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André Balogh
Rudolf A. Treumann

Physics of Collisionless Shocks

Space Plasma Shock Waves

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Space Plasma Shock Waves



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Cover illustration: (see p. 474). Sketch of the upwind heliosphere in the magnetised interstellar wind blowing against the solar wind, creating a Bow Shock and Ion Wall in the LISM, as well as the Heliopause. SPS is the solar planetary system showing the two orbits of the Voyager 1 (V1) and Voyager 2 (V2) spacecraft. V1 is shown to cross the Termination Shock (TS) at larger northern distance than V2 in the south. The Termination Shock is the red ring that confines the solar wind. The (exaggerated) asymmetry of the TS and location of SPS in the solar wind is caused by the interstellar magnetic field. Outside of the Termination Shock and the heliopause is the Heliosheath. Colour symbolises density, dark blue the low LISM density up to blue the low solar wind density. Density gradients are not included except for the compression of the upwind heliosheath which is drawn in lighter green. The pickup ion component in the outer upwind solar wind is not made visible.

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Preface

The present volume is the result of an attempt to collect our contemporary knowledge about the formation, structure and internal physics of *nonrelativistic collisionless* shocks of the kind we encounter in the near Earth environment where they are accessible to men with the help of spacecraft. This accessibility has the indisputable advantage of allowing to perform observations and measurements *in situ* the shock as well as the shock environment. This advantage elevates nonrelativistic collisionless shock theory from the state of theoretical speculation to the level where it becomes subject to tests and experiments. It moreover opens the window to the discovery of a variety of effects which the speculative researcher cannot be aware of without experimental guidance.

This volume has been structured into two parts. Part I deals with the basics of collisionless shock physics in as far as it can be treated in the nonrelativistic approximation. Fortunately all shocks in Earth's environment fall into this class of shocks. From an astrophysical point of view this fact might be unfortunate. Astrophysical shocks are manifest in the various forms of electromagnetic radiation they emit, and shocks that produce this wide range of radiation to the amounts seen must be relativistic. Relativistic shocks are not accessible *in situ*. Nevertheless, the detection of nonrelativistic shocks in space and the investigation of their structure, dynamics and physical processes serve as a guide for the investigation of astrophysical shocks and thus cannot be denied or diminished in their importance, even though it would be highly advantageous to observe a relativistic shock *in situ* as well (which, on the other hand, with high probability would cause fatal hazards for mankind or even inhibit any intelligent life on Earth and, if it existed, as well nearby).

Nonrelativistic collisionless shocks in highly dilute ordinary matter of the kind we encounter in the interplanetary space and heliosphere are mediated by medium and long range interactions between charged particles introduced by the electromagnetic field that acts on the free charges via the Lorentz force. In order to produce free charges, the matter must be hot. Thus collisionless shocks are high-temperature plasma shocks. They are subject to the different though correlated dynamics of ions and electrons. The difference of three orders of magnitude in the masses of these two constituents brings with it a severe complication and an enormous richness of effects, the current knowledge (as of April 2008 when Part I was finished) of which is described in some detail in Part I.

Part II deals with applications of the theory and experiments that have been reviewed in Part I to two of the four classes of collisionless shocks present in the heliosphere, planetary and cometary bow shocks, and last not least the heliospheric termination shock. We are in the fortunate position to have one of these shocks invariably available for experiment: Earth's bow shock wave, which is the result of the supercritical solar wind plasma flow that interacts with the geomagnetic field. The Earth's bow shock has not only the advantage of being continuously at hand, it also is bent concomitantly providing quasi-perpendicular and quasi-parallel sections where quasi-perpendicular and quasi-parallel shocks can almost continuously be studied *in situ* and in detail. The heliospheric termination shock is of particular interest because it is a concave shock and is probably mediated

by the presence of the anomalous cosmic ray component in the heliosphere. Moreover, it is the sole stellar wind termination shock that has been accessible to *in situ* observations. As such its properties are of invaluable astrophysical significance. They are large scale and closer to planar shock properties and thus represent one or the other type. One might have wished to also include applications to interplanetary shocks and solar coronal shocks. The latter are, however, not directly accessible, while the radiation they emit can be studied in relation to effects which are observable *in situ* in interplanetary space. We feel that their inclusion requires a different, more astrophysical treatment. Interplanetary shocks are occasionally accessible. However, for reasons of time, we were not in the position to include a review of them here. Both types of shocks will be the subject of related but separate reviews in Space Science Reviews.

There is no book without omissions, misinterpretations, misconceptions and even errors. We are very well aware of the fact that presumably the present volume also suffers from these diseases. We have tried to do our best in the attempt to avoid them as far as possible. Nevertheless, several of the presentations and opinions expressed in this volume are coloured by our personal knowledge and our personal ignorance and also by the general lack of insight into the enormous complexity of the physics of collisionless supercritical shocks, by the limited amount of data that are available for checking diverging theories, by the opacity of collisionless theories which are obscured by opinions and the very complex and occasionally non-transparent or sometimes even opaque approaches by the various kinds of numerical simulations, as well as by the omissions, misinterpretations and as well over-interpretations of some of these approaches.

One substantial hurdle in the completion of collisionless shock theory is the lack of extended full particle PIC code simulations in large three-dimensional systems, for full ion-to-electron mass ratios and with sufficiently large particle numbers. The simulation runs available to us have been critically included, it is however unknown how representative they are for constructing a global, much less an ultimate collisionless shock theory. Experience shows that about every ten years a further step ahead is done that forces one to abandon earlier conclusions and results. Such it was in the early seventies when magnetohydrodynamic shock theory turned out to be too simple an approximation, which did not abandon it yet for another decade. Then, in the eighties, hybrid theories using few macro-particles came into fashion providing interesting results about shock reformation as the last word at that time, which ultimately had to be given up when, in the nineties, full particle PIC simulations took over and showed that the entire reformation theory worked completely differently for quasi-perpendicular and quasi-parallel shocks. The rise of real-mass-ratio PIC simulations promises further fundamental changes in shock theory as the interaction between electrons and ions and the different instabilities in which both particles are involved will presumably cause another revolution—in particular in the highly nonlinear, non-homogeneous shock state that cannot be treated analytically and for which imagination is not strong enough to construct a consistent picture in mind that is neither guided by experiment nor by simulations.

All high resolution simulations are restricted to small systems. They just provide local information on the physics of shocks. Anticipating a complete theory requires embedding them into large-scale systems, either building up such systems out of many small-scale

systems, or embed the latter into MHD simulations. Such approaches are still completely illusionary from a practical point of view, because the fluid and the microscopic approaches have technically little in common. Maybe a kind of renormalisation procedure must first be developed before this problem can be treated even on supercomputers, as the differences in scales between the microscopic and macroscopic domains are so huge. In such a procedure a shock would consist of many microscopic domains which, in order to yield the macroscopic shock conserving the important microscopic effects, should be collected to blocks, retaining only those microscopic effects that survive the averaging procedure. One of the main problems here is buried in the proper consideration of the boundary conditions at the bounds of the many microscopic domains.

A tantalising key question concerns shock particle acceleration which is now more than half a century in work without having led to a complete satisfactory picture of how charged particles can be accelerated by shocks out of the thermal background to cosmic ray energies, illuminating the poverty of our imagination and intellectual capabilities. Hopefully, the next decade of numerical simulation will provide the key to the solution of this problem.

In this volume, we were only able to draw an intermediate and incomplete though contemporary (at the time of writing) picture of the physics involved in collisionless shocks (Part I was finished in April 2008, Part II in July 2008). This picture is subject to corrections, additions and modifications. It is our hope that it will just serve as a reference frame for further successful research that will be going far beyond what is contained herein.

We are indebted to a small number of colleagues who have contributed by discussions. Their efforts are gratefully acknowledged at this place. Finally, we express our gratitude to Dr. Harry Blom, Mrs. Jennifer Satten and Felix Portnoy of the New York Springer office for their professional and sensitive handling of the final edition of this book. Last not least, we thank the Springer Production Office at Vilnius for their excellent work in type setting and production of this volume to our great satisfaction.

Bern, Switzerland

André Balogh
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Part I

Collisionless Shock Theory

Introduction

Part I of this book deals with the history, theory and simulation of non-relativistic collisionless shocks like those encountered in space plasma in the close and more distant spatial environment of the Earth.¹

Chapter 2 poses the general problem of a collisionless shock as the problem of understanding how in a completely collisionless but streaming high-temperature plasma shocks can develop at all, forming discontinuous transition layers of thickness much less than any collisional mean free path length. The history of shock research is briefly reviewed. It is expressed that collisionless shocks as a realistic possibility of a state of matter have been realised not earlier than roughly half a century ago. The basic properties of collisionless shocks are noted in preparing the theory of collisionless shocks and a classification of shocks given in terms of their physical properties, which are developed in the following chapters. Naturally, due to limitations of space, the historical overview is cursory, describes the early history, gasdynamic shocks, the realisation of the collisionless shock problem, and its investigation over three decades in theory and observation until the advent of the numerical simulation age. Types of collisionless shocks are identified, and the distinction is made between electrostatic and magnetised shocks, MHD shocks, evolutionarity and coplanarity, switch-on and switch-off shocks. The criticality of shocks is defined as the transition from subcritical dissipative to supercritical viscous shocks, and it is noted that supercritical shocks are particle reflectors.

Chapter 3 develops the basic sets of equations which lead to the conservation laws describing collisionless plasma shock waves. The evolution of shock waves by wave steepening is discussed as the fundamental nonlinear process. Once the shock exists, one defines the boundary (Rankine-Hugoniot) conditions in the case that the shock can be described by fluid theory, i.e. for magnetogasdynamic shocks. Various analytical models of shock formation are presented, and the basic waves and instabilities are discussed which may become important in collisionless shock physics. The general dispersion relation describing low- β -shocks is derived and the wave modes (whistlers, Alfvén waves, magnetosonic waves) are given. A survey of the theory of anomalous resistivity is presented in the quasi-linear limit and beyond for the various shock-relevant instabilities, and the mechanisms of shock particle reflection is discussed. Finally, a briefing on numerical simulation techniques is added, providing a short idea on this important field and its methods.

Chapter 4 separates out the subclass of collisionless subcritical shocks which emerge at Mach numbers below the critical Mach number. In space such shocks are quite rare. This

¹A concise version of this Part I with emphasis on astrophysical applications has previously been published in Treumann RA (2009) *Astron Astrophys Rev* 19. Its extension into the relativistic domain reviewing the current state of relativistic and ultra-relativistic shock research can be found in Bykov AM, Treumann RA (2011) *Astron Astrophys Rev* 21, see the Reference list of Chapter 2.

chapter introduces them as shocks which are determined by the competition of nonlinear steeping of a low frequency wave and dispersive or dissipative processes which come into play under different dispersive conditions. The former, when first order gradients affect the dispersive properties of the wave, the latter when the gradient, i.e. the shock thickness, is of the order of the dissipative scale. The subcritical cross-shock potential is determined, and an expression for the dissipation scale is given.

Chapter 5 reviews the theory and simulation of quasi-perpendicular super-critical shocks by formulating the quasi-perpendicular shock problem as being different from the quasi-parallel shock. Distinction between both types becomes important at supercritical Mach numbers where the angle Θ_{Bn} between the shock normal and the ambient upstream (pre-shock) magnetic field becomes important. This requires the investigation of shock-particle reflection which is a function of this angle. Its kinematic formulation leads to the clear distinction of quasi-perpendicular and quasi-parallel shocks for angles $\Theta_{Bn} < \frac{1}{4}\pi$ and $\Theta_{Bn} > \frac{1}{4}\pi$, respectively. Quasi-perpendicular shocks develop a well expressed foot whose physics is described. Furthermore, based on numerical simulation the shock and foot structure is discussed in depth for different electron-to-ion mass ratios, foot and ramp scales are determined, and a substantial part of the chapter is devoted to shock reformation for the case when the ion-plasma β_i is large and for the other case, when oblique whistlers stabilise the shock in two dimensions. High-Mach number shocks will always become non-stationary in a certain sense. The relevant instabilities are investigated: Buneman, two-stream, modified two stream, electron tail generation and particle heating.

Chapter 6 turns to the description of quasi-parallel shocks which in a certain sense form the majority of observed shocks in space. These shocks possess an extended foreshock region with its own extremely interesting dynamics for both types of particles, electrons and ions, reaching from the foreshock boundaries to the deep interior of the foreshock. Based on a variety of numerical simulations the properties of the foreshock are investigated and its importance in shock physics is elucidated. In fact, it turns out that the foreshock is the region where dissipation of flow energy starts well before the flow arrives at the shock. This dissipation is caused by various instabilities excited by the interaction between the flow and the reflected particles that have escaped to upstream from the shock. Interaction between these waves and the reflected and accumulated particle component in the foreshock causes wave growth and steeping, formation of shocklets and pulsations and causes continuous reformation of the quasi-parallel shock that differs completely from quasi-perpendicular shock reformation. It is the main process of maintaining the quasi-parallel shock which by its nature principally turns out to be locally nonstationary and, in addition, on the small scale making the quasi-parallel shock close to becoming quasi-perpendicular for the electrons. This process can be defined as turbulent reformation. Various effects like hot flow anomalies and generation of electromagnetic radiation are discussed. The importance of this kind of physics for particle acceleration is noted.

Chapter 7 is a first application of collisionless shock theory. It provides a discussion of some ideas on particle acceleration at non-relativistic supercritical shocks mainly pointing out that it is the quasi-parallel shocks which are capable of providing the seed particle population needed for entering the first-order Fermi acceleration cycle in shock acceleration. In general, collisionless shocks in the heliosphere cannot accelerate particles to very

high energies. The energies that can be reached are limited by several bounds: the non-relativistic nature of the heliospheric collisionless shocks to which this review restricts, the finite size of these shocks, the finite width of the downstream region, and by the nature of turbulence which in the Fermi process is believed to scatter the particles on both sides of the shock. As a fundamental problem of the acceleration mechanism the injection of seed particles is identified. Some mechanisms for production of seed particles are invoked but the most probable mechanism is the interaction between the quasi-parallel shock and the energetic tail of the reflected population during quasi-parallel shock reformation. Multiple reflection of particles in this process may produce a seed particle component that overcomes the threshold for the Fermi process. Also, acceleration of electrons begins to uncover its nature due to the nearly quasi-perpendicular nature of the quasi-parallel shock on the small scale, i.e. the electron scale, as described in Chapter 6. Energetic electrons are required for the observed radiation from shocks both in space as in astrophysics. Noting that most of the acceleration physics under astrophysical conditions is focussing on shock acceleration, the ultimate clarification of these processes is important. It requires further investigation and the transition to much higher Mach number shocks than have been possible to simulate so far. This shifts collisionless shock physics naturally into the relativistic domain.

The Shock Problem

Abstract. This introductory chapter provides a brief historical account of the emerging interest in shock waves. It started from consideration of hydrodynamic flows but was stimulated by military purposes, with Ernst Mach being the first defining its physical properties. The problem sharpened after realising that shocks in plasmas predominantly evolve under the collisionless conditions encountered in space. The different stages of shock exploration are briefly described and the conditions for plasma shock are listed. We define the main terms used in shock physics and distinguish between the different types of collisionless shocks, electrostatic, magnetised, MHD-shocks, switch-on and switch-off shocks, and evolutionarity and criticality of shocks are briefly discussed.

2.1 A Cursory Historical Overview

The present book deals with the physics of shock waves in our heliosphere only, a very particular class of shocks: shocks in collisionless high-temperature but non-relativistic plasmas. The physics of shocks waves is much more general, however, covering one of the most interesting chapters in many-particle physics, from solid state to the huge dimensions of cosmic space. Shock waves were involved when stars and planets formed and when the matter in the universe clumped to build galaxies, and they are involved when the heavy elements have formed which are at the fundament of life and civilisation. Mankind, however, has become aware of shocks only very lately. The present section gives a concise account of its history in human understanding.

2.1.1 The Early History

Interest in shocks has arisen first in gasdynamics when fast flows came to attention not long before Ernst Mach's realisation [*Mach & Wentzel*, 1884, 1886; *Mach & Salcher*, 1887; *Mach*, 1898] that in order for a shock to evolve in a flow, an obstacle must be put into the flow with the property that the relative velocity V of the obstacle with respect to the bulk flow exceeds the velocity of sound c_s in the medium, leading him to define the critical velocity ratio, $\mathcal{M} = V/c_s$ (see Figure 2.1), today known as the famous “Mach number” and contributing even more to his public recognition and scientific immortality than his other great contribution to physics and natural philosophy, the preparation of the path to General Relativity for Albert Einstein by his fundamental analysis of the nature of gravity and the equivalence of inertial and gravitational masses.

Mach was experimenting at this time with projectiles that were ejected from guns using the new technique of taking photographs of cords hat are generated by projectiles when traversing a transparent though very viscous fluid (for an example see Figure 2.2). Even

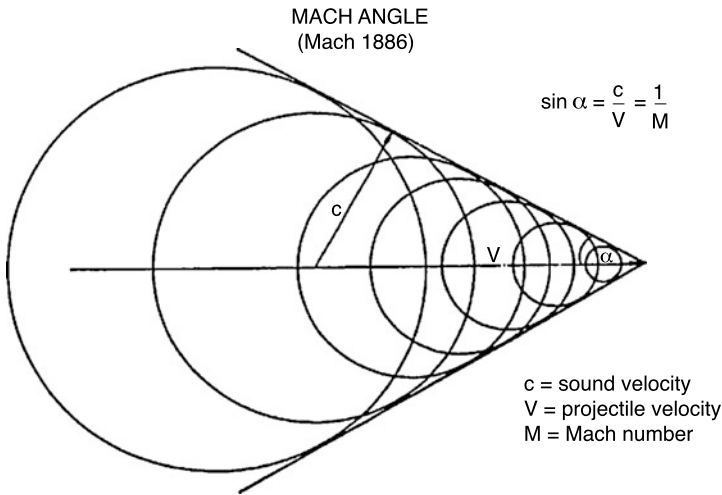


Figure 2.1: Mach's 1886 drawing of the form of a projectile generated shock of Mach-opening angle $\alpha = \sin^{-1}(1/M)$. The projectile is moving to the right (or the air is moving to the left) at speed V . The shock front is formed as the envelope of all the spherical soundwave fronts excited along the path of the projectile in the air when the sound waves expand radially at sound speed c .

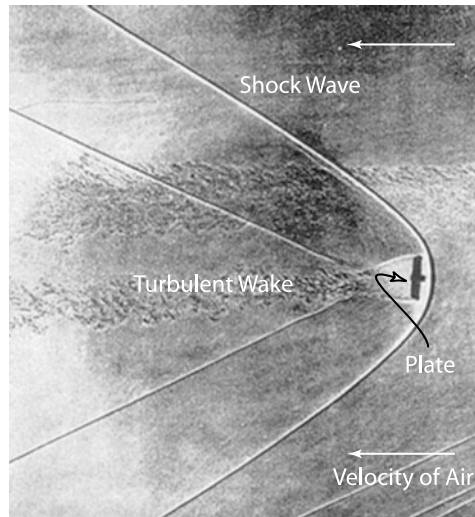


Figure 2.2: A cord photograph of a blunt plate moving at supersonic velocity in a transparent viscous medium causing a blunt thin shock wave of hyperbolic form and leaving a turbulent wake behind (photograph taken after H. Schardin, Lilienthal-Gesellschaft, Report No. 139). This figure shows nicely that blunt objects cause blunt nosed shocks standing at a distance from the object. This is similar to what the Earth's dipole magnetic field does in the solar wind.

though Mach was not working for the Austrian Ministry of War, these endeavours were in the very interest of it. Mach mentioned this fact, somewhat ironically, in his famous later publication [*Mach*, 1898] where he was publicly reviewing his results on this matter and where he writes:¹

“As in our todays life shooting and everything connected to it under certain circumstances plays a very important, if not the dominant, role you may possibly turn your interest for an hour to certain experiments, which have been performed not necessarily in view of their martial application but rather in scientific purpose, and which will provide you some insight into the processes taking place in the shooting.”

The physical interpretation of Mach’s observations was stimulated by the recognition that, no matter how fast the stream would be as seen from a frame of reference that is at rest with the projectile, the projectile being an obstacle in the flow causes disturbances to evolve in the flow. Such disturbances are travelling waves which in ordinary gasdynamics are sound waves. In flows with Mach number $\mathcal{M} < 1$, sound waves can reach any upstream position thereby informing the flow of the presence of the obstacle and leaving the flow sufficient time for rearranging and changing its direction in order for turning itself smoothly around the obstacle. In this case no shock will evolve.

In the opposite case $\mathcal{M} \geq 1$, the stream remains uninformed until it hits the obstacle, being completely unprepared for its presence, and something catastrophic happens, i.e. *the flow is shocked*. In this case the flow is too fast for the disturbances generated by the obstacle to propagate large distances upstream. They can propagate only a finite distance from the obstacle up to a certain point in a time that is shorter than the flow needs from this point to arrive at the obstacle. Hence, it is clear that from the farthest upstream position reached by the waves – i.e. from the shock to the obstacle – the flow velocity must have dropped below the velocity of sound, which requires the shock to be a thin discontinuity surface at which the flow has been braked and a substantial part of its directed motional energy is transformed into heat, i.e. into disordered motion. The shock discontinuity surface thus separates the undisturbed cooler upstream flow from the disturbed warmer downstream flow between shock and obstacle that contains all the obstacle-excited sound waves, and the temperature and density of the flow must both increase across the shock from upstream to downstream, which implies that the pressure, which is the product of density and temperature, increases as well.

2.1.2 Gasdynamic Shocks

The geometry of the shock can be constructed from the characteristic curves of sound wave propagation around the obstacle. In the downstream region the subsonic flow bends around the obstacle. Since retardation implies that the volume of a flow element is reduced, the downstream density is increased above upstream and flow energy is converted into thermal energy, giving the region downstream of the shock a higher temperature than upstream. Higher temperature means increased disorder and thus implies that entropy is generated at the shock making the whole process irreversible, which can be done only when the shock

¹The original text is in German and can be found in a recent recollection by *Pohl* [2002].

front supports dissipation. Finally, due to the presence of the broad spectrum of sound waves the downstream region supports irregular motion and is to some extent turbulent.

This is the simple global physics implied. It can be described by the equations of compressible gasdynamics where the emphasis is on compressible, because sound waves are fluctuations in density. The microscopic processes are, however, more complicated even in ordinary gasdynamics, implying some knowledge about the relevant dissipative transport coefficients heat conduction and viscosity, which determine the physical thickness of the shock front, the shock profile, and the process of heating and generation of entropy. A first one dimensional theory goes back to *Becker* [1922]. Later comprehensive reviews can be found in *Landau & Lifshitz* [1959] and *Zeldovich & Raizer* [1966]. An extended and almost complete review of the history of shock waves has been given only recently by *Krehl* [2007].

Gasdynamic shocks are dominated by binary collisions and thus are collisional. Binary collisions between the molecules of the fluid respectively the gas are required, since collisions are the only way the molecules interact among each other as long as radiative interactions are absent. They are required for the necessary heating and entropy generation. There is also some dissipative interaction with the sound waves; this is, however, weak and usually negligible compared to the viscous interaction, except under conditions when the amplitudes of the sound waves at the shortest wavelengths become large. This happens to be the case only inside the shock front, the width of which is of the order of just a few collisional mean free paths $\lambda_{\text{mfp}} = (N\sigma_c)^{-1}$. The latter is defined as the inverse product of gas number density N and collisional cross section $\sigma_c \simeq \text{a few} \times 10^{-19} \text{ m}^2$. Thus, for the mean free path to be small, it requires large gas densities which are rather rarely found in interstellar or interplanetary space, the subject in which we are interested in this volume. Shocks developing under such – practically collisionless – conditions are called collisionless shocks, the term “collisionless” implying that binary collisions can completely be neglected but should be replaced by other non-collisional processes that are capable to warrant the production of entropy. The widths of collisionless shocks are much less than the theoretical collisional mean free path, and any dissipative processes must be attributed to mechanisms based not on collisions but on collective processes.

In gases collisionless conditions evolve naturally when the temperature of the gas rises. In this case the gas becomes dilute with decreasing density and, in addition, the (generally weak) dependence of the collisional cross section on temperature causes a decrease in σ_c . This effect becomes particularly remarkable when the temperature substantially exceeds the ionisation energy threshold. Then the gas makes the transition to a plasma consisting of an ever decreasing number of neutral molecules and an increasing number of electron-ion pairs. At temperatures far above ionisation energy the density of neutrals can be neglected compared to the density of electrons and ions. These are electrically charged, with the interaction between them dominated not anymore by short-range collisions but by the long-range Coulomb force that decays $\propto r^{-2}$, with interparticle distance $r = |\mathbf{r}_2 - \mathbf{r}_1|$. The collisional mean free path $\lambda_{\text{mfp}} = (N\sigma_c)^{-1} \equiv \lambda_c$ now contains the Coulomb-cross section $\sigma_c = \sigma_C$, thereby becoming the Coulomb mean free path λ_C . In plasma the two components, electrons and ions, are about independent populations that are coupled primarily through the condition of maintaining the quasi-neutrality of the plasma. The Coulomb

cross section itself is inversely proportional to the fourth power of the particle velocity, $\sigma_C = (16\pi N^2 v^4 / \omega_{pe}^4)^{-1} \propto v^{-4}$ and, for fast particles, decreases rapidly with v . Accordingly, $\lambda_C \propto v^3$ increases, readily becoming as large as the macroscopic extension of the entire plasma. For thermal particles it increases with temperature as $\lambda_C \propto T^{\frac{3}{2}}$, thus becoming very large as well. In interplanetary space this length is of the order of several AU. Any shock of lesser width must therefore be completely collisionless. This inescapable conclusion causes severe problems in the interpretation and physical understanding of collisionless shocks in space. It is such shocks, which are frequently encountered in interplanetary and interstellar space and which behave completely different from their collisional counterparts in fluids and gases. The present volume deals with their properties as they are inferred from direct observation and interpreted in theory and simulation.

2.1.3 Realising Collisionless Shocks

The possibility of collisionless shocks in an ionised gaseous medium that can be described by the equations of magnetohydrodynamics was anticipated first by *Courant & Friedrichs* [1948]. The first theory of magnetised shocks was given by *De Hoffman & Teller* [1950] in an important paper that was stimulated by after-war atmospheric nuclear explosions. This paper was even preceded by Fermi's seminal paper on the origin of cosmic rays [Fermi, 1949] where he implicitly proposed the existence of collisionless shocks when stating that cosmic rays are accelerated in multiple head-on collisions with concentrated magnetic fields (plasma clouds or magnetised shock waves), in each collision picking up the difference in flow velocity between the flows upstream and downstream of the shock. The De Hoffman & Teller paper ignited an avalanche of theoretical investigations of magnetohydrodynamic shocks [Helfer, 1953; Lüst, 1953, 1955; Marshall, 1955; Syrovatskij, 1958; Shafranov, 1957; Vedenov et al, 1961; Vedenov, 1967; Davis et al, 1958; Gardner et al, 1958; Syrovatskii, 1959; Bazer & Ericson, 1959; Ludford, 1959; Montgomery, 1959; Kontorovich, 1959; Liubarskii & Polovin, 1959; Colgate, 1959; Sagdeev, 1960, 1962a, b; Fishman et al, 1960; Germain, 1960; Kulikovskii & Liubimov, 1961; Morawetz, 1961, and others]. Early on, *Friedrichs* [1954] had realised already that the method of characteristics could be modified in a way that makes it applicable as well to magnetohydrodynamic shocks. This allowed for the formal construction of the geometrical shapes of magnetohydrodynamic shocks developing in front of given obstacles of arbitrary profile.

Production of collisionless shocks in the laboratory encountered more severe problems, as the dimensions of the devices were small, temperatures comparably low, and densities high such that collisional effects could hardly be suppressed. Nevertheless, first preliminary experimental results on various aspects of the structure of nearly collisionless shocks were reported by *Bazer & Ericson* [1959], *Patrick* [1959], *Wilcox et al* [1960, 1961], *Auer et al* [1961, 1962], *Keck* [1962], *Camac et al* [1962], *Fishman & Petschek* [1962], *Brennan et al* [1963] and others. The first successful production of collisionless shocks in laboratory experiments [Kurtmullaev et al, 1965; Paul et al, 1965; Eiselevich et al, 1971] revealed that the collisionless shocks investigated were highly nonstationary and, depending on Mach number, exhibited complicated substructuring, including strongly heated electrons [Paul

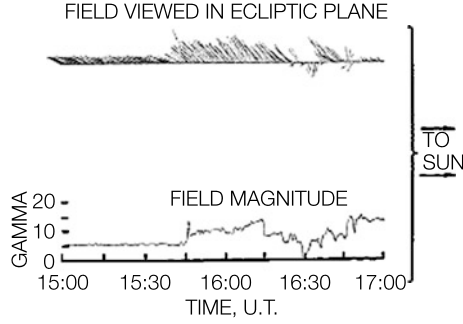


Figure 2.3: The original magnetic field recordings by Mariner II of October 7, 1962. The passage of Earth’s bow shock wave occurs at 15:46 UT and is shown by the spiky increase in the magnetic field followed by a period of grossly enhanced magnetic field strength [after *Sonett & Abrams, 1963*]. The upper part shows the direction of the magnetic field vector in the ecliptic plane and its sudden change across the bow shock.

et al, 1967] and electric potential jumps that were extended over $\sim 100\lambda_D$, many Debye lengths λ_D inside the shock transition [for a review see *Eselevich, 1982*]. The existence of heated electrons and electric potentials already led *Paul et al [1967]* to speculate that, at higher Mach numbers, shocks could be *in principle* non-stationary. Further laboratory studies by *Morse et al [1972]* and *Morse & Destler [1972]* with the facilities available at that time seemed to confirm this conjecture.

However, the first indisputable proof of the real existence of collisionless shocks in natural plasmas came from spacecraft observations, when the Mariner (see Figure 2.3) and IMP satellites passed the Earth’s bow shock, the shock wave that is standing upstream of Earth’s magnetosphere in the solar wind [*Sonett & Abrams, 1963; Ness et al, 1964*]. That such shocks should occasionally exist in the solar wind had already been suggested about a decade earlier by *Gold [1955]* who concluded that the sharp sudden-commencement rise in the horizontal component of the geomagnetic field at the surface of the Earth that initiates large magnetic storms on Earth not only implied an impinging interplanetary plasma stream – as had been proposed twenty years earlier by *Chapman & Ferraro [1930, 1931]* – but required a very high velocity $\mathcal{M} > 1$ solar wind stream that was able to (at that time believed to occur just temporarily) create a bow shock in front of the Earth when interacting with the Earth’s dipolar geomagnetic field. Referring to the then accepted presence of the stationary and super-Alfvénic solar wind flow, *Zhigulev & Romishevskii [1960]*, and somewhat later *Axford [1962]* and *Kellogg [1962]*, simultaneously picking up this idea, suggested that this shock wave in front of Earth’s dipole field should in fact represent a *stationary bow shock* that is standing in the solar wind, being at rest in Earth’s reference frame.

The above mentioned seminal spacecraft observations *in situ* the solar wind by *Sonett & Abrams [1963]* and *Ness et al [1964]* unambiguously confirmed these claims, demonstrating that the super-magnetosonic solar wind stream had indeed traversed a thin discontinuity surface, downstream of which it had entered a plasma that was in another highly disturbed irreversibly “turbulent” thermodynamic state. The transition was monitored mainly

in the magnetic field that changed abruptly from a relatively steady solar wind value around $B_{\text{sw}} \simeq 5\text{--}10\text{ nT}$ to a high downstream value that was fluctuating around an average of $\langle B \rangle \simeq 15\text{--}30\text{ nT}$ and had rotated its direction by a large angle. This shock transition surface could not be in local thermodynamic equilibrium. It moreover turned out to have thickness of the order of $\sim$ few 100 km being comparable to the gyro-radius of an incoming solar wind proton and thus many orders of magnitude less than the Coulomb mean free path. For a measured solar wind-plasma particle density of $N \sim 10\text{ cm}^{-3}$ and an electron temperature of $T_e \sim 30\text{ eV}$ the Coulomb mean free path amounts roughly to about $\lambda_C \sim 5\text{ AU}$, vastly larger than the dimension of the entire Earth's magnetosphere system which has an estimated linear extension in the anti-sunward direction of $\lesssim 1000 R_E \approx 3 \times 10^{-3}\text{ AU}$. Thus Earth's bow shock represents a truly collision-free shock transition.

This realisation of the extreme sharpness of a collisionless shock like the Earth's bow shock immediately posed a serious problem for the magnetohydrodynamic description of collisionless shocks. In collisionless magnetohydrodynamics there is no known dissipation mechanism that could lead to the observed extremely short transition scales $\Delta \sim r_{ci}$ in high Mach number flows which are comparable to the ion gyro-radius r_{ci} . Magnetohydrodynamics neglects any differences in the properties of electrons and ions and thus barely covers the very physics of shocks on the observed scales. In its frame, shocks are considered as infinitely narrow discontinuities, narrower than the magnetohydrodynamic flow scales $L \gg \Delta \gg \lambda_d$; on the other hand, these discontinuities must physically be much wider than the dissipation scale λ_d with all the physics going on inside the shock transition. This implies that the conditions derived from collisionless magnetohydrodynamics just hold far upstream and far downstream of the shock transition, i.e. far outside the region where the shock interactions are going on. In describing shock waves, collisionless magnetohydrodynamics must be used in an asymptotic sense, providing the remote boundary conditions on the shock transition. One must look for processes different from magnetohydrodynamics in order to arrive at a description of the processes leading to shock formation and shock dynamics and the structure of the shock transition. In fact, viewed from the magnetohydrodynamic single-fluid viewpoint, the shock should not be restricted to the steep shock front, it rather includes the entire shock transition region from outside the foreshock across the shock front down to the boundary layer at the surface of the obstacle. And this holds as well even in two-fluid shock theory that distinguishes between the behaviour of electrons and ions in the plasma fluid.

2.1.4 Early Collisionless Shock Investigations

Evolutionary models of magnetohydrodynamic shocks [Kantrowitz & Petschek, 1966] were based on the assumption of the dispersive evolution of one of the three magnetohydrodynamic wave modes, the compressive “fast” and “slow magnetosonic” modes and the incompressible “intermediate” or Alfvén wave. It is clear that, in a non-dissipative medium, the evolution of a shock wave must be due to the nonlinear evolution of a dispersive wave disturbance. In those modes where the shorter wavelength waves have higher phase velocity, the faster short-wavelength waves will overcome the slower long-wavelength waves and cause steepening of the wave. If there is neither dissipation nor

dispersion, the wave will start breaking. Wave steepening can be balanced by dispersion, which leads to the formation of large amplitude isolated solitary wave packets. When, in addition, the steepening causes shortening of the wavelength until the extension of the wave packet in real space becomes comparable to the internal dissipation scale, a shock ramp may form out of the solitary structure. This ramp separates the compressed downstream state from the upstream state.

Such processes have been proposed to take place in an early seminal review paper by Sagdeev [1966] on the collective processes involved in the evolution of collisionless shocks that was based on the earlier work of this author [Sagdeev, 1960, 1962a, b; Moiseev & Sagdeev, 1963a, b]. The ideas made public in that paper were of fundamental importance for two decennia of collisionless shock research that was based on the nonlinear evolution of dispersive waves. In particular the insight into the microscopic physical processes taking place in collisionless shock formation and the introduction of the equivalent pseudo-potential method (later called ‘Sagdeev potential’ method) clarified many open points and determined the direction of future shock research.

Already the first *in situ* observations identified the collisionless shocks in the solar wind like Earth’s bow shock as magnetised shocks. It is the magnetic field which determines many of their properties. The magnetic field complicates the problem substantially by multiplying the number of possible plasma modes, differentiating between electron and ion dynamics, and increasing the possibilities of nonlinear interactions. On the other hand, the presence of a magnetic field introduces some rigidity and ordering into the particle dynamics by assigning adiabatically invariant magnetic moments $\mu = T_{\perp}/|\mathbf{B}|$ to each particle of mass m , electric charge q , and energy $T_{\perp} = \frac{1}{2}mv_{\perp}^2$ perpendicular to the magnetic field $\mathbf{B}$.

Because of the obvious importance of dispersive and dissipative effects in shock formation, for more than one decade the theoretical efforts concentrated on the investigation of the dispersive properties of the various plasma modes, in particular on two-fluid and kinetic modes [Gary & Biskamp, 1971, for example] and on the generation of anomalous collision frequencies [Krall & Book, 1969; Krall & Liewer, 1971; Biskamp & Chodura, 1971, for example] in hot plasma even though it was realised already by Marshall [1955] that high-Mach number shocks cannot be sustained by purely resistive dissipation like anomalous resistivity and viscosity alone. This does not mean that the investigation of anomalous resistivity and viscosity by itself would make no sense. At the contrary, it was realised very early that wave-particle interactions replace binary collisions in collisionless plasmas, thereby generating anomalous friction which manifests itself in anomalous transport coefficients, and much effort was invested into the determination of these coefficients [e.g. Vekshtein *et al.*, 1970; Bekshtein & Sagdeev, 1970; Liewer & Krall, 1973; Sagdeev, 1979, and others].

These anomalous transport coefficients do in fact apply to low-Mach number shocks. However, when the Mach number exceeds a certain – surprisingly low and angular dependent – critical limit, the anomalous resistive or viscous time scales that depend on the growth rates of instabilities become too long in order to generate the required dissipation, heating and increase in entropy fast enough for maintaining a quasi-stationary shock. The critical Mach number estimated by Marshall [1955] was $\mathcal{M}_{\text{crit}} = 2.76$ for a perpendicu-

lar shock. This follows from the condition that $\mathcal{M} \leq 1$ right downstream of the shock, implying that the maximum downstream flow speed $V_2 = c_{s2}$ should just equal the downstream sound speed in the shock-heated flow [Coroniti, 1970a]. At higher Mach numbers the solution Nature finds is that the shock ramp specularly reflects back upstream an ever increasing portion of the incoming plasma corresponding to the fraction of particles whose excess motional energy the shock is unable to convert into heat.

Shock reflection had first been suggested and inferred as an important mechanism for shock dissipation by Sagdeev [1966]. In a magnetised perpendicular shock like part of Earth's bow shock, the shock-reflected ions generate a magnetic foot in front of the shock ramp, as has been realised by Woods [1969, 1971]. This foot is the magnetic field of the current carried by the reflected ion stream that drifts along the shock ramp in the respective crossed magnetic field and magnetic field gradients. It is immediately clear that the efficiency of reflection must depend on the angle $\Theta_{Bn} = \cos^{-1}(\mathbf{B}_1, \mathbf{n})$ between the shock ramp, represented by the shock normal vector $\mathbf{n}$, and the upstream magnetic field $\mathbf{B}_1$.

This angle allows to distinguish between perpendicular $\Theta_{Bn} = 90^\circ$ and parallel $\Theta_{Bn} = 0^\circ$ magnetised shocks. In the ideal reflection of a particle, it is its flow-velocity component V_n parallel to the shock normal that is inverted. Since the particles are tied to the magnetic field by gyration, they return upstream with velocity that is (at most) the projection $v_{\parallel} = -V_n \cos \Theta_{Bn}$ of this component onto the upstream magnetic field. In a perpendicular shock $v_{\parallel} = 0$ and the reflected particles do not really return into the upstream flow but perform an orbit of half a gyro-circle upstream extension around the magnetic field $\mathbf{B}_1$. In the intermediate domain of quasi-perpendicular $\Theta_{Bn} > 45^\circ$ and quasi-parallel $\Theta_{Bn} < 45^\circ$ shocks, particles leave the shock upstream along the magnetic field $\mathbf{B}_1$. While for quasi-perpendicular shocks all particles after few upstream gyrations return to the shock and pass it to downstream, the efficiency of reflection decreases with decreasing angle Θ_{Bn} such that, theoretically, for parallel shocks no particles are reflected at all.

The reflection process depends on the shock potential, height of magnetic shock ramp, shock width and plasma-wave spectrum in the shock. Because of the complexity of the equations involved, its investigation requires extensive numerical calculations. Such numerical simulations using different plasma models have been initiated in the early seventies [Forslund & Shonk, 1970; Biskamp & Welter, 1972a, b] and have since become the main theoretical instrumentation in the investigation of collisionless shocks, accompanying and completing the wealth of data obtained from *in situ* measurements of shocks in interplanetary space and from remote observations using radio emissions from travelling interplanetary shocks or the entire electromagnetic spectrum as is believed to be emitted from astrophysical shocks from the infrared through optical, radio and X-rays up to gamma rays which have been detected from highly relativistic shocks in astrophysical jets.

2.1.5 Three Decades of Exploration: Theory and Observation

During the following three decades many facets of the behaviour of collisionless shocks could be clarified with the help of these observations and support from the first numerical simulations. For the initial period the achievements have been summarised at different

stages by *Anderson* [1963], who reviewed the then known magnetohydrodynamic shock wave theory which still was not aware of the simulation possibilities which opened up by the coming availability of powerful computers and, in the first and much more important place, by *Sagdeev* [1966], who gave an ingenious summary of the ideas that had been developed by him [*Sagdeev*, 1960, 1962a, b] and his coworkers [*Vedenov et al*, 1961; *Vedenov*, 1967; *Kadomtsev & Petviashvili*, 1963; *Moiseev & Sagdeev*, 1963a, b; *Galeev & Karpman*, 1963; *Galeev & Oraevskii*, 1963; *Karpman*, 1964a, b; *Karpman & Sagdeev*, 1964] on the nonlinear evolution of collisionless shocks from an initial disturbance growing out of a plasma instability. *Sagdeev* [1966] demonstrated particularly clearly the physical ideas of dispersive shock-wave formation and the onset of dissipation and advertised the method of the later-so called Sagdeev potential.² *Sagdeev's* seminal paper laid the foundation for a three decade long fruitful research in nonlinear wave structures and shock waves.

This extraordinarily important step was followed by the next early period that was represented by the review papers of *Friedman et al* [1971], *Tidman & Krall* [1971], *Biskamp* [1973], *Galeev* [1976], *Formisano* [1977], *Greenstadt & Fredricks* [1979], again *Sagdeev* [1979] and ultimately two comprehensive review volumes edited by *Stone & Tsurutani* [1985] and *Tsurutani & Stone* [1985].

These last two volumes summarised the state of the knowledge that had been reached in the mid-eighties. This knowledge was based mainly on the first most sophisticated observations made by the ISEE 1 & 2 spacecraft which for a couple of years regularly traversed the Earth's bow shock wave at many different positions and, in addition, crossed a number of interplanetary travelling shocks. Great discoveries made by these spacecraft were – among others – the division of the bow shock into quasi-perpendicular and quasi-parallel parts, the identification of the upstream structure of the bow shock, which was found to be divided into the undisturbed solar wind in front of the quasi-perpendicular bow shock and the foot region, the evolution of the ion distribution across the quasi-perpendicular bow shock [*Paschmann et al*, 1982; *Skopke et al*, 1983], and the extended foreshock region upstream of the quasi-parallel part of the bow shock, thereby confirming the earlier theoretical claims mentioned above that supercritical shocks should reflect part of the solar wind back upstream.

Much effort was invested into the investigation of the foreshock which was itself found to be divided into a narrow electron foreshock and an extended ion foreshock, the former being separated sharply from the undisturbed solar wind. In the broad ion foreshock region, on the other hand, the ion distribution functions were found evolving from a beam-like distribution at and close to the ion foreshock boundary towards a nearly isotropic diffuse

²Historically it is interesting to note that the Sagdeev potential method was in fact independently used already by *Davis et al* [1958] in their treatment of one-dimensional magnetised magnetohydrodynamic shocks, who numerically calculated shock solutions but missed the deeper physical meaning of the Sagdeev potential which was elucidated in the work of *Sagdeev*. Of course, the method of transforming an arbitrary second order one-dimensional (ordinary) differential equation into the equation of a particle moving in an equivalent potential well was known for long in classical mechanics and the solution of the telegraph equation (at least since the work of *Gustav R. Kirchhoff* around 1850), but its extraordinarily fruitful application to the nonlinear equations of plasma physics that led to the understanding of soliton and shock formation was entirely due to *Sagdeev*.

distribution deep inside the ion foreshock [Gosling *et al*, 1978, 1982; Paschmann *et al*, 1981].

An important observation was that the reflected ions become accelerated to about four times solar wind energy when being picked-up by the solar wind and coupling to the solar wind stream, a fact being used later to explain the acceleration of the anomalous component of cosmic rays in the heliosphere. It was moreover found that the reflected ions strongly interact with the upstream solar wind [Paschmann *et al*, 1979] via several types of ion beam interactions [Gary, 1981], first observed a decade earlier in the laboratory by Phillips & Robson [1972]. These instabilities were found to generate broad spectra of low-frequency electromagnetic modes that fill the foreshock with an intense spectrum of low-frequency electromagnetic fluctuations that propagate in all directions with and against the solar wind.³ Some of the lowest frequency fluctuations can even steepen and generate large amplitude low frequency or quasi-stationary wave packets in the solar wind resembling small, spatially localised and travelling magnetic ramps or shocklets. When the solar wind interacts with these shocklets, the solar-wind stream becomes already partially retarded long before even reaching the very shock front. In other words, the entire foreshock region is already part of the shock transition.⁴

Another important observation concerned the evolution of the electron distribution in the vicinity of the bow shock and, in particular, across the quasi-perpendicular bow shock [Feldman *et al*, 1982, 1983]. It was, in fact, found that the electron distribution evolved from the nearly Boltzmannian plus halo distribution in the solar wind to make the transition into a flat-top heated electron distribution, when crossing the shock ramp. The flat top suggested that strong electron heating takes place inside the shock on a short time-scale and that the shock ramp contains a stationary electric field which is partially responsible for the reflection of particles back upstream, while at the same time heating the inflowing electrons. Moreover, the distributions showed that electrons were also reflected at the shock, escaping in the form of narrow-beam bundles into the upstream solar wind along the solar wind magnetic field line (or flux tube) that is tangential to the bow shock. Upstream of the shock, these electron beams were strong enough to generate plasma fluctuations around the plasma frequency and to give rise to radio emission at the harmonic of the plasma frequency. Plasma wave observations [Rodriguez & Gurnett, 1975; Anderson *et al*, 1981] were confirmed by these observations as being excited by the shock reflected electron beams.

The period following the comprehensive reviews by Stone & Tsurutani [1985] and Tsurutani & Stone [1985] were devoted to further studies of the bow shock by other spacecraft like AMPTE, and of other interplanetary shocks accompanying solar ejection events like CMEs by Ulysses and other spacecraft, in particular by the Voyager 1 & 2 satellites who investigated travelling shocks and corotating (with the sun) interacting regions in interplanetary space, heading for encounters with bow shocks of the outer planets and, ultimately, the heliospheric termination shock. To everyone's excitement the heliospheric termination

³A professional timely review of the observations of all waves observed in the foreshock was given by Gurnett [1985] and is contained in the above cited volume edited by Tsurutani & Stone [1985].

⁴The importance of this insight has for long time, actually up to this time, not been realised, even though it was suggested by the ISEE observations and later by AMPTE IRM observations long ago.

shock was crossed in a spectacular event in December 2004 by Voyager 1 at a distance of ~ 94 AU, confirming the prediction of its existence and at the same time opening up a large number of new problems and questions that had not been expected or anticipated before. Ten years after the above two reviews, *Russell* [1995] edited a proceedings volume collecting the accumulated research of this period. And another ten years later *Li, Zank & Russell* [2005] organised a conference that was devoted to reviewing the more recent achievements in shock research.

2.1.6 The Numerical Simulation Age

The advent of powerful computer resources in the mid-sixties completely changed the research attitudes also in collisionless-shock physics. Starting from the idea that all the physics that can be known is contained in the basic physical laws, which can be represented by a set of conservation equations,⁵ the idea arose that, searching the domain of solutions of these conservation laws in the parameter space prescribed by observations with the help of the new computing facilities would not only liberate one from the tedious burden of finding mathematically correct analytical solutions of these laws, but would also expand the accessible domain of solutions into those directions where no analytical solutions could be found. With this philosophy in mind a new generation of researchers started developing numerical methods for solving the basic nonlinear equations of plasma physics with the help of powerful computer facilities circumventing the classical methods of solving partial differential equations and enabling attacking any nonlinear problem by sometimes straight forward brute force methods. In fact, collisionless shocks are particularly well suited for the application of such methods just because they are intrinsically nonlinear. The decades since the late sixties were thus marked by an inflation of numerical approaches, so-called computer simulations, to shock physics thereby paralleling similar developments in all fields of exact scientific research. The ever-increasing capacities of the computer have been very tempting for performing simulations. However, the capacities are still suited only for very well-tailored simulation problems.

It is, moreover, clear that numerical simulations *per se* do not allow to make fundamental discoveries which go beyond the amount of information that is already contained in the equations one is going to solve on the computer. Still, this is an infinite number of problems out of which the relevant and treatable must be carefully and insightfully extracted. Any simulation requires, in addition, the application of subsequent data analysis which closely resembles the analysis of observational data obtained in real space or laboratory experiments. Because of this second reason one often speaks more correctly of *computer experiments* or numerical experiments instead of numerical simulations.

There are two fundamentally different directions of simulations [for a collection of methods, see, e.g. *Birdsall & Langdon*, 1991]. Either one solves the conservation laws which have been obtained to some approximation from the fundamental Liouville equation, or one goes right away back to the number of particles that is contained in the simulation volume and solves for each of them the Newtonian equation of motion in the self-

⁵Below in Chapter 3 we will provide a number of such model conservation laws.

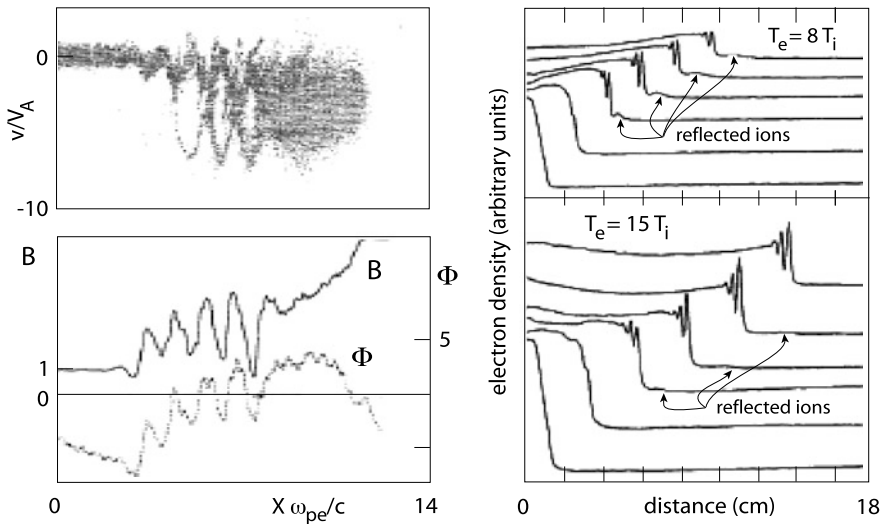


Figure 2.4: One-dimensional numerical simulations by (right) Taylor *et al* [1970] and (left) Biskamp & Welter [1972a] of shocks with reflected ions. In Taylor *et al*'s simulation the shock ramp is shown in time-stacked profiles with time running from left corner diagonally upward. The density ramp is shown. The initial profile has been assumed as a steep ramp. Reflection of ions leads with increasing time to its oscillatory structure and generation of a foot as had been suggested by Woods [1969]. Biskamp & Welter's simulation shows the phase-space (upper panel), the evolution of reflected (negative speed) ions and heating behind the shock ramp. In the lower panel the evolution of the magnetic field with distance in magnitude and direction is shown. The field exhibits strong undulations caused by the gyrating reflected ions in the solar wind in front of the ramp. A rotation of the field angle through the shock is also detected.

generated fields. Both methods have been applied and have given successively converging results. Both approaches have their advantages and their pitfalls. Which is to be applied depends on the problem which one wants to solve. For instance, in order to determine the global shape of the Earth's bow shock it makes no sense to refer to the full particle approach; a fluid approach is good enough here. On the other hand if one wants to infer about the reflection of particles from the shock in some particular position on the smaller scale, a full particle code or also a Vlasov code would possibly be appropriate.

Since the beginning, computer experiments have proven very valuable in collisionless shock physics. In order to elucidate the internal physical structure of shocks one is, however, directed to particle codes rather than fluid codes. First electrostatic simulations of one-dimensional collisionless shocks have been performed with very small particle numbers and in very small simulation boxes (spatio-temporal boxes with one axis the space coordinate, the other axis time) by Dawson & Shanny [1968], Forslund & Shonk [1970] and Davidson *et al* [1970]. They observed the expected strong plasma heating (see Figure 2.5). Slightly later Biskamp & Welter [1972a, b], from similar one-dimensional simulations, found the first indication for reflection of particles in shock formation. With larger simulation boxes a decade later, the same authors [Biskamp & Welter, 1982] confirmed

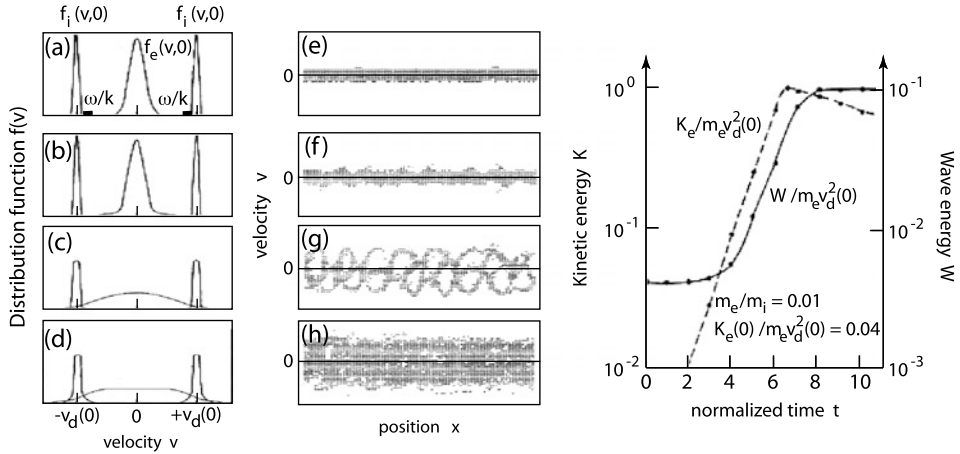


Figure 2.5: A one-dimensional purely electrostatic particle simulation of shock plasma heating as observed by Davidson *et al* [1970]. The (unrealistic model) simulation in the shock reference frame consists of a thermal electron background and two counter-streaming ion beams (*left column*) of exactly same temperature and density modelling the inflow of plasma (forward ion beam) and reflected ions (backward beam). The electrons are assumed hotter than ions neglecting their bulk motion but cold enough for allowing the Buneman instability [Buneman, 1959] to develop which requires a strong shock. The *left panel* shows the evolution of the distribution functions: broadening of all distributions in particular of electrons indicating the heating. The *middle column* is the electron phase space representation showing the disturbance of electron orbits, increase of phase space volume and final thermalisation. Probably the most interesting finding is in the third phase space panels which shows that as an intermediary state the electron orbits show formation of voids in phase space. These holes contain strong local electric fields but were not recognised as being of importance at this time. They were independently discovered theoretically only a few years later [Schamel, 1972, 1975, 1979; Dupree *et al*, 1975; Dupree, 1982, 1983, 1986; Berman *et al*, 1985]. The *right panel* shows the evolution of electron kinetic and electric field energy during the evolution.

that, in the simulations, high Mach number shocks indeed reflected ions back upstream when the Mach number exceeded a certain critical value.

These first full particle simulations were overseeded in the eighties by hybrid simulations, which brought with them a big qualitative and even quantitative step ahead in the understanding of collisionless shocks [Leroy *et al*, 1982; Leroy, 1983, 1984; Leroy & Mangeney, 1984; Scholer, 1990; Scholer & Terasawa, 1990, and many others], a method where the ions are treated as particles, while the electrons are taken as a fluid. These were all one-dimensional simulations where an ion beam was allowed to hit a solid wall until being reflected from the wall and returning upstream. Electrons were treated as a *massless isothermal fluid*. The perpendicular shock that was investigated by these methods, evolved in the interaction of the incoming and reflected ion beams and, in the frame of the fixed reflecting wall, moved upstream at a certain measurable speed. The free energy required for the shock and entropy production resided in the velocity difference between the upstream and reflected ion beams.

The first of these simulations concerned subcritical perpendicular shocks, showing the effect of shock steepening. It included some numerical resistivity. Increasing the Mach

number, monitored the transition to ion reflection and formation of the predicted shock foot in front of the shock. The later one-dimensional simulations included oblique magnetic fields and started investigating the reflection process in dependence on the shock normal angle Θ_{Bn} . They, moreover, focussed on particle acceleration processes which are of primary interest in astrophysical application.

The hybrid simulations reproduced various properties of collisionless shocks like the ion-foot formation in perpendicular and quasi-perpendicular shocks, shock reformation in quasi-perpendicular shocks and ion acceleration in a combination of hybrid and test-particle simulations. They also allowed to distinguish between subcritical and supercritical shocks. They allowed for the investigation of the various low-frequency waves which are excited in the ion foreshock by the super-critical shock-reflected ion distribution. Since, in such simulations, the electrons act passively being a neutralising fluid, their dynamics is completely neglected. Nevertheless, a number of important properties of collisionless shocks was recovered. In addition it was found that, contrary to pure fluid theory, supercritical shocks turned out to possess a self-generated non-coplanar magnetic field component. And it was also found that the shock ramp itself was nonstationary and behaved stationary only in the average over time scales long compared with the reformation-time scale and spatial scales extended along the shock front. Clearly, hybrid simulations can be used to describe the meso-scale structure of shocks, but most of the interesting shock physics takes place on shorter and smaller scales, when the electron dynamics is included. This can be done by various approximations, relaxing the severe assumption of massless Boltzmannian electrons.

Currently the investigation of collisionless shocks has achieved a certain state of saturation. Analytical theories have basically been exhausted for the last fifteen years. The complexity of shock structure has set a fairly strict bound to it unless completely new mathematical methods will be brought up allowing for taking this complexity into account without further increasing the lack of mathematical simplicity and transparency.

Measurements *in situ* have as well been pushed to their bounds. Improved instrumentation and the increase in on-board data storage capabilities and speed of on-board data analysis and data reduction provided an enormous data pool that is already available and which is by far not yet exhausted. However, most of the data obtained are from single spacecraft observations and thus suffer from being single point measurements such that it is difficult or even impossible to distinguish between spatial effects and temporal evolution.

On the other hand, the multi-spacecraft observations like combinations of different spacecraft suffer from the accidental character of their connectedness with the bow shock or other shocks in interplanetary space. Specially designed multi-spacecraft missions like CLUSTER suffer from the inflexibility of tuning the sub-spacecraft separations. Some of these pitfalls can be overcome by sophisticated data analysis methods [Paschmann & Daly, 2000] but the progress is not as overwhelming as one would hope. Moreover, between design and launch of such multi-spacecraft lies roughly one decade such that the instrumentation cannot keep track with the development in technology which is available when the spacecraft is put into orbit.

Finally, even though the computing capabilities increase almost exponentially, the currently available computers still cannot manage to solve the orbits of the vast numbers of particles in a realistic volume in space and phase space. For instance the spatial volume at the perpendicular bow shock wave of Earth amounts to roughly $V_{BS} \sim 50 R_E^3 \approx 2 \times 10^{19} \text{ m}^3$ which for an average density of $N \sim 10^7 \text{ m}^{-3}$ implies just $NV_{BS} \sim 10^{26}$ protons plus the same number of electrons in the volume and passing through the shock at a speed of a few 100 km/s. The orbits of such numbers of particles in their self-excited and external fields cannot be calculated even by the most powerful computers.

Hence, in particle codes, particles must be grouped together into macro-particles neglecting their short-distance interactions and short-distance fields. As long as this grouping affects small particle numbers its effects on the final results are certainly negligible. However, since computers so far can manage to solve only up to the order of few times $\sim 10^9$ particle orbits simultaneously, investigation of the full perpendicular shock with particle codes requires a macro-particle size collection of the order of $\sim 10^{17}$ individual particles into one macro-particle, which is not small anymore but forms super-clusters of particles with all their unknown internal and external dynamics. Since at a density of $N \sim 10^7 \text{ m}^{-3}$ each particle occupies an elementary volume of $N^{-1} \sim 10^{-7} \text{ m}^3$, each of these macro-particles occupies a volume of $\sim 10^{10} \text{ m}^3$ or a plasma blob of 10 km^3 volume.

For comparison, in the solar wind the Debye length is $\lambda_D \sim 23 \text{ m}$; the *real* Debye volume is thus six orders of magnitude smaller than the macro-particle volume, negligibly small compared with the macro-particle volume and cannot be resolved in three-dimensional full particle simulations. Any resolvable linear spatial scale must be larger than $\ell_{MP} \sim 5 \text{ km}$. The electron gyro-radius is about 200 km, the electron inertial length about 1700 km such that these lengths can be resolved. Similarly, any proton scale can be resolved as well. However care must be taken when simulating problems where the electron Debye scale is involved like in radiation problems. Moreover since the scale of macro-particles is a real-space scale, the validity of any one- or two-dimensional simulations in the directions of the neglected coordinates is restricted to scales $\gg 5 \text{ km}$ which in the solar wind is not a severe restriction. In order to resolve the Debye scale in radiation problems the system is assumed to be homogeneous over $\sim 500\text{--}1000 \text{ km}$ in two-dimensional and $50,000\text{--}100,000 \text{ km}$ in one-dimensional simulations in the directions on which coordinates the simulation is independent. At the same time, however, in the remaining simulation direction the resolution can be made sufficiently high enough to resolve the relevant times and lengths.

However, the main problem is not the resolution itself. It is the neglect of the interactions between the particles constituting the macro-particles, since these are dynamical systems which continuously exchange particles, momentum, and energy among themselves which may or may not introduce systematic errors. The limitations on the simulation results are very difficult to estimate, however, and in most cases are believed to be unimportant. Conservation of macro-particle number and total energy are usually good measures of this validity of the simulations, which are usually stopped when these conservations becomes violated in the course of the runs, and the results are taken valid for simulation times only where the energy and particle numbers were conserved within a few per cent. However, even in this case it must be stressed that, from simulations, it can only

be inferred what is already contained in the equations. They are not suited for making any real physical discoveries, while they can make predictions that can be checked against observation, thereby validating or falsifying the underlying theory.

2.2 When Are Shocks?

At a first glance the notion of a collisionless shock seems nonsensical⁶ [for a popular review see *Sagdeev & Kennel, 1991*]. Indeed, one intuitively imagines that a shock requires that something is shocked by collisions. However, Nature is not always organised the way one naively believes. Though ordinary gasdynamic shocks or shocks in condensed matter are indeed shocks which cannot be thought of when ignoring the high collisionality,⁷ collisionless systems like high-temperature dilute plasmas take advantage of nonlocal, non-binary, ‘anomalous collisions’ between particles and the existing external and selfconsistently-generated electromagnetic fields. These ‘anomalous collisions’ are in fact long-range *collective interactions* between groups of particles and fields; they lead to correlations between these groups, such that the particles do not anymore behave like freely moving ballistic particles. Because of the correlations the particles, together with the fields, organise themselves to form structures. The possibility of such structure formation in a completely collision-free plasma has been an extraordinarily important insight. It leads to close interaction and the generation of irreversible dissipation processes in plasma that result in heating, acceleration of groups of particles, generation of entropy and emission of radiation, which all are different forms of energy distribution in and energy loss from the plasma. It is the very meaning of these collective processes, to make the otherwise completely dissipation-free plasma capable of returning to a thermodynamic state of destruction of the available free energy and transforming it into heat and other less valuable forms of energy. One of the structures, which contribute to this kind of dissipation, are collisionless shock waves.

Collisionless shock waves form when a large obstacle is put into a plasma flow that is either super-Alfvénic or super-magnetosonic in the frame of the obstacle. Super-Alfvénic flows have Mach numbers $\mathcal{M}_A = V/V_A > 1$, where $\mathbf{V}_A = \mathbf{B}/\sqrt{\mu_0 m_i N}$ is the Alfvén velocity in the magnetised plasma of density N and magnetic field $\mathbf{B}$, and m_i is the ion mass. Correspondingly, super-magnetosonic flows have Mach numbers $\mathcal{M} \equiv \mathcal{M}_{ms} = V/c_{ms} > 1$, where $c_{ms}^2 = V_A^2 + c_s^2$ is the square of the magnetosonic speed, and $c_s^2 = \partial P / \partial \rho$ is the square of the ordinary sound speed, with P the isotropic pressure, and $\rho = m_i N$ the mass density. In plasma of respective electron and ion temperatures T_e and T_i (in energy units) and adiabatic indices $\gamma_{e,i}$ for electrons and ions, the latter is (with sufficient accuracy) given by $c_s^2 = (\gamma_e T_e + \gamma_i T_i) / m_i$. We should, however, note that this discussion is based on fluid considerations, and even then it is valid in this form only for the so-called fast-magnetosonic

⁶This is also the first reaction one encounters when talking to either laymen or even physicist working in a different area of physics. Usually one earns a forgiving smile from either of them.

⁷This means, ignoring the short range forces involved in binary collisions when the particles (or solid objects) literally touch and hit each other.

mode perpendicular to the magnetic field $\mathbf{B}$. The magnetosonic speed depends on the wave propagation angle θ with respect to the magnetic field. Its general fluid expression

$$c_{ms}^2(\theta) = c_{ms}^2 \pm [(V_A^2 - c_s^2)^2 + 4V_A^2 c_s^2 \sin^2 \theta]^{\frac{1}{2}} \quad (2.1)$$

where c_{ms} is the angle-independent expression given above, shows that the magnetosonic velocity contains two branches, c_{ms}^+ related to the *fast* magnetosonic wave mode with the positive sign in front of the root, and c_{ms}^- related to the *slow* magnetosonic wave mode with negative sign in front of the root. In addition there is the Alfvén wave with speed V_A that is independent of the sound velocity c_s . Of these three modes only the fast mode propagates perpendicular to the magnetic field. However, in all other directions all three modes can exist and, therefore, each of them may form a shock, if only the flow speed exceeds its velocity. Hence there can, in principle, three different kinds of shocks exist, the *fast*, *slow*, and *Alfvénic* (or *intermediate*, because its speed is intermediate between the fast and slow waves) shocks, respectively.

Collisionless shocks are macroscopic phenomena in which very many particles are involved. The requirement on the obstacle is that its diameter D in the two directions perpendicular to the flow must be very large compared to the intrinsic scales of the flow while at the same time being much less than the collisional mean free path. The largest intrinsic scale of a magnetised plasma is the ion gyro-radius $r_{ci} = V_{i\perp}/\omega_{ci}$, where $V_{i\perp}$ is the ion flow velocity perpendicular to $\mathbf{B}$, and $\omega_{ci} = eB/m_i$ is the ion-cyclotron angular frequency. Hence, with the Coulomb mean free path $\lambda_C = (N\sigma_C)^{-1}$ defined earlier, the above condition is simply that

$$r_{ci} \ll D \ll \lambda_C \quad (2.2)$$

The shock forms an extended surface that is bent around the obstacle. An example of the form of such a shock is drawn in Figure 2.6, showing (among various other of its properties) the average shape of the Earth's bow shock in an ecliptic cross section that is caused in the interaction of the solar wind with the – approximately – dipolar geomagnetic field.

The radius of curvature $R_c \gg \Delta$ of the shock perpendicular to the flow will always be much larger than the width $\Delta \sim r_{ci}$ of the shock in the direction of flow, the latter being of the order of a few ion gyro-radii only. Collisionless shocks can in good approximation be considered as thin, locally flat surfaces of width Δ and outer shock normal $\mathbf{n}$. Given the function of the shock surface $F_S(\mathbf{r}) = r$, where r is the radius vector from the (arbitrary) centre to one point on the shock surface, the shock normal is defined as

$$\mathbf{n}(\mathbf{r}) = -\frac{\nabla F(\mathbf{r})}{|\nabla F(\mathbf{r})|} \quad (2.3)$$

the negative sign telling that it points outward (the shock being seen from the inside as a concave surface), the gradient accounting for the direction of strongest variation, and the normal is normalised to the gradient since it is a unit vector that satisfies the condition $\mathbf{n} \cdot \mathbf{n} = 1$.

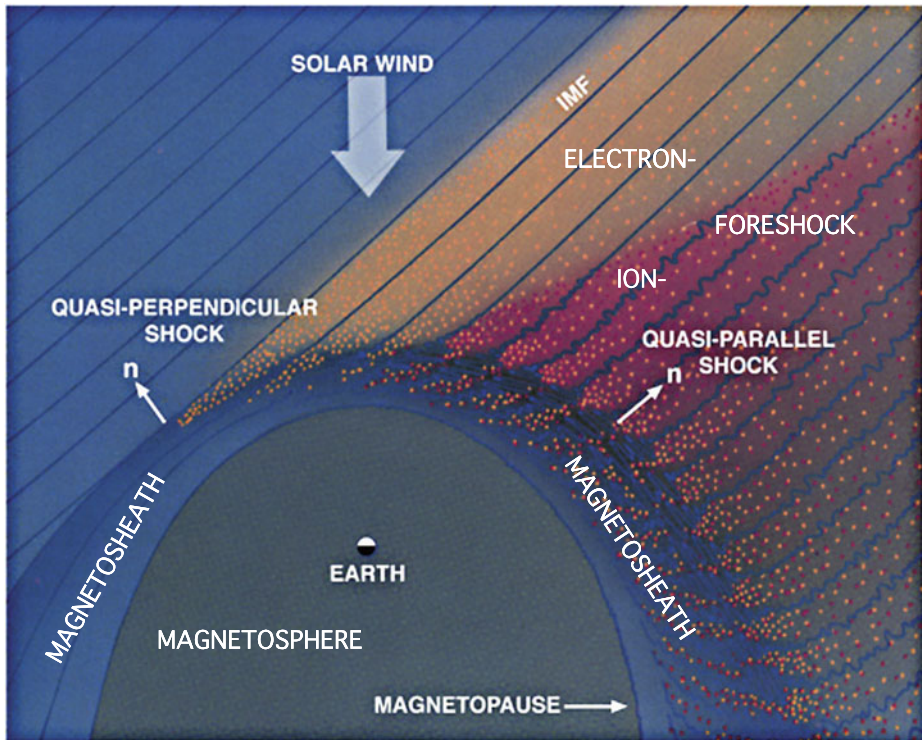


Figure 2.6: A two-dimensional schematic view on Earth's steady-state bow shock in front of the blunt magnetosphere [after Tsurutani & Stone, 1985, with permission of the American Geophysical Union] which forms when the supersonic solar wind streams against the dipolar geomagnetic field. The bow shock is the diffuse hyperbolically shaped region standing at a distance in front of the magnetopause. The inclined *blue lines* simulate the solar wind magnetic field (interplanetary magnetic field IMF). In this figure they lie in the plane. The direction of the shock normal is indicated at two positions. Where it points perpendicular to the solar wind magnetic field the character of the bow shock is perpendicular. In the vicinity of this point where the solar wind magnetic field is tangent to the bow shock the shock behaves quasi-perpendicularly. When the shock is aligned with or against the solar wind magnetic field the bow shock behaves quasi-parallel. Quasi-perpendicular shocks are magnetically quiet compared to quasi-parallel shocks. This is indicated here by the gradually increasing oscillatory behaviour of the magnetic field when passing along the shock from the quasi-perpendicular part into the quasi-parallel part. Correspondingly, the behaviour of the plasma downstream of the shock is strongly disturbed behind the quasi-perpendicular shock. Finally, when the shock is super-critical, as is the case for the bow shock, electrons and ions are reflected from it. Reflection is strongest at the quasi-perpendicular shock but particles can escape upstream only along the magnetic field. Hence the upstream region is divided into an electron (yellow) and an ion foreshock accounting for the faster escape speeds of electrons than ions.

The important consequence of the above scaling is that, locally, of all spatial derivatives $\partial/\partial x \sim \partial/\partial y \ll \partial/\partial n$ only the derivative across the shock front counts. The gradient operator ∇ thus reduces to the derivative in the direction opposite to the local shock normal $\mathbf{n}$ or, with coordinate n ,

$$\nabla \simeq -\mathbf{n} \frac{\partial}{\partial n} \sim -\mathbf{n} \frac{[\dots]}{\Delta} \quad (2.4)$$

where the brackets $[\dots] = (\dots)_2 - (\dots)_1$ stand for the difference of the values of the quantity under consideration downstream in Region 2 behind the shock minus the value upstream in Region 1 in front of the shock. Formal theory is extensively making use of this fact. Clearly, however, it loses its importance in full particle simulations and in particular in simulations in more than one dimension. In those simulations the shock turns – to a certain degree trivially – out to be nonstationary, i.e. even in one-dimensional simulations the shock steepness (and thus the local shape of the shock) varies with time.

Rather than its shape, the most important physical property of the shock is its capability of slowing down the upstream flow from super-magnetosonic to magnetosonic or even sub-magnetosonic speed across the shock, i.e. from $\mathcal{M} > 1$ to $\mathcal{M} \leq 1$, on the extraordinarily short distance Δ of the shock width. Similar as for gasdynamic shocks, this implies that a substantial amount of upstream flow energy must be converted into compression of the plasma, i.e. into enhanced pressure over the narrow length Δ . This is possible only if, in addition, the plasma is heated over the same distance. The shock thus creates a downstream region of high pressure, producing entropy, and separates it from the upstream lower-pressure region. From this point of view, collisionless shocks are alike to gasdynamic shocks, even though there are no collisions between particles. Other mechanisms are required bridging the lack of collisions and providing the necessary dissipation. Moreover, being a permeable boundary between two regions of different temperatures and pressures, the shock is *not in thermal equilibrium*. Thermal non-equilibria in closed systems cannot survive, however. They have the tendency of evolving towards thermal equilibrium. It is therefore important to realise that *a shock cannot be stationary*; in order to be maintained over long periods *it must be continuously reformed*. This is indeed the case with all the collisionless shocks observed in the heliosphere. Collisionless shocks as quasi-stationary, non-transient phenomena occur only in open systems – like in the solar wind-magnetosphere interaction – and are continuously reformed on the expense of the energy and momentum density of the plasma inflow.

2.3 Types of Collisionless Shocks

When speaking about shock waves the implication is that, like an ordinary wave, the shock propagates on the background plasma. This propagation is not obvious when looking for instance at Earth's bow shock which, when neglecting its irregular change of position, is about stationary in the long-term average, standing in front of the magnetosphere in the solar wind. However, the more correct view is when looking at the shock in the frame of the moving solar wind. Seen from there the shock is propagating in the solar wind upstream towards the Sun at velocity that is approximately comparable to the solar wind speed. In the solar wind frame, the shock appears as a compressive upstream moving wave front (or wave ramp), like a tsunami in the ocean following a seaquake. Later we will illuminate the question, how such a large amplitude shock ramp can form. In preparation of a deeper

discussion we, in this section, provide a rough classification of the various types of shocks that can develop in a collisionless plasma.

2.3.1 Electrostatic Shocks

Shocks can be classified from various points of view. The first and simplest classification is with respect to the electrodynamic properties of shocks. Plasmas consist of electrically charged particles which under normal conditions for maintaining overall charge neutrality occur in about equal numbers per volume element. They have same number densities, $N_e = N_i = N$. When the different charges do in the average not possess different bulk velocities the plasma is free of electric currents $\mathbf{j} = Ne(\mathbf{V}_i - \mathbf{V}_e) = 0$ and, in the absence of an external magnetic field $\mathbf{B}_0$, the plasma is free of magnetic fields. It behaves purely electric. In this case, a shock wave which occurs in the flow is called an *electrostatic shock*.

In the heliosphere such electrostatic shocks are rare, because most moving plasmas are magnetised. They may, however, occur under certain very special conditions even in the strongly magnetised plasmas in the auroral zones of magnetised planets and in the particle acceleration zones in the solar corona during particular flare events. In these cases, electrostatic shocks are strictly one-dimensional however and occur only on very small scales where they contribute to the generation of magnetic-field aligned electrostatic fields. These fields can promptly accelerate particles to energies of the order of the total macroscopic electrostatic potential drop. We will return to this problem in this book in Chapter 7 on particle acceleration. However, this kind of shocks does not belong to the regular large-scale genuine shocks that lie at the focus of this text. Nevertheless, the theoretical and numerical investigation of electrostatic shocks was instrumental for the understanding of shock physics.

2.3.2 Magnetised Shocks

The vast majority of collisionless shocks in the heliosphere – and as well in astrophysics – belongs to a different class of shocks known as *magnetised shocks* simply, because the plasma in the heliosphere is magnetised and allows for electric currents to flow across and along the magnetic field. For instance, when a magnetised super-magnetosonic moving plasma is shocked, it is quite natural that the different drift motions of particles of opposite charges generated in the plasma gradient of the shock ramp cause electric drift currents to flow across the magnetic field. These currents are accompanied by proper secondary magnetic fields and, in addition, cause other effects like anomalous transport and plasma heating. Magnetised shocks, therefore, behave quite differently from electrostatic shocks. They are of macroscopic nature and, because of their abundance, are much more important.

Since the presence of the magnetic field is the first main distinction between electrostatic and magnetised shocks, it is reasonable, for a first classification, to distinguish between shocks in which the upstream magnetic field is tangential to the shock surface and those shocks, where the magnetic field is perpendicular to the shock surface. The reason for such a distinction is that in the former case the upstream flow velocity, $\mathbf{V}_1 \perp \mathbf{B}_1$, is perpendicular to the upstream magnetic field while, in the second case, the upstream

flow is parallel to the magnetic field, $\mathbf{V}_1 \parallel \mathbf{B}_1$. This has the consequence that, in the perpendicular case, the magnetic term $\mathbf{V}_1 \times \mathbf{B}_1 \neq 0$ in the Lorentz force is finite. The magnetic field lines are convected with the flow and pile up at the shock ramp. In the parallel case, the magnetic term in the Lorentz force vanishes identically. From a naive point of view, such shocks become unmagnetised and should behave like gasdynamic shocks. This is, however, not the case, as we will see in later chapters. Shocks, where the flow is parallel to the field, behave very different from shocks where the flow is perpendicular.

In order to distinguish between these two types of shocks, one rather refers to the shock normal defined in Eq. (2.3), which gives a precise local definition of the shock surface. Defining the shock normal angle Θ_{Bn} through

$$\tan \Theta_{Bn} = \mathbf{n} \cdot \mathbf{B}_1 / |\mathbf{B}_1| \quad (2.5)$$

we can then distinguish *perpendicular shocks* with $\Theta_{Bn} = \frac{1}{2}\pi$ and *parallel shocks* with $\Theta_{Bn} = 0$. In the nature, at a large bent shock, the two extreme cases are realised only over a small portion of the entire shock surface. One therefore rather distinguishes *quasi-perpendicular* and *quasi-parallel* shocks, which are operationally defined by $\frac{1}{4}\pi < \Theta_{Bn} \leq \frac{1}{2}\pi$ and $0 \leq \Theta_{Bn} < \frac{1}{4}\pi$, respectively. In the overlap region $\frac{1}{6}\pi < \Theta_{Bn} \leq \frac{1}{3}\pi$ one speaks about *oblique* shocks keeping, however, in mind that in contrast to the distinction between parallel and perpendicular shocks the term oblique is conventional only and is not required by physics. Oblique shocks simply have mixed parallel/perpendicular shock properties.

Figure 2.6 shows the artistic drawing of an example of a supercritical collisionless shock wave, the Earth's bow shock, located in the super-magnetosonic and magnetised solar wind exhibiting regions of all three kinds of supercritical shock waves along its bent hyperbolically shaped surface. The outer normal $\mathbf{n}$ to the shock has been drawn in two point, the completely perpendicular region on the left where the interplanetary magnetic field is strictly tangential to the idealised shock surface, and in the strictly parallel region where the idealised magnetic field is about parallel to the shock normal. There is a pronounced difference in both locations shown in the drawing in reference to the observations which will be discussed in depth in Chapter 10. This difference refers to the degree of distortion of the magnetic field and shock in both positions. The perpendicular shock is considerably less disturbed than the parallel shock. The yellow and red coloured regions indicate the spatial domains filled with particles, electrons and protons, reflected from the supercritical shock. In the region along the shock surface between the perpendicular and parallel parts the shock changes its character gradually from quasi-perpendicular through oblique to quasi-parallel.

The magnetic field introduces another important property of the plasma, i.e. a pressure anisotropy. The pressures parallel $P_{\parallel}$ and perpendicular $P_{\perp}$ to the magnetic field can become different, and the pressure is in fact a tensor $\mathbf{P} = P_{\perp} \mathbf{I} + (P_{\parallel} - P_{\perp}) \mathbf{BB}/B^2$, where $\mathbf{I}$ is the unit tensor. This anisotropy has the effect that the ratio $\beta = 2\mu_0 P/B^2$ of thermal to magnetic field pressure in general becomes anisotropic as well, with $\beta_{\parallel} = 2\mu_0 P_{\parallel}/B^2 \neq \beta_{\perp} = 2\mu_0 P_{\perp}/B^2$. Owing to this we define an anisotropy factor

$$A = \frac{\beta_{\perp}}{\beta_{\parallel}} - 1 \equiv \frac{P_{\perp}}{P_{\parallel}} - 1 \equiv \frac{T_{\perp}}{T_{\parallel}} - 1 \quad (2.6)$$

which can be positive for an excess in perpendicular energy, negative for an excess in parallel energy, and it also can vanish when the anisotropy is vanishingly small.

2.3.3 MHD Shocks

The simplest and historically also the first model approach to collisionless shocks was within magnetohydrodynamic fluid theory by simply adding to the frictionless gasdynamic equations the Lorentz force

$$\mathbf{F}_L = \mathbf{j} \times \mathbf{B} \quad (2.7)$$

with current density $\mathbf{j}$ defined through Ohm's law

$$\mathbf{j} = \sigma(\mathbf{E} + \mathbf{V} \times \mathbf{B}) \quad (2.8)$$

Under collisionless ideally conducting conditions the conductivity is $\sigma \rightarrow \infty$, and Ohm's law is replaced by the ideal MHD condition

$$\mathbf{E} = -\mathbf{V} \times \mathbf{B} \quad (2.9)$$

for the relation between the electric field $\mathbf{E}$, fluid velocity $\mathbf{V}$, and magnetic field $\mathbf{B}$. These relations must be completed by some equation of state relation between the plasma pressure P , density N and temperature T (in energy units) for which usually the ideal gas law $P = NT$ is taken to hold in its adiabatic version, assuming that shock formation proceeds on such a short scale that the temperature cannot adjust. This is in fact not an unreasonable assumption as the shock cannot be in thermal equilibrium as has been argued above on different reasons.

The three possible shocks in such a case are just the *fast*, *slow* and *intermediate* shocks related to the three MHD wave modes that have been mentioned above. They may occur depending on which of the wave phase speeds is exceeded by the flow, and for the slower shocks, of course, under the additional condition that the faster waves are inhibited in the medium because otherwise, when the faster waves would be excited in the interaction of flow and obstacle they would propagate upstream at faster speed than the shock itself and inform the flow about the presence of the obstacle. No shock would be formed in this case from the slower Alfvénic or slow modes. Since this is the more realistic case the most frequently observed shocks are fast shocks. However, occasionally also slow or intermediate shocks have been claimed to have been detected in interplanetary space.

The four possible shock transitions in terms of the relations between flow and MHD wave mode speeds [De Hoffman & Teller, 1950; Balogh & Riley, 2005] are:

$$\begin{aligned} \text{trans 1 : } & V > c_{ms}^+, & \text{trans 2 : } & c_{ms}^+ > V > c_{int} \\ \text{trans 3 : } & c_{int} > V > c_{ms}^-, & \text{trans 4 : } & c_{ms}^- > V \end{aligned}$$

where the velocities $c_{ms}^\pm$ have been defined in Eq. (2.1), and $c_{int} = V_A \cos \theta$ is the angle-dependent Alfvén velocity of the intermediate wave. Not all these transitions can, however,

be realised. *Wu & Kennel [1992]* have shown that entropy considerations allow only for the transitions $1 \rightarrow 2$, $1 \rightarrow 3$, $1 \rightarrow 4$, $2 \rightarrow 3$, $2 \rightarrow 4$, and $3 \rightarrow 4$. Of these the first is a fast mode shock transition, while the transition $3 \rightarrow 4$ is a slow mode, and the remaining ones are all intermediate shock transitions which together with the slow mode might sometimes exist under the above mentioned restrictions. However, usually an obstacle will excite all three waves together, and then only the fast mode will cause a shock. The observation of slow or intermediate shocks thus requires very special conditions to exist in the plasma. In addition, any shocks will have to respect also the evolutionary condition discussed briefly below.

In the MHD frame, one can define simple relations between the parameters of the streaming gas upstream and downstream of the shock, the so-called *Rankine-Hugoniot relations*. These relations give a first idea of the conditions at the shock transition. They result from the conservative character of the MHD equations which, in fact, are conservation laws for the mass flow, momentum flow and energy density. They will be derived in the next Chapter 3. The restriction on them is that the processes, which determine the generation of entropy and causing the irreversibility of the shock, must all be strictly confined solely to the shock transition region. This is not so easy to achieve as it might seem at first glance. The shock transition might be much broader and more extended than the proper shock ramp suggests. Hence the parameter values that enter the Rankine-Hugoniot relations, are taken correctly only sufficiently far away to both sides of the shock ramp in order to assure of not mixing-in processes that are dissipative and thus are not contained in the MHD Rankine-Hugoniot conservation laws. The problem is then that the shock surface itself must be extended enough, compared with the distance from the shock where the parameter values are taken. In addition, its curvature should still be negligible in order to not destroy the assumption of shock planarity that is involved into the Rankine-Hugoniot relations.

2.3.4 Evolutionarity

In the above paragraph we mentioned that not all of the six possible shock solutions in MHD can be realised. The actual reason for this lies in the so-called conditions of evolutionarity of shock waves, which are based on the hyperbolic nature of the conservation laws which allows wave propagation only if it is in accord with causality [*Lax, 1957*].

For MHD waves with dissipation these conditions have been discussed by *Jeffrey & Taniuti [1964]*, *Kantrowitz & Petschek [1966]*, *Liberman & Velikhovich [1986]* and others. They also hold in the collisionless regime, because causality is a general requirement in Nature, meaning in this case that the drop in speed across a shock (in the wave mode of the shock) must be large enough for the normal component of the downstream flow to fall below the corresponding downstream mode velocity. For a fast shock this implies the following ordering of the normal flow and magnetosonic velocities to both sides of the shock:

$$V_{1n} > c_{1ms}^+, \quad \text{while} \quad V_{2n} < c_{2ms}^+$$

where the numbers 1, 2 refer to upstream and downstream of the fast shock wave.

The first condition is necessary for the shock to be formed at all; it is the second condition which (partially) accounts for the evolutionarity. Otherwise the small fast-mode disturbances excited downstream and moving upward towards the shock would move faster than the flow, they would overcome the shock and steepen it without limit. Since this cannot happen for a shock to form, the downstream normal speed must be less than the downstream fast magnetosonic speed. Similar conditions hold for any shock as also for large amplitude shocks. Furthermore, for fast shocks the flow velocity must be greater than the intermediate speed on both sides of the shock, while for slow shocks it must be less than the intermediate speed on both sides. These conditions hold because of the same reason as otherwise the corresponding waves would catch up with the shock front, modify and destroy it and no shock could form.

2.3.5 Coplanarity

Another observation of MHD shocks is related to the directions of the magnetic field and flows to both sides of the shock front. These directions are not arbitrary. At the contrary, it can be shown from the MHD conservation laws respectively the Rankine-Hugoniot relations that the flow and magnetic field directions in front and behind the shock front in MHD lie in the same plane, i.e. they are coplanar. This property had been realised already by *Marshall* [1955] and has been discussed in depth by *Kantrowitz & Petschek* [1966] and others [see e.g. *Burgess*, 1995]. For a stationary ideal MHD shock wave with no other wave activity or kinetic processes present outside the shock transition, such that dissipation takes place solely inside the narrow shock transition and this transition region can be considered as infinitesimally thin with respect to all other physical scales in the plasma, the electric field in the shock rest frame is strictly perpendicular to the magnetic field, given by Eq. (2.9). The equation

$$\nabla \times \mathbf{E} = 0 \quad (2.10)$$

which is the stationary Faraday's law, and the shock normal $\mathbf{n}$ defined in Eqs. (2.3), (2.4) yield that the scalar product between $\mathbf{n}$ and the difference in the tangential components of the magnetic field to both sides vanishes:

$$(V_{n2} - V_{n1})\{\mathbf{B}_{t2} \times \mathbf{B}_{t1}\} = 0 \quad (2.11)$$

The difference in the normal components across the shock does clearly not vanish, such that $V_{n2} \neq V_{n1}$. Hence

$$\mathbf{n} \cdot \{\mathbf{B}_2 - \mathbf{B}_1\} = 0 \quad (2.12)$$

which implies not only that $\mathbf{n}$ is normal to the tangential components of the magnetic field on both sides of the infinitesimally thin discontinuity, which would be a trivial conclusion, but also that the two tangential components to both sides are strictly parallel. They may – and should – have different lengths but will have same direction across the shock.

Coplanarity does not strictly hold, however. For instance, when the shock is non-stationary, i.e. when its width changes with time or in the direction tangential to the shock,⁸ the right-hand side in Faraday's law does not vanish, and coplanarity becomes violated.

Also, any upstream low frequency electromagnetic plasma wave that propagates along the upstream magnetic field, possesses a magnetic wave field that is perpendicular to the upstream field. When it encounters the shock, this tangential component will be transformed and amplified across the shock. This naturally introduces an out-of-plane magnetic field component, thereby violating the co-planarity condition. There are also other effects which at a real non-MHD shock violate coplanarity. We will not discuss them at this location. One particular case in MHD is noted in the following subsection.

2.3.6 Switch-On and Switch-Off Shocks

Parallel shocks in MHD should, theoretically, behave exactly like gasdynamic shocks, not having any upstream tangential magnetic field component and should also not have any downstream tangential field. The 'tangential field' in this case has 'no direction'.

This conclusion does not hold rigorously, however, since plasmas consist of charged particles which are sensitive to fluctuations in the field and can excite various waves in the plasma via electric currents which then become the sources of magnetic fields. We will later in this volume encounter cases when kinetic effects in parallel and quasi-parallel shocks play an important role in their physics and are well capable of generating tangential fields at least on scales shorter than the ion scale.

However, even in MHD one stumbles across the interesting fact that this kind of shocks must have peculiar properties. The reason is that they are not, as in gasdynamics, the result of steepened *sound* waves, in which case they would simply be purely electrostatic shocks. At the contrary, the waves propagating *parallel* to the magnetic field are *Alfvén* and *magnetosonic* waves. Alfvén waves contain transverse magnetic field components. These transverse wave fields, in a parallel shock, are in fact tangential to the shock. Hence, if a purely parallel shock steepens, the transverse Alfvén waves do steepen as well, and the shock after the transition from upstream to downstream switches on a tangential magnetic component which originally was not present. Such shocks are called *switch-on* shocks. Similarly one can imagine the case that a tangential component behind the shock is by the same process switched off by an oppositely directed switch-on field, yielding a *switch-off shock*. Both cases are theoretically possible and models have been provided for instance by *Kennel & Edmiston [1988]* for the resistive MHD case, even though there is only little experimental or observational evidence for the existence of such shocks in space.

The problem of whether or not such shocks exist in MHD is related to the question whether or not an Alfvén wave steepens non-linearly when propagating into a shock. To first order this steepening for an ordinary Alfvén wave is zero. However, to second order

⁸The latter is moved by MHD to processes in the internal structure of the dissipation region of width Δ , and this is determined by processes not covered in MHD which assumes that Δ is not a function of time and space, being defined only locally.

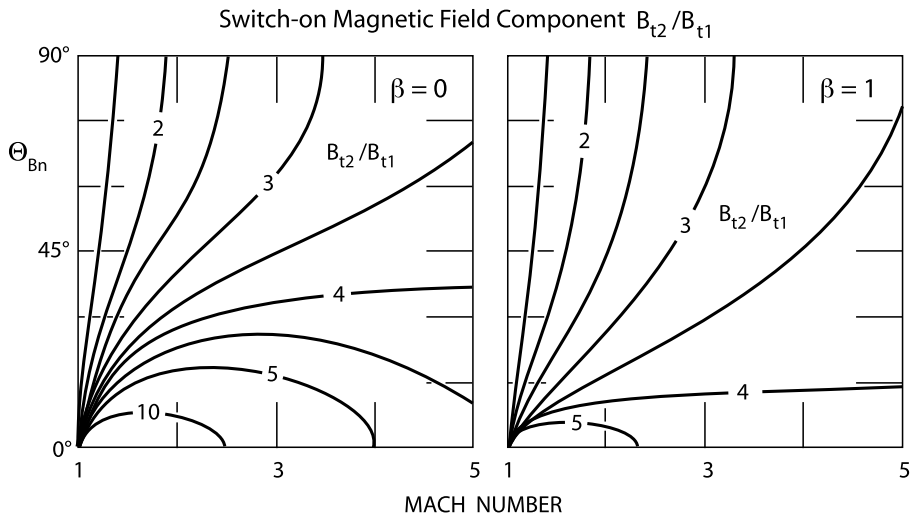


Figure 2.7: The tangential switch-on magnetic field component for an MHD shock as function of Mach number $\mathcal{M}$ and shock normal angle Θ_{Bn} for two extreme upstream plasma β values [after *Farris et al*, 1994, with permission from American Geophysical Union]. The adiabatic index is assumed as $\gamma = 5/3$.

a wave trailing the leading Alfvén wave feels its weak transverse magnetic component. This trailing wave therefore propagates slightly oblique to the main magnetic field and thus causes a second order density compression which in addition to generating a shock-like plasma compression changes the Alfvén velocity locally. In the case when the trailing wave is polarised in the same direction as the leading wave it also increases the transverse magnetic field component downstream of the compression thereby to second order switching on a tangential magnetic component. A whole train of trailing waves of same polarisation will thus cause strong steepening in both the density and tangential magnetic field.

Clearly, this kind of shocks is a more or less exotic case of MHD shocks whose importance is not precisely known [the rare case of observation of a switch-on shock has been reported by *Farris et al*, 1994]. Figure 2.7 shows the theoretical dependence of the downstream tangential magnetic field component as function of Mach number $\mathcal{M}$ and shock normal angle Θ_{Bn} for two different upstream plasma β . The downstream tangential field is given in relation to the upstream tangential magnetic field amplitude and arises here due to the breaking of the magnetic field when crossing the shock. The effect is seen to be largest for quasi-parallel shocks and increasing Mach numbers.

There are other efficient mechanisms of magnetic field generation in shock waves based on the famous Weibel instability which come into play when the shocks have high Mach numbers and the plasma is anisotropic. These mechanisms are particularly strong when relativistic effects must be taken into account which is the case under most astrophysical conditions. These effects are not anymore simple MHD but appear naturally in the kinetic treatment of shocks.

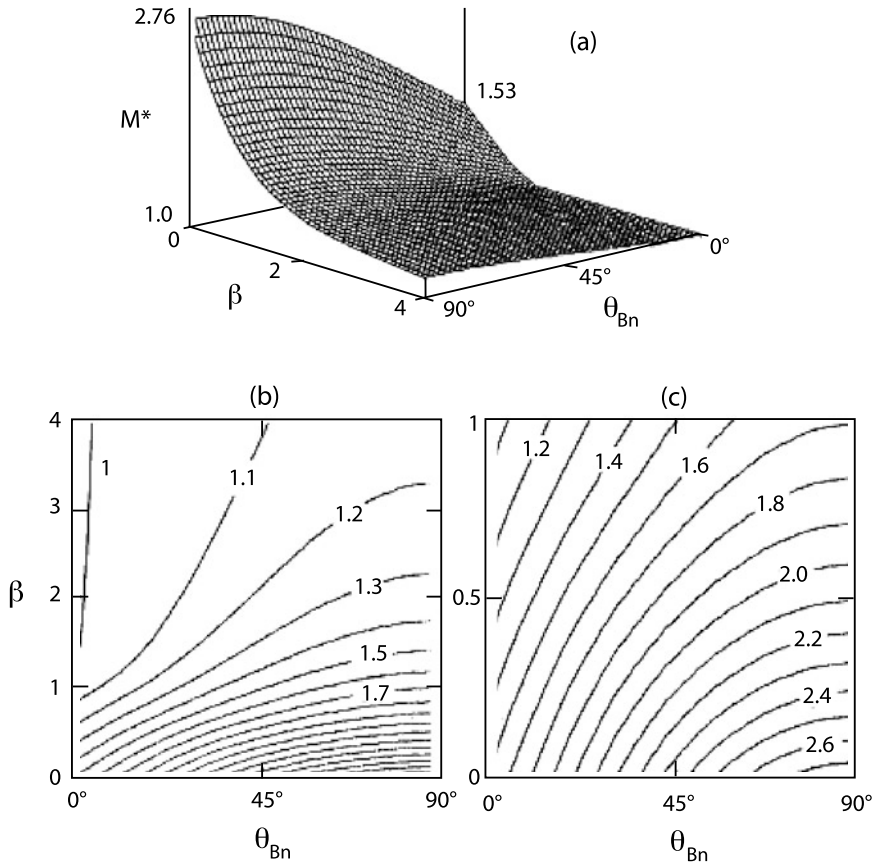


Figure 2.8: Parametric dependence of $\mathcal{M}_c$ for a fast shock on the upstream plasma- β_1 (top) and shock angle Θ_{Bn} for the special case of adiabatic $\gamma = \frac{5}{3}$. The lower panels show two β_1 ranges of critical Mach number contours [from Edmiston & Kennel, 1984]. For large $\beta_1 > 1$ the critical Mach number is close to $\mathcal{M}_c \sim 1$, while for smaller β_1 it is a strong function of Θ_{Bn} having its lowest value $\mathcal{M}_c = 1.53$ for parallel and $\mathcal{M}_c = 2.76$ for perpendicular shocks. The latter value is the same as that originally inferred already by Marshall [1955].

2.4 Criticality

In this final section of this preparatory chapter we single out a most important property of shocks which leads to another *physically* justified classification of collisionless shocks into *subcritical* and *supercritical* shocks, according to their Mach-numbers $\mathcal{M} < \mathcal{M}_c$ being smaller or $\mathcal{M} > \mathcal{M}_c$ larger than some critical Mach-number $\mathcal{M}_c$. We already mentioned that at high Mach-numbers the existence of a critical Mach number for collisionless shocks was predicted long ago from consideration of the insufficiency of dissipation in the shock to provide fast enough retardation of the inflow, plasma thermalisation, and entropy pro-

duction. For a resistive shock *Marshall* [1955] had numerically determined the critical Mach number to $\mathcal{M}_c \approx 2.76$.

Subcritical shocks are capable of generating sufficient dissipation to account for retardation, thermalisation and entropy in the time the flow crosses the shock from upstream to downstream. The relevant processes are based on wave-particle interaction between the shocked plasma and the shock-excited turbulent wave fields. These processes will be discussed in Chapter 4.

For supercritical shocks this is, however, not the case. Supercritical shocks must evoke mechanisms different from simple wave-particle interaction for getting rid of the excess energy in the bulk flow that cannot be dissipated by any classical anomalous dissipation. Above the critical Mach number the simplest efficient way of energy dissipation is rejection of the in-flowing excess energy from the shock by reflecting a substantial part of the incoming plasma back upstream. The physical processes involved into the reflection process and its effects on the structure of the shock will be discussed below in separate Chapters 5 and 6 for both quasi-perpendicular and quasi-parallel shocks.

Independent of this, the determination of the critical Mach number poses an interesting problem. This was formulated by *Kantrowitz & Petschek* [1966] who realised that the magnetic width of the shock must always exceed the magnetic Reynolds length for dissipative shocks, and that the finite magnetic field compression ratio, therefore, sets an upper limit to the rate of resistive dissipation that is possible in an MHD shock. Plasmas possess several dissipative lengths, depending on which dissipative process is considered. Any nonlinear wave that propagates in the plasma should steepen as long, until its transverse scale approaches the longest of these dissipative scales. Then dissipation sets on and limits its amplitude [*Coroniti*, 1970].

Thus, when the wavelength of the fast magnetosonic wave approaches the resistive length, the magnetic field decouples from the wave by resistive dissipation, and the wave speed becomes the sound speed downstream of the shock ramp. The condition for the critical Mach number is then given by $V_{2n} = c_{2s}$. Similarly, for the slow-mode shock, because of its different dispersive properties, the resistive critical-Mach number is defined by the condition $V_{1n} = c_{1s}$. Since these quantities depend on wave angle, they have to be solved numerically. This was done by *Edmiston & Kennel* [1984]. The critical fast-mode Mach number varies between 1 and 3, depending on the upstream plasma parameters and flow-angle to the magnetic field. It is usually called *first* critical Mach number, because there is theoretical evidence in simulations for a second critical Mach number, which comes into play when the shock structure becomes time dependent [*Krasnoselskikh et al*, 2002], whistlers accumulate at the shock front and periodically cause its reformation. The dominant dispersion is then the whistler dispersion. An approximate expression for this *second* or *whistler* critical Mach number is

$$\mathcal{M}_{2c} \propto \left(\frac{m_i}{m_e} \right)^{\frac{1}{2}} \cos \Theta_{Bn} \quad (2.13)$$

where the constant of proportionality depends on whether one defines the Mach number with respect to the whistler-phase or group velocities. For the former it is $\frac{1}{2}$, and for the latter $\sqrt{27/64}$ [*Oka et al*, 2006]. These authors have used GEOTAIL data to confirm the

existence of this whistler-critical Mach number that separates the regions of sub-critical and super-critical shocks in bulk flow velocity V_{1n} -magnetic angle Θ_{Bn} space.

It is clear that it is the smallest critical Mach number that determines the behaviour of the shock. In simple words: $\mathcal{M} > 1$ is responsible for the existence of the shock under the condition that an obstacle exists in the flow, which is disturbed in some way such that fast waves can grow, steepen and form shocks. When, in addition, the flow exceeds the next lowest Mach number for a given Θ_{Bn} the shock at this angle will make the transition into a supercritical shock and under additional conditions, which have not yet be ultimately clarified, will start reflecting particles back upstream. If, because of some reason, this would not happen, the flow might have to exceed the next higher critical Mach number until reflection becomes possible. In such a case the shock would become metastable in the region where the Mach number becomes supercritical, will steepen and shrink in width until other effects and – ultimately – reflection of particles can set on.

2.5 Remarks

In this chapter we have not yet dug into the theory and observation of collisionless shocks. Instead we have tried to provide the basic philosophy of shock formation. Recently *Balogh & Riley [2005]* gave a beautiful and comprehensive account of the complexity involved into the shock problem that is illustrated in Figure 2.9 and which we like to cite at this occasion because we cannot say it in any better way:

“... the concept [of shocks] originally inherited from collision-dominated media has been successfully extrapolated to collisionless plasmas... both qualitatively and quantitatively. There is nevertheless an inherent contradiction at the heart of the concept of collisionless shock waves. On the one hand, a discontinuous transition between two states of the plasma is assumed, represented by nominally well-defined parameters which characterise unambiguously the upstream and downstream states, and which are linked by continuity equations, augmented as necessary by additional equations. On the other hand, the necessary dissipation in the shock transition itself generates phenomena which can propagate information, primarily in the form of particle populations which have interacted with the shock, away from the shock front, upstream as well as downstream in general, and which, in turn, modify in a complex manner the states of the plasma and the temporal and spatial functions of its constituents... As a result, the collisionless shock possibly never propagates into a region which has not already been modified, to a greater or lesser extent, by phenomena which originate at the shock and which therefore affect, through a set of complex feedback processes, the shock transition itself. This picture is fully born out by numerically simulated shocks, in which the initial conditions, or ‘states’ of the upstream and downstream plasma are fully defined at the outset; but once the shock is generated, the characteristic parameters of the medium evolve as a result of particles propagating away from the shock, generating, through a set of (usually nonlinear) wave particle interactions and a range of instabilities, significantly more complex ‘states’ than was assumed at the start...”

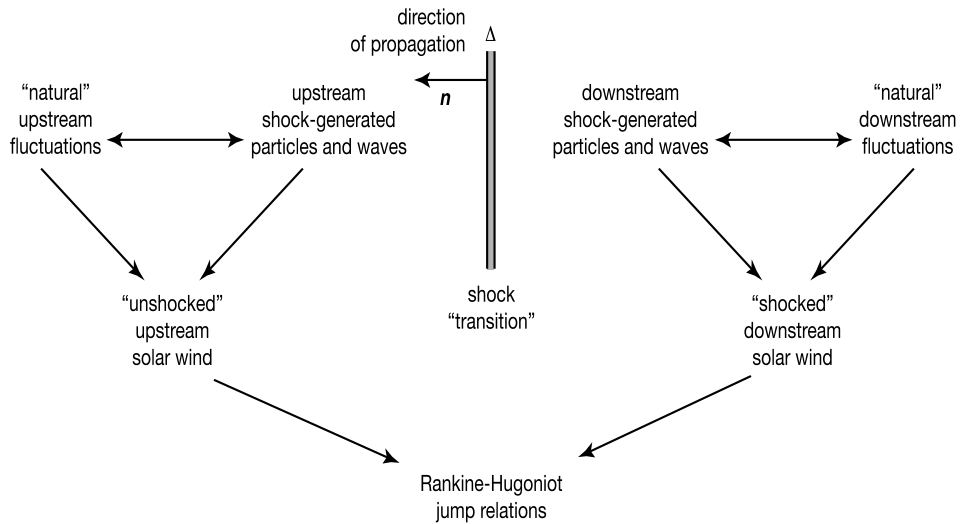


Figure 2.9: A schematic representation of the complexity of processes accompanying a collisionless supercritical shock and the relationship between the various states upstream and downstream of the shock transition [after *Balogh & Riley, 2005*, courtesy A. Balogh, ISSI]. There is a noticeable symmetry between the upstream and downstream shock regions, even though the effects and physical properties of both regions are very different belonging to two mutually coupled but different plasma states.

Clearly, the picture developed in this text by *Balogh & Riley [2005]* is that of a supercritical shock, it already refers to the complicated nonlinear set of equations that determine the behaviour of shocks, and it also anticipates much of what will follow in the coming chapters, the particle reflection process from supercritical shocks, the effects those reflected particles have on their surroundings, transportation of information, mass, energy, and fields, excitation of instabilities, waves, and the entire zoo of nonlinear interactions, convection and reaction on the shock which, as we have already clarified, is an entity that is not in thermal equilibrium and therefore evolves and reforms continuously trying to reach thermal equilibrium but being hindered to ultimately establish it by the supercritical inflow. It is no surprise that shocks not only constitute an extraordinarily interesting but also a challenging phenomenon in particular when the plasma is collisionless.

We have in this chapter listed most of the terms and classifications of shocks and the basic terms which are used in their description. We encountered several ways of classification of collisionless shocks which we will follow in the later chapters of this text. Their properties, however, must be described by models which depend on the assumption of which scales are considered to be important for the processes to be investigated. This will be done in the next Chapter 3 where a number of models are described starting from gasdynamic models and proceeding with increasing sophistication to kinetic and particle models of collisionless shocks.

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Basic Equations and Models

Abstract. This chapter presents the basic theoretical background of collisionless shock physics. It states the basic process of shock formation as the growth of a small disturbance in the plasma by the action of the intrinsic nonlinearity of flow, independent of the cause of the initial disturbance. The latter can be an external driver like a piston or a blast, it can also be an internal instability. Shocks form when nonlinearity causes steeping (or steepening) of the disturbance in space and some process exists which prevents breaking of the steep wave. Such processes are of dissipative or dispersive nature and are discussed in ascending importance. An intermediate step is the evolution of solitary waves based on the Sagdeev pseudo-potential. After this fundamental discussion, the plasma kinetic equations are given and the Rankine-Hugoniot jump conditions at shocks are derived with the shock solutions explicitly given. Critical Mach numbers are defined beyond which dissipation is unable to prevent wave breaking. The relevant wave instabilities causing initial disturbances, dispersion and dissipation are discussed at length. Transport ratios are given, and anomalous transport is reviewed. Finally, shock particle reflection is identified as the basic process of shock stabilisation preventing breaking. The last section provides a cursory and incomplete briefing on numerical simulation techniques.

3.1 Wave Steeping

Shocks have a certain width Δ and a certain jump in density N , temperature T , pressure P and magnetic field B across this width from a given upstream value to a downstream value. This jump is by no means infinitesimal. At the contrary it is usually several times the upstream value in magnitude. Thus, looked at as a wave, a shock is a highly nonlinear wave structure of wavelength Δ with amplitude that cannot be neglected compared with the upstream value. Therefore, the basic equations describing a shock cannot be linearised, as is usually done in considering wave phenomena. These equation must be solved in their full nonlinearities. This, however, is barely possible and can be done analytically only in very rare cases, which usually are not of interest. On the other hand, shocks evolve inside the plasma from small disturbances. It is thus reasonable to ask for the nonlinear evolution of such a small harmonic disturbance in order to learn, how a disturbance can evolve into a very large amplitude shock ramp.

3.1.1 Simple Waves: Steeping and Breaking

The simplest way to do this, is to consider the evolution of so-called simple waves [see, e.g., *Witham, 1966*]. Simple waves are one-dimensional sinusoidal disturbances of the

plasma velocity of the form $V(x) = A \sin kx$ moving on the plasma background at speed c . The total derivative of this disturbance is given by

$$\frac{dV}{dt} \equiv \frac{\partial V}{\partial t} + V \frac{\partial V}{\partial x} = 0 \quad (3.1)$$

and in the absence of any forces and friction is assumed to vanish. This has the meaning that, sitting on the crest of the disturbance and moving with it, one does not see any variation. Hence this is the equation which describes the evolution of the disturbance as long as no friction or other force comes into play which in the initial state is a reasonable assumption.

Let us now investigate, how such a disturbance will evolve when it propagates through the plasma. Clearly, during propagation the main effect on the shape of the disturbance arises from the second nonlinear term, which can be written as $Vk \cos kx$. Inserting for V this becomes $\frac{1}{2} \sin 2kx$. Hence, the nonlinear term in the above Eq. (3.1) generates harmonic sidebands of half wavelength and half the amplitude of the original wave. These waves, by the same mechanism, also generate sidebands on their own, now at quarter original wavelength and amplitude, and so on with increasingly shorter wavelengths. The total amplitude is the superposition of all these sideband harmonics

$$V_k(x) = \sum_l \frac{A}{2^l} \sin 2^l kx, \quad (l = 0, 1, 2, \dots) \quad (3.2)$$

All these waves and sidebands propagate at the same velocity c . This can be easily seen when, for propagating waves, replacing $kx \rightarrow k(x - ct)$ in the above expression. They locally superpose and add to the wave amplitude. Because ever shorter wavelengths contribute, the wave steepens until the gradients become so steep that other processes take over. If this does not happen, the wave will turn over and break.

This is illustrated in Figure 3.1 in the co-moving frame of a sinusoidal wave. The lower left part of the figure shows the nonlinear mechanism. Since the velocity is largest at the maxima it speeds up the motion of the maxima with respect to the remaining parts of the wave profile. Moreover, the actions on the positive and negative maxima are oppositely directed, and the wave starts forming a ramp corresponding to a shock front. This happens at time t_b . For times $t > t_b$ the wave will turn over and collapse. This can be prevented only by additional processes which set on when the wavelength of the ramp becomes so short that in Eq. (3.1) terms of higher-order gradients in the velocity must be taken into account.

Equation (3.1) can in fact be understood as the lowest order equation describing the evolution of a wave packet of wave number k . In general, in the wave frame of reference its right-hand side is a function $F(V)$ that can be expanded with respect to V . The first higher-order term in this expansion turns out to be of second order in the spatial derivative $\nabla = \hat{\mathbf{x}} \partial_x$, where $\hat{\mathbf{x}}$ is the unit vector in the direction of x . The next higher-order term is third order in ∇ , and so on. Up to third order the resulting equation then reads

$$\frac{\partial V}{\partial t} + V \frac{\partial V}{\partial x} = \frac{\partial}{\partial x} D \frac{\partial V}{\partial x} - \beta \frac{\partial^3 V}{\partial x^3} + \dots \quad (3.3)$$

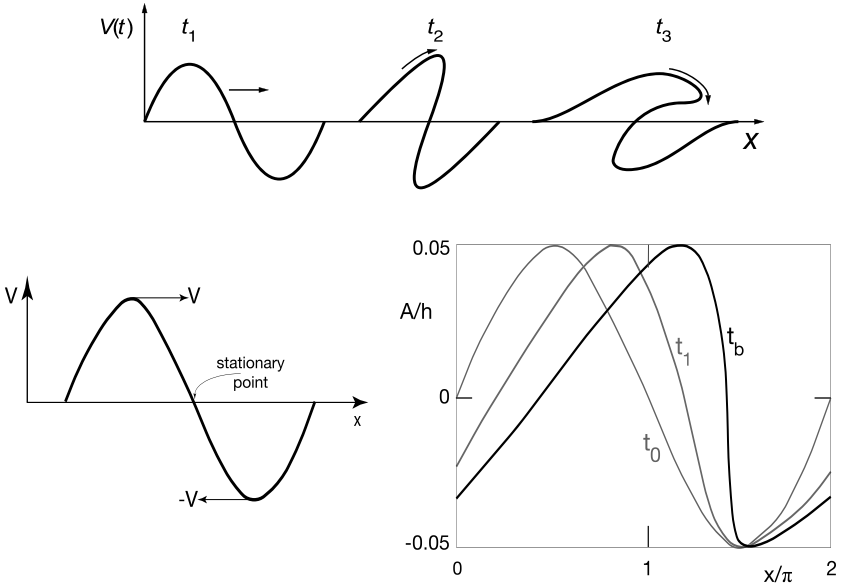


Figure 3.1: *Top:* Schematic of the steepening and wave breaking phenomenon, illustrated for three successive times t_1 , t_2 , t_3 . At t_2 the wave has steepened to maximum, and in t_3 it collapses in the absence of any retarding effects. *Bottom:* steepening and final breaking of an initial sinusoidal (simple) wave. *Bottom left:* Initial wave form in wave frame showing the nonlinear action of the wave on its own shape. *Bottom right:* Calculation of wave form steepening in a shallow fluid of depth h [after Zahibo *et al.*, 2007]. The wave profile is shown at the initial time t_0 , intermediate time t_1 , and breaking time t_b when the wave starts turning over. Steepening of the profile is well expressed.

The first term on the right is a diffusive term with diffusion coefficient $D(x)$. The third term with arbitrary coefficient β is the lowest-order contribution of wave dispersion to the evolution of the wave shape and amplitude. This can be most easily seen when taking the linearised equation, assuming V to be a small disturbance only of the wave speed c , neglecting the nonlinearity in the second term on the left by approximating it with $c\nabla V$, and subsequently Fourier analysing for a harmonic perturbation $V = A \exp i(kx - \omega t)$ with wave number k and frequency ω . This procedure yields (for constant D) the following *dispersion relation*

$$\omega - kc + k^3 \beta = -ik^2 D \quad (3.4)$$

On the left of this equation there is the relation between the frequency ω and wave number k , while on the right appears an imaginary term that depends on the diffusion coefficient D and the square of the wave number k . Imaginary terms in frequency imply damping. Hence, as we noted above, the second-order spatial derivative term in Eq. (3.3) corresponds to diffusive dissipation of flow energy, while (for real β) the third term in Eq. (3.3) causes the wave to disperse, i.e. waves of different wave-numbers, respectively different wavelengths, propagate at different phase velocities.

3.1.2 Burgers' Dissipative Shock Solution

Returning to the nonlinear equation (3.3), we can thus conclude from the above analysis that a wave, if not breaking, will steepen so long until either the diffusive or the dispersive terms on the right start competing with the nonlinearity, in which case the wave may assume a stationary shape that is balanced either by diffusion or by dispersion. Diffusion implies balance by real dissipation with energy transformed into heat, while dispersion implies that the 'dangerous' short wavelength waves, which cause the steepening, either run the wave out or do not catch up with the wave when the slope of the wave profile exceeds a certain steepness. In the dispersive case the wave will either exhibit short wavelength fluctuations in front of the steepened profile or behind it, depending on whether the shorter wavelength sidebands are retarded or accelerated. However, we can also conclude that dispersion alone should be unable to generate a shock since, in the simple form discussed here, it does not produce irreversible dissipation and hence no heating and increase in entropy. For a shock profile to be created, some kind of diffusive process will be necessary.

Stationary Burgers Equation

Equation (3.3) allows to distinguish two extreme cases. The first case is that of purely diffusive compensation of the nonlinear steepening. In this case the dispersive term can be neglected, and one obtains the Burgers equation

$$\frac{\partial V}{\partial t} + V \frac{\partial V}{\partial x} = D \frac{\partial^2 V}{\partial x^2} \quad (3.5)$$

which is a non-linear diffusion (or heat conduction) equation. In contrast to the ordinary linear heat conduction equation, the Burgers equation possesses stationary solutions due to the above mentioned compensation of diffusive spread by nonlinear steepening. These stationary solutions can be found when transforming to a coordinate system moving with the wave by introducing the new coordinate $y = x - ct$. Then Burgers' equation becomes

$$D \frac{\partial^2 V}{\partial y^2} = (V - c) \frac{\partial V}{\partial y} \quad (3.6)$$

We are interested only in solution which are regular at infinity with vanishing derivatives. Introducing the variable $V' = V - c$ the first integral is easily obtained. Integrating a second time the solution found is then

$$\frac{V}{c} = 1 - \tanh\left(\frac{x - ct}{2D/c}\right) \quad (3.7)$$

The form of this solution is a typical shock ramp which is displayed in Figure 3.2. The ramp is sitting on the wave velocity c . Its width is $\Delta = 2D/c$. The shock solutions produced by Burgers' equation are thus *propagating non-oscillatory* shocks; they are simple stationary constant amplitude ramps of the sort of tsunamis. We should, however, keep in mind that they are produced solely by nonlinear steepening and its compensation through diffusion. When the latter is large, the shock will be steep; in the opposite case it will be a flat ramp only, and its relative height is a function of its width.

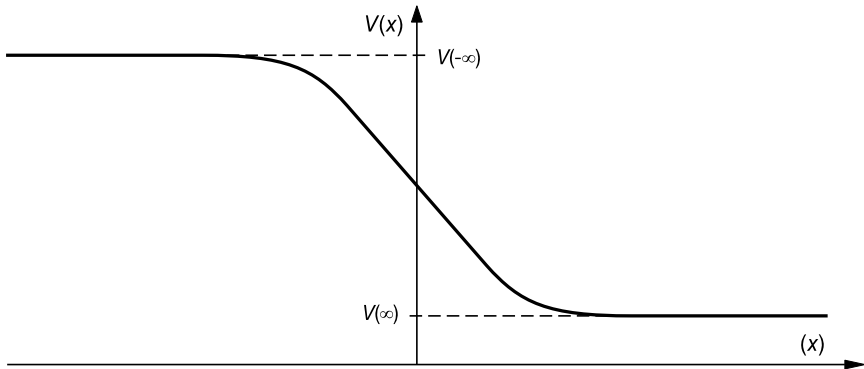


Figure 3.2: The stationary solution of the Burgers equation is a smooth shock ramp of width $\Delta = 2D/c$, depending on shock velocity c and diffusion coefficient D .

Time-Dependence

Since Burgers' equation is an ordinary diffusion equation it can be solved by the usual methods of treating heat conduction evolving from some initial state. Such an investigation is necessary in order to justify that the stationary state of the shock-ramp Burgers solution can indeed be reached by evolution out of the initial state V_0 . To this end one transforms the time-dependent Burgers equation (3.5) through introduction of a new variable ϕ via $V = -2D(\partial \ln \phi / \partial y)$ into the common form of a diffusion equation $\partial \phi / \partial t = D \partial^2 \phi / \partial y^2$, which has the commonly known solution

$$\phi(y, t) = \frac{1}{(4\pi Dt)^{\frac{1}{2}}} \int_{-\infty}^{\infty} d\eta \exp \left[-\frac{(y - \eta)^2}{4Dt} - \frac{1}{2D} \int_0^{\eta} V_0(\tau) d\tau \right] \quad (3.8)$$

The initial disturbance satisfies the condition of convergence $\int_0^y dy' V_0(y') \leq \text{const} \cdot y$ for $y \rightarrow \infty$, which yields the requirement that $\int_{-\infty}^{\infty} dy' V_0(y') = \Theta < \infty$ as well as the time-asymptotic solution

$$V(y, t \rightarrow \infty) \simeq -2D \frac{d}{dy} \ln G \left[\frac{y}{(4Dt)^{\frac{1}{2}}} \right] \quad (3.9)$$

The dummy function $G(x)$ is a transformed version of the function ϕ that is given by $\pi^{\frac{1}{2}} G(x) = e^{-\Theta/4D} \int_{-\infty}^x d\eta e^{-\eta^2} + e^{\Theta/4D} \int_x^{\infty} d\eta e^{-\eta^2}$. Figure 3.3 shows this asymptotic profile of Burgers shock-ramp solution. The characteristic shape of this solution contains a smooth wavelike increase up to a flat plateau, followed by the shock ramp and a smooth transition to the undisturbed state. The ramp is moving to the right in the direction of the original wave propagation. This is seen from the time dependence of the crest of the ramp. Clearly, most shock transitions in space do not exhibit this smooth rise inside the downstream region of the shock, indicating that the Burgers solution has pure model character which does not really confirm with the plasma reality.

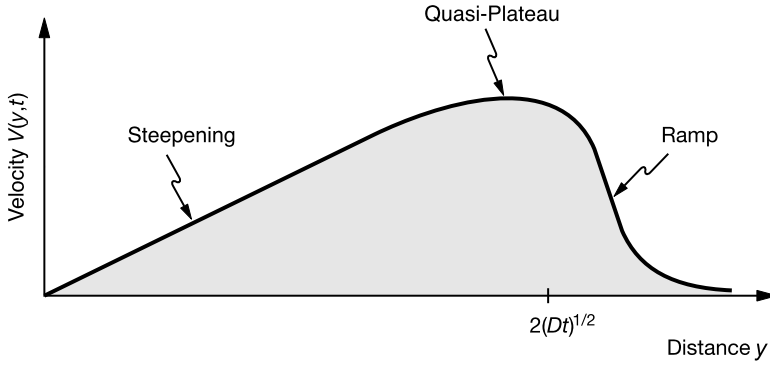


Figure 3.3: The time-asymptotic Burgers' shock solution Eq. (3.9) which evolves from the initial disturbance V_0 through steepening and dissipative ramp formation after a given diffusive time at a location which is determined by the time t and diffusion coefficient D . The characteristic shape of this solution contains a smooth wavelike increase up to a flat plateau followed by the shock ramp and a smooth transition to the undisturbed state. The ramp is moving to the right in the direction of original wave propagation. This is seen from the time dependence of the crest of the ramp. Clearly, most shock transitions in space do not exhibit this smooth rise inside the downstream region of the shock, indicating that the Burgers solution has pure model character which does not really confirm with the plasma reality.

3.1.3 Korteweg-de Vries Dispersion Effects

Balancing the nonlinearity with the help of dissipation is one possibility. The remaining possibility which in the absence of dissipation becomes the dominant, is balancing nonlinearity with dispersion. In this case we can neglect the diffusion term in Eq. (3.3) to obtain the so-called Korteweg-de Vries equation¹

$$\frac{\partial V}{\partial t} + V \frac{\partial V}{\partial x} + \beta \frac{\partial^3 V}{\partial x^3} = 0 \quad (3.10)$$

Similar to Burgers' equation, the Korteweg-de Vries equation also allows for stationary localised solutions. Such solutions are restricted to a finite spatial interval because dispersion does not cause irreversible effects. As before we assume that the stationary solution moves at speed c , and we introduce the co-moving coordinate $y = x - ct$, this time measured from the centre of the localised disturbance. The Korteweg-de Vries equation then transforms into the third order ordinary differential equation

$$(V - c) \frac{\partial V}{\partial y} + \beta \frac{\partial^3 V}{\partial y^3} = 0 \quad (3.11)$$

¹It might be of interest to note that to find a method that solves this general time-dependent equation analytically took more than a century, and for this purpose it was necessary to develop the whole apparatus of non-relativistic quantum mechanics. This was done in a seminal paper by *Gardner et al* [1967].

Solution of this stationary equation requires the prescription of boundary conditions at $y \rightarrow \pm\infty$ for which, because in the non-dissipative case no ramp can be formed, we choose $V = \partial V / \partial y = 0$. It can be shown by substitution that the solution is the function

$$V_{KdV}(x - ct) = 3c \operatorname{sech}^2 \left[\sqrt{\frac{c}{\beta}} \frac{(x - ct)}{2} \right] \quad (3.12)$$

This function is a so-called *soliton* solution; it describes a stationary bell-shaped solitary wave pulse propagating at velocity c along x without any change of form. The width of this pulse is $\Delta = 2\sqrt{\beta/c}$ and depends on the velocity c and the dispersion parameter β in such a way that the faster the pulse moves, the narrower the pulse becomes. In addition there is a distinct relation between the amplitude A of the pulse and its width

$$\Delta = 2(3\beta/A)^{\frac{1}{2}} \quad (3.13)$$

from where it follows that large amplitude Korteweg-de Vries solitons are fast and narrow. Like in the case of the Burgers equation, the Korteweg-de Vries equation is a model equation which results from the dispersive properties of nonlinear waves. However, it is interesting that it can be derived for real problems arising in plasma wave propagation, and several variants of it have in the past been applied to the plasma. Hence, in describing the nonlinear evolution of plasmas it is a more realistic model than Burgers' equation which is simply a consequence of the strongly dispersive and practically dissipation-free properties of plasmas.

The stationary Korteweg-de Vries equation can have a whole chain of such solitons with the solitons having completely different amplitudes and widths. These solitons have the interesting property that they can pass through each other during mutual encounters without having any effect on their widths and amplitudes; only the phases and spatial positions of the waves from which the solitons form will change during the collision. The question of how these chains of solitons are produced is a question that can be answered only when solving the time-dependent Korteweg-de Vries problem imposing a certain initial condition similar to that imposed above on the time-dependent Burgers equation. In the case of the time-dependent Korteweg-de Vries equation the solution cannot be found in such a simple way, however. Solving it rather constitutes a major mathematical problem which requires solving an equivalent Schrödinger equation [Gardner *et al*, 1967]. It turns out that the soliton amplitudes in the chain which solves the Korteweg-de Vries equation are related in some manner by an infinite set of invariants of the Korteweg-de Vries equation. In addition, the time-dependent Korteweg-de Vries equation also supports wave trails which accompany the solitons forming an oscillatory (turbulent) background of spatially dependent amplitudes on which the solitons propagate.

Clearly, these soliton chains are no shocks; they are wave pulses which after some steeping time evolve into stationarity and are completely reversible practically not leaving any effect on the plasma if one neglects the microscopic processes which take place in the plasma. But this is precisely the door where the speculation comes in about such solitons in collisionless plasma being the initial state of the formation of collisionless shocks. Because, if one can manage a soliton in the chain to move so fast that its width becomes

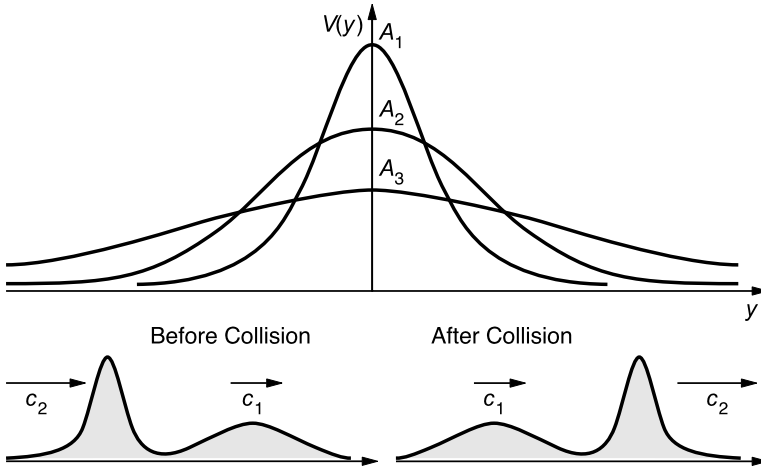


Figure 3.4: *Top:* Shape of Korteweg-de Vries solitons for different widths and amplitudes (labelled A_1, A_2, A_3). Solitons of large amplitude are narrower than those with smaller amplitude. *Below:* In the interaction between two Korteweg-de Vries solitons of different speeds only the position changes while the collision has no effect on amplitude or shape.

comparable to the intrinsic plasma scales, then the wave field of the soliton should distort the microscopic particle motion causing some kind of dissipation which necessarily will turn the soliton into a shock wave by generating entropy and producing a difference between the states upstream and downstream of the soliton. The soliton in this case borrows from Burgers shock solutions, and mathematically the Korteweg-de Vries equation regains the lost dissipative Burgers second-order term becoming a Korteweg-de Vries-Burgers equation. This was, actually, the point *Sagdeev* [1966] made intrinsically in his famous theory of shock formation in collisionless plasma.

3.1.4 Sagdeev's Pseudo-potential

The Korteweg-de Vries equation is the ideal candidate for introducing one particular notion that has become immensely important in soliton and shock research, the so-called *Sagdeev potential*. The Sagdeev potential is a pseudo-potential introduced in order to solve a certain class of nonlinear partial differential equations and to distinguish between solitary wave and shock solutions of these equations. This method takes advantage of the similarity of the first integral of the particular class of equations to the equation of motion of a hypothetical particle in classical mechanics. Knowledge of the Sagdeev potential then reduces the problem of solution to the mere discussion of the behaviour of a particle in the pseudo-potential well.

The stationary Korteweg-de Vries equation (3.11) can be directly integrated once. Applying the boundary conditions at infinity, the integration constant in the first integration becomes zero yielding the nonlinear second-order differential equation

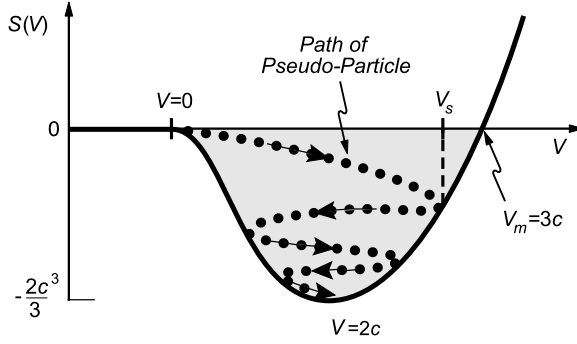


Figure 3.5: A sketch of the Korteweg-de Vries Sagdeev pseudo-potential. Solutions exist only in the region of $S(V) \leq 0$. The maximum soliton amplitude is just $V = 3c$. The minimum potential is at $V = 2c$. The *dotted line* shows the path of a pseudo-particle for shock formation in presence of dissipation when it “steps down” the potential well to approach the “minimum energy” state.

$$\beta \frac{\partial^2 V}{\partial y^2} = V \left(c - \frac{1}{2} V \right) \quad (3.14)$$

The similarity of this equation to Newton’s equation of motion of a pseudo-particle of mass β in the force field given on the right-hand side is quite obvious. Here V is the pseudo-spatial coordinate, and y is the pseudo-time. This equation can be solved by multiplying it with the pseudo-velocity $\partial V / \partial y$, after which the right-hand side can be represented as the derivative of the Sagdeev pseudo-potential, $S(V)$, which in this case is a function of the (real) velocity V . It becomes explicit in the pseudo-energy conservation law

$$\frac{\beta}{2} \left(\frac{\partial V}{\partial y} \right)^2 = \frac{V^2}{2} \left(c - \frac{1}{3} V \right) \equiv -S(V) \quad (3.15)$$

after having integrated a second time and again applied the vanishing boundary conditions at infinity. Because the left-hand side of this expression is a positive quantity, solutions exist only under the condition that the Sagdeev pseudo-potential is attractive, i.e. is negative

$$S(V) = \frac{V^2}{2} \left(\frac{V}{3} - c \right) < 0 \quad (3.16)$$

and, as we already know, solutions can exist only in the region of velocity space where $V < 3c$. Figure 3.5 shows a sketch of the Korteweg-de Vries Sagdeev potential. It vanishes at $V = 0$ and $V = 3c$ and has its minimum of $S = -3c^3/2$ at $V = 2c$. In terms of energy states a pseudo-particle (soliton) can assume any of the energy levels inside the negative portion of the Sagdeev potential. The soliton with maximum amplitude $V = 3c$ is at the “highest” level $S = 0$. But there can be many solitons at this level with amplitudes between 0 and $3c$. The “most stable” soliton at the “ground state” has minimum Sagdeev potential and amplitude $V = 2c$, and there is only one soliton with such an amplitude. In the absence of dissipation all these solitons are stable.

The actual solutions (3.12) of the Korteweg-de Vries equation can then be found from $\partial V / \partial y = \sqrt{-2S(V)/\beta}$ by simple quadrature, solving the integral

$$y - y_0 = \int_0^V \frac{dV}{[-2S(V)/\beta]^{\frac{1}{2}}} \quad (3.17)$$

We do not further discuss the stationary solution of the Korteweg-de Vries equation; it is but one example which can be solved by the Sagdeev potential method. In the literature it has been demonstrated that a very large number of other nonlinear problems in plasma related to solitary waves can be treated in the same way, sometimes under much more complicated conditions and leading to different types of solitary solutions.

All these solutions are, however, dissipation free and do not directly lead to shock wave solutions. As *Sagdeev* [1966] pointed out, they will turn into shock solutions whenever anomalous processes at short wavelengths cause the appearance of some kind of anomalous dissipation under the ideal conditions of non-collisionality.

This claim is a most important insight that can, however, be based only on the kinetic theory of the microscopic interaction between waves and particles and waves and waves in plasma far from thermal equilibrium, the so-called *collective processes* which dominate the behaviour of high temperature plasmas in which shock waves are at home. The dot-arrowred line in Figure 3.5 shows the presumable “path” of such a dissipative soliton in the Sagdeev potential “energy” space. The soliton pseudo-particle will in this case step down the potential, possibly in an oscillatory way. In the case when the system is open and energy is continuously supplied it might reach a stationary shock state with shock amplitude V_s or settle at the ultimate minimum of the Sagdeev potential. The dynamics of this depends on the microphysics.

3.2 Basic Equations

Before discussing these processes and their relevance for shock wave formation, we need to briefly introduce the equations which lie at the fundament of all these processes and to discuss their macroscopic consequences. We will, in the present chapter, distinguish between two approaches to the description of shocks, the theoretical and the numerical approaches, respectively. The former deals with the average properties of collective plasma behaviour and the investigation of wave growth from an infinitesimal perturbation up to a large amplitude shock, the latter refers to the dynamics of macro-particles (as has been described in Chapter 2) and is independent of the average equations as it simply solves Newton’s equations of motion of the many macro-particles that constitute the plasma in their self-consistent fields, where the fields are obtained from Maxwell’s equations of electrodynamics. any final shock theory must combine both approaches because the fundamental basic equations cannot be solved analytically, while the numerical approach provides data which cannot be understood without a follow-up theoretical investigation tailored to serve the effects found in the numerical simulation experiments.

3.2.1 Kinetic Plasma Equations

Collisionless shock waves represent the final result of collective interactions in which very many particles and in addition the full electromagnetic fields are involved. It is thus quite reasonable that they cannot be described by test particle theory which considers the motion of non-interacting particles separate from other particles and fields. Test particle theory can however be applied if one is not interested in the formation of shocks but instead in its effect on small numbers of particles. This is used in the shock theory of charged particle acceleration which will be the subject of Chapter 7.

Maxwell-Vlasov Equations

The basic equations on which shock physics is founded are the kinetic equations of a plasma [cf., e.g., *Montgomery & Tidman, 1964; Tidman & Krall, 1971*, and others] or, at the best, some of its simplifications, in addition to the full set of the equations of electrodynamics.

Since collisions can be neglected, and thus the Boltzmann collision term in the kinetic equations is suppressed, these equations reduce to the (non-relativistic) Vlasov-Maxwell set of equations

$$\frac{\partial F^\pm}{\partial t} + \mathbf{v} \cdot \nabla F^\pm + \frac{e_\pm}{m_\pm} (\mathbf{E} + \mathbf{v} \times \mathbf{B}) \cdot \nabla_{\mathbf{v}} F^\pm = 0 \quad (3.18)$$

where $F^\pm(\mathbf{v}, \mathbf{x}, t)$ are the electron and ion phase space distributions, distinguished by the respective $+$ and $-$ signs, which depend on the six-dimensional phase space composed of velocity, $\mathbf{v}$, and real space, $\mathbf{x}$, coordinates. $m_+ \equiv m_i$ and $m_- \equiv m_e$ are the ion and electron masses, respectively; $e_+ = e$ is the ion charge, $e_- = -e$ electron charge, e the elementary charge, and $\mathbf{E}(\mathbf{x}, t)$, $\mathbf{B}(\mathbf{x}, t)$ are the electromagnetic fields which are independent on velocity while being functions of space and time. Finally, $\nabla_{\mathbf{v}} \equiv \partial/\partial \mathbf{v}$ is the velocity gradient operator acting on the phase space distributions. These two Vlasov equations (3.18) are coupled mutually and to the electromagnetic fields through Maxwell's equations

$$\left. \begin{aligned} \nabla \times \mathbf{B} &= \mu_0 \epsilon_0 (\partial \mathbf{E} / \partial t) + \mu_0 \sum_{\pm} e_{\pm} \int d\mathbf{v}^3 F^{\pm} \mathbf{v}, & \nabla \cdot \mathbf{B} &= 0 \\ \nabla \times \mathbf{E} &= -(\partial \mathbf{B} / \partial t), & \nabla \cdot \mathbf{E} &= \epsilon_0^{-1} \sum_{\pm} e_{\pm} \int d\mathbf{v}^3 F^{\pm} \end{aligned} \right\} \quad (3.19)$$

The second term on the right in the first of these equations is the electric current density; the term on the right in the last of these equations is the electric space charge density (divided by the dielectric constant of vacuum, ϵ_0). These equations already account for the coupling of the field to the particles through the definition of the electric current and particle densities as zero and first moments of the one-particle phase space distributions.

Shocks evolve from infinitesimal wave disturbances; one hence considers two different states of the plasma with the physics of both of them contained in the above equations. These two states are first the final average slowly evolving state of the fully developed shock, and second the strongly time-dependent evolution of the infinitesimal disturbance from the thermal level where it starts up to the formation of the shock. In the first state the

shock possesses a distinct shock profile while in the second state one deals with initially infinitesimal fluctuations. When the fluctuation amplitude approaches the shock strength the two different ways of looking at the shock should ideally lead to the same result. According to this distinction one divides all field and plasma quantities, A , into their slowly varying averages, $\langle A \rangle$, and fast fluctuations, δA , superimposed on the averages according to the prescription

$$A = \langle A \rangle + \delta A, \quad \langle \delta A \rangle = 0$$

The second part of this prescription breaks down when the fluctuations become very large or non-symmetric, but it makes life easier to deal with zero averages of small fluctuations as long as the fluctuation amplitude remains to be small. Frequently it is the only way of extracting a solution from the above nonlinear and complex set of equations (3.18)–(3.19).

In what follows we will apply different simplifications to all these basic equations referring to the last conditions of simplification. Only the last section of the present chapter will, finally, deal with the numerical simulation technique which, in fact, will become the most important tool in the investigation of shocks in the remaining three chapters of this first part of the book.

Equations for Averages and Fluctuations

Let us for simplicity temporarily indicate the averages $\langle \dots \rangle$ of the distribution functions and fields by the subscript 0 on the unbraced quantities, and the fluctuations by small letters f , $\mathbf{e}$, $\mathbf{b}$. Then, on applying the above prescription of averaging to the Vlasov equation, we obtain the kinetic equation for the average distribution functions $F_0^\pm(\mathbf{v}, \mathbf{x}, t)$ in the form

$$\frac{\partial F_0^\pm}{\partial t} + \mathbf{v} \cdot \nabla F_0^\pm + \frac{e_\pm}{m_\pm} (\mathbf{E}_0 + \mathbf{v} \times \mathbf{B}_0) \cdot \nabla_{\mathbf{v}} F_0^\pm = -\frac{e_\pm}{m_\pm} \langle (\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} f^\pm \rangle \quad (3.20)$$

Here the average quantities are assumed to vary on much longer spatial and temporal scales than the fluctuation scales such that the condition of averaging $\langle f, \mathbf{e}, \mathbf{b} \rangle = 0$ remains valid. This average Vlasov equation contains a non-vanishing pseudo-collision term on its right which accounts for the effect of the correlations between the fluctuations and particles on the average distribution. In contrast to the Vlasov equation, the Maxwell equations (3.19) retain their form with the sole difference that the full distribution functions $F^\pm$ appearing in the expression for the electric current density in Ampère's law and in the space charge term in Poisson's equation are to be replaced by their average counterparts $F_0^\pm$, yielding

$$\left. \begin{aligned} \nabla \times \mathbf{B}_0 &= \mu_0 \epsilon_0 (\partial \mathbf{E}_0 / \partial t) + \mu_0 \sum_{\pm} e_{\pm} \int d\mathbf{v}^3 F_0^\pm \mathbf{v}, & \nabla \cdot \mathbf{B}_0 &= 0 \\ \nabla \times \mathbf{E}_0 &= -(\partial \mathbf{B}_0 / \partial t), & \nabla \cdot \mathbf{E}_0 &= \epsilon_0^{-1} \sum_{\pm} e_{\pm} \int d\mathbf{v}^3 F_0^\pm \end{aligned} \right\} \quad (3.21)$$

In order to obtain equations for the fluctuations one subtracts the set of averaged equations from the full set of equations and orders the terms for the fluctuation quantities $f = F - F_0$, $\mathbf{e} = \mathbf{E} - \mathbf{E}_0$, $\mathbf{b} = \mathbf{B} - \mathbf{B}_0$. This procedure leaves the Maxwell equations unchanged when all

quantities appearing in them are replaced by the fluctuating quantities, and the fluctuating Vlasov equation becomes

$$\begin{aligned} \frac{\partial f^\pm}{\partial t} + \mathbf{v} \cdot \nabla f^\pm + \frac{e_\pm}{m_\pm} (\mathbf{E}_0 + \mathbf{v} \times \mathbf{B}_0) \cdot \nabla_{\mathbf{v}} f^\pm = & - \frac{e_\pm}{m_\pm} (\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} F_0^\pm \\ & - \frac{e_\pm}{m_\pm} (\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} f^\pm + \frac{e_\pm}{m_\pm} \langle (\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} f^\pm \rangle \end{aligned} \quad (3.22)$$

Up to this stage the fluctuations are allowed to have arbitrarily large amplitudes; it is only their scales which must be much shorter than the scales of the average field quantities. This means for instance that the width of the shock transition regions should be much larger than the wavelengths of the fluctuations.

The last equation is in fact the equation that describes the evolution of fluctuations. however the coupling to the average quantities is still so strong in this equation that it can be solved only together with the average equation. In particular the average “collision term” appearing on its right provides the greatest complications. It will therefore be simplified considerably in treating real problems.

On the other hand, the “collision term” in the average equation is the term that is responsible for anomalous dissipation and is thus the most interesting term in any theory that deals with the evolution of shock waves. For a spectrum of properly chosen fluctuations this term prevents large amplitude waves from indefinite steeping and breaking and provides the required dissipation of kinetic energy, entropy generation, and shock stabilisation. In its general version given above it should contain the whole physics of the shock including the complete collective processes which occur before real particle collisions come into play.

However, the complexity of these equations is still too large for solving them. So one needs further simplifications in order to infer about the behaviour of shocks. The simplest and at the same time very effective simplification is to ask for the macroscopic conservation laws and the conditions of change of the plasma quantities across the shock transition layer which are in accord with the above fundamental kinetic equations. These are the magnetogasdynamical equations and the Rankine-Hugoniot jump relations.

3.2.2 Conservation Laws

Following the philosophy of simplification we will first, before asking for the internal processes taking place in the shock transition, the generation of dissipation, particle reflection, entropy production etc., look into the global – i.e. large-scale – structure of a shock. In order to do this we need consider only the global plasma and field quantities, density $\langle N \rangle = N_0$, flow velocity $\langle \mathbf{V} \rangle = \mathbf{V}_0$ respectively momentum density $\langle N\mathbf{V} \rangle = N_0\mathbf{V}_0$, pressure $\langle P \rangle = P_0$, magnetic field $\langle \mathbf{B} \rangle = \mathbf{B}_0$, electric field $\langle \mathbf{E} \rangle = \mathbf{E}_0$, current density $\langle \mathbf{j} \rangle = \mathbf{j}_0$, entropy $\mathcal{S}$ and so on.

These quantities are all averages or result from average moments over the global distribution function $\langle F \rangle = F_0$. Since we will be dealing in the following only with average

moments we suppress both the angular brackets and index 0. The prescription of taking moment of order i is

$$M^i = \int d\mathbf{v}^3 \mathbf{v}^i F \quad (3.23)$$

where $\mathbf{v}^i = \mathbf{v} \dots \mathbf{v}$ is understood as the i -fold dyadic product. The first three moments are $N = \int d\mathbf{v}^3 F$, $N\mathbf{V} = \int d\mathbf{v}^3 \mathbf{v} F$, $\mathbf{P} = m \int d\mathbf{v}^3 (\mathbf{v} - \mathbf{V})(\mathbf{v} - \mathbf{V}) F$. Clearly, the diagonal of the pressure tensor $\mathbf{P}$ gives the average energy density and also defines the local temperatures $T_{\parallel}, T_{\perp}$ parallel and perpendicular to the average magnetic field. Operating in the usual way with these definitions on the average Vlasov equation (3.20) produces the well-known full – i.e. infinite – set of magnetogasdynamic equations for the infinite chain of moments of $F^{\pm}$ for each particle species $\pm = e, i$. The first two of them are

$$\frac{\partial N^{\pm}}{\partial t} + \nabla \cdot (N\mathbf{V})^{\pm} = 0 \quad (3.24)$$

$$\frac{\partial (N\mathbf{V})^{\pm}}{\partial t} + \nabla \cdot (N\mathbf{V}\mathbf{V})^{\pm} + \frac{1}{m_{\pm}} \nabla \cdot \mathbf{P}^{\pm} = \frac{e_{\pm} N^{\pm}}{m_{\pm}} (\mathbf{E} + \mathbf{V}^{\pm} \times \mathbf{B}) + \int d\mathbf{v}^3 \mathbf{v} \mathcal{C}^{\pm} \quad (3.25)$$

where by $\mathcal{C}$ the pseudo-collision term on the right of Eq. (3.20) is meant. Because this term conserves particle number (or mass) the zero moment of it vanishes identically and does not contribute to the first (zero-order) of the above moment equations. Wave particle interaction neither changes particle number nor mass density. These are strictly conserved as is shown by the above particle number conservation equation.

In the first order moment equation it produces a wave friction term that has the explicit form

$$-\frac{1}{m_{\pm}} \left\{ \frac{1}{\mu_0} \frac{\partial}{\partial t} \langle \mathbf{e} \times \mathbf{b} \rangle + \nabla \cdot \left[\left(\frac{\epsilon_0}{2} \langle \mathbf{e}^2 \rangle + \frac{1}{2\mu_0} \langle \mathbf{b}^2 \rangle \right) \mathbf{I} - \left(\epsilon_0 \langle \mathbf{e}\mathbf{e} \rangle + \frac{1}{\mu_0} \langle \mathbf{b}\mathbf{b} \rangle \right) \right] \right\} \quad (3.26)$$

All these terms are in fact of nature ponderomotive force-density terms contributed by the average wave pressure gradients; the first term results from the wave Poynting moment, the second is the gradient of a pure isotropic wave pressure, the third is related to wave pressure anisotropy. The inverse proportionality of this entire expression to the mass shows that the main contribution is due to the electron momentum density equation. The effect on the ions can be neglected as they are (in the non-relativistic case considered here) insensitive to ponderomotive effects. We note in passing, that it is this term which while affecting the motion of the electron gas will be responsible for the appearance of anomalous collisions, anomalous resistivity and viscosity, which we will discuss at a later occasion.

The two above equations do not form a complete system of equations. The first contains number density flux, the main constituent of the second equation which, as a new entity, contains the pressure. For $\mathbf{P}$ one, in principle, can derive an energy conservation (heat conduction) equation which would contain the new quantity of heat flux, the next higher moment. On the other hand, one can replace the pressure equation that follows from the energy conservation law, by equations of state, $\mathbf{P}(N, \gamma, T_{\parallel}, T_{\perp})$, which express the pressure tensor components through density, temperature, adiabatic coefficient γ etc. This

is the usual procedure applied when investigating shock solutions. One should, however, be aware of the fact that equations of state in non-equilibrium are merely approximations which hold under certain conditions of either isothermality – which does not apply to shocks as they are not in thermal equilibrium – or adiabaticity. The latter condition is quite reasonable in dealing with the fast processes taking part in the shock environment when the flow passes across the shock front in a time so short that thermalisation becomes impossible.

The idea is to apply the momentum equations to an extended shock that represents a ramp in real space. In the spirit of our discussion in Chapter 2, the first step to do this is assuming that the shock is a thin planar discontinuity that moves at a certain shock velocity U in the shock normal direction $\mathbf{n}$ across the plasma. If we confine all the micro-processes to the interior of the shock plane, i.e. if we go far enough away from the shock plane upstream and downstream, then we can apply the above dissipation-free average conservation laws to the shock and ask only for the differences in the plasma and field parameters between downstream and upstream of the shock, trying to express the downstream values in terms of the undisturbed upstream flow and field values. In doing this, we completely neglect the “pseudo-collision” terms on the right of these equations, since all physics that is going on will be confined to the transition region as wide as it can be. For a plane rigid stationary shock surface this assumption is good enough. However, when doing so, with the above separate conservation laws for electrons and ions, we immediately run into severe problems even in the simplest completely interaction-free case. The reason is that electrons and ions because of their different mass behave completely differently while at the same time cannot be treated separately as they are coupled through charge conservation and electrical neutrality and through their unequal contributions to the electric current density and therefore to the fields, a difficulty that has been discussed by *Woods* [1971].

3.3 Rankine-Hugoniot Relations

In order to overcome this difficulty one is forced to further simplification of the conservation equations by adding up the electron and ion equations [cf., e.g., *Baumjohann & Treumann*, 1996, as for one of the many accounts available in the literature]. To this purpose one must define new centre-of-mass variables

$$m = \sum_{\pm} m_{\pm} = m_i \left(1 + \frac{m_e}{m_i} \right), \quad N = \frac{\sum_{\pm} m_{\pm} N^{\pm}}{\sum_{\pm} m_{\pm}}, \quad \mathbf{V} = \frac{\sum_{\pm} m_{\pm} (N\mathbf{V})^{\pm}}{\sum_{\pm} m_{\pm} N^{\pm}} \quad (3.27)$$

This leads to the magnetogasdynamic or MHD equations for a single-fluid plasma. Since the equation of continuity remains unchanged from Eq. (3.24) it suffices to write down the momentum conservation equation

$$\frac{\partial(mN\mathbf{V})}{\partial t} + \nabla \cdot (mN\mathbf{V}\mathbf{V}) = -\nabla \cdot \mathbf{P} + \rho \mathbf{E} + \mathbf{j} \times \mathbf{B} \quad (3.28)$$

where $\mathbf{P} = \mathbf{P}_e + \mathbf{P}_i$ is the total pressure tensor, and ρ is the electric charge density $\rho = e(N_i - N_e)$ which in quasi-neutral plasmas outside the shock is assumed to be zero such

that the second term on the right containing the average electric field vanishes outside the shock ramp. The last term in this equation is the Lorentz force written in terms of the average current from Ampère's law

$$\mu_0 \mathbf{j} = \nabla \times \mathbf{B} \quad (3.29)$$

The displacement current can be safely neglected because these equations hold only for very slow variations with frequency smaller than the ion cyclotron frequency $\omega \ll \omega_{ci} = eB/m_i$, scales much larger than the ion gyro-radius $L \gg r_{ci} = V_{i\perp}/\omega_{ci}$, and wave speeds much less than the speed of light. Note that this equation is completely collisionless. If we would have retained the pseudo-collision term on the right in the electron equation this would simply have added an electron ponderomotive force term on the right.

In fact, together with Maxwell's equations these equations are not yet complete in a double sense. They have to be completed with appropriate equations of state for the pressure components, as has been mentioned above, and they have to be completed by a relation between the current and the average electric field that appears in Maxwell's equations, i.e. with an appropriate Ohm's law. This is found by subtracting the electron and ion momentum conservation equations and turns out to be quite complicated [cf., e.g., *Krall & Trivelpiece, 1973*]. In slightly simplified form the collisionless Ohm's law reads

$$\mathbf{E} + \mathbf{V} \times \mathbf{B} = \frac{1}{eN} \mathbf{j} \times \mathbf{B} - \frac{1}{eN} \nabla \cdot \mathbf{P}_e + \frac{m_e}{e^2 N} \frac{\partial \mathbf{j}}{\partial t} \quad (3.30)$$

Note that on the right only electron terms are contained in this expression. Also, an electron ponderomotive term – responsible for anomalous transport effects – would appear on the right if we would retain the pseudo-collision term.

However, even in this form even though the system is non-collisional Ohm's law is still too complex for treating the conservation laws at a shock transition. The reason is that the right-hand side introduces second order spatial derivatives into Faraday's law through the pressure gradient and current expressions. One therefore argues that for sufficiently flat shock transitions the terms on the right can be neglected. This argument implies that one must go far enough away from the shock into a region where any shock excited turbulence has decayed away in order to apply global conservation laws to the shock. This can be done when only the left-hand side in Ohm's law is retained and the ideal MHD frozen-in condition holds:

$$\mathbf{E} = -\mathbf{V} \times \mathbf{B} \quad (3.31)$$

Assuming that the shock is plane and narrow as shown in Figure 3.6 such that any variations along the shock can be ignored and the sole variation is along the shock normal, Eqs. (3.28)–(3.31), the continuity equation and Maxwell's equations become all one-dimensional and can be integrated along $\mathbf{n}$ across the shock transition (with regular boundary conditions at $x = \pm\infty$). Applying the definition of the shock normal (2.3) in Chapter 2 and the prescription for the ∇ -operator in Eq. (2.4) transforms these equations into a non-linear algebraic system of equations for the jumps [...] of the field quantities

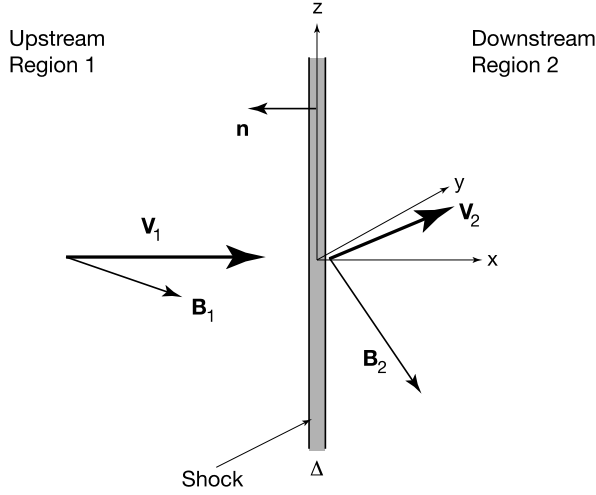


Figure 3.6: A sketch of the planar thin (width Δ) shock geometry with $\mathbf{n}$ shock normal, upstream and downstream regions, bulk flow velocities and magnetic field vectors.

$$\begin{aligned}
 \mathbf{n} \cdot [N\mathbf{V}] &= 0 \\
 \mathbf{n} \cdot [mN\mathbf{V}\mathbf{V}] + \mathbf{n} \cdot \left[P + \frac{B^2}{2\mu_0} \right] - \frac{1}{\mu_0} \mathbf{n} \cdot [\mathbf{B}\mathbf{B}] &= 0 \\
 [\mathbf{n} \times \mathbf{V} \times \mathbf{B}] &= 0 \\
 \mathbf{n} \times [\mathbf{B}] &= 0 \\
 \left[mN\mathbf{n} \cdot \mathbf{V} \left\{ \frac{V^2}{2} + w + \frac{1}{mN} \left(P + \frac{B^2}{\mu_0} \right) \right\} - \frac{1}{\mu_0} (\mathbf{V} \cdot \mathbf{B}) \mathbf{n} \cdot \mathbf{B} \right] &= 0
 \end{aligned} \tag{3.32}$$

Here, for simplicity, the pressure has been assumed isotropic. The last equation is the rewritten energy conservation equation, where $w = c_v P / k_B N$ is the ideal gas enthalpy density, c_v the specific heat, and k_B Boltzmann's constant. This system of equations is the implicit form of the *Rankine-Hugoniot* conservation equations in ideal magnetogasdynamics (ideal MHD). In this version it contains all ideal MHD discontinuities of which shock waves are a subclass, the class of solutions with a finite flow across the discontinuity, compressions (in density), and increases in temperature T , pressure P , and entropy $\mathcal{S}$ across the discontinuity in the transition from upstream to downstream.

3.3.1 Explicit MHD Shock Solutions

We are not interested in the full set of solutions of the above system of jump conditions (3.32). We rather look for genuine shock conditions. This requires finite mass flux $\mathcal{F} = NV_n$ across the shock in the normal direction. The first of the Rankine-Hugoniot relations (3.32) tells that the jump $[\mathcal{F}] = 0$. Hence $\mathcal{F} = \text{const}$, and we must sort for solutions with $\mathcal{F} \neq 0$, or $N_1 V_{n1} = N_2 V_{n2}$, in order to be dealing with a shock.

Introducing the specific volume $\mathcal{V} = (mN)^{-1}$, the whole system of jump conditions can be factorised [Baumjohann & Treumann, 1996] and can be written in the form

$$\mathcal{F} \left(\mathcal{F}^2 - \frac{B_n^2}{\mu_0 \langle \mathcal{V} \rangle} \right) \left\{ \mathcal{F}^4 + \mathcal{F}^2 \left(\frac{[P]}{[\mathcal{V}]} - \frac{\langle \mathbf{B} \rangle^2}{\mu_0 \langle \mathcal{V} \rangle} \right) - \frac{B_n^2}{\mu_0 \langle \mathcal{V} \rangle} \frac{[P]}{[\mathcal{V}]} \right\} = 0 \quad (3.33)$$

In the one-fluid approximation magnetogasdynamic shock waves with $\mathcal{F} \neq 0$ are contained in the expression in curly braces which still depends on the jumps in pressure $[P]$ and specific volume $[\mathcal{V}]$ and thus on the energy conservation equation respectively the equation of state. We will not discuss this equation further as in the following more insight can be gained from explicit consideration of a few particular cases.

Under the special condition that the flow in the upstream Region 1 is along x (anti-parallel to $\mathbf{n}$) and the upstream magnetic field $\mathbf{B}_1 = (B_{1x}, 0, B_{1z})$ is in the (xz) -plane, and assuming $\mathcal{F} \neq 0$, the jump conditions Eqs. (3.32) simplify. Since V_1, B_{1x}, B_{1z}, P_1 are known quantities, it is convenient to introduce normalised variables for the corresponding downstream values

$$\frac{N_2}{N_1} \rightarrow N_2, \quad \frac{\mathbf{V}_2}{V_1} \rightarrow \mathbf{V}_2, \quad \frac{T_{1,2}}{mV_1^2/2} \rightarrow T_{1,2}, \quad \frac{\mathbf{B}_{1,2}}{\sqrt{\mu_0 m N_1 V_1^2}} \rightarrow \mathbf{B}_{1,2}$$

where the temperature is taken in energy units. Instead of it we may also use the corresponding thermal speeds $v_{1,2}$ which by the above normalisation are normalised to V_1 . This then yields the following normalised Rankine-Hugoniot relations, in which $B_n = \text{const}$ as a consequence of the vanishing divergence of the magnetic field,

$$\begin{aligned} N_2 V_{n2} &= 1 \\ V_{n2} B_{z2} - V_{z2} B_n &= B_{z1} \\ B_{z2} B_n - V_{z2} &= B_{z1} B_n \\ 2N_2 (v_2^2 + V_{n2}^2) + B_{z2}^2 &= 2(1 + v_1^2) + B_{z1}^2 \\ V_{n2}^2 + V_{z2}^2 + 2B_{z2} B_{z1} + 5v_2^2 &= 1 + 2B_{z1}^2 + 5v_1^2 \end{aligned} \quad (3.34)$$

The energy conservation equation yields the last in these expressions. There the enthalpy is taken into account giving the factor 5 in front of the thermal velocities. These five equations can be combined into a third-order equation for one of the downstream unknown quantities, for instance V_{n2} , expressed in terms of the upstream values

$$a_3 V_{n2}^3 + a_2 V_{n2}^2 + a_1 V_{n2} + a_0 = 0 \quad (3.35)$$

where $a_0 = -B_n^2[B_{z1}^2 + B_n^2(1 + 5v_1^2)]$, $a_1 = 2B_n^2(1 + 2B_{z1}^2 + 5v_1^2) - \frac{1}{2}B_{z1}^2$, $-a_2 = 1 + 5v_1^2 + 8B_n^2 + \frac{5}{2}B_{z1}^2$, $a_3 = 4$. Below we discuss a few simple illustrative solutions of this equation.

3.3.2 Perpendicular Shocks

For strictly perpendicular shocks we have $B_n = 0$, $B_{z1} = B_1$, $a_0 = 0$, $a_1 = -\frac{1}{2}B_1^2$, $-a_2 = 1 + 5v_1^2 + \frac{5}{2}B_{z1}^2$, $a_3 = 4$. Equation (3.35) turns into a quadratic equation yielding the solution

$$V_{n2} = \frac{1}{8} \left\{ 1 + \left(1 + \frac{5}{2}\beta_1 \right) B_1^2 + \left[\left(1 + \left(1 + \frac{5}{2}\beta_1 \right) B_1^2 \right)^2 + 2B_1^2 \right]^{\frac{1}{2}} \right\} = \frac{1}{N_2} = \frac{1}{B_2} \quad (3.36)$$

Since the condition for a shock to exist is that the normal velocity $V_{n2} < 1$ in Region 2, we immediately conclude that in a perpendicular shock the density and tangential magnetic field components in Region 2 increase by the same fraction as the normal velocity drops, and this fraction is determined by the plasma- β ratio $\beta_1 = 2\mu_0 N_1 T_1 / B_1^2$ in Region 1, where $T_1 = T_{e1} + T_{i1}$ is the total temperature. The condition on V_{n2} implies that the Mach number takes the form (now in physical units)

$$1 < \mathcal{M} = \frac{1}{1 + 5\beta_1/6} \frac{V_1}{V_{A1}} = \frac{1}{1 + 5\beta_1/6} \mathcal{M}_A \quad (3.37)$$

Here $\mathcal{M}_A$ is the Alfvén-Mach number which is the flow to Alfvén velocity ratio. In cold plasmas or plasmas containing strong magnetic fields $\beta \ll 1$, and the Mach number is simply the Alfvén-Mach number. Conversely, in hot plasmas the Mach number becomes about the ordinary gasdynamic Mach number.

For the increase in normalised temperature one finds accordingly

$$\frac{T_2}{T_1} = 1 + \frac{4}{5T_1} \left[\frac{N_2^2 - 1}{2N_2^2} + \frac{1 - N_2}{\mathcal{M}_A^2} \right] > 1 \quad (3.38)$$

This is always larger than one. Perpendicular shocks cause plasma heating during shock transition time and thus cause also increase in entropy

$$\Delta \mathcal{S} \propto \ln \left[\frac{1}{N_2} \left(\frac{T_2}{T_1} \right)^{\frac{1}{\gamma-1}} \right] \quad (3.39)$$

which holds under the ideal gas assumption.

The Critical Mach Number for Perpendicular Shocks

To show how the critical Mach number of a shock arises from the Rankine-Hugoniot relations we consider the strictly perpendicular case with vanishing upstream pressure $P_1 = 0$. The explicit jump conditions become very simple in this case:

$$\begin{aligned} N_1 V_1 &= N_2 V_2 \\ V_1 B_1 &= V_2 B_2 \\ N_1 V_1^2 + \frac{B_1^2}{2\mu_0 m} &= N_2 V_2^2 + \frac{P_2}{m} + \frac{B_2^2}{2\mu_0 m} \\ \frac{1}{2} V_1^2 + \frac{B_1^2}{\mu_0 m N_1} &= \frac{1}{2} V_2^2 + \frac{\gamma}{\gamma-1} \frac{P_2}{m N_2} + \frac{B_2^2}{\mu_0 m N_2} \end{aligned} \quad (3.40)$$

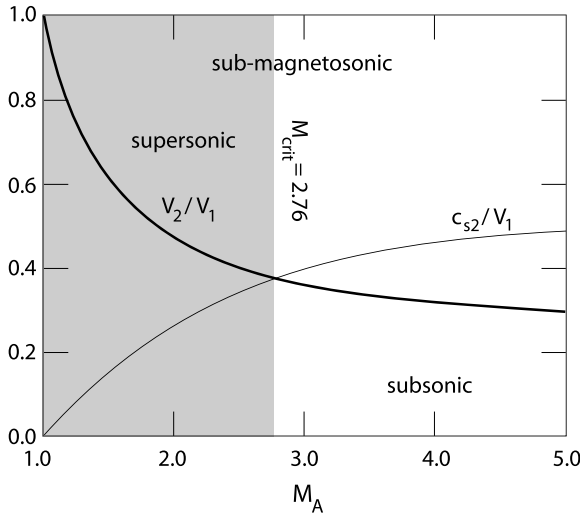


Figure 3.7: Dependence of the downstream normalised flow V_2/V_1 and sound c_{s2} velocities on the upstream Alfvénic Mach number for an ideal MHD perpendicular shock with zero upstream pressure $P_2 = 0$. The crossing of the two curves defines the critical Mach number which is $\mathcal{M}_A = 2.76$.

where B is the only existing tangential component of the magnetic field here, and $\gamma = \frac{5}{3}$ is the adiabatic index (valid for fast = adiabatic transitions across the shock). This is the simplest imaginable case of an MHD shock, and it is easy to solve these equations. Figure 3.7 shows the resulting relation between the normalised downstream flow V_2/V_1 and downstream sound speed $c_{s2}/V_1 = \sqrt{5P_2/3mN_2}/V_1$ as function of the upstream Alfvénic Mach number $\mathcal{M}_A = B_1/V_1\sqrt{\mu_0 m N_1}$.

The two curves in the figure cross each other at the critical Mach number which in the present case is $\mathcal{M}_{crit} = 2.76$ and where the downstream sound speed exceeds the flow speed. Below the critical Mach number the downstream flow is still supersonic (though clearly sub-magnetosonic!). Only above the critical Mach number the downstream flow velocity falls below the downstream sound speed. There is thus a qualitative change in the shock character above it that is not contained in the Rankine-Hugoniot conditions.

3.3.3 Parallel Shocks

This case is not well treated in magnetogasdynamics conditions as we have explained earlier. Since the magnetic field is normal to the shock it is theoretically unaffected by the presence of the shock which therefore should become purely gasdynamic. In the above perpendicular shock jump conditions one can for this case simply delete the magnetic terms. However, this does not cover the real physics involved into parallel shocks which must be treated on the basis of kinetic theory and with the simulation tool at hand.

3.3.4 High Mach Numbers

This limit applies when the ram pressure of the flow is very high and exceeds the thermal pressure. Then all terms including $v_{1,2}$ can be neglected. Moreover, one usually also neglects the magnetic field in this case, and the shocks become then purely flow determined with $V_{n2} = N_2^{-1} \sim B_2^{-1} \simeq \frac{1}{4}$, suggesting that both, the magnetic field and density should not increase by more than a factor of 4.

In fact, the observations in interplanetary space indeed confirm that all shocks that have been observed there are weaker than this. However, again, this reasoning does not really apply in plasmas because at very high Mach numbers other effects come into play which are connected to the kinetic nature of a plasma, electrodynamic effects, and the differences in the electron and ion motions. Ultimately relativistic effects must be taken into account. These become susceptible first for electrons, increasing their mass but at the same time distinguishing them even stronger from the inert ions, because the electron dynamics changes completely in the relativistic domain.

In addition, high Mach number shocks are supercritical and even though it seems that one could treat them in the simple way as has been done here, the kinetic effects involved into their physics inhibit such a simplistic interpretation of high-Mach number shocks. High-Mach number shocks readily become turbulent, exciting various kinds of waves which grow to large amplitudes and completely modify the environment of the shock which cannot be treated any more as quiet. In such a turbulent environment shocks assume intermittent character losing stationarity or even identity as a single ramp which the flow has to surpass when going from Region 1 to Region 2. Occasionally the distinction between two regions only may become obsolete. There might be more than one transition regions, subshocks form, the ramp will evolve its own structures. And these structures come and go, are temporarily created and damp away to make space for the evolution of other new structures. Probably, high Mach number shocks exist only temporarily at one and the same spatial location. They are highly dynamical, changing their nature, structure, shape, steepness and intensity along the surface of the shock such that they strongly deviate from one-dimensionality and even from two-dimensionality. They are time-dependent, reforming themselves continuously in different regions of space and thus cannot be described by a simple plane shock geometry of the kind we have assumed. Later in this book at the appropriate place we will consider moderately high Mach number shocks when dealing with the extended class of supercritical shocks.

Oblique Shocks

Real shocks do not belong neither to the very particular classes of parallel nor perpendicular shocks. Real shocks are oblique in the sense that the upstream magnetic fields $\mathbf{B}_1$ are inclined with respect to the shock normal $\mathbf{n}$. As mentioned earlier, one distinguishes between *quasi-parallel* and *quasi-perpendicular* shocks depending on the shock normal angle Θ_{Bn} being closer to 0° or 90° . Since we will treat the properties of these shocks separately in some following chapters, we are not going to discuss them at this place.

Shocks around finite size obstacles will never be really plane. The best approximations to plane shocks are interplanetary shocks. Bow shocks in front of magnetised planets, comets or other bodies are always curved. They assume all kinds of shock properties along their surfaces reaching from perpendicular to oblique and parallel. Curving their surface implies that the shock normal changes its angle with respect to the direction of the upstream flow $\mathbf{V}_1$.

Since shocks evolve for Mach numbers $\mathcal{M}_{ms} = V_1/c_{ms} > 1$ they occupy a finite volume of space only. For shock formation the Mach number based on V_{1n} normal to the shock is relevant. Defining the flow normal angle $\cos \Theta_{Vn} = |\mathbf{V}_1 \cdot \mathbf{n}|$, formation of a bow shock in front of a finite sized object like a magnetosphere is restricted to flow normal angles

$$\Theta_{Vn} < \cos^{-1}(\mathcal{M}_{ms}^{-1}) \quad (3.41)$$

For instance, at a nominal Mach number $\mathcal{M}_{ms} \sim 8$ shocks exist for angles $\Theta_{Vn} < 82.8^\circ$. Thus the opening angle of a shock decreases with Mach number. The shock surface around blunt obstacles like the magnetospheres of a Earth and Jupiter has hyperbolic shape away from the shock nose, thereby keeping the angle θ_{Vn} about constant. In fact, observations in space near Earth and Jupiter have shown the shock to exist very far outside the ecliptic plane.

3.4 Waves and Instabilities

It has been mentioned several times that shocks evolve from waves mainly through nonlinear wave steepening and the onset of dissipation and dispersion. Moreover, it is the various modes of waves that are responsible for the generation of anomalous dissipation, shock ramp broadening, generation of turbulence in the shock environment and shock ramp itself, as well as for particle acceleration, shock particle reflection and the successive effects. The idea is that in a plasma that consists of electro-dynamically active particles the excitation of the various plasma wave modes in the electromagnetic field as collective effects is the easiest way of energy distribution and transport. There is very little momentum needed in order to accelerate a wave, even though many particles are involved in the excitation and propagation of the wave, much less momentum than accelerating a substantial number of particles to medium energy. Therefore any more profound understanding of shock processes cannot avoid bothering with waves, instabilities, wave excitation and wave particle interaction.

3.4.1 Dispersion Relation

Waves are a very general phenomenon of most media. However, they do not fall from sky. Instead, they evolve from small thermal fluctuations in the medium. Such fluctuations are unavoidable. In order for a wave to propagate in the medium a number of conditions need to be satisfied, however. The first is that the medium allows for a particular range of frequencies ω and wave-vectors $\mathbf{k}$ to exist in the medium; i.e. it allows for eigen-oscillations

or eigenmodes. These ranges are specified by the dispersion relation $D(\omega, \mathbf{k}, \dots) = 0$ which formulates the condition that the dynamical equations of the medium possess small-amplitude solutions. This dispersion relation is usually derived in the linear infinitesimally small amplitude approximation. However, nonlinear dispersion relations can sometimes also be formulated in which case $D(\omega, \mathbf{k}, |\mathbf{a}|)$ depends on the fluctuation amplitude $|\mathbf{a}|$ as well.

Plasmas are electromagnetically active media with the electromagnetic field governed by Maxwell's equations. Since there the plasma properties enter only through the material equations (i.e. current density $\mathbf{j}$, space charge ρ), the dispersion relation is most easily obtained from them. Moreover both, the current and space charge in a plasma, depend on the number densities in the plasma; i.e. the space charge variation can be included into the current variation. It is then simple matter to derive the general electromagnetic wave equation for the fluctuating fields, $\mathbf{e}$, on a much slower evolving background, $\langle \mathbf{B}, \mathbf{E} \rangle$,

$$\nabla^2 \mathbf{e} - \nabla(\nabla \cdot \mathbf{e}) - \epsilon_0 \mu_0 \frac{\partial^2 \mathbf{e}}{\partial t^2} = \mu_0 \frac{\partial \mathbf{j}}{\partial t} \quad (3.42)$$

The magnetic fluctuation field $\mathbf{b}$ is completely determined from Maxwell's equations, and the current appearing on the right is expressed conveniently through the space-time dependence of the fluctuation-conductivity tensor $\boldsymbol{\sigma}(\mathbf{x}, t)$ as

$$\mathbf{j}(\mathbf{x}, t) = \int d\mathbf{x}' \int_{-\infty}^t dt' \boldsymbol{\sigma}(\mathbf{x} - \mathbf{x}', t - t') \cdot \mathbf{e} \quad (3.43)$$

an expression that implicitly accounts for causality due to the integration over the entire past of the current up to the observation time t . Since the fluctuation current is a functional of the complete set of particle distribution functions (through the zero, N , and first, $N\mathbf{V}$, moments) the complete evolution of the fluctuations up to a large amplitude shock is contained in these expressions. However, for practical purposes one linearises this equation by assuming that the fluctuation-conductivity tensor is, to first order, independent of the fields $\mathbf{e}, \mathbf{b}$. The above equation (3.42) becomes linear under this assumption and can be Fourier analysed, with wave vector $\mathbf{k}$ and frequency ω , yielding (with $c^2 = 1/\mu_0 \epsilon_0$ the square of the speed of light)

$$\left[\left(k^2 - \frac{\omega^2}{c^2} \right) \mathbf{I} - \mathbf{k}\mathbf{k} - i\omega\mu_0 \boldsymbol{\sigma}(\omega, \mathbf{k}) \right] \cdot \mathbf{e}(\omega, \mathbf{k}) = 0 \quad (3.44)$$

The quantities in this expression satisfy the following symmetry relations $\mathbf{e}(-\omega, -\mathbf{k}) = \mathbf{e}^*(\omega, \mathbf{k})$, $\boldsymbol{\sigma}(-\omega, -\mathbf{k}) = \boldsymbol{\sigma}^*(\omega, \mathbf{k})$. Setting the expression in brackets to zero yields the equation for the linear eigenmodes $\omega = \omega(\mathbf{k})$. For convenience we define the dielectric tensor

$$\boldsymbol{\epsilon}(\omega, \mathbf{k}) \equiv \mathbf{I} + \frac{i}{\omega\epsilon_0} \boldsymbol{\sigma}(\omega, \mathbf{k}) \quad (3.45)$$

which satisfies the same conditions as the fluctuation conductivity. Then we can finally write the general dispersion relation in the compact form as the determinant of the bracketed expression

$$D(\omega, \mathbf{k}) \equiv \text{Det} \left[\left(k^2 - \frac{\omega^2}{c^2} \right) \mathbf{I} - \mathbf{k}\mathbf{k} + \boldsymbol{\epsilon}(\omega, \mathbf{k}) \right] = 0 \quad (3.46)$$

The particular linear physics of the plasma is contained in the dielectric tensor through the conductivity. (We note in passing that for any classical medium the above conditions together with the dispersion relation are the equivalent of the well-known Kramers-Kronig relations of causal fluctuations in quantum mechanics [see, e.g., *Landau & Lifshitz, 1980*].) In addition to the solution of the dispersion equation (3.46) the first problem consists in the determination of the linear fluctuation conductivity tensor which enters the dielectric tensor (3.45). For this one needs to go to the appropriate plasma model. However, we repeat that without this linear step one cannot obtain any susceptible information about the nature of a shock wave. This we have explained in breadth in Chapter 2 and the preceding sections.

The linear dispersion relation Eq. (3.46) has plane wave mode solutions of the form $\propto \exp i(\mathbf{k} \cdot \mathbf{x} - \omega t)$, which are eigenmodes of the particular plasma model which is described by the kinetic equations (or appropriate simplifications of the kinetic equations) of the plasma. The dynamics of the plasma enters through the wave conductivity tensor which can be determined from the Fourier transformed expression for the current density

$$\mathbf{j}(\mathbf{k}, \omega) = \sum \pm e_{\pm} \int d\mathbf{v}^3 \mathbf{v} f^{\pm}(\mathbf{v}, \mathbf{k}, \omega) = \boldsymbol{\sigma}(\mathbf{k}, \omega) \cdot \mathbf{e}(\mathbf{k}, \omega) \quad (3.47)$$

as the first moment of the fluctuating part $f^{\pm}$ of the distribution function $F^{\pm}$. Solving for the integral and expressing $f^{\pm}$ through the electric wave field $\mathbf{e}$ yields the wanted form of the wave conductivity tensor $\boldsymbol{\sigma}$. The problem is thus reduced to the determination of $f^{\pm}$ from Eq. (3.22) where we drop the average terms on the right-hand side retaining only terms linear in the fluctuations:

$$\frac{\partial f^{\pm}}{\partial t} + \mathbf{v} \cdot \nabla f^{\pm} + \frac{e_{\pm}}{m_{\pm}} (\mathbf{E}_0 + \mathbf{v} \times \mathbf{B}_0) \cdot \nabla_{\mathbf{v}} f^{\pm} = -\frac{e_{\pm}}{m_{\pm}} (\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} F_0^{\pm} \quad (3.48)$$

Operating with a Fourier transform on this equation then yields the following expression

$$\left[1 - \frac{ie_{\pm}}{m_{\pm}(\mathbf{k} \cdot \mathbf{v} - \omega)} (\mathbf{E}_0 + \mathbf{v} \times \mathbf{B}_0) \cdot \nabla_{\mathbf{v}} \right] f^{\pm} = \frac{ie_{\pm}}{m_{\pm}} \frac{i(\mathbf{e} + \mathbf{v} \times \mathbf{b}) \cdot \nabla_{\mathbf{v}} F_0^{\pm}}{\mathbf{k} \cdot \mathbf{v} - \omega} \quad (3.49)$$

which determines $f^{\pm}$ in terms of the average and fluctuating field quantities, which is just what we want. One can now make assumptions about the average fields and distribution function in order to explicitly calculate $f^{\pm}$. Usually these assumptions are $\mathbf{E}_0 = 0$, $\mathbf{B}_0 = B_0 \hat{\mathbf{z}}$ with $B_0 = \text{const.}$ Then the operator $\nabla_{\mathbf{v}} = -\hat{\mathbf{z}} \partial / \partial \phi$ on the left simplifies to a mere derivative with respect to the gyration angle ϕ of the particles. Further assumptions on $F^{\pm}$ are that the average distributions are gyro-tropic, in which case the integration with respect to ϕ becomes trivial.

With such assumptions it is not difficult so still tedious to solve for $f^\pm$ and finally get the conductivity tensor. The respective expressions have been given in various places and will not be repeated here. Good references among others are *Montgomery & Tidman* [1964], *Gary* [1993], *Baumjohann & Treumann* [1996]. One can even include weak inhomogeneity [cf., e.g. *Krall & Trivelpiece*, 1973] in the average distribution function and fields (which is the eikonal approximation holding under the assumption that the gradient scale is much longer than the wavelength, or $|\nabla/\mathbf{k}| \ll 1$). This is needed when considering an inhomogeneous initial state like a given soliton or shock structure and investigating its prospective stability. In this case the plasma background is not homogeneous anymore because a soliton or shock has already evolved in it and has locally modified the plasma. Any wave modes which will be excited on this modified background will then not be influenced only by sideband formation, steeping, nonlinearity and dispersion but also by the change of the plasma properties from location to location. This implies that the waves themselves change character and properties across a shock.

Damping/Growth Rate

The solutions of the dispersion relation are in most cases complex, and for real wave vector $\mathbf{k}$ can be written as $\omega(\mathbf{k}) = \omega_r(\mathbf{k}) + i\gamma(\omega_r, \mathbf{k})$, where the index r indicates the real part, and γ is the imaginary part of the frequency which itself is a function of the real frequency and wave number, because each mode of given frequency can behave differently in time, and the wave under normal conditions will be dispersive, i.e. it will not be a linear function of wave number. In most cases the amplitude of a given wave will change slowly in time, which means that the imaginary part of the frequency is small compared to the real frequency. If this is granted, then γ can be determined by a simple procedure directly from the dispersion relation $D(\omega, \mathbf{k}) = D_r(\omega, \mathbf{k}) + iD_i(\omega, \mathbf{k})$, which can be written as the sum of its real D_r and imaginary D_i parts because a small imaginary part γ in the frequency changes the dispersion relation only weakly, and it can be expanded with respect to this imaginary part. Up to first order in γ/ω one then obtains

$$D_r(\omega_r, \mathbf{k}) = 0, \quad \gamma(\omega_r, \mathbf{k}) = -\frac{D_i(\omega_r, \mathbf{k})}{\partial D_r(\omega_r, \mathbf{k})/\partial \omega|_{\gamma=0}} \quad (3.50)$$

The first of these expressions determines the real frequency as function of wave number $\omega_r(\mathbf{k})$ which can be calculated directly from the real part of the dispersion relation. The second equation is a prescription to determine the imaginary part of the frequency, i.e. the damping or growth rate of the wave.

Remarks

Two remarks on the dispersion relation are in place. First, the weak damping/growth rate solutions $\omega = \omega(\mathbf{k})$ of the above general dispersion relation – themselves called dispersion relations – are also of use in the weakly nonlinear case. They can be understood as the lowest-order expansion term of a more general nonlinear dispersion relation $\omega = \omega(\mathbf{k}, |\mathbf{e}, \mathbf{b}|^2)$ which depends weakly on the wave amplitude or wave energy $|\mathbf{e}, \mathbf{b}|^2$.

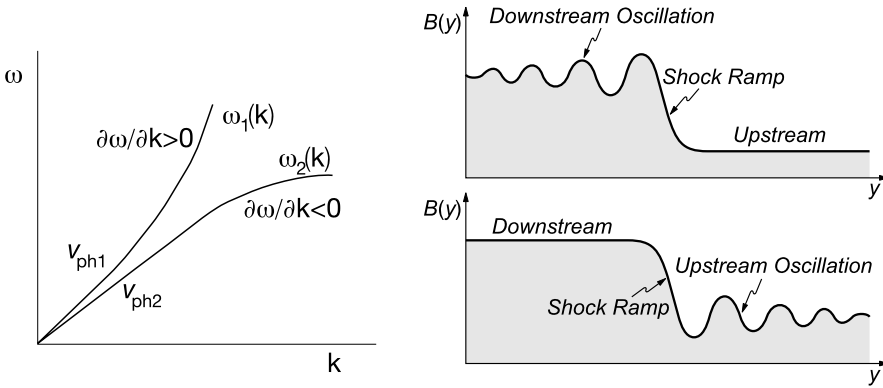


Figure 3.8: *Left:* Two different types of dispersions in the (real) $(\omega, \mathbf{k})$ -plane. Short waves with concave dispersion have slower group velocity than long waves and form a trail on the long wave. Short waves with convex dispersion move faster than long waves. *Right:* The effect of this difference in sideband wave velocity on a laminar subcritical shock wave. Convex dispersions produce run-away waves which appear as spatially damped upstream oscillations (in the B-field, for instance). The trailing waves resulting from concave dispersion occur as downstream spatially damped oscillations. Maximum wave amplitudes are observed near the shock ramp in both cases.

When taking into account higher order expansion terms in the wave amplitude it produces other non-linear equations which govern the amplitude evolution of the wave under consideration. Such an equation is the non-linear Schrödinger equation we will get familiar with when discussing transport processes.

Second, from the dispersion relation $\omega_r = \omega_r(\mathbf{k})$ one can infer in which way steepening of a wave is compensated by the dispersion of the wave. Figure 3.8 on its left shows two typical cases of (real) dispersion curves of low frequency waves in the (ω_r, k) -plane from which shock waves could evolve [after Sagdeev, 1966]. Both curves have in common that they exhibit linear dispersion at long wavelengths, i.e. at small wave-numbers k , with slope giving the phase velocities of the waves. In this region all nonlinearly generated sidebands have same phase and group velocities causing broadening of the wave spectrum and steepening. However, at higher wave-numbers the dispersion curves start diverging from linear slope, one of the waves turning convex, the other concave. These turnovers imply a change in phase and group velocities. The convex dispersion implies that shorter wavelengths generated in the convex part of the dispersion curve move faster than the long waves. They will thus catch up with the long wavelength wave and run away ahead of the wave forming upstream precursors of the wave as shown for the shock in the lower part on the right. On the other hand, for the concave dispersion shorter wavelength waves fall behind the long waves. They represent a wave trail following the large amplitude long wave as is shown for the shock in the upper part on the right. Hence a simple glance at the dispersion curves already dismantles the possible properties of the expected nonlinearity and the structure of the shock.

A word of caution is in place here, however. This reasoning does not hold for all shocks but for subcritical laminar shocks only. Supercritical higher Mach number shocks

will behave in a more complicated way being much less dependent on dissipation and dispersion.

3.4.2 The MHD Modes – Low- β Shocks

The waves from which a shock forms are the lowest-frequency plasma modes that are excited under the particular conditions of the shocked plasma. In low- β (cold) plasma these are the three fundamental MHD modes. Since all collisionless shocks in the heliosphere are magnetised the magnetic field has to be included, and the global shocks are not purely electrostatic even though subshocks developing in them can well behave approximately electrostatic.

We are already familiar with the three low- β magnetogasdynamic modes, the fast and slow magnetosonic and the intermediate Alfvén waves. These are the lowest frequency eigenmodes of a homogeneous not necessarily isotropic plasma, i.e. when a small disturbance is present in the plasma it will propagate in one or all of these modes. Their dispersion relation follows from

$$D_{\text{mhd}}(\omega, \mathbf{k}) = \begin{pmatrix} \omega^2 - V_A^2 k_{\parallel}^2 - c_{ms}^2 k_{\perp}^2 & 0 & -c_s^2 k_{\parallel} k_{\perp} \\ 0 & \omega^2 - V_A^2 k_{\parallel}^2 & 0 \\ -c_s^2 k_{\parallel} k_{\perp} & 0 & \omega^2 - c_s^2 k_{\parallel}^2 \end{pmatrix} = 0 \quad (3.51)$$

which depends on the parallel and perpendicular components of $\mathbf{k}$ only in a very simple way. Moreover, it is a purely real dispersion relation lacking any imaginary part and therefore also any damping which is of course typical for a low frequency dissipation-free plasma. It is a different question of how these modes can be excited, and we will come to this at a later stage.

The phase velocity c_{ms} of these modes has been dealt with already in Eq. (2.1). Since these waves are linear waves with no dispersion, their dispersion relation is simply $\omega = kc_{ms}(\theta)$ with θ the angle between the wave vector $\mathbf{k}$ and magnetic field $\mathbf{B}$. Figure 3.9 shows the real space angular dependence of these three phase velocities for two special cases. Clearly in the direction perpendicular to the magnetic field only the fast mode propagates and, hence, strictly perpendicular MHD shocks are fast shocks as has been noted. In the direction parallel to $\mathbf{B}$ all three waves can propagate.

Magnetosonic Solitons in Cold Plasma

We will now show that the method of the Sagdeev pseudo-potential can be used to understand the formation of a fast mode solitary wave (or soliton) propagating strictly perpendicular to the magnetic field $\mathbf{B}$. This has first been shown by *Davis et al* [1958]. What results from this procedure will not yet be a shock, because, as we have noted before, shock formation requires the presence of some kind of dissipation, while the equations on which the present theory is based are strictly dissipation-free (remember that we have dropped the correlation terms on the right-hand side of the kinetic equation before deriving the moment equations and that we have not yet discussed any way of how dissipation occurs when these terms are taken into account).

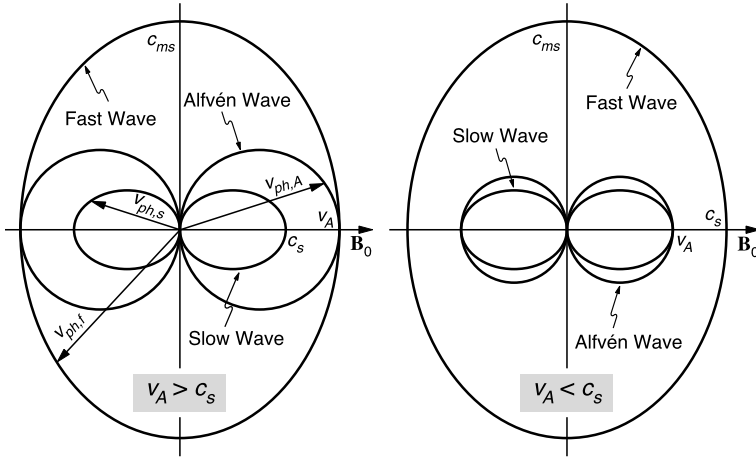


Figure 3.9: Wave vector diagram of two cases of MHD waves in the plane of the magnetic field.

In order to find the stationary solutions we are looking for, one must retain the non-linearity in the stationary one-dimensional quasi-neutral magnetogasdynamic equations. This nonlinearity appears in the convective term $V_n \nabla_n V_n$ in the equation of motion. From constancy of normal flux $[\mathcal{F}] = 0$ one has $NV_n = N_1 V_1$ and $E_y = N_1 B_1$ where the index 1 means undisturbed values far upstream. Since only electrons contribute to the current by their drift in the crossed electric and magnetic fields $\mathbf{E} = -\mathbf{V} \times \mathbf{B}$, $\mathbf{B} = B\hat{\mathbf{z}}$ in the shock frame, we must retain a small component $E_n = -BV_y$ across the shock, effectively produced by the difference in electron and ion motion and causing the shock current to flow in y -direction in the shock transition. It occurs in the stationary equation of motion on the scale of the shock transition, i.e. on the scale of the ion gyro-radius, and as current $\mathbf{j}_y = -eNV_y\hat{\mathbf{y}}$ in Ampère's law

$$mNV_n \frac{dV_n}{dx} = -eNBV_y, \quad \frac{dB}{dx} = -\mu_0 eNV_y \quad (3.52)$$

Combining these equations yields the normal fluid velocity as function of the magnetic field and the initial bulk flow velocity V_1 at infinity

$$V_n = V_1 \left(1 - \frac{B^2 - B_1^2}{2\mu_0 m N_1 V_1^2} \right) = V_1 \left(1 - \frac{1}{2} \frac{V_A^2 - V_{A1}^2}{V_1^2} \right) \quad (3.53)$$

showing that the bulk velocity decreases from V_1 when the magnetic field B increases. The second term in the first parentheses is the difference in the ratio of magnetic pressures $B^2/2\mu_0$ at the location under observance and $B_1^2/2\mu_0$ at infinity upstream to the kinetic pressure mNV_1^2 at infinity. This is written in terms of Alfvén velocities in the second parentheses on the right.

With the help of this expression, E_y and Ampère's law we obtain

$$\mu_0 e N_1 V_y = \left[\frac{B^2 - B_1^2}{2\mu_0 m N_1 V_1} - 1 \right] \frac{dB}{dx} \quad (3.54)$$

which can be used to eliminate the velocity components and obtain an equation for the variation in the magnetic field given in the form of the energy conservation equation of a pseudo-particle of mass 1 and velocity B in the Sagdeev pseudo-potential $S(B)$:

$$\frac{1}{2} \left(\frac{dB}{dx} \right)^2 = -S(B), \quad S(B) = -\frac{(B - B_1)^2 [(B + B_1)^2 / 4\mu_0 m N_1 V_1^2 - 1]}{2\lambda_e^2 [1 - (B^2 - B_1^2) / 2\mu_0 m N_1 V_1^2]^2} \quad (3.55)$$

where $\lambda_e = c/\omega_{pe}$ is the electron inertial length (electron skin depth). The electron inertial length is the only length scale that appears in the above equation when we perform a dimensional analysis. It therefore turns out that the characteristic width, $\Delta \sim \lambda_e$, of the magnetosonic solitons is of the order of the electron skin depth. This can also be seen when for small amplitude disturbances the last expression is expanded with respect to B . Defining $b = (B - B_1)/B_1$, $b_m = (B_m - B_1)/B_1$, $\xi = x/\Delta$ yields to first order

$$\frac{db}{d\xi} \simeq \pm b(b_m - b)^{\frac{1}{2}} \quad \rightarrow \quad b \simeq \frac{4b_m}{(1 + \exp - |\xi|)^2}, \quad \Delta = \frac{\lambda_e}{\sqrt{b_m}} \quad (3.56)$$

giving the above scaling of the soliton width. In addition the inverse scaling of Δ with the maximum soliton amplitude b_m is reproduced.

As discussed before, $S < 0$ is required for solutions to exist. It is clear that in the absence of dissipation no shock can emerge from these stationary waves. They are solitary waves, stationary wave structures of finite spatial extensions and amplitude.

Returning to the original variables, the maximum soliton amplitude is obtained from $S(B_m) = 0$ as $B_m = B_1(2V_1/V_{A1} - 1)$ which together with Eq. (3.53) yields that $B_m < 3B_1$, and consequently the Alfvénic Mach number for the solitons to exist is $\mathcal{M}_A < 2$. The closer the Mach number approaches the maximum Mach number 2 the narrower the solitons become. This means that they steepen and, for Mach numbers $\mathcal{M} > 2$, will overturn and break, because the dispersion does no longer balance the nonlinearity. Solitons cannot exist anymore at Mach numbers such high.

When a way can be found to generate dissipation in the region occupied by the soliton, then a magnetosonic soliton can evolve into a shock wave. Sagdeev's idea [Sagdeev, 1966] was that this can happen when the soliton becomes large amplitude and narrow enough such that in the steep rise of its crest sufficient dissipation could be generated by anomalous collisions and anomalous friction. These anomalous collisions would generate sufficient entropy that the states on the two sides of the soliton would differ from each other and the flow across the soliton would irreversibly change. In this case the soliton would turn into a dissipative subcritical laminar shock. Such shocks will be discussed in Chapter 4. Formally we may, of course, model the dissipative effect by simply defining some "collision frequency ν " and introducing a "collision term" on the right of Eq. (3.55).

In order to do this we go one step back to the equation from which (3.55) has been obtained and add the collision term there:

$$\frac{d^2B}{dx^2} = -\frac{\partial S(B)}{\partial B} - \nu \frac{dB}{dx} \quad (3.57)$$

This equation is modelled exactly after the equation of a damped oscillator, where dB/dx is the velocity. We should note here that this modelling has not been justified yet and in fact is not justified by any of our arguments yet. It not only requires the proof of the existence of an anomalous collision frequency ν , it also requires the proof that from the kinetic equations containing the correlation terms an equation of the above structure can be derived.

Ignoring these objections and exploiting the analogy with the damped oscillator we may conclude from the equation (3.57) for the damped oscillator that the inclusion of anomalous “collisions” will dissipate the kinetic energy of motion of the pseudo-particle during its oscillation in the Sagdeev pseudo-potential $S(B)$ until the particle will finally come to rest at the bottom of the potential well. This is the case we have discussed earlier in connection with the Korteweg-de Vries-Burgers equation. It is drawn schematically in Figure 3.5 when a shock wave forms from the soliton. The value of B at minimum in the Sagdeev pseudo-potential where the pseudo-particle ultimately settles is the magnetic field level

$$B_2 = \frac{1}{2}B_1 [(8\beta_1 + 1)^{\frac{1}{2}} - 1] \quad (3.58)$$

far downstream of the shock that has formed in this dissipation process from the magnetosonic soliton. This value is determined by the upstream plasma- β value.

Taking this for granted, we can conclude that the damped oscillations the pseudo-particle performs on its damped downward path in the Sagdeev pseudo-potential are the spatially damped oscillations of the field $B(x)$ downstream of the shock. Moreover, the shock possesses an overshoot in B at shock position $B_2 < B_{ov} < B_m$ which is smaller than B_m but larger than B_2 . The existence of damped downstream oscillations is in agreement with the concave shape of the dispersion relation of magnetosonic waves which, for large k , become dispersive and approach the lower-hybrid branch. This can be seen directly from the dispersion relation for perpendicular ($\theta = 90^\circ$) propagating magnetosonic waves which reads

$$\omega_{ms} = V_A k_\perp (1 + k^2 \lambda_e^2)^{-\frac{1}{2}} \quad (3.59)$$

For $k^2 \lambda_e^2 \ll 1$ the wave has constant phase velocity and is non-dispersive, becoming gradually dispersive with increasing k when the effective phase velocity decreases. Hence, $\partial \omega / \partial k < 0$ and the dispersion is concave, for very large $k \lambda_e \gg 1$ approaching the lower hybrid frequency $\omega_{lh} = \sqrt{\omega_{ce} \omega_{ci}} = \omega_{ce} \sqrt{m_e / m_i}$ where it flattens out, as shown in Figure 3.10. At oblique angles $90^\circ - \delta \gg \sqrt{m_e / m_i} \approx 1/43$ the dispersion is inverted, and $\partial \omega / \partial k > 0$. Here the shorter waves run the soliton out and appear on the upstream side as spatially damped oscillations. Now their scale is the ion inertial length $\lambda_i = c / \omega_{pi}$. However, shocks with convex dispersion where the shorter waves outrun the soliton will not

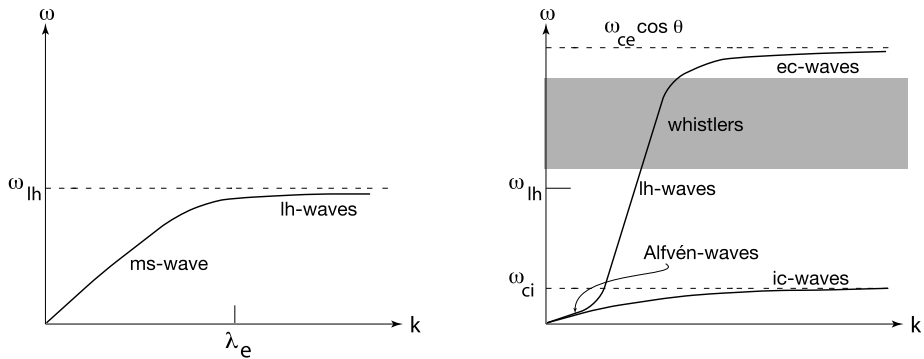


Figure 3.10: *Left:* Dispersion relation of perpendicular magnetosonic waves with concave dispersion at large k , where the wave effectively becomes a lower-hybrid wave with frequency close to the lower-hybrid frequency ω_{lh} . *Right:* Dispersion relation in cold plasma at nearly parallel propagation. The ion branch starts from the left-handed Alfvén modes and goes in resonance at the ion-cyclotron frequency ω_{ci} with concave dispersion. The electron branch starts from the right-handed Alfvén mode, has convex dispersion, passes through the lower hybrid mode into the whistler mode $\omega_{lh} \ll \omega \ll \omega_{ce}$, assumes concave dispersion, and finally goes in resonance at the electron-cyclotron frequency ω_{ce} , where it becomes the electron-cyclotron wave.

exhibit a sharp shock profile. Rather they will be oscillating shocks with smoothed out ramp.

The theory presented above applies to a low-frequency plasma of velocity $V_1 \gg v_e, v_i$ larger than the thermal velocities of the plasma components. When the relation between V_1 and the thermal velocities changes, one must take into account thermal effects. These change the nature of the solitons and shock substantially. These changes will be discussed in greater depth in Chapter 4.

3.4.3 Whistlers and Alfvén Shocks

Many waves are capable of dispersively evolving into solitons or other similar stationary wave pulses if only their dispersion relation allows it. Among those waves we here consider only two particular cases, whistlers in cold plasma and Alfvén waves in low but finite β conditions.

Whistler Solitons in Cold Plasma

We now return to the cold plasma dispersion relation including both kinds of particles, electrons and ions. For parallel propagation $\mathbf{k} = k_{\parallel} \hat{\mathbf{z}}$, where $\mathbf{B} \parallel \hat{\mathbf{z}}$, and the angle $\theta = 0$, the dispersion relations of linear plasma waves are

$$\frac{k^2 c^2}{\omega^2} = 1 - \frac{\omega_{pe}^2}{\omega(\omega \mp \omega_{ce})} - \frac{\omega_{pi}^2}{\omega(\omega \pm \omega_{ci})} \quad (3.60)$$

The frequencies in the nominators are the electron and ion plasma frequencies, $\omega_{pe, pi}$, respectively. In the denominators appear the electron and ion cyclotron frequencies, $\omega_{ce, ci}$

respectively. These relations describe right-hand and left-hand polarised waves according to the upper and lower signs. Basically two branches coming from the resonances in the denominators are described by this relation. Figure 3.10 on the right shows a plot of the two dispersion curves. For the evolution of shocks the most interesting part is the shaded whistler-mode dispersion relation.

The whistler dispersion relation is the (upper sign) electron part of the above dispersion relation. Neglecting the non-resonant ion contribution it reads

$$\frac{k^2 c^2}{\omega^2} = 1 + \frac{\omega_{pe}^2}{\omega(\omega_{ce} - \omega)} \quad (3.61)$$

with the second term being large because of the resonance in the denominator. Thus one can also neglect the 1 on the right finding that solutions exist only for $\omega < \omega_{ce}$ as is also seen in the above drawing. Then the dispersion relation becomes

$$k^2 \lambda_e^2 \simeq \omega / (\omega_{ce} - \omega) \quad \rightarrow \quad \omega \simeq \omega_{ce} \left(1 + \frac{1}{k^2 \lambda_e^2} \right)^{-1} \quad (3.62)$$

which exhibits its concave character confirming that short wavelengths whistlers will fall behind the main shock pulse. Nonlinear analysis of these waves goes back to *Montgomery* [1959], *Sagdeev* [1966], and *Kakutani* [1966] and is based on the fluid equations we used before. Let the plasma again be moving in x -direction antiparallel to the shock normal $\mathbf{n}$ and write the magnetic field in polar coordinates as $\mathbf{B} = B_\perp (0, \cos \theta, \sin \theta)$; then one again obtains the canonical Sagdeev form of the first integral of the equation of motion of a pseudo-particle at pseudo-position $B_\perp$ and pseudo-velocity $dB_\perp/dx$ as

$$\frac{1}{2} \left(\frac{dB_\perp}{dx} \right)^2 = -S(B_\perp) \quad (3.63)$$

The Sagdeev pseudo-potential is a complicated expression which simplifies considerably for a uniform upstream state. We introduce the normalised variables $b_\perp = B_\perp/B_{\perp m}$, $\xi = x/\lambda_e$ and $\beta_\perp = B_\perp^2/\mu_0 m N_1 V_1^2$ writing

$$\frac{1}{2} \left(\frac{db_\perp}{d\xi} \right)^2 = -S(b_\perp) = -\frac{1}{8\beta_{\perp m}} \frac{b_\perp(b_\perp^2 - 1)}{(1 - 1/2\beta_\perp)^2} \quad (3.64)$$

Solitons exist for $B_\perp \leq B_{\perp m}$ which is the maximum whistler soliton amplitude, and for $\beta_\perp > \frac{1}{2}$. For the maximum amplitude we have

$$B_{\perp m} < \sqrt{2\mu_0 m N_1 V_1^2} \sim \sqrt{m_i/m_e} B_1 \approx 43 B_1 \quad (3.65)$$

and the whistler soliton velocity $V_1 \gg V_{A1}$ yielding a soliton Mach number range of

$$\sqrt{m_i/4m_e} < \mathcal{M}_A < \sqrt{m_i/2m_e} \quad \rightarrow \quad 22 < \mathcal{M}_A < 30 \quad (3.66)$$

which identifies the whistler solitons as being high-Mach number solitons, indeed. In case they evolve into shocks, these shocks are high-Mach number as well. This might cause

other effects which have not been considered so far. Hence the present formal theory must be taken with caution in application to real problems of much lower Mach numbers. One of the neglected conditions is quasi-neutrality, which demands that $\epsilon_0 E / e N \Delta x \ll 1$. This leads to the further restriction on $B_\perp$ and the pulse width Δ

$$B_\perp / B_1 \ll (2\pi c^2 / V_{A1}^2)^{\frac{1}{4}} \sim 1.5c / V_{A1}, \quad \Delta \sim \Delta x \sim \lambda_i (B_1 / B_\perp) \quad (3.67)$$

where $\lambda_i = c / \omega_{pi}$ is the ion skin depth (ion inertial length). It follows that these whistler pulses should have quite large characteristic widths and, moreover, characteristic frequencies $\omega \sim \frac{1}{2} \omega_{lh} (B_\perp / B_1)$, far below the lower hybrid frequency. These properties identify the whistlers as right-handed (rotating clockwise along x) high-frequency Alfvén wave pulses. These low frequency whistler/high frequency whistler-Alfvén waves can indeed be excited by a cold shock-reflected ion beam as will be shown below in the section on ion beam instabilities.

In view of our remarks we do not give rigorous derivations of these approximate formulae. It will turn out later when discussing numerical simulations that whistlers do indeed occur at shocks and have been observed early on as well in laboratory experiments on collisionless shocks [e.g., *Decker & Robson, 1972*] that, however, other fluctuations driven by ion reflection are of greater importance in structuring supercritical Mach number shocks. Still it is highly probable that whistlers are excited in shock waves as the conditions will be in favour of them when sufficient free energy is available in the shock front because of several reasons, one of them electron heating in the perpendicular direction [early observations suggested their presence near shocks in space, see, *Rodriguez & Gurnett, 1975*]. In this case whistlers become very important for producing dissipation via a short-wavelength instability, called *decay instability*, which had been predicted by *Galeev & Karpman [1963]* and for which evidence has been found in the above laboratory observations by *Decker & Robson [1972]*.

Alfvén Solitons at Finite- β

Alfvén waves are non-dispersive. However, when the plasma temperature increases, dispersion in the direction perpendicular to the magnetic field sets on. From the general dispersion relation in the very low-frequency limit one then obtains [e.g., *Baumjohann & Treumann, 1996*] for the frequency of the Alfvén wave

$$\omega^2(k_\parallel, k_\perp) = k_\parallel^2 V_A^2 \frac{1 + k_\perp^2 r_{ci}^2}{1 + k_\perp^2 \lambda_{pe}^2} \quad (3.68)$$

where the ion gyro-radius is slightly modified with temperature according to $r_{ci}^2 \rightarrow r_{ci}^2 (\frac{3}{4} + T_e / T_i)$, and $V_A^2 = B^2 / \mu_0 m N$ is the square of the Alfvén speed. This dispersion relation describes two kinds of Alfvén waves depending on $k_\perp \sim r_{ci}^{-1}$ or $k_\perp \sim \lambda_e^{-1}$. These modes become important when the plasma- $\beta < 1$. The phase velocities of the two modes together with the ordinary Alfvén wave are shown in Figure 3.11. The modes differ in their dispersive properties from ordinary Alfvén waves as they propagate oblique to the magnetic field,

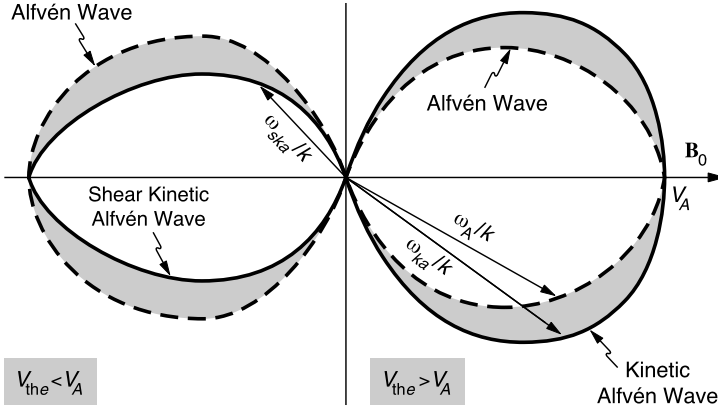


Figure 3.11: Phase velocity diagram of the three Alfvén wave modes in the $(\mathbf{B}, \mathbf{k})$ -plane of the magnetic field and wave vector for the two cases when the alfvén velocity is larger *left* or smaller *right* than the electron thermal speed v_{the} . The ordinary Alfvén wave describes a circle in this plane. The two other phase velocities describe deformed curves.

i.e. the wave energy propagates under an angle to the magnetic field for $k_{\perp}$ is independent of $k_{\parallel}$.

At finite temperatures $1 > \beta > m_e/m_i$ the wave is called *kinetic* Alfvén wave. Its perpendicular wavelength becomes the order of the ion gyro-radius, and the phase speed parallel to the magnetic field increases. At low temperatures $\beta < m_e/m_i$ the wave is called *shear* or better *inertial* Alfvén wave with perpendicular wavelength comparable to the electron skin depth, and the effective parallel phase velocity decreases. In terms of prospective Alfvénic shocks this means that a kinetic Alfvén shock in direction parallel to the external field will support oscillations upstream of the shock ramp, while an inertial Alfvén shock will support downstream oscillations and thus possess a sharp shock ramp.

The dispersive properties of both kinetic Alfvén modes enable the existence of stationary wave pulses. These have been discovered first by Hasegawa & Mima [1976] for the *kinetic* mode. For the low Alfvén frequencies quasi-neutrality is a good assumption. Moreover, the electrons have time enough to behave Boltzmann-like. Hence with the parallel electric potential $\phi_{\parallel}$ we have

$$N_e = N_0 \exp(e\phi_{\parallel}/T_e) \quad (3.69)$$

In the perpendicular direction we use the electric potential $\phi_{\perp}$. Maxwell's equations, the nonlinear ion momentum conservation, and Poisson's equation then reduce to

$$\begin{aligned} \frac{\partial B_{\perp}}{\partial t} &= \nabla_{\perp} \nabla_{\parallel} (\phi_{\perp} - \phi_{\parallel}) \\ \nabla_{\perp}^2 \nabla_{\parallel}^2 (\phi_{\perp} - \phi_{\parallel}) &= \mu_0 \nabla_{\parallel} \frac{\partial j_{\parallel}}{\partial t} \\ \frac{\partial N_i}{\partial t} &= \frac{1}{B_0 \omega_{ci}} \nabla_{\perp} \cdot \left(N_i \nabla_{\perp} \frac{\partial \phi_{\perp}}{\partial t} \right) \end{aligned} \quad (3.70)$$

The field-aligned current $j_{\parallel}$ is carried by the hot electron component, such that its divergence is given by $\nabla_{\parallel} j_{\parallel} = e(\partial N_e / \partial t)$. Introducing dimensionless variables $\xi = x\omega_{ci}/c_{ia}$, $\zeta = z\omega_{pi}/c$, $\tau = \omega_{ci}t$, $N' = N/N_0$ and measuring the potentials in units of e/T_e , then transforming to a comoving coordinate system $\eta = \kappa_{\perp}\xi + \kappa_{\parallel}\zeta - \tau$, the whole system of equations is reduced to the nonlinear equation

$$\kappa_{\perp}^2 \kappa_{\parallel}^2 N' \frac{d^2 \ln N'}{d\eta^2} = (1 - N') (N' - \kappa_{\parallel}^2) \quad (3.71)$$

which is in the form suited for the Sagdeev pseudo-potential method, yielding $(dN'/d\eta)^2 = -S(N', \kappa_{\parallel}, \kappa_{\perp})$ with

$$S(N', \kappa_{\parallel}, \kappa_{\perp}) = -\frac{2N'}{\kappa_{\parallel}^2 \kappa_{\perp}^2} [(1 - N') (N' + \kappa_{\parallel}^2) + (1 + \kappa_{\parallel}^2) N' \ln N'] < 0 \quad (3.72)$$

for soliton solutions to exist. These solutions give the density as function of the linear coordinate η . Interestingly, there are solutions which are dilutions and solutions which are compressions. The condition for existence of soliton solutions is independent of the perpendicular wavenumber. Hence it is the parallel electric field that is responsible for the formation of solitons and balance of the nonlinear steeping. Solitons form only in parallel direction with the magnetic field being inclined to the soliton which in the perpendicular direction is flat. When such a soliton attains dissipation and turns into a shock, it will become a quasi-parallel shock preceded by damped upstream waves that have outrun the shock ramp. That the shock will be quasi-parallel can be easily seen from the fact that $k_{\perp} \gg k_{\parallel}$, the shock front will be perpendicular to the external field, and therefore $b_{\perp} \ll b_{\parallel}$ as required for a quasi-parallel shock.

Inspection of Eq. (3.72) shows that the Sagdeev potential vanishes at $N' = 0$, $N' = 1$, and N'_m . Compressive (rarefaction) solitons occur at $N' > 1$ ($N' < 1$). Only compressive solitons are of interest in shock formation. The maximum amplitude N'_m of compressive solitons follows from setting the bracket to zero. It is approximately given by the solution of $N' + \kappa_{\parallel}^2 \approx (1 + \kappa_{\parallel}^2) \ln N'$ which, for $\kappa_{\parallel}^2 = 1$ is $N'_m \approx 3$. The minimum of $S(N')$ for compressive solitons is found by taking the derivative of the bracket and putting it to zero. Setting $N' = 1 + n'$ and expanding the logarithm one finds the minimum of the Sagdeev pseudo-potential trough at

$$N'_{KAS} = 1 + n' \approx \frac{1 + 3\kappa_{\parallel}^2}{1 + \kappa_{\parallel}^2} < 3 \quad (3.73)$$

This is the prospective maximum amplitude of a kinetic Alfvén shock (KAS) evolving in the presence of dissipation from a compressive kinetic Alfvén soliton. The compression at a KAS ramp is thus limited to a factor < 3 . Finally, the value of the Sagdeev pseudo-potential at its absolute minimum gives the steepest gradient of the density in the shock ramp, $(dN'/d\eta)_{\max} = |S(N'_{KAS})|$ at the turning point of the density in the ramp. This ratio provides an estimate of the width of the KAS-shock pulse

$$\Delta_{KAS} \sim N'_{KAS} / |S(N'_{KAS})| > |3 - \kappa_{\parallel}^2|^{-1} \quad (3.74)$$

in normalised units (where in the final estimate at the right we have used the maximum shock amplitude). It depends on the given values of $\kappa_{\parallel}$ while being independent of $\kappa_{\perp}$, in agreement with the fact that only values parallel to the field will affect the shock structure.

We will not investigate the case of *inertial* Alfvén wave solitons. In the heliosphere these shocks occur only in the auroral zones of strongly magnetised planets and, possibly, also deep in the solar corona in the region of very strong magnetic field and comparably small temperature only. The related solitons have been discussed elsewhere [Treumann *et al*, 1990; Berthomier *et al*, 1999]. Shocks forming there should always be rarefactions, containing very little plasma and having highly oscillating wakes on their downstream side, containing short wavelength inertial Alfvén waves which have been left behind the faster shock ramp. In fact, such very diluted plasmas have indeed been inferred in the auroral zones of planets like Earth and Jupiter. They contain large electric potential drops; however the interpretation of these as shock potentials is quite unusual and not generally accepted.

Remarks on the Generation of Dissipation

Sagdeev [1966] gave an idea of how dissipation can be generated in a shock ramp, realising that shocks must contain a – large scale – electric potential drop (or its equivalent as, for instance, an equivalent electric field corresponding to the shock ramp density gradient) at which low energy ions will be reflected. (Note that electrons will instead become accelerated by this potential across the shock.) No matter how few those ions are, they will return into the upstream medium, where they (as Sagdeev [1966] had already noted) become accelerated tangentially along the (perpendicular) shock surface by the upstream convection electric field until gaining enough kinetic energy to overcome the shock potential barrier, passing the shock, and escaping downstream. These ions form a current in front of the shock that carries free energy and will ultimately become unstable with respect to the *two-stream* instability, scattering upstream electrons and in this way cause dissipation. The physical mechanism of this process will be discussed later in this chapter in the section on anomalous transport. First we need to be informed about the instabilities and waves that are relevant with respect to shock formation. To these mechanisms we will continuously return when discussing the different types of shocks and the corresponding numerical simulations. Without them neither the existence nor the structure of shocks in collisionless plasmas can be understood. The models presented so far cannot give more than hints in which direction one has to pursue. Shock physics is too complicated for analytical theory.

3.4.4 Instabilities

An instability arises when an active medium like a warm plasma reacts to the presence of an amount of available free energy that exceeds the thermal energy. Since any form of available free energy keeps the plasma away from thermodynamic equilibrium, restoration of equilibrium becomes necessary. This is most easily done by exciting fluctuations out of the thermal fluctuation background and let them grow to amplitudes large enough for either causing dissipation or transporting the energy away to a location where it can be

dissipated by other processes. If no fluctuation is injected from externally, the instability selects one or more particular eigenmodes out of the thermal fluctuations which as a consequence causes a wave to start growing and propagating across the plasma. The selection of the unstable frequency and wavelength ranges is usually ruled by resonance of the fluctuation with the number of particles which carry the free energy, with the fastest growing fluctuation winning the race for consuming the free energy. However, other ways of exciting waves are also possible when several waves are sharing the amount of free energy. In all cases the waves that survive are those with the fastest growth and least damping. [Theoretical overviews of instabilities and dissipation related to shocks have been given by *Galeev, 1976; Sagdeev, 1979; Wu et al, 1984; Papadopoulos, 1985; Winske, 1985b.*]

Since most instabilities follow this recipe of growing out of the thermal background they start as infinitesimal disturbances which can be described by linear dispersion theory. Hence instability theory in plasma can be based on the dispersion relation Eq. (3.46) and for weakly growing waves on the expression for the growth rate (3.50). Plane linear wave modes possess a phase factor $\exp i(\mathbf{k} \cdot \mathbf{x} - \omega t)$. It is clear that instability sets on whenever the imaginary part of the frequency becomes positive; $\gamma(\omega, \mathbf{k}, \dots) > 0$ is the growth rate of the instability of the particular wave mode that becomes unstable.

Another view is that a wave grows along its path of propagation. In this case the growth is spatial and, at real frequency ω , is described by the imaginary part of the wave number $k_i < 0$, which must be negative for a wave growing when propagating into direction $+\mathbf{x}$. For propagating waves there is a natural relation between the temporal and spatial growth rates that is given by $k_i = -\gamma/(|\partial\omega/\partial k|)$ (here written just for one component of the spatial growth vector and given in terms of the group velocity $\partial\omega/\partial k$ of the wave in this particular direction).

In the context of collisionless shocks the instabilities of interest can be divided into two classes. The first class contains those waves which can grow themselves to become a shock. It is clear that these waves will be of low frequency and comparably large scale because otherwise they would not evolve into a large macroscopic shock. We have already discussed a few candidates and their nonlinear evolution in the previous sections, among them magnetosonic, Alfvénic and whistler modes. In this section we will investigate a number of waves which form secondarily after an initial seed shock ramp has grown in some way out of one of these wave modes, these are ion modes which have now been identified to be responsible for structuring, shaping and reforming the shock. In fact real oblique shocks – which are the main class of shocks in interplanetary space and probably in all space and astrophysical objects – cannot survive without the presence of these ion waves which can therefore be considered the wave modes that really produce shocks in a process of taking and giving between shock and waves.

The second class are waves that accompany the shock and provide anomalous transport coefficients like anomalous collision frequencies, friction coefficients, heat conductivity and viscosity. These waves are also important for the shock as they contribute to entropy generation and dissipation. However, they are not primary in the sense that they are not shock-forming waves.

Among them there is another group that only carries away energy and information from the shock. These are high-frequency waves, mostly electrostatic in nature, produced

by electrons, or when electromagnetic they are in the free-space radiation modes. In the latter case they carry the information from remote objects as radiation in various modes, radio or x-ray to Earth, informing of the existence of a shock. In interplanetary space it is only radio waves which fall into this group as the radiation measure of the heliospheric shocks is too small to map them into x-rays. These groups of waves we will briefly mention below; they will however play a more important role in this text when discussing measurements and observations of the various types of shock waves that are encountered in the heliosphere.

Ion-Beam Driven Instabilities – $\omega \lesssim \omega_{ci}$

The shock waves in the heliosphere are magnetised. As long as we are interested in their formation and properties we can restrict to low frequency electromagnetic waves in warm plasma. Such waves are excited by plasma streams or kinetic anisotropies in one or the other way. The simplest instability known which distorts the magnetic field by exciting Alfvén waves that are propagating along the magnetic field is the firehose mode.

Firehose Mode. The firehose mode is the result of a pressure (or temperature) anisotropy in plasma with the parallel pressure $P_{\parallel}$ exceeding the perpendicular $P_{\perp}$ and magnetic $B^2/2\mu_0$ pressures. *Sagdeev [1966]* gave a simple intuitive explanation of this instability based on the insight that the parallel thermal motion of the adiabatic magnetised ions along the magnetic field exerts a centrifugal force on the field lines. When this force exceeds the restoring forces of the magnetic pressure and perpendicular plasma pressure, the centrifugal force wins and a small excursion of the magnetic field starts growing and propagates as a wave along the magnetic flux tube like on a string. The condition for instability is

$$P_{\parallel} - P_{\perp} > B^2/\mu_0 \quad (3.75)$$

Since the pressure anisotropy on the left means that there is an excess in parallel energy in the plasma, the plasma possesses free energy which by the instability is fed into the excitation of Alfvén waves with frequencies $\omega_A \ll \omega_{ci}$, transported away with Alfvén speed and ultimately dissipated in some way – as expected. The waves excited are ordinary Alfvén waves, however, and not suited for shock formation. Below we will once more encounter this mode in discussing ion beam instabilities.

Kinetic Alfvén Waves. Excitation of kinetic Alfvén waves requires $\beta < 1$ and a different process. In the solar wind the β -condition is barely satisfied except possibly in the very strong coronal magnetic fields or locally (possibly in Corotating Interaction Region boundaries when the magnetic field may become compressed without just forming a shock). Kinetic Alfvén waves possess a finite electric field component parallel to the magnetic field which can accelerate electrons. However, the inverse mechanism is also possible that electrons moving along the magnetic field in the opposite direction, become retarded by this field component and feed their energy into the kinetic Alfvén wave. A process similar to this has been suggested by *Hasegawa [1979]* in different context for bouncing electrons

in a locally inhomogeneous magnetic field represented as $B(z) = B_0(1 + az^2)$. The electron beam conserves the magnetic moment when moving along the magnetic field, interacting adiabatically with the parallel wave electric field for long wavelength

$$E_{\parallel} = \frac{k_{\perp}^2 r_{ci}^2}{1 + k_{\perp}^2 r_{ci}^2} \left| \frac{\omega B_{\perp}}{k_{\perp}} \right| \quad (3.76)$$

The condition that the electrons form a beam is that at the resonance with the wave the derivative of the electron distribution with respect to the resonant parallel electron energy of motion $\varepsilon_{\parallel} = m_e \omega_b^2 / 2a$ is positive, $\partial F_e / \partial \varepsilon_{\parallel} > 0$. Here the square of the bounce frequency is $\omega_b^2 = 2\mu \omega_{ce} / m_e$, and $\mu = m_e V_{\parallel}^2 / 2B$ is the magnetic moment. One can calculate the growth rate of the wave as

$$\gamma_{\text{KAW}} \simeq (k_{\perp} \lambda_e)^{-2} (\omega^4 / a k_{\parallel}^3 v_e^3), \quad \omega = k_{\parallel} V_A \quad (3.77)$$

Thus, an electron beam moving along an inhomogeneous magnetic field can excite kinetic Alfvén waves.

Probably more important than this is, however, the interaction of ions which are reflected from a solitary pulse and move back upstream ahead of the pulse, as had been suggested by *Sagdeev* [1966]. The reflected ions will represent a beam that is moving against the initial plasma inflow which by itself is another ion beam neutralised by the comoving electrons. This configuration leads to a ion beam-ion beam interaction and should cause an instability because free energy is present in the two counter-streaming beams. The various instabilities this process may cause have been reviewed by *Gary* [1993].

Kinetic Growth Rate. Before coming to discussing the relevant instabilities we should briefly mention the waves which can be driven by them. We already noted that in thermal plasma most waves will rest in thermal fluctuations. Once a wave which is an eigenmode of the plasma is injected it will experience thermal damping until it disappears in the background fluctuations. Hence, a wave that is assumed to grow must overcome this damping which for propagation parallel to the average magnetic field $\mathbf{B}_0$ in a uniform plasma is given by

$$\gamma(\omega, \mathbf{k}) \simeq \sqrt{\frac{\pi}{2}} \sum_s \frac{\omega_s^2}{2\omega v_s} \left(V_s - \frac{\omega}{k_{\parallel}} \right) \exp \left[\frac{(\omega \pm \omega_{cs} - k_{\parallel} V_s)^2}{2k_{\parallel}^2 v_s^2} \right] \quad (3.78)$$

where the index s identifies the species, $v_s = \sqrt{2T_s/m_s}$ is the thermal velocity of species s , V_s its average parallel bulk drift velocity, $\omega_s \equiv \omega_{ps}$ its plasma frequency, and we have dropped the index r (for real) at ω which in this expression is understood as real anyway. Note that we are going to take into account several different ion species and thus need an extra index to distinguish between them all. The simplified cold dispersion relation is

$$\omega^2 - k^2 c^2 - \sum_s \frac{\omega_s^2 (\omega - k_{\parallel} V_s)}{\omega - k_{\parallel} V_s \pm \omega_{cs}} = 0 \quad (3.79)$$

which determines the approximate real frequency. However, when thermal effects are included, then there is no way to avoid the numerical solution of the full kinetic dispersion relation. In the following we will generally refer to such numerical solutions.

It is obvious from the expression for the damping rate, that for sufficiently large average drift velocities V_s of species s larger than the phase velocity $\omega/k_{\parallel}$ of the wave this particular species contributes a positive term to the damping rate γ which, when large enough can dominate the entire damping rate. In this way streaming is one way to cause instability. In the absence of streaming γ is independent of V_s and is negative for a thermally isotropic plasma. In the presence of a temperature anisotropy, however, this may change as we have seen for the firehose mode. The above damping rate (3.78) does not account for thermal anisotropy which is, however, simple matter to include [cf., e.g. Gary, 1993].

We note finally that γ is a resonant damping/growth rate yielding resonant instability driven by small groups of resonant particles. The firehose mode is non-resonant since all particles contribute to it. Generally most non-resonant instabilities can only be found by solving the full dispersion relation numerically.

At low frequencies it suffices for our purposes of understanding shock physics to deal with a three-component plasma consisting of two ion species and one neutralising electron component which we assume to follow a Maxwellian (thermal) velocity distribution. Moreover, we assume that the drifting ion components are Maxwellians as well. In conformity with the above remarks on a resonant instability we assume that the dominant ion component has large density $N_i \gg N_b$, and the second component represents a weak fast beam of density N_b propagating on the ion-electron background with velocity $V_b \gg V_i \approx 0$. Following Gary [1993] it is convenient to distinguish the three regimes: cool beams ($0 < v_b < V_b$), warm beams ($v_b \sim V_b$), and hot beams ($v_b \gg V_b$). Figure 3.12 shows the beam configurations for these three cases and the location of the wave resonances respectively the position of the unstable frequencies.

Cool Ion Beam: Right-Hand Instability. Assume that the ion beam is thermally isotropic and cool in the above sense, i.e. its velocity relative to the bulk plasma is faster than its thermal speed. In this case a right-handed resonant instability occurs. In the absence of a beam $V_b = 0$ the parallel propagating mode is a right-circularly polarised magnetosonic wave propagating on the lowest frequency whistler dispersion branch with $\omega \approx k_{\parallel} V_A$. In presence of a drift this wave becomes unstable, and the fastest growing frequency is at frequency $\omega \simeq k_{\parallel} V_b - \omega_{ci}$. This mode propagates parallel to the beam, because $\omega > 0$, $k_{\parallel} > 0$, and $V_b > 0$. The numerical solution of this instability for densities $0.01 \lesssim N_b/N_i \lesssim 0.10$ at the wave-number $k_{\parallel}$ of fastest growth rate identifies a growth rate of the order of the wave frequency $\gamma \sim \omega$ and

$$\gamma_m \simeq \omega_{ci} (N_b/2N_e)^{\frac{1}{3}} \quad (3.80)$$

for the maximum growth rate γ_m , where $N_e = N_i + N_b$ is the total density from quasi-neutrality. This instability drives waves propagating together with the beam in the direction of the ion beam on the plasma background which has been assumed at rest. If applied,

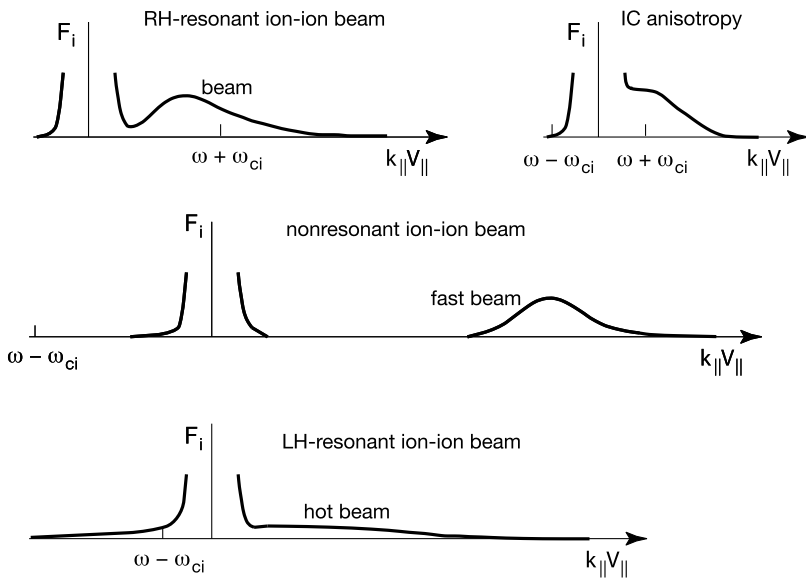


Figure 3.12: The three cases of ion beam-plasma interaction and the location of the unstable frequencies. Shown is the parallel (reduced) distribution function $F_i(k_{\parallel}v_{\parallel})$, where for simplicity the (constant) parallel wave-number $k_{\parallel}$ has been included into the argument. Right handed resonant modes (RH) are excited by a cool not too fast beam. When the beam is too fast the interaction becomes nonresonant. When the beam is hot, a resonant left hand mode (LH) is excited. In addition the effect of temperature anisotropy is shown when a plateau forms on the distribution function [after Gary, 1993].

for instance, to shock reflected ions then for 2% reflected ions the maximum growth rate is $\sim 0.2\omega_{ci}$, and $V_b \sim 1.2\omega_{ci}/k_{\parallel}$, and $k_{\parallel} \sim 0.2\omega - ci/V_A$ which gives $V_b \sim 6V_A$. In the solar wind the Alfvén velocity is about $V_A \approx 30$ km/s. Hence the velocity difference between shock reflected ions and solar wind along the magnetic field should be roughly ~ 180 km/s. The thermal velocity of the ion beam must thus be substantially less than this value, corresponding to a thermal beam energy less than $T_b \ll 100$ eV which in the solar wind, for instance, is satisfied near the tangential field line. The solar wind travels at 300–1200 km/s. Complete reflection should produce difference speeds twice these values. The above value is thus not unreasonable for travelling shocks, but for bow shock reflected ions applies to the quasi-perpendicular portion of the bow shock only. We may thus conclude that this wave mode could be excited in the solar wind by shock reflected ion beams near quasi-perpendicular shocks.

Warm Ion Beam: Left-Hand Instability. The above instability is present when the ion beam is rather cold. When the temperature of the ion beam increases and the background ions remain to be cold, then beam ions appear on the negative velocity side of the bulk ion distribution and go into resonance there with the left-hand polarised ion-Alfvén wave. Their maximum growth rate is a fraction of the growth rate of the right-hand low

frequency whistler mode. Nevertheless it can excite the Alfvén-ion cyclotron wave which also propagates parallel to the beam. For this instability the beam velocity must exceed the Alfvén speed $V_b > V_A$.

At oblique propagation both the right and left hand instabilities have smaller growth rates. But interestingly, it has been shown [Goldstein *et al*, 1985] that the fastest growing modes then appear for oblique $\mathbf{k}$ and harmonics of the ion cyclotron frequency $\omega \sim n\omega_{ci}$, with $n = 1, 2, \dots$.

Nonresonant Ion Instability: Firehose Mode. When the ion beam is fast and cold it does not go into resonance because its velocity is too high. In this case all ions participate in a nonresonant instability which in fact is a thermal firehose mode where the ion beam has sufficient energy to shake the field line. This mode propagates antiparallel to the ion beam, has small phase speed and negative helicity. This mode has large growth rate for large N_b/N_e and V_b/V_A simply because then there are many beam ions and the centrifugal force is large while the beam velocity lies outside any resonant wave speed. It is trivial that this instability becomes stronger when the ion beam is composed of heavier ions as the larger mass of these increases the centrifugal force effect.

Electron Instabilities and Radiation – $\omega \sim \omega_{pe}$

Other than ion beam excited instabilities electron-beam instabilities are not involved in direct shock formation (unless the electron beams are highly relativistic which in the entire heliosphere is not the case; we do not investigate relativistic shocks in this text anyway). The reason is that the frequencies of electron instabilities are high. However, just because of this reason they are crucial in anomalous transport being responsible for anomalous collision frequencies and high frequency field fluctuations. The reason is that the high frequency waves lead to energy loss of the electrons retarding them while for the heavier ions they represent a fluctuating background scattering them. In this way high frequency waves may contribute to the basic dissipation in shocks even though this dissipation for supercritical shocks will not be sufficient to maintain a collisionless shock or even to create a shock under collisionless conditions. This is also easy to understand intuitively, because the waves need time to be created and to reach a substantial amplitude. This time in a fast stream is longer than the time the stream needs to cross the shock. So waves will not accumulate there; rather the fast stream will have convected them downstream long before they have reached substantial amplitudes for becoming important in scattering.

When we are going to discuss electromagnetic waves which can be excited by electrons we also must keep in mind that such waves can propagate only when there is an electromagnetic dispersion branch in the plasma under consideration. These electromagnetic branches in $(\omega, \mathbf{k})$ -space are located at frequencies below the electron cyclotron frequency ω_{ce} . The corresponding branch is the whistler mode branch. Electrons will (under conditions prevailing at shocks) in general not be able to excite electromagnetic modes at higher frequencies than ω_{ce} . We have seen before that ion beams have been able to excite whistlers at low frequencies but above the ion-cyclotron frequency. This was possible only because of the presence of the high frequency electron whistler branch as a channel for

wave propagation. In fact, ion cyclotron waves can for this reason also have higher electromagnetic harmonics. For electrons under conditions in the shock environment it is generally not possible to directly excite free space electromagnetic modes. Electromagnetic waves excited by them propagate on the whistler branch or its low frequency Alfvénic extension.² Electron beams thus excite electromagnetic whistlers and right-handed Alfvén waves. They also excite a variety of electrostatic emissions which we will mention later as well.

Whistlers. Whistlers can be driven in two ways, either by an electron temperature anisotropy [Kennel & Petschek, 1966], or by electron beams (or heat fluxes) [Gary, 1993]. In the former case the conditions for instability are that the perpendicular electron temperature $T_{e\perp} > T_{e\parallel}$ exceeds the parallel electron temperature, and that the parallel energy of the resonant electrons $\mathcal{E}_{\parallel} = \frac{1}{2}m_e V_{\parallel}^2 > B^2/2\mu_0 N_e$ exceeds the magnetic energy per electron.

Gary [1993] has investigated the more relevant case of whistler excitation by an electron beam. He finds from numerical solution of the full dispersion relation including an electron beam in parallel motion that with increasing beam velocity V_b the real frequency of the unstable whistler decreases, i.e. the unstably excited whistler shifts to lower frequencies on the whistler branch while remaining in the whistler range $\omega_{ci} < \omega < \omega_{ce}$. Both, the background electrons and beam electrons contribute resonantly. The most important finding is that the whistler mode for sufficiently large $\beta_i \sim 1$ (which means low magnetic field), N_b/N_e and T_b/T_e has the lowest beam velocity threshold when compared with the electrostatic electron beam instabilities as shown in Figure 3.13. This finding implies that in a relatively high- β plasma a moderately dense electron beam will first excite whistler waves. In the shock environment the conditions for excitation of whistlers should thus be favourable whenever an electron beam propagates across the plasma along the relatively weak magnetic field. The electrons in resonance satisfy $V_{\parallel} = (\omega - \omega_{ce})/k_{\parallel}$ and, because $\omega \ll \omega_{ce}$ the resonant electrons move in the direction opposite to the beam. Enhancing the beam temperature increases the number of resonant electrons thus feeding the instability.

On the other hand, increasing the beam speed shifts the particles out of resonance and decreases the instability. Hence for a given beam temperature the whistler instability has a maximum growth rate a few times the ion cyclotron frequency.

3.4.5 “Transport Ratios”

Measured wave spectra are complex and opaque, i.e. it is very difficult from an inspection of their shape to identify the wave modes that are present in the plasma volume under investigation. In some rare clean cases one can conclude from the observation of a particular maximum peak in the wave power or the observation of only one single field component which wave has been detected. In the general case of broad spectra or mixed

²In fact there is one exception to this statement. There exist free space electromagnetic mode branches (radiation) above and even below the electron cyclotron frequency on which electron excited electromagnetic waves could in principle propagate. The mechanism to excite them is the electron cyclotron maser instability [for a recent review see, e.g. Treumann, 2006], which is a very particular instability that becomes awakened under conditions which to our knowledge are barely satisfied in the non-relativistic shock environment.

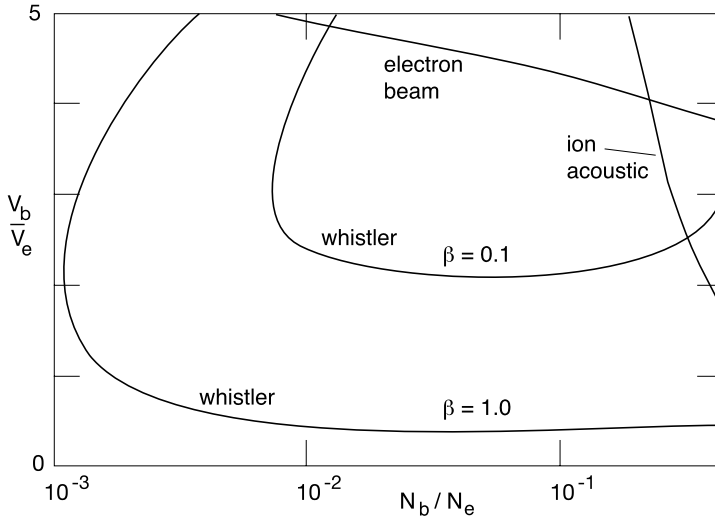


Figure 3.13: The regions of instability of the electron beam excited whistler mode in density and beam velocity space for two different β compared to the ion acoustic and electron beam modes. Instability is above the curves. The whistler instability has the lowest threshold in this parameter range [after Gary, 1993].

spectral peaks and various field components lacking clear dominance of one field component it becomes nearly impossible to decide about the waves. Clearly, when the spectrum is shapeless power-law one, in most cases, is dealing with developed turbulence or localised structures in which case it makes no sense to distinguish and search for single modes. Then one must seek shelter among the well developed methods of analysing turbulence.

On the other hand, if the spectra indicate the presence of single waves, one would like to have some quantities at hand which help identifying which modes one is dealing with. It would be helpful if one could measure simultaneously both, the wave frequency and the wave number spectra. This is possible, however, only with sophisticated multi-spacecraft constellations. And even then only in the rarest cases the determination of the dispersion relation from experiment will be possible.

In application of these theoretical arguments to real observations one therefore has defined some quantities, called “transport ratios”, which have turned out to be quite valuable in helping identifying some of the wave modes. Such transport ratios for electromagnetic waves have been given by Gary [1993]. Below we list the for our purposes most interesting of them.

Polarisation. The polarisation of a wave magnetic field with respect to wave number $\mathbf{k}$ is given by

$$P = ib_S/b_A \quad (3.81)$$

where b_S , b_A are the components of the magnetic fluctuation field $\mathbf{b}$ in the directions $\mathbf{S}$, $\mathbf{A}$ of magnetosonic and Alfvén waves, respectively, i.e. the vector $\mathbf{A} = \mathbf{k} \times \mathbf{B}_0$ is perpendicular

to the wave vector and the ambient magnetic field, while the vector $\mathbf{S}$ is perpendicular to $\mathbf{k}$ (because of the vanishing divergence $\nabla \cdot \mathbf{b} = 0$) in the plane $(\mathbf{k}, \mathbf{B}_0)$. The waves are more magnetosonic or more Alfvénic whether $|P| > 1$ or $|P| < 1$, respectively. For $\text{Re } P > 0$ (< 0) the waves are right-hand (left-hand) polarised.

Compression. The magnetic compression of the wave measures the relative variation in the parallel magnetic fluctuation field

$$C_B = \langle b_{\parallel}^2 \rangle / \langle |\mathbf{b}|^2 \rangle \quad (3.82)$$

where the fluctuations are taken at a given pair $(\omega, \mathbf{k})$. The angular brackets $\langle ab \rangle$ mean taking the real part of the correlation function of the two bracketed quantities.

Parallel Compressibility. This ratio together with the compression ratio provides a tool for estimating how compressive a wave is. It is defined for species s as

$$C_{\parallel s} = \frac{B_0^2}{\langle b_{\parallel}^2 \rangle} \frac{\langle b_{\parallel} \Delta N_s \rangle}{N_s B_0} \quad (3.83)$$

Non-coplanarity Ratio. This ratio measures the fluctuating field component out of the plane $(\mathbf{k}, \mathbf{B}_0)$, and is given by

$$C_c = \langle |b_A|^2 \rangle / \langle |\mathbf{b}|^2 \rangle \quad (3.84)$$

Alfvén Ratio. Defining $\Delta \mathbf{V}_A = \mathbf{b} / \sqrt{\mu_0 m_i N}$, where N is the total plasma density, the Alfvén ratio is defined as

$$R_{As} = \langle |\Delta \mathbf{V}_s|^2 \rangle / \langle |\Delta \mathbf{V}_A|^2 \rangle \quad (3.85)$$

Here $\Delta \mathbf{V}_s$ is the flow velocity of species s . An Alfvén wave has $\Delta \mathbf{V}_i = \pm \mathbf{V}_A \mathbf{b} / B_0$, and its own Alfvén ratio is $R_{Ai} = 1$. The Alfvén ratio thus measures the fraction of Alfvén waves contained in the near-zero frequency fluctuations.

Cross-Helicity. Helicity of a wave is another identifier of the wave mode, it is in particular useful for determining the direction of propagation of the wave by considering its sign. One has therefore defined a ratio which provides a measure of it:

$$\frac{1}{2} H_{cs} = \langle \mathbf{b} \cdot \Delta \mathbf{V}_s \rangle / [\langle |\Delta \mathbf{V}_s|^2 \rangle + \langle |\mathbf{b}|^2 \rangle] \quad (3.86)$$

Parallel propagating Alfvén waves have helicity $H_{ci} = -1$. Fast magnetosonic nearly parallel propagating modes have also $|H_{ci}| \sim 1$, and $H_{ci} = 0$ for perpendicular propagation.

3.5 Anomalous Transport

It has been mentioned several times that the evolution of shocks requires the generation of some kind of dissipation. Under the conditions of non-collisionality the generation of dissipation must be intrinsic to the plasma. In fact, assume that a source of free energy is switched on in the plasma. It is then quite natural to imagine that this available free energy will act on the plasma in a way to dissipate itself, distribute itself all over the plasma and ultimately transform itself into heat, entropy and create a new thermal equilibrium. Seen from this point of view the occurrence of instability is the first step in this chain of processes directed toward thermal equilibrium. Putting an obstacle into a fast but otherwise thermalised plasma stream is clearly a way of providing free energy, because seen from the frame of the obstacle the plasma is not in equilibrium; there is a large velocity difference between plasma and obstacle and thus a large amount of free latent energy available in the system. Dissipation, however, requires small scale interactions between the plasma constituents. Macroscopic motions with their long scales contribute to large scale structures but do not directly act on the microscopic scales.

In order to affect the particle motion and contribute to friction among the particles small scale processes have to be called for. These processes take mostly place on electron scales. Moreover, since it is much easier to excite fluctuations in the electric field than in the magnetic field, these interactions are electrostatic. In the following we consider the electrostatic fluctuations which are expected to contribute to the generation of anomalous dissipation. These processes can be divided into those which are not affected by the presence of an external magnetic field called unmagnetised, and those where the external magnetic field must be taken into account in the particle motion, i.e. magnetised processes. The distinction is made by the relation between plasma and cyclotron frequencies.

When $V_A \ll c\sqrt{m_e/m_i}$, the electron cyclotron frequency is much less than the electron plasma frequency, $\omega_{pe} \gg \omega_{ce}$ and one is dealing with an unmagnetised case. Also, when the entire dynamics is restricted to the direction parallel to the magnetic field, the problem can be considered to be unmagnetised. The complete unmagnetised dispersion relation including all species s and their drifts V_s is

$$1 + \sum_s K_s(\omega, \mathbf{k}) = 0, \quad K_s(\omega, \mathbf{k}) = \frac{\omega_s^2}{N_s} \frac{\partial}{\partial \omega} \int \frac{d\mathbf{v}^3 F_{0s}(\mathbf{v})}{\omega - \mathbf{k} \cdot \mathbf{v}} \quad (3.87)$$

The function K_s is the susceptibility contribution of species s with average distribution function F_{0s} . In a Maxwellian component plasma $K_s = -(1/2k^2 \lambda_{Ds}^2) Z'(\zeta_s)$ can be expressed through the plasma dispersion function $Z(\zeta_s)$, with λ_{Ds} the Debye-length, $\zeta_s = (\omega - \mathbf{k} \cdot \mathbf{V}_s)/\sqrt{2}k v_s$, and thermal speed v_s of component s . Note that here the sign of charge is included in the cyclotron frequency, i.e. for electrons $-\omega_{ce}$, for ions $+\omega_{ci}$.

When the plasma is magnetised, which applies to all other cases, the susceptibility becomes more involved. For Maxwellian components it reads

$$K_s(\omega, \mathbf{k}) = \frac{1}{k^2 \lambda_{Ds}} \left[1 + \frac{\omega e^{-\eta_s}}{\sqrt{2}|k_{\parallel}|v_s} \sum_{l=-\infty}^{\infty} I_l(\eta_s) Z(\zeta_s^l) \right], \quad \zeta_s^l \equiv \frac{\omega + l\omega_{cs}}{\sqrt{2}|k_{\parallel}|v_s} \quad (3.88)$$

Here $\eta_s = (k_{\perp} r_{cs})^2$ with $I_l(\eta_s)$ the order- l Bessel function of imaginary argument.

This last expression suggests that magnetised electrons will support *electron cyclotron harmonics* $l\omega_{ce}$, which, when purely perpendicular, are *Bernstein modes*. There are also ion Bernstein modes $l\omega_{ci}$, but (with the exception of injection of a localised perpendicular ion beam) they do not play any susceptible role in transport. Exciting both electron and ion Bernstein modes requires beams perpendicular to the magnetic field, in which case when applied to shocks the *lower hybrid* instability will become the most important agent in generating anomalous dissipation.

As with the electromagnetic waves discussed in the previous sections the electrostatic dispersion relations can be solved analytically in closed form only in particular simple cases, and one has to retreat to a numerical approach. In the following we present the few most important cases and their effect on the generation of anomalous dissipation keeping in mind, however, that most collisionless shocks in the heliosphere are supercritical, and anomalous dissipation does not contribute substantially to their evolution and maintenance. Dealing in Chapter 6 with supercritical shocks, particle reflection is the dominant dissipation process, and we will refer to the instabilities of the previous section.

3.5.1 Electrostatic Wave Particle Interactions

The wave friction term Eq. (3.26) requires the determination of the fluctuation amplitudes of the wave modes that are responsible for causing anomalous friction. Again, the first step is to identify the unstable wave modes. In the second step we will then either have to determine the saturation level amplitudes of these waves or to consider their further interaction with particles or other waves.

Here we list instabilities of interest in the anomalous dissipation process only. When discussing application to observations we will later in passing also mention instabilities which are involved in radiation from shocks. However, radiation provides no substantial energy loss, and the dissipation caused by radiation under the conditions of the nonrelativistic shocks in the heliosphere is completely negligible and does neither affect shock formation nor shock structure. One possible exception are shocks in the solar atmosphere which sometimes are accompanied by x-ray emission which, however, is not a genuine unstable plasma process in this case. Such processes are believed to take place in solar coronal shocks.

Unmagnetised Electron and Ion Instabilities

The unmagnetised dispersion relation in a Maxwellian component plasma consisting of one electron and one ion component has two solution, electron plasma waves or Langmuir waves $\omega^2 = \omega_{pe}^2 + 3k_{\parallel}^2 v_e^2$ at $k_{\parallel} \lambda_{De} \ll 1$ and ion-acoustic waves. The Langmuir wave can be driven unstable by a parallel electron beam of velocity $V_e \gtrsim 3v_e$.

Ion-Acoustic Waves. More interesting are ion-acoustic waves. In the absence of any difference velocity between electrons and ions these are strongly damped plasma waves propagating along the magnetic field with dispersion

$$\frac{\omega^2}{k_{\parallel}^2} \simeq \frac{3T_i}{m_i} + \frac{T_e}{m_i} \frac{1}{1 + k_{\parallel}^2 \lambda_{De}^2}, \quad k_{\parallel} \lambda_{Di} \ll 1 \quad (3.89)$$

These waves become dispersive at larger $k_{\parallel}$ and for large electron temperatures with concave dispersion curve. At small parallel wave numbers their dispersion is linear resembling sound waves with velocity $c_{ia}^2 \simeq (T_e + 3T_i)/m_i$. Obviously for dispersion to compensate nonlinearity large electron temperatures $T_e > 3T_i$ are required in which case also the damping is small. Ion acoustic waves if managing to overcome damping are therefore a candidate for electrostatic unmagnetised shock formation: Their dispersion favours shock ramps with slow moving wave trails and, as we will see later, they also can contribute to dissipation thus satisfying all conditions as candidates for shock formation. However, such shocks are purely electrostatic and do not affect the magnetic field. They will therefore only be of interest in sub-structuring magnetised shocks, possibly contributing to the formation of subshocks of short wavelength of the order of several Debye lengths.

In fact going from the kinetic description to the fluid description and writing down the continuity and momentum conservation equations for ion acoustic waves for parallel propagation and one-dimensionality and combining it with Poisson's equation

$$\frac{\partial N}{\partial t} + \nabla_{\parallel}(NV) = 0, \quad \frac{\partial V}{\partial t} + V \nabla_{\parallel} V = -\frac{e}{m_i} E_{\parallel}, \quad E_{\parallel} = -\nabla_{\parallel} \phi \quad (3.90)$$

under the well justified assumption that for the low ion-acoustic frequencies the electrons behave as thermalised hot Boltzmannians with density $N_e = N \exp(e\phi/T_e)$ depending exponentially on the electrostatic potential ϕ , assuming quasi-neutrality and localised stationary solutions, we manipulate all these equations into the Sagdeev pseudo-potential form

$$\frac{1}{2}(\nabla_{\parallel} \phi)^2 = -S(\phi), \quad S = -\frac{m_i N_1 V_1^2}{\varepsilon_0} \left[\left(1 - \frac{2e\phi}{m_i V_1^2} \right)^{\frac{1}{2}} + \frac{T_e}{m_i V_1^2} \exp \frac{e\phi}{T_e} \right] \quad (3.91)$$

As usual the subscript 1 refers to values far upstream of the localised solution. In the absence of dissipation this solution for $S < 0$ yields solitons of maximum potential amplitude ϕ_m found from setting the bracket to zero. ϕ_m corresponds to a maximum compressive amplitude $N_m = N \exp(e\phi_m/T_e)$. With ion acoustic Mach number $\mathcal{M}_{ia} = V_1/c_{ia}$ it is found that solitons exist only in the supersonic regime $\mathcal{M}_{ia} > 1$. The soliton speed can be expressed through the maximum potential (or density via Boltzmann's expression) as

$$\frac{V_1}{c_{ia}} = \frac{1}{\sqrt{2N}} \frac{N_m - N}{\{N_m - N[1 + \ln(N_m/N)]\}^{\frac{1}{2}}}, \quad m_i V_1^2 > 2e\phi_m$$

The condition on the maximum compression amplitude N_m in this expression simply requires that the potential energy must be less than the initial flow energy. This sets a limit on the possible Mach numbers $\mathcal{M}_{ia}^2 > 2e\phi_m/T_e$ or, when combined with the definition of the latter, $\mathcal{M}_{ia}^2/2 > \ln(N_m/N)$.

Since the denominator in the former equation must be real, this condition requires that

$$\ln(N_m/N) + 1 - N_m/N < 0$$

Expressing herein the density ratio through the Mach number, one finds that for ion-acoustic solitons to exist the Mach number is limited to values below a surprisingly small critical Mach number

$$\mathcal{M}_{ia} < \mathcal{M}_{ia}^{\text{crit}} \simeq 1.6$$

Ion acoustic solitons cannot exist for Mach numbers exceeding $\mathcal{M}_{ia}^{\text{crit}}$. For higher inflow velocities V_1 the ion acoustic soliton will either not evolve or break down.

The range of possible Mach numbers is rather limited which is simply due to the fact that for higher speeds the dispersion is unable to sustain a stationary state. Any solution will be non-stationary, wave like or unstable. Concerning the formation of shocks one in addition to dispersion requires that dissipation is produced. Since it is known that ion acoustic waves are Landau damped with damping rate $\gamma_{L,ia}$, one can argue that Landau damping will cause a shock profile on the ion acoustic soliton with downstream state different from the upstream state [Ott & Sudan, 1969; Tidman & Krall, 1971]. The downstream density is then found to be $N_2 \simeq N_1 \exp[2(\mathcal{M}_{ia} - 1)]$ and will exhibit trailing oscillations, as has been discussed above. However, Landau damping takes time, and therefore the general argument applies to this kind of shock formation that the damping will not have time to work for large Mach numbers. Thus the damping argument applies only to subcritical shocks of Mach numbers smaller than $\mathcal{M}_{ia} \simeq 1.6$. Such weak (electrostatic) ion acoustic shocks can indeed evolve and may contribute to sub-structuring of stronger supercritical shocks in the region where the Mach number has already dropped to values below the critical.

Sagdeev [1966] has favoured reflection of inflowing particles from the leading edge of the soliton [Moiseev & Sagdeev, 1963] over Landau damping. This reflection affects ions with energy less than the soliton potential ϕ_m and causes oscillations of long wave length. More important is that the reflected ions form an ion-ion beam configuration and are thus subject to the ion instabilities discussed previously yielding waves which may generate dissipation but do also propagate upstream of the shock where they cause wave particle interactions and retard the inflow ahead of the shock.

Electron Current Driven Ion Acoustic Instability. So far we have not asked for the reason of an ion acoustic wave to grow. This can be achieved in the simplest way by letting one of the plasma components drift with respect to the other. If in a two-component plasma the electron drift with respect to the ions they effectively carry a current $j_{\parallel} = -eNV_e$ which is in most cases – but not necessarily – along the magnetic field. Here we assumed again quasi-neutrality $N_e = N_i = N$ which dispenses us from considering space charges and solving Poisson's equation. In this case assuming weak growth such that we can apply the general instability theory with $\gamma \ll \omega$ the growth rate of the ion acoustic wave Eq. (3.89) becomes

$$\frac{\gamma_{ia}}{\omega} = \sqrt{\frac{\pi}{2}} \frac{\omega^2}{k^3 c_{ia}^2} \frac{\mathbf{k} \cdot \mathbf{V}_e - \omega}{2v_e} \exp \left[-\frac{(\omega - k_{\parallel} V_e)^2}{2k_{\parallel}^2 v_e^2} \right] \quad (3.92)$$

Instability sets on for $V_e > \omega/k_{\parallel}$ when the electron velocity, which is the current drift velocity, exceeds the phase velocity of the ion acoustic wave, i.e. when – approximately –

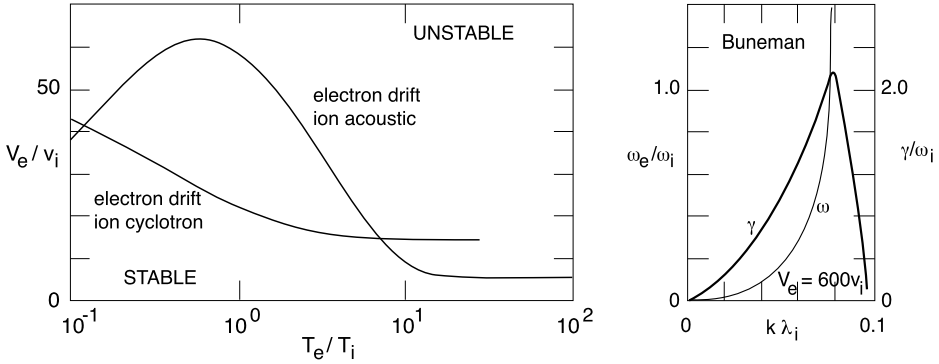


Figure 3.14: *Left:* The threshold drift speed for the electron current driven ion-acoustic instability as function of the electron to ion temperature ratio. For comparison the threshold for the parallel electron current driven ion cyclotron instability in a magnetised plasma is shown for a dense plasma with $\omega_{pe}/\omega_{ce} = 10$ [after *Kindel & Kennel, 1971*, courtesy American Geophysical Union]. *Right:* Buneman electron drift-ion two-stream instability, frequency and growth rate as function of wave number at large electron current drifts when the ion-acoustic instability has changed to the Buneman mode.

$V_e > c_{ia}$. These waves have relatively long wavelength $k_{\parallel}\lambda_{De} \ll 1$. The threshold for marginal stability of these electron current driven ion acoustic waves can be obtained from setting $\gamma_{ia} = 0$, yielding

$$V_e \simeq \frac{\omega}{k_{\parallel}} \left[1 + \left(\frac{m_i}{m_e} \right)^{\frac{1}{2}} \left(\frac{T_e}{T_i} \right)^{\frac{3}{2}} \exp \left(-\frac{3}{2} - \frac{T_e}{2T_i} \right) \right] \quad (3.93)$$

The second term in the brackets results from the Landau damping of the ion acoustic waves. It is seen that this term disappears for hot electrons with $T_e \gg T_i$ thus lowering the threshold for instability to its marginally smallest value $V_e = c_{ia}$. The threshold is shown graphically as function of the temperature ratio in Figure 3.14. The threshold is measured in ion thermal speeds v_i and is quite high. Moreover, the electron temperature must be high implying that the electron distribution must be much hotter than the ion distribution. In the solar wind this is usually satisfied but fails downstream of the shock. The physical reason for the electron temperature to be high for instability is that the distribution must have a positive slope in v in the region of overlap with the cold ion distribution for resonant instability, $\partial F_{0e}/\partial v_{\parallel}|_{\omega/k_{\parallel}} > 0$. There must be more fast than slow electrons in the phase velocity frame of the wave in order to push the wave to higher momentum and energy, i.e. causing instability.

It is clear that this kind of interaction between the ion acoustic wave goes on the expense of the motional energy of the resonant electrons. Hence one expects that ion acoustic waves retard and scatter the current electrons thereby reducing the current flow, implying that first an electron current along the magnetic field has to be generated. This resembles collisional friction which the resonant electrons experience and can thus be interpreted as the production of anomalous resistance in the plasma. An interpretation like

this has been put forward by *Sagdeev* [1966] and has been elaborated in depth afterwards. Below we will return to this theory.

Buneman Electron Current Two Stream Instability. Returning to Figure 3.14 we observe a change in the threshold curve for $V_e/v_i > \sqrt{m_i/m_e T_e/T_i}$. Here the electrons become thermally slow with respect to their drift velocity, $v_e < V_e$. The ion-acoustic instability under this condition changes into the Buneman instability which is an electron current fluid instability and is also known under the name electron current drift or electron-ion two stream instability. It has been discovered by *Buneman* [1958] and favoured for application in shock physics already by *Sagdeev* [1966]. It should be noted that the transition from ion-acoustic to two-stream instability has been investigated in depth by *Dum & Chodura* [1979].

Now, treating the electrons as cold the kinetic effects disappear, and the complex dispersion relation of the Buneman instability becomes

$$1 - \frac{\omega_{pi}^2}{\omega^2} - \frac{\omega_{pe}^2}{(\omega - k_{\parallel} V_e)^2} = 0 \quad (3.94)$$

Note that under these conditions the weak growth theory cannot be applied anymore. Instead one must find the growth rate from the complex solutions of this quartic expression. Fortunately, this equation can be solved since for resonant electrons the third term becomes dominant. The instability has real frequency $\omega \sim k_{\parallel} V_e$ and maximum growth rate of the order of the ion plasma frequency

$$\gamma_{\text{Bun,max}} \simeq \frac{\sqrt{3}}{16^{\frac{1}{3}}} \left(\frac{m_i}{m_e} \right)^{\frac{1}{6}} \omega_{pi} \sim 2.4 \omega_{pi} \quad (3.95)$$

Figure 3.14 on its right shows the Buneman frequency and growth rate for a case of very large electron current drift $V_e = 600 v_i$ in dependence on the wave number $k_{\parallel} \lambda_{Di}$. The maximum growth of the instability is close to $k_{\parallel} \lambda_{Di} \sim 0.9$ at short wavelengths slightly larger than the Debye length. The growth rate of this instability is very large. This implies that the instability is very strong and grows very fast thereby consuming a substantial fraction of the current streaming energy. The nonlinear treatment of this instability is of particular interest for shock physics. Since the instability grows so fast it makes little sense to treat it analytically for reasons which will become clear when dealing with the application of numerical simulations to shocks. The physics involved into the two stream instability can be described as follows. Both the electrons carrying the current and the ions are cold. The situation is thus two stream, and the instability is not resonant but reactive with all particles participating. This is the reason for its strength and rapidity. Because it consumes a fraction of the bulk flow energy of the electron current, the current becomes decelerated, and the energy is going mainly into the electrons which are heated by the instability until the instability stabilises when $V_e < v_e$. Then the ion-acoustic instability takes over. The Buneman two stream instability is thus accompanied by a burst in electron temperature and a rapid decrease in current. However, the final state of the instability is not a stationary

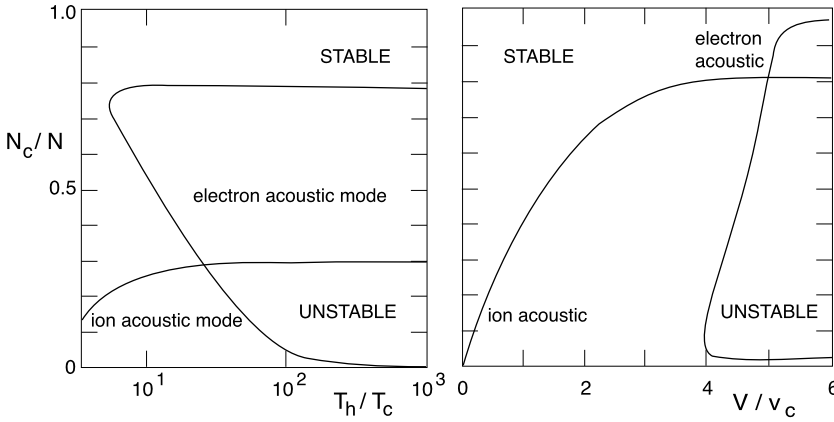


Figure 3.15: *Left:* Regions of existence of electron and ion acoustic modes in the density/temperature plane. Here $N = N_c + N_h = N_i$ is the total quasi-neutral density. The temperatures given as the hot to cold temperature ratio. The electron acoustic mode extends to larger cold density and higher hot temperatures than the ion acoustic mode. *Right:* Unstable versus stable domains for electron and ion acoustic modes in the density/drift velocity plane for the case when the electron acoustic mode is destabilised by a drifting hot plasma carrying a current.

state because the ion acoustic instability continues to grow with its own dynamics, possibly ending up in the formation of solitons when the Mach number has sufficiently decreased by current relaxation and heating, the former decreasing V_e , the latter increasing c_{ia} . But even during the blow-up phase of the two-stream instability structuring similar to soliton formation occurs. This can only be inferred from numerical simulation.

Modified Two-Stream Instability. This is a variant of the ordinary two-stream instability driven by a relative drift between electrons and ions but this time with the drift direction across the ambient magnetic field $\mathbf{B}$ [McBride *et al*, 1972] thereby accounting for differences in the cross field particle drifts which may lead to cross-field (e.g. diamagnetic) current flow. The dispersion relation for the modified two-stream instability is

$$\left[1 - \frac{\omega_{pi}^2}{(\omega - kV_i)^2} - \frac{\omega_{pe}^2}{\omega^2} \right] \left[1 - \frac{\omega_{ce}^2 \cos^2 \Theta_{Bn}}{\omega^2 (1 + k^2 \lambda_D^2)} \right] = \frac{\omega_{ce}^2 \sin^2 \Theta_{Bn}}{\omega^2 (1 + k^2 \lambda_D^2)} \left[1 - \frac{\omega_{pi}^2}{(\omega - kV_i)^2} \right] \quad (3.96)$$

This expression is written here in the electron frame of reference and with angle Θ_{Bn} . One recognises that the first term is the ordinary Buneman two-stream term. However, for oblique propagation $\Theta_{Bn} \neq 90^\circ$ and $\Theta_{Bn} \neq 0^\circ$ the two-stream mode couples to the whistler mode. It is this coupling which makes the modified two-stream instability interesting for shocks. Dispersion curves and growth rates are shown in Figure 3.16 for $\Theta_{Bn} = 60^\circ$, and $V_i = V_A$, and an artificial mass ratio $m_i/m_e = 80$ which has been taken in view of numerical simulations to be discussed later.

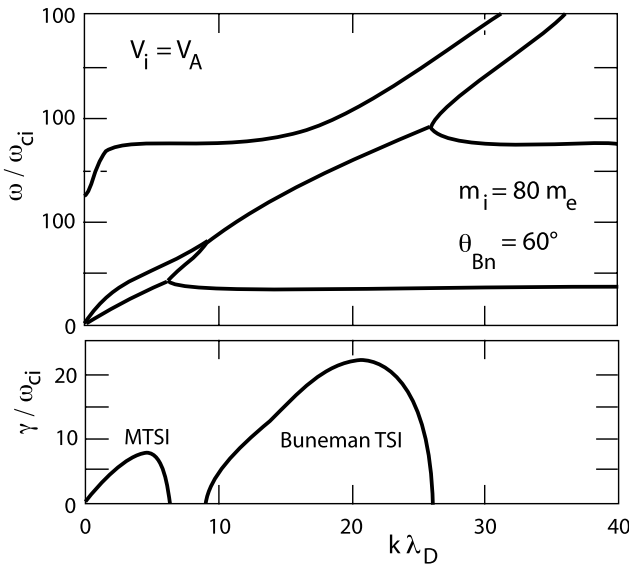


Figure 3.16: *Top:* Dispersion curves for the MTSI and Buneman TSI showing the coupling of the modes in dependence on wave number [after Matsukiyo & Scholer, 2003, courtesy American Geophysical Union]. *Bottom:* Growth rates for the Buneman TSI and MTSI. The Buneman TSI is at shorter wavelengths and higher frequencies but has larger growth rate while the MTSI has longer wavelengths, lower frequency. However, even though the MTSI does not exist for $\Theta_{Bn} = 0^\circ$ and $\Theta_{Bn} = 90^\circ$, it has much lower threshold than the Buneman TSI and therefore is much more likely to be excited than the Buneman TSI.

The modified two-stream instability (MTSI) operates also for relative drifts smaller than the electron thermal but larger than the ion-acoustic velocity and even for $T_i \sim T_e$ which makes it potentially important if only such perpendicular drifts can be generated. The unstable frequency is in the range of the lower-hybrid frequency. Hence the ions can be taken unmagnetised with strongly magnetised electrons. However, it requires oblique relative electron drifts since for perpendicular drift the instability disappears meaning that the unmagnetised ions propagate under an angle to the magnetic field while the magnetised electrons move only parallel to the magnetic field.

On the other hand, for unmagnetised ions and in the presence of a perpendicular electric field $\mathbf{E} \perp \mathbf{B}$ the electrons may perform a $\mathbf{V}_{de} = \mathbf{E} \times \mathbf{B} / B^2$ drift perpendicular to both fields relative to the ions which corresponds to a perpendicular drift current $\mathbf{J}_{de\perp} = -eN\mathbf{V}_{de}$, in which case the modified two-stream instability will become unstable if only $V_{de} > c_{ia}$ exceeds the ion acoustic velocity. This might easily be the case inside the shock ramp where the shock produces an electric field the action of which is threefold: it retards the inflowing ions and reflects the lower energy ions back into the upstream flow, and it gives rise to the electron drift current on the transverse shock scale on which the ions are unmagnetised, i.e. under shock conditions this is the ion inertial length λ_i .

Electron-Beam Electron-Acoustic Instability. A mode very similar to the ion-acoustic mode is the electron acoustic mode. Here the role of the ions is played by the cold (subscript c) electron background. In presence of another hot (subscript h) electron component the physics becomes very similar to the ion-acoustic wave, and a weakly damped resonant mode with real dispersion relation

$$\omega^2 \simeq \omega_{pe,c}^2 \frac{1 + 3k_{\parallel}^2 \lambda_{De,c}^2}{1 + 1/k_{\parallel}^2 \lambda_{De,h}^2}, \quad k_{\parallel} \lambda_{DE,c} < 1 < k_{\parallel} \lambda_{DE,h} \quad (3.97)$$

can propagate in the range of wave numbers indicated on the right. This weak damping can, like in the case of the ion acoustic instability, be overcome if the hot electron component performs a drift or if another electron beam is injected into the plasma. In the former case with ions and cold electrons at rest the hot electron drift implies current flow. The regions of existence and stability of the electron acoustic instability are shown in Figure 3.15 in comparison with the ion acoustic instability. Growth rates have been found numerically to be quite large, the order of the ion plasma frequency and thus similar to the Buneman mode. On the other hand, numerical experiments have not detected strong electron acoustic waves.

Inspection of the electron acoustic dispersion relation shows that the mode is dispersive as well. For very low cold electron temperatures this dispersion is similar to ion acoustic waves suggesting that electron acoustic solitons could be formed in a similar way. In fact, such solitons have been calculated analytically for different parameter regimes [Dubouloz *et al*, 1991]. However, observations do not seem to support their existence. Also numerical simulations [Matsukiyo *et al*, 2004] have not shown the formation of such solitons. Clearly, electron acoustic waves can have a different dynamics because of the fast reaction times of the electrons, leading to rapid thermalisation of the plasma. These questions still remain to be open. In any case, if electron acoustic waves exist, the electron acoustic mode is quite well suited for plasma heating in shocks where plasmas of different temperatures mix. Its advantage is also that it proceeds on a very fast time scale close to the inverse of the plasma frequency. It is hence well suited for fast production of anomalous dissipation of energy. Moreover, since it very obviously damps rapidly it may act as an agent of about immediate transformation of excess energy in the electrons into heating electrons. None of these question has been understood nor answered properly at the time of writing.

Ion-Beam Ion-Acoustic Instability. An instability very similar to electron beam driven electron acoustic instability is its ion pendant when a cold ion core and hot ion beam in presence of a neutralising electron background become unstable [Gary & Omidi, 1987]. The mode excited in this case is again the ion acoustic mode, this time propagating at very low phase velocities $\omega/k \simeq c_{ia}(N_{i,c} - N - i, b)/N_e$ less than c_{ia} . In this range the waves have no dispersion. This wave is, however, very easy to destabilise because of its low beam velocity threshold which can lie even below the thermal speed of the ion core component. On the other hand the growth rates of this instability are very low. Measured as before for the electron-acoustic, ion-acoustic and Buneman instabilities in terms of the

ion plasma frequency, its maximum growth rate is just of the order of a few per cent of $\gamma_m \sim 0.3\omega_{pi}$, very small compared to the growth rate of the electron-current driven ion acoustic instability which was of the order of $\gamma_m \sim \omega_{pi}$. It should thus be less important in anomalous processes.

Electrostatic Ion-Cyclotron Harmonic Instabilities. So far we have considered only unmagnetised instabilities. We now turn to listing the most important – in view of the anomalous processes we have in mind when considering shock formation – *magnetised electrostatic* instabilities. These instabilities occur when a magnetic field is present in the plasma and the electrons must be considered to be magnetised.

The frequencies of the waves in question will therefore be well below the electron cyclotron frequency $\omega \ll \omega_{ce}$ falling into the range of and above the ion cyclotron frequency $\omega \gtrsim \omega_{ci}$. According to linear theory the magnetised modes in this range occur at harmonics $\omega_l \sim l\omega_{ci}$ of the ion cyclotron frequency. We are thus dealing here with electrostatic ion-cyclotron harmonics.

Following a suggestion by *Drummond & Rosenbluth* [1962] ion cyclotron instabilities have been proposed as generating anomalous collisions by *Kindel & Kennel* [1971] who advocated their importance because of their apparently lower instability threshold than the ion acoustic instability for electron current shown in Figure 3.14. In fact ion-cyclotron harmonic waves oblique to the magnetic field can become unstable in the presence of an electron current flowing along the magnetic field. The reason is that – in contrast to the earlier mentioned strictly perpendicular Bernstein mode resonances – such oblique waves possess a field-aligned projection component of their electrostatic field which can resonate with the parallel current electrons via the Doppler-shifted resonance condition $k_{\parallel}V_{\parallel e} = \omega \pm l\omega_{ci}$ of which only Landau resonance $l = 0$ is effective for $|\omega - k_{\parallel}V_{e\parallel}| < \sqrt{2}|k_{\parallel}|v_e$. The ions contribute only a weak resonant damping of the waves for $l = -1$. Note that the obliqueness of propagation of these modes implies in contrast to Bernstein modes that the resonance does not occur precisely at the harmonics but lies somewhere in between in the harmonic bands where the dispersion curves have particular geometrical forms [cf., any book on basic plasma wave physics, e.g., *Baumjohann & Treumann*, 1996]. Its precise location depends heavily on the exact prescribed conditions of the plasma and current velocity $V_{e\parallel}$, and no analytic expression can be provided.

Strongest weak instability occurs in the harmonic range $1 < l < 2$ with growth rate $\gamma \ll \omega$ and $k_{\parallel}r_{ci} \sim 1$, i.e. wavelengths comparable to the thermal ion gyro-radius in the frame of the ions and propagation angles $\theta < 85^\circ$. The velocity threshold decreases with increasing electron temperature T_e/T_i simply because more electrons go into resonance then. But for large ratios it is taken over by the ion acoustic instability as seen from Figure 3.14. (Note that in the solar wind/bow shock system, for instance, the ratio is about $T_e/T_i \approx 10$ changing across the bow shock to become $T_e/T_i \lesssim 1$; hence one may expect the first electrostatic ion-cyclotron harmonic to be present wherever parallel electron currents flow in the shock on the upstream side, while downstream neither current driven ion cyclotron nor ion acoustic instabilities should exist.)

We have already noted that ion beams can also excite ion cyclotron harmonic waves. Parallel beams excite similar waves with similar properties like parallel electron currents

for propagation angles $0^\circ \ll \theta < 90^\circ$ rather close to perpendicular. Perpendicular ion beams, on the other hand excite a broad spectrum of ion cyclotron harmonics on the background ion component depending on beam energy and background temperature. These excitations are restricted to a distance of the order of the ion gyro-radius $r < r_{ci,b}$ of the beam ions only, because at larger perpendicular scales the beam ions are themselves magnetised. Thus this kind of excitation is of importance merely when heavy ion beams penetrate the plasma, for instance when a heavy ion beam component of large perpendicular energy is reflected from a supercritical shock wave back upstream, or when ions become highly accelerated in interaction with the shock and penetrate across the shock onto the downstream side where they have much larger gyro-radii than background ions. On the other hand, in the shock foot of a quasi-perpendicular supercritical shock a large number of electrostatic ion cyclotron harmonics may be excited as on this scale the reflected ions are effectively unmagnetised.

Electrostatic Electron-Cyclotron Instability. This instability is the pendant to the former ion cyclotron instability at the much higher frequency perpendicular (or oblique) electron cyclotron harmonics (for purely perpendicular propagation these are the Bernstein modes). It is driven by the coupling between a sufficiently fast ion beam and the electron cyclotron harmonics at perpendicular wave numbers $k_\perp > 0$ and can also be driven unstable when reflected ions are present for instance in the foot of the supercritical perpendicular shock. When driven unstable by an oblique current it becomes the electrostatic modified two stream instability gaining high importance thereby.

Electrostatic Drift Instabilities in Inhomogeneous Plasma

The last group of instabilities we will refer to in the context of shock physics are instabilities that are caused in presence of plasma inhomogeneity. Three basic kinds of plasma inhomogeneities can be identified: density $N(\mathbf{x})$, magnetic $\mathbf{B}(\mathbf{x})$, and temperature $T(\mathbf{x})$ real space inhomogeneities. The dependencies of these (average) quantities on space may in most cases not be independent. In the following, for the purposes of investigation of their effects on wave excitation, we will take them as being in fact mutually independent. For simplicity we will take into account only density inhomogeneities perpendicular to $\mathbf{B}_0$ on scales much larger than the wave length. We also assume linearity, which is justified because under the assumption of weak gradients the effects of the inhomogeneity will be weak as well and thus cause only slow wave growth. In this case one can expand the density with respect to the perpendicular direction x up to first order and write

$$N_s(x) = N_{s0}(1 + \varepsilon_N x), \quad \text{where} \quad \varepsilon_N \equiv [\nabla_x N_s(x)]_{x=0}$$

The effect of the inhomogeneity is that the magnetised particle component behaves adiabatically and starts performing a diamagnetic drift motion $V_{Ns} \hat{\mathbf{y}} = (\varepsilon_N v_s^2 / \pm \omega_{cs}) \hat{\mathbf{y}}$ perpendicular to the magnetic field and density gradient into $\pm y$ direction, depending on the sign of the particle charge. The $\pm$ -sign in the denominator indicates that the cyclotron frequency is taken here including the sign of the charge. Drift motions of this kind cause a perpendicular drift current j_{dy} to flow in the plasma because particles of different charge

sign move in opposite directions. The waves excited under such conditions satisfy the modified dispersion relation

$$1 - \varepsilon_N + \sum_s K_s(\mathbf{k}, \omega, \varepsilon_N) = 0 \quad (3.98)$$

The new susceptibility $K_s(\varepsilon_N)$ is still local, i.e. it changes on scales larger than the gradient scale, but locally depends only on ε_N . It is of the same form as Eq. (3.88) with two differences: because of the occurrence of the finite perpendicular drift $V_{dy}\hat{\mathbf{y}}$ the frequency in the factor in front of the sum in (3.88) is exchanged with the Doppler-shifted frequency $\omega \rightarrow \omega - k_{\perp} V_{Ns}$. In addition a new drift term

$$+ \frac{\omega V_{Ns}}{k_{\perp} v_s^2} e^{-\eta_s} \sum_{l=-\infty}^{\infty} \frac{l \omega_s I_l(\eta_s)}{\omega + l \omega_{cs}}$$

must be added inside the brackets. Instabilities resulting from this dispersion relation in the weak instability limit bear the general name of *drift* or *universal* instabilities. They are resonant with the drift motion and have frequencies $\omega \simeq k_{\perp} V_{Ni}$ and long wave lengths satisfying $k_{\perp} r_{ci} \lesssim 1$. For smaller wavelengths these waves are highly dispersive and thus can form nonlinear structures.

This mode becomes particularly interesting and important in shock physics when the drift speed is so large that the ions can be considered as unmagnetised. This happens because the frequency of the drift mode increases with V_N and quickly exceeds the ion cyclotron frequency. In this case the frequency of maximum growth is close to the lower-hybrid frequency ω_{lh} and the drift mode becomes a lower-hybrid drift wave.

Lower-Hybrid Drift Instability. For unmagnetised ions, when the frequency of the drift wave is $\omega \gg \omega_{ci}$, the ion cyclotron frequency can be neglected and the susceptibility simplifies

$$K_s(\omega, \mathbf{k}) = \frac{1}{k^2 \lambda_{Ds}^2} \left[1 + \frac{\omega - \mathbf{k}_{\perp} \cdot \mathbf{V}_{Ns}}{\sqrt{2} k v_i} Z\left(\frac{\omega}{\sqrt{2} k v_i}\right) \right] \quad (3.99)$$

The positive slope on the distribution function which is responsible for instability is in this case on the perpendicular part which depends on the drift velocity. There the maximum of the distribution is shifted out of the origin to the location of the drift velocity. The real frequency and growth rates are given by [Gary & Sanderson, 1979]

$$\frac{\omega}{k_{\perp} V_{Ne}} \simeq - \left[\left(1 + \frac{T_e}{T_i} \right) \frac{e^{\eta_e}}{I(\eta_e)} - 1 \right]^{-1} \quad (3.100)$$

$$\frac{\gamma}{\omega} \simeq \sqrt{\frac{\pi}{2}} \frac{T_e}{T_i} \frac{V_{Ne}}{v_i} \frac{e^{\eta_e}}{I_0(\eta_e)} \left[\left(1 + \frac{T_e}{T_i} \right) \frac{e^{\eta_e}}{I(\eta_e)} - 1 \right]^{-2} \quad (3.101)$$

The unstable wave propagates antiparallel to the direction of the electron gradient drift, i.e. in the direction of the electric drift current. Maximum growth of this lower hybrid wave

has been found at long wavelengths $k_{\perp,m} r_{ci} \sim \sqrt{m_i/m_e}$ over a relatively broad frequency range close to the lower hybrid frequency $\omega_{lh} \simeq \omega_{ce} \sqrt{m_e/m_i}$.

The reason for the lower-hybrid drift instability to just excite the lower-hybrid frequency is that it is a perpendicular two-stream instability similar to the Buneman mode that is, however, driven by the bulk velocity difference between electron and ion gradient drifts. It thus propagates on the background ion component. In fact in the $(\omega, k_{\perp})$ -plane the drift-beam mode $\omega_d \simeq k_{\perp} V_{Ne}$ couples to many ion cyclotron harmonics thereby exciting almost all of them. However, largest growth occurs in the harmonic dispersion band that contains the lower hybrid frequency which is a strong plasma resonance.

The lower-hybrid drift instability is the strongest in the family of the electrostatic ion cyclotron instabilities. As a two stream instability its maximum growth rate remains still modestly large being of the order of a few ion cyclotron frequencies, $\gamma_{lh,m} \simeq (1-3)\omega_{ci}$. In a plasma leaving sufficient time $\tau_{nl,sat} \gg \omega_{ci}^{-1}$ for growth and saturation it may well play a substantial role in generating dissipation. We will see later that this instability indeed provides the highest so far inferred from instabilities anomalous collision frequency which turns out to be of the order of the lower-hybrid frequency itself. Its relevance in application to collisionless shocks is however questionable, because of the above argument.

In addition, the lower-hybrid drift instability appears to stabilise under $\beta > 1$ conditions [Davidson et al, 1977]. In application to shocks this restriction, if it translates into the nonlinear regime, presents a severe barrier to the use of the lower hybrid drift instability as generator of anomalous resistance, dissipation and entropy generation.

3.5.2 Anomalous Resistivity

Resistivity is defined via the Drude formula $\eta = v/\epsilon_0 \omega_{pe}^2$, with $v = \sigma_e N v_e$ the collision frequency. The latter, under collisionless conditions, becomes the anomalous collision frequency ν_a and is the quantity containing the interaction between electrons and the nonlinear wave fluctuations.

This becomes obvious when realising that the Spitzer-Coulomb collision frequency $\nu_C \sim \omega_{pe}/N \lambda_{De}^3$ is proportional to the ratio of the *plasma wave fluctuation level in thermal equilibrium* $W_{th} = \frac{1}{2} \epsilon_0 \langle \mathbf{e}_{th}^2 \rangle$ to thermal energy, $\nu_C \sim \omega_{pe} W_{th}/NT_e$. Under saturated instability conditions it is then reasonable to assume that the actual fluctuations $\langle \mathbf{e}^2 \rangle$ replace the thermal fluctuations in this expression which yields the Sagdeev formula

$$\nu_a \simeq \frac{W_{sat}}{NT_e} \omega_{pe} \quad (3.102)$$

The problem is thus reduced to the determination of the *nonlinear saturation level of the unstable wave spectrum*. Its determination requires knowledge of the electric current $\mathbf{j} \simeq -e \langle N \mathbf{V}_e \rangle$ as a functional of the electric wave fluctuation field $\mathbf{e}$. (We are speaking here only of electrons since in collisionless plasma electrons – because of their much faster mobility than ions – are the particles which carry the electric current. The electrons feel the friction of the waves and become retarded by anomalous collisions thereby dissipating the kinetic energy of the current and contributing to collisionless Joule heating of the plasma.) The evolution of the electron current is – in principle – given by the electronic

part of Eq. (3.24), respectively Ohm's law (3.30), if on the right-hand sides the average anomalous electronic friction terms (3.26) are added, since these are the crucial terms containing the wave-particle interactions.

We are interested only in the parallel collision frequency here. In the nonlinear stationary state the time derivatives can be neglected. In order to obtain a first expression for the parallel anomalous collision frequency ν_a we assume that the last term in Eq. (3.24) is of the form

$$-\frac{\nu_a}{m_e} N V_{\parallel} = \int d^3v v_{\parallel} \mathcal{C}_e$$

When inserting from Eq. (3.26), keeping only the parallel component and the electric part which in the microscopic interactions dominates over the electromagnetic fluctuations, since these affect only frequencies below ω_{ci} , one obtains for the anomalous collision frequency

$$\nu_a \simeq \frac{1}{N V_{\parallel}} \nabla_{\parallel} \langle W_{sat} \rangle \quad (3.103)$$

which is general but still preliminary. It requires knowledge of the average power density of the electric wave field $\langle W_{sat} \rangle = \frac{1}{2} \epsilon_0 \langle \partial_{\omega}(\omega \mathcal{E}) | \mathbf{e}^2 | \rangle$ which here can be of arbitrarily large amplitude and arbitrary spatial structure. Remember that the only condition implied in deriving (3.26) was that the fluctuations were fast both in space and time compared to the slow changes in the plasma background quantities. In this sense the parallel gradient operator $\nabla_{\parallel}$ acts on the slow variability of the wave power.

In order to proceed, another equation is required which determines the evolution of the wave power. This lacking equation can only be formulated in Fourier space $(\omega, \mathbf{k})$ and should contain all the nonlinear interactions and thus cannot be of general nature. A simplifying assumption is that it describes the nonlinear evolution according to a kind of wave-kinetic equation

$$\frac{d}{dt} W_k \simeq \frac{\partial W_k}{\partial t} + \left(\mathbf{V} - \frac{\partial \omega}{\partial \mathbf{k}} \right) \cdot \nabla W_k = 2\gamma(\omega, \mathbf{k}, W_k) W_k + \dots \quad (3.104)$$

We assume the system has reached stationarity such that the wave spectrum W_k does not evolve with time anymore. In this case the left-hand side simplifies, and we can express the spatial derivative of the stationary wave spectrum as

$$\nabla_{\parallel} W_k \simeq \frac{2\gamma(\omega, \mathbf{k}, W_k)}{|V_{\parallel} - \partial \omega / \partial k_{\parallel}|} W_k \quad (3.105)$$

which after transforming from Fourier into real space and averaging over the fast fluctuations can be used in the above expression (3.103) to express the anomalous collision rate through the average wave power. Such expressions will be given below. Usually the current-drift speed $|V_{\parallel}| \gg |\partial \omega / \partial k_{\parallel}|$ is much larger than the wave group velocity, and the latter can be neglected. This yields the inverse square dependence of $\nu_a \propto |V_{\parallel}|^{-2}$ on the current drift velocity.

The derivation of an anomalous collision rate contains a number of crucial assumptions of which the most severe concern the simplifications in the terms in the kinetic wave equation (3.104) where we suppressed a refraction term $(\nabla\omega) \cdot (\partial W_k / \partial \mathbf{k})$ – which becomes important in strongly inhomogeneous plasmas like in a shock ramp – and neglected all terms on the right with the exception of the growth rate term. Since the dependence of γ on the wave amplitude has been retained, some generality in the nonlinearity is still retained. Stronger simplifications are made when restricting to pure quasilinear theory. In this theory the nonlinear dependence of the growth rate on the wave amplitude is dropped. This confirms to the conventional approach to anomalous dissipation. The most elaborated *quasilinear* theory (valid for any direction of propagation including electromagnetic contributions) can be found in Yoon & Lui [2007] where also the application to one particular mode (the lower-hybrid drift mode discussed below) is given and a rudimentary contribution of the neglected Coulomb collision term is retained.

In any case the mechanism of saturation of the nonlinear wave field must be known in order to obtain a useful practical expression for v_a . In the following we review only the three wave modes that contribute strongest to anomalous resistance, ion-acoustic, Buneman two-stream, and lower-hybrid drift modes.

Nonlinear Evolution of Waves

Quasilinear theory [Sagdeev, 1979; Yoon & Lui, 2007] is so far the simplest and most effective approach to the calculation of anomalous collision frequencies. This approach uses the linear growth rate $\gamma(\mathbf{k})$ of the instability yielding the simplified formula

$$v_a \simeq \frac{1}{m_e N V_{\parallel}^2} \int \frac{d^3 k}{8\pi^3} \frac{\mathbf{k} \cdot \mathbf{V}_e}{\omega(\mathbf{k})} \gamma(\mathbf{k}) W_{sat}(\mathbf{k}), \quad W_{sat}(\mathbf{k}) \equiv \frac{\varepsilon_0}{2} \frac{\partial \omega \varepsilon(\omega, \mathbf{k})}{\partial \omega} \langle |\mathbf{e}^2| \rangle \quad (3.106)$$

It requires in addition knowledge of $\gamma(\mathbf{k})$ and the drift velocity $\mathbf{V}_e = V_{\parallel} \hat{\mathbf{z}}$.

Anomalous Ion Acoustic Collision Frequency. For instance, from this expression, assuming $V_{\parallel} > c_{ia}$ and $k_{max} \lambda_{De} \sim 1$, and $\gamma_{ia} \sim \omega V_{\parallel} / v_e$ which holds for the ion acoustic instability in the large drift limit, one just obtains the above Sagdeev expression (3.102) for the anomalous collision rate which is good for application when the saturation level of the instability is measured.

Ion acoustic waves saturate by scattering off thermal plasma ions according to the resonant scattering process $\omega - \omega' = (\mathbf{k} - \mathbf{k}') \cdot \mathbf{v}_i$ where the prime ' indicates the frequency and wave numbers of the scattered wave. This process has been used by Sagdeev [1966] to explicitly calculate the ion acoustic anomalous collision frequency

$$v_{a,ia} \simeq 0.01 \omega_{pi} \frac{V_{\parallel}}{c_{ia}} \frac{T_e}{T_i} \theta^{-2} \quad (3.107)$$

which holds for large electron current drifts $V_{\parallel} \gg c_{ia}$ and for narrow angles $\theta > 0$. Actually, experiments have shown that this expression overestimates the anomalous resistance suggesting that anomalous collisions are less effective than theory predicts. More precise

theories than the above estimate have been developed by *Vekshtein et al* [1970], and have been reviewed by *Sagdeev* [1979].

Anomalous Two-Stream Collision Frequency. As noted earlier, the two-stream instability is a very strong instability causing large current momentum losses and rapid plasma heating. It switches itself off during evolution and will therefore not be a stationary instability. It causes, in fact, different effects which probably destroy its direct importance in collision generation.

It has phase speed $\omega/k \sim v_e \sqrt{m_e/m_i}$ substantially below electron thermal speed implying that it stays relatively long in the volume of excitation which supports its effect on the local plasma. Its wave energy density, from simple consideration is less than ion energy $W_{ts} \leq T_i$. It yields a large theoretical collision frequency

$$v_{a,ts} \simeq \omega_{pi} \gg v_{a,ic} \quad \text{with} \quad v_{a,ic} \sim 0.3 k_{\parallel} v_e < \omega_{ci} \quad (3.108)$$

of the order of the ion plasma frequency and several orders of magnitude larger than Spitzer-Coulomb collision frequency. As such it is a serious candidate for generating anomalous dissipation and heating in shocks whenever a two-stream situation is encountered. This is indeed frequently the case as we will see in the supercritical shock Chapters 5 and 6.

If in the above Eq. (3.108) we compare the two-stream collision rate, for instance, with the ion-cyclotron collision rate $v_{a,ic}$ that had been favoured by *Kindel & Kennel* [1971], we find that it has far larger growth rate than the ion cyclotron wave and will thus always dominate when the current is strong, $V_{\parallel} > v_e$. At weaker currents the ion-acoustic instability will be faster (because of the slowness of the ion-cyclotron instability) as long as the magnetic field is weak and $V_{\parallel} \gg c_{ia}$.

Anomalous Lower-Hybrid Drift Collision Rate. As noted earlier, the lower-hybrid drift instability is particularly important as it is the exceptional representative of a fast growing (universal) drift wave instability which in the presence of gradients in plasma can always be expected to grow. Clearly shocks are a particularly good candidate for such an instability because of the steep density and magnetic field gradients occurring in compressive shocks. Moreover, in application to shocks, other than at pressure equilibrium boundaries like the magnetopause [see, e.g., *Treumann et al*, 1991; *Winske & Omidi*, 1995], the magnetic gradient adds positively to the growth rate.

In calculating the quasilinear saturation level of this instability makes use of the wave number at maximum growth $k_{max}^2 \lambda_{Di}^2 = 2/(1 + \omega_{pe}^2/\omega_{ce}^2)$ and $\partial(\omega\epsilon)/\partial\omega = 1 + \omega_{pe}^2/\omega_{ce}^2 \gg 1$ in dense plasma like the vicinity of the bow shock, for instance. The saturation wave level [Davidson, 1978] then follows after solving the quasilinear diffusion equation to be fraction of the drift energy

$$W_{sat,th} \simeq \frac{m_e N V_{de}^2}{8(1 + \omega_{pe}^2/\omega_{ce}^2)} \quad (3.109)$$

From here the anomalous collision rate follows as

$$v_{a,lh} \simeq \sqrt{\frac{\pi}{2}} \left(\frac{\omega_{pe}^2}{\omega_{lh}} \right) \frac{W_{sat,lh}}{NT_i} \approx \sqrt{\frac{\pi}{2}} \left(\frac{r_{ci}}{4L_N} \right)^2 \omega_{lh} \quad (3.110)$$

proportional to the ratio of ion-gyroradius to density gradient scale length $L_N = |\nabla \ln N|^{-1}$. This growth rate is large for steep gradient scales close to the ion gyro-radius, a condition that holds in the shock ramp. The propagation of electrons in the lower-hybrid drift case is perpendicular to the magnetic field since the electrons perform a diamagnetic drift which constitutes the electric current. In the above one-dimensional theory everything has been reduced to the coordinate parallel to the current. In a shock wave this current will flow along the shock surface, perpendicular to the magnetic field and shock ramp density/field gradient.

The anomalous lower-hybrid drift collision frequency is large and renders the lower-hybrid drift instability a viable candidate for generation of anomalous dissipation in shocks if only the condition that maximum growth is found for $\beta < 1$ can be circumvented. Shocks in the heliosphere in most cases satisfy the inverse condition $\beta \gtrsim 1$. For the same reason anomalous lower-hybrid drift collision frequencies are not believed to be involved in reconnection, for instance, except in the presence of a very strong guide field. Whether this condition on β indeed provides a serious restriction has not yet been clearly verified, however, neither theoretically nor in numerical simulation or observation.

Recently *Yoon & Lui* [2007] reviewed the theory of anomalous resistivity for the lower-hybrid drift and two-stream instabilities and derived some (slightly) more precise (but considerably more involved) expressions than the formulae given above. They also included electromagnetic effects and arbitrary directions of propagation to maintain that the lower-hybrid drift anomalous resistance can be very large indeed, up to a factor of 10^{10} larger than Spitzer-Coulomb resistance.

Runaway Effects. Since the collision frequency is an inverse function of the particle current drift velocity it allows for the interesting effect that particles of sufficiently high speed cannot be captured by the electric field. They instead get another push and escape as so-called runaway electrons. This effect (the physical reason of which is momentum balance) is known since long and applies to some of the anomalous collision processes as well as to Spitzer collisions. The reason is that the wave level saturates yielding a constant collision frequency for every mode in question. Hence, fast particles do simply not interact but escape like in free ballistic flight. Hence there will always be some energetic particles which behave like freely streaming particles. These, when flowing along the magnetic field, constitute a moderately energetic particle beam and may provide a seed particle population for further acceleration even in the presence of anomalous collisions.

Other Effects. Several other effects have not been mentioned here in relation to anomalous effects. These are wave decay processes, generation of radiation, and resonant wave absorption processes in inhomogeneities. It is interesting to note that, historically, *Sagdeev* [1966] in following calculations of *Moiseev & Sagdeev* [1963] proposed that such wave

decay processes would contribute substantially to anomalous dissipation in subcritical shocks by enhancing the number of waves present and thus enhancing the probability of particle scattering by waves.

Wave Decay Processes. The first of these belong to the class of wave-wave interaction which are the pendant of collisions between particles in the wave picture without invoking particles. Such interactions must satisfy the wave momentum and wave energy conservation laws which can be written simply as

$$\sum_{\alpha} \mathbf{k}_{\alpha} = 0, \quad \sum_{\alpha} \omega_{\alpha}(\mathbf{k}) = 0 \quad (3.111)$$

where the index α marks the particular wave mode, and each pair $[\omega_{\alpha}(\mathbf{k}), \mathbf{k}_{\alpha}]$ satisfies its own linear (or nonlinear) dispersion relation. According to the number of waves involved, the smallest possible number is 3, there are three-wave interactions, four-wave interactions, ... with for weak interactions decaying probability. For the three-wave case one important case is that a strong wave (one that has become very strong due to linear wave growth) such a three wave process can lead to wave decay into two other weaker modes of different dispersion due to the relations

$$\omega_0 \rightarrow \omega_1 + \omega_2, \quad \mathbf{k}_0 + \mathbf{k}_1 + \mathbf{k}_2 = 0, \quad \omega_0 \gg (\omega_1, \omega_2) \quad (3.112)$$

We do not go into the details of these process; they are rather involved because of the complexity of the three dispersion relations which have to be taken into account and which can be different when waves from one branch jump over into another branch feeding a wave mode there. These processes have, however, a number of consequences of which four are important:

- They contribute to the excitation of wave modes in a plasma which are not directly driven by an instability but result from the decay of an instability-driven large amplitude intense wave, in this way contributing to the production of a broad spectrum of turbulence that consists of many different and even possibly weakly damped modes in the plasma which otherwise would not exist, when only the decay is stronger than the natural damping of the wave.
- Decay processes limit the intensity of an instability-driven mode and reduce it substantially to the advantage of other modes. In this way they weaken the contribution of the particular mode to anomalous collisions while they might enhance collisionality by producing a broad turbulent background fluctuation spectrum.
- By generating other waves in the plasma they may provide a background from which other instabilities may grow which are driven by sources which otherwise would not overcome the instability threshold.
- The broad spectrum produced in plasma wave decay processes may move upstream of the shock and modify the upstream conditions in a way not foreseen in the Rankine-Hugoniot relations. Hence such processes cannot be handled in simple

plasma modelling of shock wave generation. They can only be investigated by particularly tailored numerical simulation techniques.

Radiation. Only radiation generated from plasma waves is of interest in heliospheric shocks because the densities are generally far too low for reaching a substantial emission measure in synchrotron or x-ray radiation. Radiation can then be generated predominantly by mode conversion or mode coupling. The difference between the two is that in mode conversion an intense high frequency plasma wave propagating up a density gradient gradually transforms into a free space mode while mode coupling or mode decay involves at least three modes in a stimulated wave collision process which also works in homogeneous plasma.

The more interesting case is a special case of the wave decay process, it is in fact its inversion, when two plasma waves join in interaction to inject their energies into a high frequency radio wave that is able to escape from the plasma. Radiation production is thus always a process of energy loss that leads to cooling of the plasma. However, in the heliosphere this cooling is weak and can be safely neglected. this kind of radiation is in fact degraded to an energetically completely unimportant process that has only indicative power.

The free space modes can be either an ordinary or an extraordinary wave, both propagating above their upper cut-off frequencies which for the left-hand polarised ordinary (LO) mode with dispersion $\omega^2 = \omega_{pe}^2 + k^2 c^2$ is the plasma frequency, and for the right-hand extraordinary (RX) mode is a cut-off frequency slightly higher than the upper-hybrid frequency $\omega_{uh}^2 = \omega_{pe}^2 + \omega_{ce}^2$.

The following wave-wave processes are of interest in generating radiation:

- the interaction of two counter-streaming electron plasma (Langmuir) waves into the transverse (T) electromagnetic wave, following the symbolic process $L + L \rightarrow T$, where the symbols L, T mean the tuples (ω_L, k_L) and (ω_T, k_T) , respectively. The energy and momentum conservation equations have been given above in Eq. (3.112). This process produces a transverse wave with nearly perpendicular propagation $k_T \ll k_L$ in the RX-mode and of frequency $\omega_T \simeq 2\omega_{pe}$, which can clearly propagate above the cut-off in weakly magnetised plasma,
- the process $L + L' \rightarrow T$ of interaction of a Langmuir wave with another Langmuir wave that has been scattered off thermal ions (i) according to the process $L + i \rightarrow L' + i^*$, where the prime indicates the scattered Langmuir wave, and the star on the ion the excited ion. The wave interaction in this case causes a lower frequency transverse wave still above the plasma frequency but closer to the cut-off of the RX-mode,
- a process similar to the one under the first item in which harmonic Langmuir waves have been generated in L–L interaction. This yields weak higher plasma harmonic radiation at frequencies $\omega \lesssim l\omega_{pe}$ with harmonic numbers $l = 3, \dots$ in the RS-T mode with intensity that decreases steeply with increasing harmonic number l ,

- merging of a Langmuir and an ion acoustic wave (S) according to $L + S \rightarrow T$, which produces a transverse wave with frequency above but very close to ω_{pe} . Whether this wave can escape from the plasma depends on whether its frequency is above the cut-off of the RX-mode, which is possible only for high frequency S waves, or whether it can tour into the OL mode for which generally the conditions of being selected in the merging process are worse than for the RX mode,
- merging of another electrostatic plasma wave like, for instance, an electron acoustic mode (EA) with a Langmuir mode according to $L + EA \rightarrow T$. Since the EA wave is near $\frac{1}{2}\omega_{pe}$, this produces a T wave at frequency $\omega \sim \frac{3}{2}\omega_{pe}$ well above the cut-off still even though still counted as fundamental mode radiation,
- then, the lower-hybrid mode can also merge with the Langmuir mode. The product of this process can, however, not escape in the RX mode as it usually does not exceed the cut-off. On the other hand, it could excite the LO mode,
- finally several nonlinear processes exciting transverse free space electromagnetic waves have been proposed like Langmuir wave collapse.

Collapse is very attractive because Langmuir waves become trapped in this process inside deep density depressions where they bounce back and forth between the walls. Thus in quite a natural way collapse generates counter-streaming Langmuir waves in highly localised places. These waves are particularly well configured for merging and escaping in the RX mode.

Moreover, during collapse the internal plasma frequency decreases rapidly, causing as well a decrease of the radiation frequency. This allows only the higher frequency part of the radiation with frequency just above the plasma frequency of the environment to escape. Radiation occurs as intense fundamental but highly bursty radiation at the plasma frequency because at the end of the collapse the intensity of the waves explodes and therefore the emitted power also explodes.

However, experimentally this process could never been proven. It has been replaced by another mechanism known as ‘stochastic growth’ which itself is doubtful as well but very popular. It takes into account the stochastic modulation of Langmuir growth in a medium of spatially fluctuating density, i.e. containing a broad spectrum of weak ion acoustic turbulence. This causes the growth rate to experience spatial modulations leading to exponential modulation of the Langmuir amplitude. Contribution to the intensity comes only from localised places, and thus the volume contributing to radiation is a fraction only of the total volume while locally radiation may be rather intense.

All these radiation processes are nonlinear of second order in the amplitudes of the involved waves. A direct process would be the linear unstable excitation of one of the free-space modes. This process, called electron-cyclotron maser, is well know to work in a highly underdense plasma with $\omega_{pe} \ll \omega_{ce}$ of the kind present in the magnetospheres of the magnetised planets under auroral conditions [for a review see *Treumann, 2006*]. In the shock environment we have instead $\omega_{pe} \gg \omega_{ce}$ and this condition is strongly violated.

On the other hand, it is not unlikely that in the shock ramp strong electric fields may accelerate electrons such that they either evolve into a loss cone distribution or form a ring distribution in phase space, both distributions favouring conditions for the electron cyclotron maser to work. In a dense plasma the maser can excite free space modes only in oblique direction at higher electron cyclotron harmonics. At Earth's bow shock we have $\omega_{pe}/\omega_{ce} \sim 50-100$. This implies that in order to reach the free space X-mode the maser should work at a very high harmonic, which is very unlikely as the maser growth rate readily decreases with harmonic number. However, excitation of the O-mode and Z-mode – the quasi-electrostatic lower frequency X-mode branch $\omega < \omega_{uh}$ –, the latter near the upper hybrid frequency ω_{uh} , may become possible. The growth rate for the O-mode is small while that for the Z-mode may be higher.

If the conditions for excitation of these modes are satisfied then it might be possible that the maser will also work in a supercritical shock wave in the presence of a strong shock ramp electric field. *Yoon et al* [2007] have recently, in a brute force numerical solution of the relevant loss-cone maser dispersion relation, tried to determine conditions when the growth rates become sufficiently high for maser excitation. While the O-mode can in principle leave the plasma directly, the Z-mode must find a way to escape from the dense plasma to free space. If excited in the shock ramp, this way may be provided by mode conversion from the shock Z-mode to the upstream X-mode or, otherwise, by coupling across the boundary of the Z-X-mode wave guide. The advantage of this kind of maser excitation, if realised in shocks, would be that it would be capable of explaining the so-called shock backbone radiation in solar and interplanetary type II radio bursts.

3.5.3 Shock Particle Reflection

The process of particle reflection from a shock wave is one of the most important processes in the entire physics of collisionless shocks, as we have noted already in several places. The mechanism of particle reflection has not yet been fully illuminated, however.

Particle reflection is required in supercritical shocks as, to our knowledge, it is the only process that can compensate for the incapability of dissipative processes inside the shock ramp to digest the fast inflow of momentum and energy into the shock. Shock particle reflection is not dissipative by itself even though in a fluid picture which deals with moments of the distribution function it can be interpreted as kind of an ion viscosity *Macmahon* [1965], i.e. it generates an anomalous viscosity coefficient which appears as a factor in front of the second derivative of the ion velocity in the ionic equation of motion. As such it also appears in the ion heat-transport equation. The kinematic ion viscosity can be expressed as

$$\mu_{vis} = m_i N v_i \lambda_{mfp} \simeq P_i / 2 \omega_{ci} \quad (3.113)$$

through the ion pressure P_i and the ion-cyclotron frequency ω_{ci} when replacing the mean free path through the ion gyro-radius. In this sense shock particle reflection constitutes by itself a very efficient non-dissipative dissipation mechanism. However, its direct dissipative action is to produce real dissipation as far as possible *upstream of the shock* in order to dissipate as much energy of motion as remains to be in excess after formation of a shock

ramp, dissipation inside the ramp, and reflection of ion back upstream. The shock does this by inhibiting a substantial fraction of inflow ions to pass across the shock from upstream into the downstream region. It is sending these ions back into the upstream region where they cause a violently unstable upstream ion beam-plasma configuration which excites a large amplitude turbulent wave spectrum that scatters the uninformed plasma inflow, heats it and retards it down to the Mach number range that can be digested by the shock. In this way the collisionless shock generates a shock transition region that extends far upstream with the shock ramp degrading to the role of playing a subshock at the location where the ultimate decrease of the Mach number from upstream to downstream takes place.

Shock reflection has another important effect on the shock as the momentum transfer from the reflected particle component to the shock retards the shock in the region of reflection thereby decreasing the effective Mach number of the shock.

The outcome of the previous paragraphs is that shock particle reflection is of incomparable importance in shock formation and in the understanding of collisionless shock physics. On the other hand it is barely understood and can, in principle, be treated only by numerical simulations. Before, in the next chapters, coming to discuss those problems in greater depth we will present below a few attempts to understand shock reflection.

Specular Reflection

Specular reflection of ions from a shock front is the simplest case to be imagined. It requires that the ions experience the shock ramp as an impenetrable wall. This can be the case when the shock itself contains a positive reflecting electric potential which builds up in front of the approaching ion. Generation of this electric potential is not clarified yet. In the soliton picture the shock potential is related to the density gradient, however, dissipation processes are also involved. Understanding its formation requires understanding the entire collisionless shock physics. In a very naive approach one assumes that in flowing magnetised plasma a potential wall is created as the consequence of charge separation between electrons and ions in penetrating the shock ramp. It occurs over a scale typically of the spatial difference between an ion and an electron gyro-radius, because in the ideal case the electrons, when running into the shock ramp, are held temporarily back in the steep magnetic field gradient over this distance while the ions feel the magnetic gradient only over a scale longer than their gyro-radius and thus penetrate deeper into the shock transition.

Reflection from Shock Potential. Due to this simplistic picture the shock ramp should contain a steep increase in the electric potential $\Delta\phi$ which will reflect any ion which has less kinetic energy $m_i V_N^2/2 < e\Delta\phi$ (see Figure 3.17). This condition contains the perpendicular ion velocity component along the shock normal. Since the ion gyrates it depends on the instantaneous angle the ion velocity has with respect to the magnetic field at the location of the ramp. In this reflection the ion velocity component $-V_N \rightarrow +V_N$ simply changes sign. For the gyrating particles this component adds up of the normal components of the bulk flow velocity V_N^f and the microscopic particle speed $\mathbf{v} = (v_\perp \cos \alpha, v_\perp \sin \alpha, v_\parallel)$, with α the gyration angle, projected on the direction of the shock normal $\mathbf{n}$. This yields

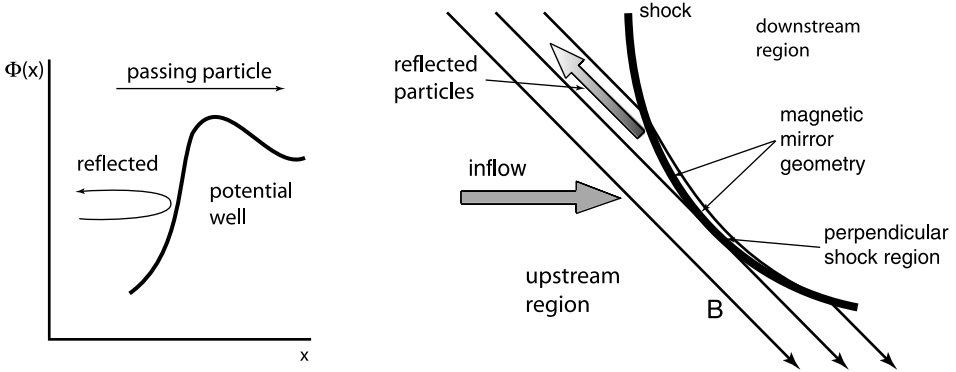


Figure 3.17: The two cases of shock reflection. *Left:* Reflection from a potential well $\Phi(x)$. Particles of energy higher than the potential energy $e\Phi$ can pass while lower energy particles become reflected. *Right:* Reflection from the perpendicular shock region at a curved shock wave as the result of magnetic field compression. Particles move toward the shock like in a magnetic mirror bottle, experience the repelling mirror force and for large initial pitch angles are reflected back upstream.

$$-V_N = -V_N^f + v_{\parallel} \cos \Theta_{Bn} + v_{\perp} \cos \alpha \sin \Theta_{Bn}$$

and the condition for specular reflection can be written as

$$(-V_N^f + v_{\parallel} \cos \Theta_{Bn} + v_{\perp} \cos \alpha \sin \Theta_{Bn})^2 < 2e\Delta\phi/m_i$$

This is a condition on the gyration angle α restricting the gyration phases of the reflected particles. For a gyrotropic upstream distribution one can average over all gyration angles from 0 to $\pi/2$ since only upstream directed velocity components reduce the inflow velocity to values below the reflection threshold, obtaining

$$(v_{\parallel} \cos \Theta_{Bn} - V_N^f)^2 + \frac{4}{\pi} (v_{\parallel} \cos \Theta_{Bn} - V_N^f) v_{\perp} \sin \Theta_{Bn} + \frac{1}{\pi} v_{\perp}^2 \sin^2 \Theta_{Bn} < \frac{2e\Delta\phi}{m_i}$$

This condition must be used on one of the velocity components $v_{\parallel}$, $v_{\perp}$ when determining the number of specularly reflected particles from the upstream ion distribution function $F_i^{up}(v_{\parallel}, v_{\perp})$.

Mirror Reflection. Another simple possibility for particle reflection from a shock ramp in magnetised plasma is mirror reflection. An ion approaching the shock has components $v_{i\parallel}$ and $V_{\parallel} = V_1 \cos(\mathbf{V}_1 \cdot \mathbf{B}_1)$ along the magnetic field. Assume a curved shock like Earth's bow shock. Close to its perpendicular part where the upstream magnetic field becomes tangential to the shock the particles approaching the shock with the stream and moving along the magnetic field with their parallel velocities experience a mirror magnetic field configuration that results from the converging magnetic field lines near the perpendicular point (compare the right part of Figure 3.17). Conservation of the magnetic moment $\mu = T_{i\perp}/B$ implies that the particles become heated adiabatically in the increasing field;

they also experience a reflecting mirror force $-\mu\nabla_{\parallel}B$ which tries to keep them away from entering the shock along the magnetic field. Particles will mirror at the perpendicular shock point and return upstream when their pitch angle becomes 90° at this location [this theory has been developed in detail by *Leroy & Mangeney, 1984; Wu, 1984*]. We will return to this mechanism in Chapter 7 when describing shock particle acceleration.

A rough estimate of the marginal upstream pitch angle for mirror reflection can be given from conservation of μ . Since the increase in field strength is according to the magnetic gradient across the shock ramp one has roughly at the perpendicular shock point $B = B_1 + (\nabla_n B)\Delta$ where Δ is the shock width and ∇_n the field gradient across the shock. Hence, to lowest order, $B = B_1 + \delta B$ with δB the magnetic compression. This yields for the upstream pitch angle at reflection $\sin^2 \alpha_1 > (1 + \delta B/B_1)^{-1}$. With compression factor $\delta B/B_1 \sim 3$ particles of upstream pitch angles $\alpha_1 > 60^\circ$ will become mirror reflected from the perpendicular shock area due to the action of the mirror force, a condition which has to be used upstream in the inflowing distribution if one wants estimating the fraction of reflected particles. This requires knowledge of the compression factor, however. The compression factor and the number of reflected particles are not independent. Hence, a selfconsistent determination requires developing the full shock theory. This can be done only by numerical simulations.

Of course, the above estimate is very crude. It demonstrates, however, that a fraction of upstream particles can, in principle, become reflected from a curved shock surface by mirroring in the converging magnetic field geometry around the perpendicular area of the shock. For fast flows reflection will always be located on the nose inflow side of the shock. This holds for ions as well as for electrons. Reflection of both sorts of particles has continuously been observed at the bow shock as is shown schematically, for instance, in Figure 2.6 which has been drawn for condition in front of a perpendicular shock at a distance inside the foot, roughly within $< 1 r_{ci}$ from the shock ramp.

Consequences of Shock Reflection. How far the reflected ions return upstream depends on the direction of the magnetic field with respect to the shock, i.e. on the shock normal angle Θ_{Bn} . For perpendicular shocks the reflected ions only pass just one gyro-radius back upstream. Seeing the convection electric field $|E_y| = |V^f B|$ they become accelerated along the shock forming a current, the velocity of which in any case exceeds the inflow velocity (which is zero in the perpendicular direction) and for sufficiently cold ions also the ion acoustic velocity c_{ia} in which case the ion-beam plasma instability will be excited in the shock foot region where the ion current flows. This may generate and anomalous collision frequency in the shock foot region. Moreover, since the excited waves accelerate electrons along the magnetic field other secondary instabilities can arise as well.

In quasi-perpendicular and oblique shocks the ions can escape along the magnetic field. In this case an ion-beam/ion beam situation arises between the upstream beam and the plasma inflow with the consequence of excitation of a variety of instabilities, electromagnetic and electrostatic. In addition, however, an ion beam-electron beam two stream situation is caused between the upstream ions and the inflow electrons which because of

the large upstream electron temperatures probably excites mainly ion-acoustic modes but can also lead to Buneman two-stream mode excitation. These modes contribute to turbulence in the upstream foreshock region creating a weakly dissipative state in the foreshock where the plasma inflow becomes informed about the presence of the shock. The electromagnetic low frequency instabilities on the other hand, which are excited in this region, will grow to large amplitude, form localised structures and after being convected by the main flow towards the shock ramp interact with the shock and modify the shock profile or even contribute to shock formation and shock regeneration.

Specular reflection from shocks is the extreme case of shock particle reflection. Other mechanisms like turbulent reflection are, however, not well elaborated and must in any case be investigated with the help of numerical simulations. In the following subsection we will in passing encounter one such mechanism in phase space hole formation in a two-stream unstable collisionless plasma.

Figure 3.18 shows the expected particle distributions for shock reflection of ions in a perpendicular shock right in the foot region. The upper part of this figure is perpendicular to the shock in flow direction. The lower part is along (tangential to) the shock surface. The different configurations of the distributions in this region can lead to the excitation of ion acoustic and two-stream instabilities. Ion beam-ion beam interaction is expected in the direction perpendicular to the shock in addition to ion-acoustic instability between the reflected ions and the main electron component in the region of positive gradient on the electron distribution. In the direction parallel to the shock surface, on the other hand, one expects an ion beam-ion interaction and a two stream instability of the main electron component with the accelerated along the shock surface reflected ion beam. Inside the ramp conditions are more involved and will be described in more detail in the context with observations in the respective chapters for both quasi-perpendicular and quasi-parallel shocks.

We note in addition that there is a peculiarity concerning quasi-parallel supercritical shocks. Due to the presence of an intense reflected and transformed ion component in the foreshock of a quasi-parallel shock there is a broad spectrum of large amplitude low frequency electromagnetic waves which are convected by the inflow stream towards the shock ramp, steepen and interact with the shock. These waves are predominantly transverse having components tangential to the shock on scales of the order and shorter than the ion gyro-radius. Consequently, quasi-parallel shocks remain to be quasi-parallel for ions, in particular for the more energetic accelerated ions. However, for the electrons all quasi-parallel supercritical shocks become quasi-perpendicular in the vicinity of the shock ramp transition such that for electrons no quasi-parallel supercritical shocks exist. This has the consequence that electrons will become reflected and accelerated all-over the shock front independent of its quasi-parallel character. We will later prove this statement by referring to observations and simulations.

Hole Formation

As the last item in this section we consider the stability of an high velocity $V_b > v_e$ electron beam or electron current. In the fluid picture we have found that such high speed

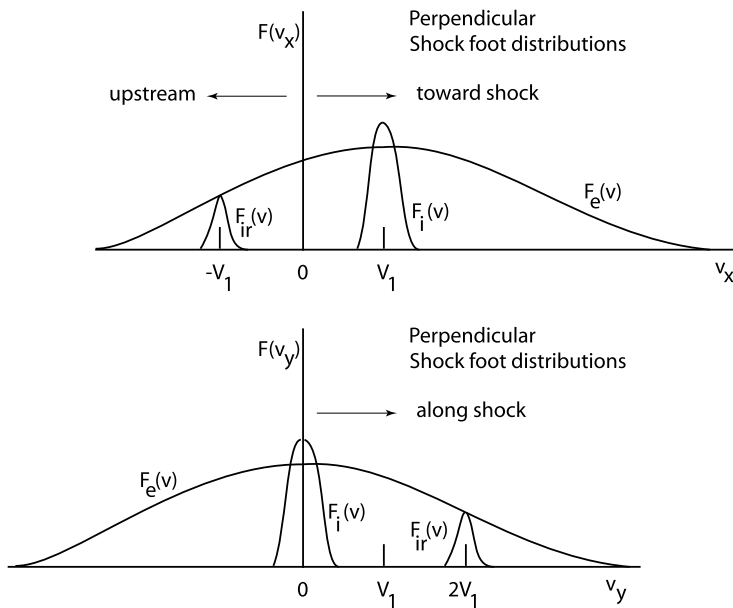


Figure 3.18: Expected particle distributions in the shock foot of a supercritical perpendicular shock which reflects ions back upstream. *Top:* Distribution functions in flow direction. The ion and electron distributions flow with velocity V_1 in shock direction. For the electrons with their high thermal speed the flow velocity is practically negligible. Reflected ions in the foot occur at velocity $-V_1$ on the scale of about one ion gyro-radius. This yields a two-ion beam configuration which is electrostatically unstable. Moreover, ion acoustic waves are excited to the right of the reflected ion distribution by the ion-acoustic instability in the positive slope of the hot electron distribution. *Bottom:* Same in the direction along the shock front. Inflowing ions and electrons have only thermal velocities in this direction. The reflected ions are accelerated to about twice the inflow velocity in the tangential inflow electric field $|\mathbf{E}| = |\mathbf{V}_1 \times \mathbf{B}_1|$. This causes an unstable ion beam-ion configuration and a two-stream configuration between accelerated reflected ions and electrons.

currents flowing through the plasma along the magnetic field or (in the presence of steep density gradients) perpendicular to both, the gradient and the magnetic field, will excite the fast growing Buneman two-stream instability. It has, however, been predicted early [Schamel, 1972; Dupree et al, 1975] that currents of this strength will undergo a kinetic instability which structures the electron and ion phase space into so-called phase space holes, which are regions of lacking particles localised in phase space while in real space represent localised electric fields, trapped particles and particle acceleration. Such holes have meanwhile been found to exist all-over in collisionless space plasmas in relation to spatially localised strong current flow as in reconnection, auroral phenomena, and also in shocks [Bale et al, 2002]. Since strong currents are expected in shocks as well in the ramp as in the foot, as we have discussed above, it is not unreasonable to assume that phase space holes might form under shock conditions as well.

Hole formation follows a nonlinear interaction known as Bernstein-Green-Kruskal (BGK) mode formation [cf., e.g., Davidson, 1972] and is based on the splitting of the

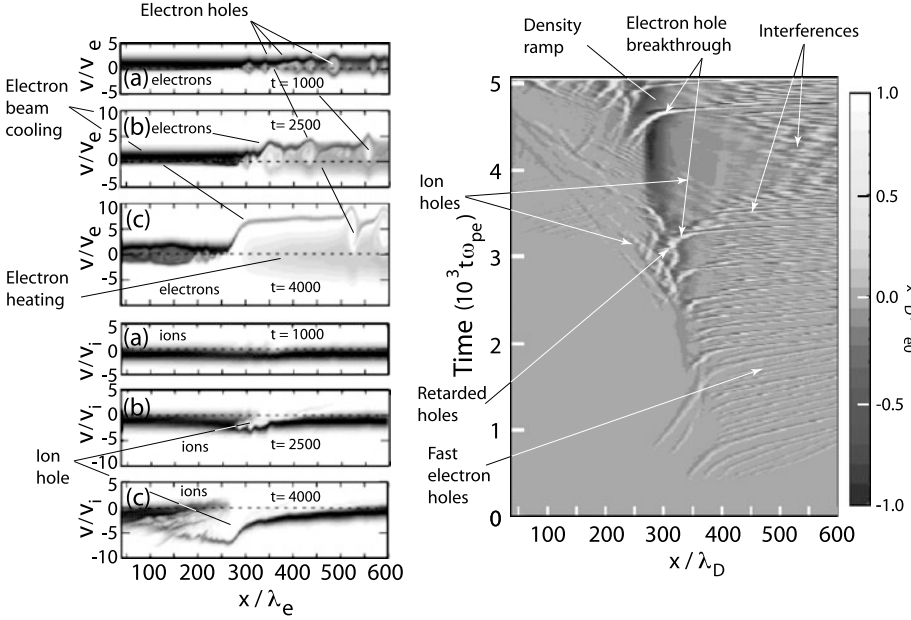


Figure 3.19: *Left:* Electron and ion phase space plots at three simulation times of the interaction of a marginally two-stream unstable plasma with a shock ramp. The ramp has been modelled as a density dip (potential wall for electrons) in the centre of the simulation box. The box has one space and one velocity coordinate. The fast electron beam current is injected into the quiet ion plasma and [after Newman *et al.*, 2001]. *Right:* Time history of the electric field during the interaction with the shock ramp showing the evolution of electron and ion phase space holes and their interaction [after Newman *et al.*, 2002].

phase space distribution function into two components, particles that are energetic enough to surpass the electric potential of a localised electric field inside a soliton, for instance, and particles of lesser energy that become either trapped or reflected from these potentials depending on the sign of particle and potential. Even if the potential is repulsive and ejects, say, electrons from the region, some trapped electrons will remain there performing oscillations and become heated up to a certain energy that leaves them still trapped. These electrons are in disordered motion and are assuming a high temperature, while the rejection and expulsion of other particles from the potential site causes their acceleration. This mechanism is quite complicated and has been treated analytically only up to a certain approximate degree in the above cited papers. In order to investigate it one better performs numerical simulations.

Figure 3.19 shows the example of a simulation of hole formation in interaction of a marginally two-stream stable current (of electron and ion of opposite bulk velocity and same initial temperature $T_e = T_i$ with a localised inhomogeneity [from Newman *et al.*, 2001, 2002]. The localised plasma inhomogeneity has been modelled as a simple density dip $\propto -\cos^2(x - x_0)$ at the centre. This is not a shock, it is, however, a potential wall which should reflect one sort of particles, in this case electrons.

The figure shows on the left phase space plots at three different simulation times for electrons and ions when due to electron reflection at the potential ramp a strong two-stream instability evolves. The holes formed on the electron distribution appear early in the simulation as egg-like distortions of the distribution in the reflected electron component. Widening of the distribution indicates the strong heating of the electrons. The holes contain dilute trapped electron, and some part of the beam becomes accelerated. At later times the heating becomes violent with a strong broadening of the electron distribution when the ion hole starts forming in the lower panels. Strong acceleration of a narrow and thus very cool electron beam is also observed in the final state. In addition the holes move along the beam, while the ion hole moves in the opposite direction.

Most interesting is the time history shown on the right in the figure. It shows the initial evolution of many small amplitude electron holes moving at fast speed to the right away from the potential ramp. At later times the ramp steepens, and the electron holes start interacting with the ion hole which moves slowly to the left. The holes intensify and finally can break through the potential ramp to escape to the right where a whole fabric of interfering holes evolves.

The importance of this observation is that the two-stream instability can form as a cause of reflection at a potential ramp. This is expected for shocks as well. Moreover, the instability causes electron and ion phase space holes to evolve and leads to completely collisionless heating due to electron trapping inside the holes, i.e. it causes irreversible heating and entropy which is needed for shocks, and it also generates a very cool electron beam to escape from the holes by continuous acceleration and collimation of cool but fast electrons. This is a very interesting and important mechanism which is capable of injecting a fast seed particle population into shock acceleration. As a side product it is a method of cooling electrons and generating cold electron beams.

3.6 Briefing on Numerical Simulation Techniques

The modern age of physics is to a large degree determined by the availability of high speed and high capacity computer systems. The use of these computing facilities for performing numerical experiments on collisionless plasmas covers now almost half a century of experience. It started with the introduction of Fermi's Monte-Carlo method and blossomed after the formulation of the Fermi-Ulam 1961 numerical model approach to cosmic ray acceleration [cf., e.g., *Lichtenberg & Lieberman, 1992*] which was based on nonlinear particle motions in electromagnetic fields. The modern state of the art in application to plasma physics has been formulated in several textbooks [e.g., *Birdsall & Langdon, 1985*] and review articles [e.g., *Dawson, 1983, 1995*]. Many problems in plasma and in particular space plasma physics with their enormous complexity could not have been solved or even attacked without computers and numerical simulations. Also, most of the discussion on shocks in the following chapters will be based on such numerical simulations which must accompany observation and experiment in order to understand what is going on in the shock environment. A brief discussion about numerical methods is therefore not only unavoidable but even necessary.

3.6.1 Basic Equations

Computers are capable of dealing simultaneously with the dynamics of many particles as we have already described in Chapter 2. The most fundamental approach in numerical simulations is hence based on the full Newton-Maxwell set of microscopic equations of which we write down here only the Newtonian subset

$$\frac{d\mathbf{v}_{is}}{dt} = \frac{e_s}{m_s}(\mathbf{E} + \mathbf{v}_{is} \times \mathbf{B}), \quad \frac{d\mathbf{x}}{dt} = \mathbf{v} \quad (3.114)$$

The connection to Maxwell's equations is done by the microscopic definition equations of the space charge $\rho(\mathbf{x}, t)$ and conduction current $\mathbf{j}(\mathbf{x}, t)$ according to

$$\rho(\mathbf{x}, t) = \sum_{is} e_s \delta(\mathbf{x} - \mathbf{x}_{is}), \quad \mathbf{j}(\mathbf{x}, t) = \sum_{is} e_s \mathbf{v}_s \delta(\mathbf{x} - \mathbf{x}_{is}) \quad (3.115)$$

In fact, these are the most general equations of a classical plasma consisting of i point-like space charges e_s of species s with mass m_s and momentum $m_s \mathbf{v}_{is}$ located at time t at location $\mathbf{x}_{is}(t)$. The point-like character is taken care of by the δ -functions. This whole set is the set of Liouville's equations [which could also formally be combined into one single *exact* equation in phase space $(\mathbf{x}, \mathbf{v})$ by introducing an *exact phase space distribution* function $F(\mathbf{x}, \mathbf{v}) = \sum_{is} \delta(\mathbf{x} - \mathbf{x}_{is}) \delta(\mathbf{v} - \mathbf{v}_{is})$. Such an equation is known as the Liouville equation].

These equations can be simplified depending on the nature of the problem. For instance, when electromagnetic effects are not of interest, the magnetic field will drop out and one uses only Poisson's equation for the electrostatic potential and the electric field in Ampère's law. When the plasma is collisionless one can use another simplification, i.e. replace the Liouville distribution with the one-particle distribution and switch to the Vlasov equation. This then produces Vlasov-codes. When one is interested only in low frequency responses of the plasma, the electrons can be treated as Boltzmannian electrons, and the electron equations are replaced by the Boltzmann dependence of the electron density on the electric potential with the ions being treated as single particles. The corresponding codes that have found extensive application in shock physics are the hybrid codes [for a brief description see, e.g., Winske, 1985a]. An even stronger simplification is the assumption of quasi-neutrality when the Poisson equation is replaced by the condition $N_e = N_i$ and the electrons are merely an instantaneously reacting neutralising background. The strongest simplification consists in the recourse to the replacement of the above kinetic equations by hydrodynamic equations, i.e. by treating electrons and ions as fluids respectively even as one single fluid. Such codes are not numerical experiments in the above sense of the words, rather they are methods to solve the hydrodynamic (or magneto-hydrodynamic) system of equations. They are not suited for the investigation of the evolution of collisionless shocks except for determining their global shapes and the gross features of the plasma fluid flow behind them around the obstacle.

3.6.2 General Methods

Either of the resulting sets of equations must be represented in digitised form in order to be prepared for treatment on a computer. One represents the spatial coordinates as a discrete grid and advances the equations over discrete time intervals. The choice of space and time steps is prescribed by the necessary accuracy of the result and by the stability of the code. The particles in the code are, however, no more point particles but of finite size, i.e. the Delta-functions assume bell-function shapes. This has the consequence that particle experience only small angle collisions well suited for studying collisionless plasmas. The integration of the equations is then performed in a series on a large number of such discrete steps of the finite sized particles the Lagrangean positions of which are used to deposit the charges and currents onto the fixed discrete grid points, followed by solving the field equations on these grid points of the many spatial cells. This produces self-consistent fields which are used in the next time step and which are smeared out over the volume by interpolation in order to construct a field continuum in which the particle orbits are advanced further. One is thus working with space and time differences and an interpolation from charge to grid and subsequently of the fields from grid to particle. In view of application to the heliospheric shocks in the following we briefly discuss the discretisation only for the electromagnetic case.

Units. Numerical methods require that all quantities including coordinates space and time are pure numbers. One thus needs to normalise them by introducing constant reference values, a density N_0 , charge-to-mass ratio e/m_0 (here m_0 is the particle rest mass either for electrons or for ions), velocity of light c (for instance, another choice is the upstream Alfvén velocity). Time is then normalised, for instance, to the plasma frequency expressed in these reference units $\omega_p^{-1} = (e^2 N_0 / \epsilon_0 m_0)^{-\frac{1}{2}}$, space in inertial lengths c/ω_p (either for ions or electrons), the electric potential for instance in $m_0 c^2 / e$ (or in units of electron temperature T_e / e), the magnetic vector potential in $m_0 c / e$. This choice of units is one of many possibilities only.

Discretisation. Both the field and particle equations must be discretised. The idea of discretisation is quite simple. One returns in history to the time just one step before Newton and Leibniz. Differential quotients become quotients of differences, higher order differential quotients become the corresponding powers of quotients of differences, mixed differential quotients become products of quotients of differences. The only trick is to assign the results to some location inside the difference interval, not necessarily the centre (!), and to do this properly. Also, time runs only in one direction: forward. Applying such a scheme one arrives at recursive equations which can be solved on a sufficiently powerful computer.

Let us assume that we reduce the electromagnetic set of equations to the electromagnetic wave equation for the vector potential component $\mathbf{A}$. This must be written in difference form

$$\frac{1}{4} \nabla^2 (\mathbf{A}^{(\frac{3}{2})} + 2\mathbf{A}^{(\frac{1}{2})} + \mathbf{A}^{(-\frac{1}{2})}) - \frac{1}{c^2 \Delta t^2} (1 + \beta D^2 \nabla^2) (\mathbf{A}^{(\frac{3}{2})} - 2\mathbf{A}^{(\frac{1}{2})} + \mathbf{A}^{(-\frac{1}{2})}) = -\mathbf{j}_T^{(\frac{1}{2})}$$

The superscripts indicate the time levels of the various terms, the ad hoc parameter β is introduced to modify the dispersion at short wavelengths, and ∇^2, D^2 are defined as

$$\nabla^2 A = \frac{A_{j+1,m} - 2A_{j,m} + A_{j-1,m}}{\Delta x^2} + \frac{A_{j,m+1} - 2A_{j,m} + A_{j,m-1}}{\Delta y^2}, \quad \frac{1}{D^2} = \frac{1}{\Delta x^2} + \frac{1}{\Delta y^2}$$

The transverse current density in the Coulomb gauge is $\mathbf{j}_T = \mathbf{j} - \nabla \partial \phi / \partial t$. Since $\nabla \cdot \mathbf{j}_T = 0$, the additional equation $\nabla^2 (\partial \phi / \partial t) = \nabla \cdot \mathbf{j}$ must be solved. The electrostatic potential $\phi^{(1)}$ is to be taken at full time step (1), while $\mathbf{j}^{(\frac{1}{2})}$ belongs to each half time step ($\frac{1}{2}$). The former is obtained from $\nabla^2 \phi^{(1)} = -\rho^{(1)}$. Finally, the fields follow from

$$\mathbf{E}^{(1)} = -\nabla \phi^{(1)} - \frac{\mathbf{A}^{(\frac{3}{2})} - \mathbf{A}^{(\frac{1}{2})}}{\Delta t}, \quad \mathbf{B}^{(1)} = \frac{1}{2} \nabla \times (\mathbf{A}^{(\frac{3}{2})} + \mathbf{A}^{(\frac{1}{2})})$$

Fields, charge and currents are defined at the centre of the cells. These equations must be solved with a so-called Poisson solver. In addition one needs to specify appropriate boundary conditions at the boundaries of the simulation box.

In a similar way one discretises the equation of motion of the particles. For this one defines $h = e\Delta t/m$ and obtains a centred form of the equation of motion as a recursion equation with unknown $\mathbf{v}^{(1)}$

$$\mathbf{v}^{(\frac{3}{2})} = \mathbf{v}^{(\frac{1}{2})} + h(\mathbf{E}^{(1)} + \mathbf{v}^{(1)} \times \mathbf{B}^{(1)})$$

In order to determine $\mathbf{B} = \frac{1}{2}(\mathbf{v}^{(\frac{3}{2})} + \mathbf{v}^{(\frac{1}{2})})$ the former equation is solved implicitly taking the scalar and vector products of the former equation with $\mathbf{B}^{(1)}$. Up to terms second order in Δt this yields the expression

$$\mathbf{v}^{(\frac{3}{2})} = \mathbf{v}^{(\frac{1}{2})} \left(1 - \frac{h^2}{2} B^2 \right) + h(\mathbf{E} + \mathbf{v}^{(\frac{1}{2})} \times \mathbf{B}) + \frac{h^2}{2} (\mathbf{E} \times \mathbf{B} + \mathbf{B} \mathbf{B} \cdot \mathbf{v}^{(\frac{1}{2})})$$

for use in the expression for $\mathbf{v}^{(1)}$ only. The set of equation obtained is then ready for computing.

However, there are two ways of computing, so-called explicit and implicit techniques. In the *explicit* technique one solves the equations as they are obtained after discretisation. In such an approach the internal errors will necessarily grow and at some stage become unstable such that the calculation must be truncated. One can artificially introduce some kind of damping of these growing error modes in order to suppress them. Justification for this is discussed in the literature [Birdsall & Langdon, 1985]. The *implicit* technique solves the problem by calculating *backward* in time [Friedman et al, 1981] such that the error modes decay away when time runs negative. We do not describe this technique here. It suffices to note that in this approach the dangerous unstable short-wavelength error modes disappear by definition and become partially eliminated from the system. The most efficient ways of calculation are the combinations of both methods.

All these methods work on a fixed grid. In application to shock ramp research the grid mesh has to be chosen refined enough for resolving shock structure. Recently *Karimabadi et al* [2005] developed a self-adaptive technique which takes care of the steeping and narrowing of a shock front in order to resolve its substructure. This is a significant progress in shock simulation technique.

We do not go into detail of the various methods and refinements of simulation techniques. Those readers who are interested in and prepare themselves for doing simulation research in collisionless shocks we rather refer to the mentioned basic literature on the techniques of numerical simulation.

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Collisionless Subcritical Shocks

Abstract. Subcritical shocks have Mach numbers below critical. This chapter gives a brief overview as a preparatory for the discussion of collisionless supercritical shocks in the next chapters. For subcritical shocks (anomalous) resistive or viscous dissipation in combination with dispersion is high enough to stabilise the shock transition. Observations of subcritical shocks in space are sparse, mostly restricted to interplanetary space. We discuss the theory of subcritical shocks, define the dissipation length and discuss the cold plasma subcritical shock model. Finally, the reasons for collisionless dissipation in subcritical shocks are illuminated.

4.1 Introduction

Subcritical shocks have Mach numbers $\mathcal{M} < \mathcal{M}_{\text{crit}}$, where the critical Mach number has been defined earlier and is determined by the condition that the flow speed in the downstream plasma behind the shock transition does not exceed the downstream sound velocity $c_{s,2}$ [Coroniti, 1970], a condition that has been used by Edmiston & Kennel [1984] and Kennel *et al* [1985] to calculate the angular dependence of the critical Mach number. For strictly perpendicular shocks it was found that the critical Mach number is $\mathcal{M}_{\text{crit}} = 2.76$ and gradually decreases with the shock becoming more and more quasi-parallel. The exact value depends, in addition, on the upstream plasma β . For large $\beta > 1$ the critical Mach number approaches $\mathcal{M} \sim 1$, while for low $\beta < 1$ its limiting value is $\mathcal{M} \sim 1.53$. Thus subcritical shocks are low Mach number shocks in the above sense that their Mach numbers are less than the critical Mach number at the given shock-normal angle Θ_{Bn} and upstream β . One could, of course, in the same way also define a critical shock-normal angle $\Theta_{Bn}^{\text{crit}}$ such that, for given plasma β and Mach number $\mathcal{M}$, shocks with $\Theta_{Bn} > \Theta_{Bn}^{\text{crit}}$ are subcritical while in the opposite case they are supercritical. Therefore relatively low Mach number curved shocks are vulnerable against consisting of regions where they behave supercritical and other regions where they behave subcritical. The calculation [Kennel *et al*, 1985] shows that quasi-perpendicular low Mach number shocks are more likely to behave subcritical than quasi-parallel shocks. Figure 4.2 shows how narrow the range is of collisionless subcritical shocks in the Mach number-plasma beta space. For parallel shocks a noticeable subcritical range exists only for low upstream beta $\beta \lesssim 0.5$. Between perpendicular and parallel shocks there is a broader transition region where quasi-perpendicular shocks are still subcritical while quasi-parallel shocks are already supercritical. Laminar quasi-perpendicular shocks can presumably be found only at small β . At larger β they become more and more fluctuating.

Of the known collisionless shocks in the heliosphere planetary bow shocks are in general not expected to exhibit any subcritical behaviour. The exceptions are cometary bow

shocks and interplanetary shocks. Cometary bow shocks are low Mach number shocks because of the strong mass loading of the solar wind in the cometary environment with heavy ions from the atmosphere of the comet, which substantially reduces the upstream flow velocity. Forward interplanetary shocks are convected with the solar wind mostly outward into the outer heliosphere and, in addition, are close to being perpendicular such that at their already relatively low Mach numbers in the solar wind frame their nearly perpendicular nature prepares them for more easily entering into the subcritical regime. Reverse shocks which occur in corotating interaction regions are instead mostly supercritical as their flow direction is sunward. Cometary shocks will be discussed in the chapter in the context of planetary bow shocks (Chapter 10). Neither coronal nor interplanetary shocks are treated in this volume, both still being subject to intense observational investigation.

Theoretically, subcritical shocks are by no means subject to the presence of an obstacle, nor are they subject to the necessity of a violent blast. As outlined in earlier places (see Chapter 3) many subcritical shocks, in particular those of very low Mach number $\mathcal{M} \gtrsim 1$ marginally entering the shock regime, are produced by the nonlinear steepening of dispersive waves. This happens to be the most frequent case in collisionless plasmas where shocklets and extraordinarily weak shock limited to spatially localised regions are very frequently observed. Examples are found in interplanetary shocks. Other examples are found throughout the Earth's Magnetosheath region, the transition region behind the Bow Shock (see Chapter 10 on Planetary Bow Shocks). Many of these shock-like structures can and have been interpreted not in the shock language but rather as nonlinear wave structures. Nevertheless, if their Mach numbers do even slightly exceed the marginal shock Mach number, then we are indeed dealing with shocks, this time with strongly subcritical shocks.

Such subcritical shocks are not subject to upstream ion reflection since the expected dissipation inside the shock front and wave dispersion should be sufficient for providing the necessary dissipation/dispersion of sustaining the shock transition from the upstream to the downstream flow. As we have discussed in the previous chapter, the kind of dispersion will either allow for upstream phase locked precursor fluctuations or downstream wake fluctuations. For illustration, Figure 4.1 shows two schematic subcritical shock profiles, one of them a dispersive shock transition resulting from nonlinear steepening of whistlers, the other the extreme case of a purely resistive shock transition.

It is, however, clear that dispersion alone cannot create a shock transition. It produces some kind of localised waves like the various kinds of solitary waves and BGK modes. These are all structures that are connected with electric potential wells. Depending on the polarity of these potentials they reflect one sort of particles out of the upstream low energy component while they trap and accelerate particles from the other component. Hence, in the complete absence of any kind of dissipation a subcritical shock can also exist only when it reflects and traps some particles.

If this is not the case, the subcritical shock must definitely be capable of generating some kind of anomalous collisionless dissipation. Thus the question arises, what kind of dissipation this might be and how it is produced. Processes that are capable of doing this have been proposed for long time [Karpman & Sagdeev, 1964; Karpman, 1964; Sagdeev, 1966]. In the following we give a brief discussion of our knowledge about subcritical

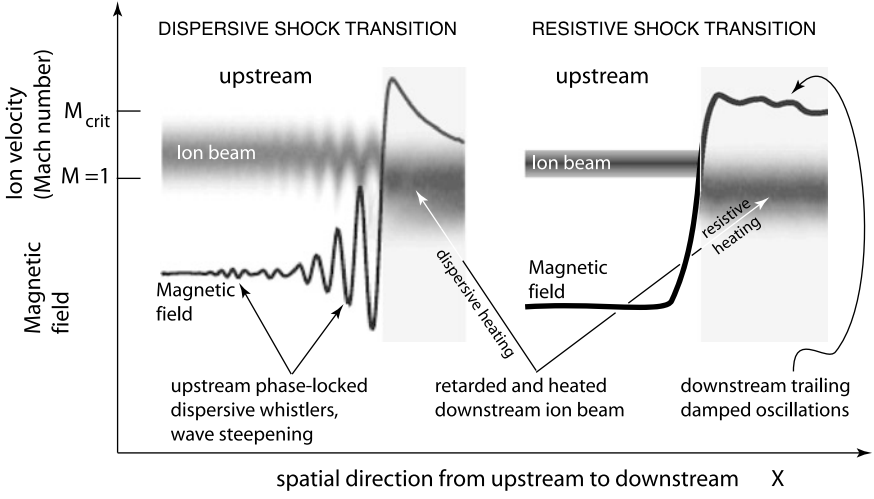


Figure 4.1: Schematic of the two types of subcritical shocks in ion phase space and magnetic field profile. In the dispersive shock transition the upstream whistler steepens nonlinearly while being phase-locked to the shock and causing coherent oscillations in the upstream ion beam. The upstream ion beam is subcritical. Heating and retarding in the shock is caused by dispersive effects. In resistive shock transition the shock forms a steep ramp lacking upstream oscillations. Heating and retarding is due to resistive Joule heating. Downstream the magnetic field may evolve trailing oscillations. Other shock transitions based on the Korteweg-deVries and related equations are also possible.

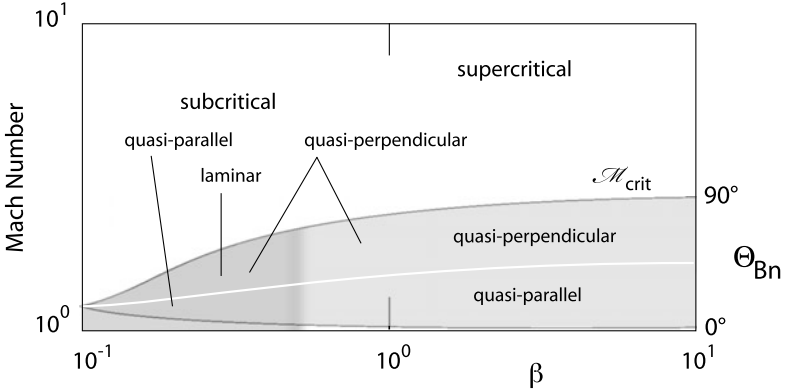


Figure 4.2: The range of supercritical and subcritical shocks in the $(\mathcal{M}_{ms}, \beta)$ -plane. The critical Mach number range is shown in shading varying between perpendicular and parallel shocks. The *white line* is the demarcation line between quasi-perpendicular and quasi-parallel shocks. The *lower dark line* belongs to strictly parallel shocks. At low β the parameter range of laminar shocks is shown in fuzzy shading. The transition to higher beta is blurred as it is not sharply defined. For a given shock normal angle Θ_{Bn} the *region below the curve* is subcritical. Quasi-parallel shocks can be subcritical only below the *white line*, depending on their shock normal angle.

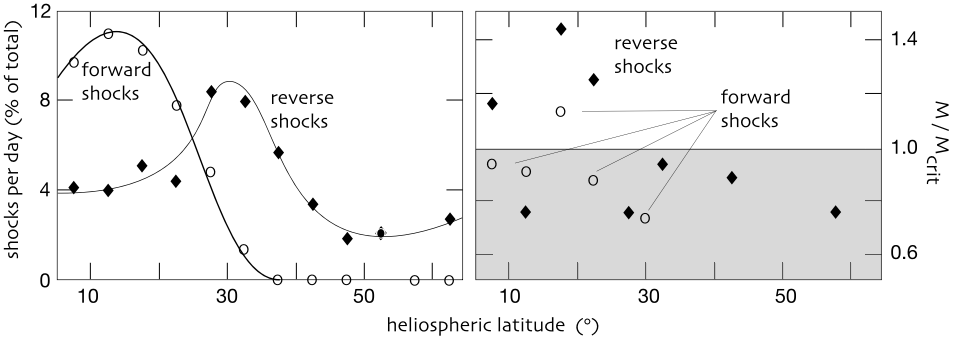


Figure 4.3: The distribution of shocks observed by the Ulysses spacecraft at a distance of ~ 5 AU as function of heliographic latitude. *Left:* The total number of corotating shocks (in %) divided into reverse and forward shocks. Forward shocks dominate at low latitudes while reverse shocks show a maximum at around 30° latitude. *Right:* Dependence of the relative shock Mach number on heliographic latitude separated into forward and reverse shocks. Shaded is the presumably subcritical domain. The low latitude forward shocks are preferentially subcritical. However subcritical corotating reverse shocks occur also on higher latitudes but are found more rarely [data taken from *Burton et al*, 1996].

shocks in collisionless plasma. Further details about the specific behaviour of shocks under cometary and interplanetary conditions can be found in the above mentioned chapters. We do not refer to simulations in this case as there are no new simulations of subcritical shocks after about 1985. Most of the investigations date back to the seventies and have been reviewed by *Biskamp* [1973], *Galeev* [1976], and *Mellott* [1985].

4.2 Observations

Mellott [1985] has reviewed the then available observations of subcritical shocks in space. Observing subcritical shocks in near-Earth space is quite rare. This might, however, be due to the difficulty of distinguish between very weak subcritical shocks and steep non-linear waves. Observational suggest that some interplanetary shocks are subcritical. These shocks are either forward or reverse shocks all being convected with the solar wind. The reverse shocks propagate against the solar wind and thus are preferably high Mach number. Forward shocks propagate with the solar wind and therefore sometimes are of sufficiently low Mach number to become subcritical. We note, however, that *Bavassano-Cattaneo et al* [1986] in an analysis of 34 interplanetary low heliographic latitude shocks found that most of the shocks were supercritical. They attributed this fact to the sensitive dependence of the critical behaviour of a shock on the plasma parameters of the highly fluctuating shock environment in the low latitude solar wind.

Burton et al [1996] analysed all out-of-ecliptic interplanetary shock observations by Ulysses in view of their properties, finding that the rate of occurrence of corotating shocks decreased dramatically with latitude, and the few shocks observed at high latitude were reverse and characterised by low Mach numbers, low β . The essence of these observations is shown in Figure 4.3 where we give the frequency of occurrence of corotating

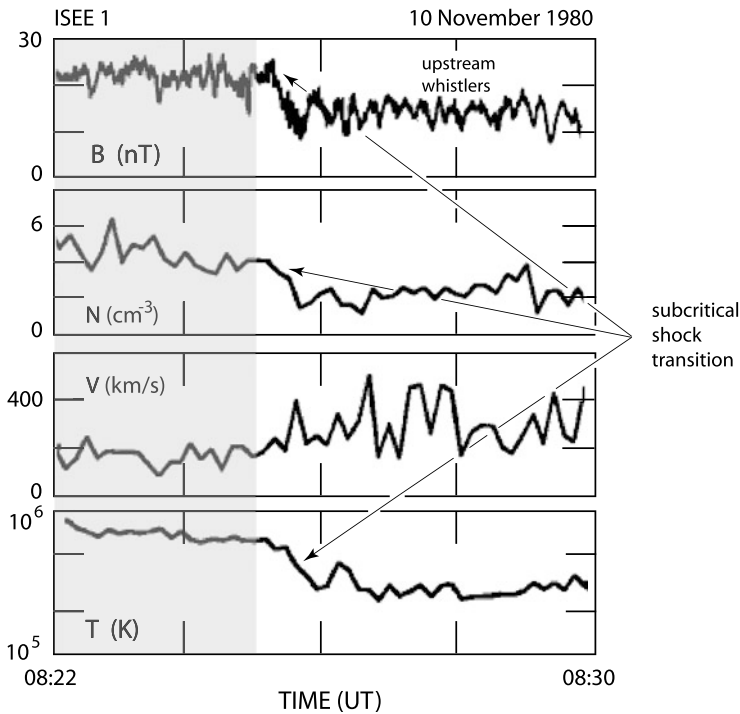


Figure 4.4: ISEE 1 crossing of a subcritical bow shock on 10 November 1980 [data taken from *Farris et al, 1994*, with permission of the American Geophysical Union]. Note the upstream standing whistlers which are responsible for the deceleration of the solar wind. The shock transition is quite smooth and of small compression ratio lacking an overshoot.

shocks at 5 AU distance as observed by *Ulysses* in dependence on heliographic latitude and Mach number ratio $\mathcal{M}/\mathcal{M}_{crit}$. Forward corotating shocks clearly dominate at low latitudes while reverse shocks have maximum occurrence around 30° heliographic latitude. Forward shocks are preferentially subcritical. However, high latitude reverse shocks seem also to be subcritical, however, occur at very low rate only.

A rare case of a subcritical Earth bow shock has been reported by *Farris et al [1994]* who identified it from the ISEE 1 recordings on 10 November 1980. Figure 4.4 shows a selection of data during this ISEE 1 outbound shock crossing coming from the magnetosheath and entering the solar wind. Clearly the shock is very low Mach number showing practically no remarkable deceleration in the streaming velocity. There are some large amplitude oscillations in V upstream of the shock but only a very smooth transition in speed from upstream to downstream. The shock is mainly seen in the magnetic field, density and plasma temperature both exhibiting only gradual changes in the shock ramp. The shock ramp is not particularly steep with compression ratio ~ 1.3 . Remarkably, no overshoot is seen in the magnetic profile. Behind the ramp the magnetic field fluctuates while the average field is still slowly increasing, which is the effect of approaching the obstacle

(magnetosphere in this case). There are magnetic fluctuations upstream and downstream. The upstream fluctuations have relatively long wavelength and are probably phase locked standing upstream whistlers which, in the solar wind frame, move in upstream direction but cannot escape from the shock. These whistlers are responsible for the gradual slowing down of the solar wind upstream of the shock transition thus providing part of the required dissipation of upstream motional energy. Downstream the fluctuations are of shorter wavelength and are irregular. Not only the absolute amplitudes but also their relative amplitudes b/B_2 are less than the relative amplitudes b_w/B_1 of the large amplitude upstream whistlers. From the density and temperature profiles the shock is identified as an entropy generator, heating and compressing the plasma.

The majority of shock observations in space is from crossings of the Earth's bow shock. However, the bow shock is in the overwhelming number of cases a highly supercritical shock. Observations will be discussed in the following chapters and in Chapter 10 on Planetary Bow Shocks. Nevertheless, in a few cases of bow shock crossings the bow shock has been identified as a locally subcritical shock when the solar wind was slow, cold (low $\beta \sim 0.5$) and dense (i.e. when the Earth was embedded into a slow solar wind stream) such that the (magnetosonic) Mach numbers were moderate, of the order of $\mathcal{M}_{ms} \sim 2$ which for nearly perpendicular shocks of $\Theta_{Bn} \sim (60-80)^\circ$ falls into the range of subcritical shocks [cf., e.g., Kennel *et al.*, 1985]. These shocks had, in general, two types of upstream waves attached to the shock ramps, a nearly coherent low frequency precursor oscillation of frequency $\omega/2\pi < 0.5$ Hz which rapidly damped away with increasing distance from the shock, and a higher frequency $\omega/2\pi > 1$ Hz and less damped oscillation. The former waves are the theoretically expected phase locked standing precursor whistlers with rest-frame frequencies of the order of $\omega \sim 10 \omega_{ci}$ and wavelengths of the order of $\lambda \sim 100$ km, in rough agreement with the estimated shock width, given in Eq. (4.2) below, which theoretically is expected to be of the order of the ion inertial length λ_i .

Figure 4.5 shows on its left the reasonably good correlation between the measured whistler wavelengths in the ISEE observations of a few subcritical bow shocks and the predicted precursor wavelengths and shock widths. The shock widths are of about the same order as the whistler precursor wavelengths, indicating that the shocks that contribute to this figure are evolving almost dispersively out of a precursor whistler, in rough agreement with the theoretical predictions [Sagdeev, 1966; Kennel & Sagdeev, 1967; Galeev, 1976]. On the right of the figure it is seen that the observations of most of the shocks are found to fall into the dispersive ion-acoustic regime, assuming that shocks may become resistive under the ion-acoustic instability. Such ion-acoustic resistive shocks should not exhibit precursor whistler waves. Some of the observations do indeed lack whistler precursors and are located in the resistive regime in the figure. We note, however, that it is not conclusive whether or not ion-acoustic waves are capable of generating sufficient resistance in subcritical shocks as the threshold for exciting ion-acoustic waves is high and presumably will barely be overcome in a low Mach number subcritical shock (cf. the discussion in Chapter 3). It is more probable that other, easier to excite instabilities and their related nonlinear effects take over in order to provide the required dissipation, and then the whistlers will become rather independent of the presence of ion-acoustic waves. In addition, the cross-shock potential might be high enough even in a subcritical shock to

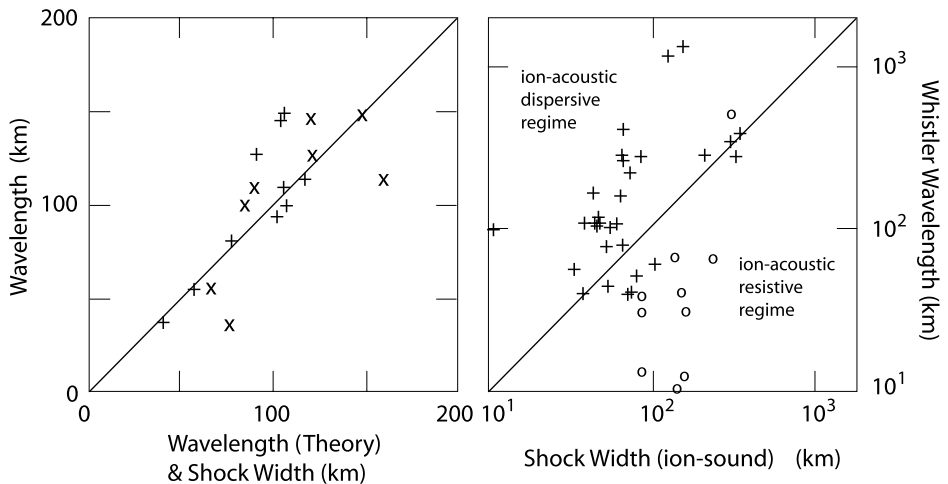


Figure 4.5: ISEE observations of subcritical shocks [data taken from *Mellott & Greenstadt, 1984*, with permission of the American Geophysical Union]. *Left:* The excellent correlation between the observed and theoretically predicted whistler precursor wavelengths during crossings of subcritical shocks (plus signs +) and the observed shock widths (crosses X). *Right:* The relation between the precursor whistler wavelengths and the shock-width as predicted from nonlinear ion-acoustic wave theory. The diagonal separates the dispersive ion-acoustic regime from the nonlinearly saturated resistive ion-acoustic regime. Under the assumption that ion-acoustic waves are responsible for dissipation in a subcritical shock precursor whistlers should be found predominantly in the dispersive region above the diagonal (plus signs +) as is approximately the case. In the case of dominance of resistance whistler precursors should be suppressed (open circles).

reflect some particles and produce small amounts of ion viscosity which are sufficient to maintain shock dissipation and shock structure. Finally, the differences between ion and electron dynamics in a subcritical shock have not yet been explored sufficiently.

Another marginally subcritical case of a bow shock crossing by ISEE 1 on 26 November 1977 at 0610 UT has been analysed by *Zilbersher [1998]*. The shock Mach number at this crossing was $\mathcal{M}_A \sim 2.7$ for a shock normal angle $\Theta_{Bn} \approx 67^\circ$ and $\beta_i \approx 0.16$, $\beta_e \approx 0.36$. According to theory such a shock should be just supercritical. The magnetic profile of the shock did not exhibit an overshoot which is what is expected in subcritical shocks. The ramp crossing was 10 s long at a shock velocity which was determined to be only $V_{sh} \approx (5-6)$ km/s at an ion inertial length $\lambda_i \approx 60$ km. The idea behind the analysis was to determine the ion behaviour in the smoothed shock. Smoothing has been done by using a wavelet analysis removing the six finest scales of a ten Daubechies wavelet decomposition [*Donoho, 1993*] of the magnetic variation. In addition the electric field was determined from the collisionless Ohm's law. As expected, the normal electric field obtained showed a bipolar structure which was restricted to the very shock ramp. Its amplitude was of the same order as the upstream convection electric field, strong enough to retard almost all the flow. Fields of this strength could just reflect some part of the ions. Using these data, the ion dynamics in the smoothed shock magnetic field and electric field was calculated with the interesting result that in this only marginally subcritical shock only a very small number

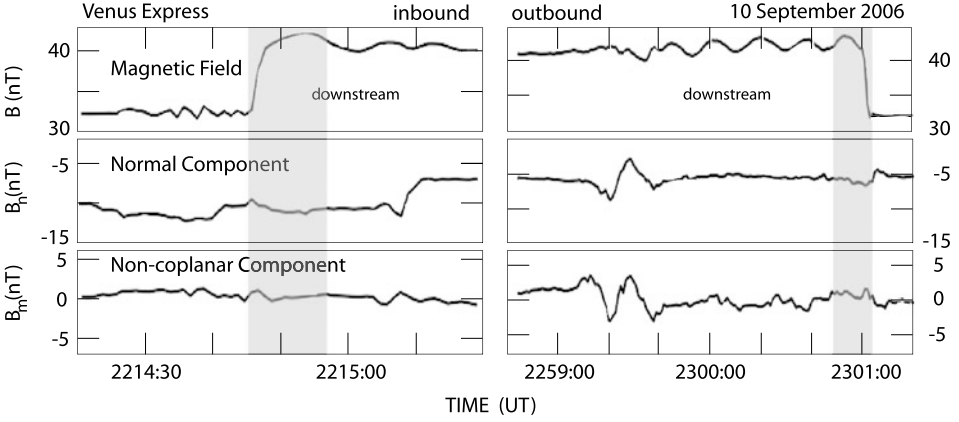


Figure 4.6: Venus Express observation of a subcritical bow shock near Venus [data taken from *Balikhin et al., 2008*, with permission of the American Geophysical Union]. The inbound and outbound crossings of the subcritical shock on 10 September 2006 are shown. The three panels give the modulus of the magnetic field, the component of the field normal to the shock, and the non-coplanar component. Grey shading shows the shock transition including the “overshoot”. The upstream field in this case was exceptionally high at 32 nT, and the shock presumably a low Mach number low- β quasi-perpendicular shock with $\Theta_{Bn} = (71-79)^\circ$ inbound and $\Theta_{Bn} \approx 80^\circ$ outbound and measured compression ratio $B_2/B_1 \approx 1.3$ corresponding to a magnetic field increase of only a factor 0.3.

of ions was reflected. The ions performed circular motions in phase space in the shock but only very few of them returned upstream. Clearly such circular motions in the ion phase space imply ion heating and ion hole formation with some of the ions becoming trapped. The transmitted ions started performing a gyration as a whole which led to the generation of finite amplitude about monochromatic oscillations of the magnetic field downstream of the shock which were clearly seen in the magnetic recordings.

The last finding is of interest in view of the recently reported observation of a subcritical bow shock structure when the Venus Express spacecraft approached the planet Venus on 10 September 2006 [*Balikhin et al., 2008*]. Two observed inbound and outbound shock crossings are shown in Figure 4.6.

The inbound crossing was at a quasi-perpendicular shock of shock-normal angle $71^\circ < \Theta_{Bn} < 79^\circ$. The magnetic field was exceptionally high near ~ 32 nT at Venus orbit, the Mach number was small, and β was small as well. The outbound crossing occurred at $\Theta_{Bn} \approx 80^\circ$ indicating that in both cases the shock was quasi-perpendicular and low Mach number. Both crossings lacked an overshoot which is also typical for subcritical shocks.

The measured compression ratio was low, $B_2/B_1 \approx 1.3$, yielding a magnetic field increase across the shock of only a factor ~ 0.3 . As expected, the normal component of the magnetic field (second panels) was practically constant across the shock. The non-coplanar component was zero both upstream and downstream of the shock but had a small peak in the shock ramp, which supports the conclusion that the shock was subcritical. For such very low Mach number subcritical shocks the shock behaves like a low amplitude magnetosonic wave [e.g. *Kennel & Sagdeev, 1967*; *Gedalin, 1998*]. This relates the veloc-

ity variations and magnetic field variations to the Mach number, yielding the Mach number as function of the compression ratio

$$\mathcal{M} - 1 \approx (B_2/B_1 - 1) \sin^2 \Theta_{Bn} \quad (4.1)$$

from which in this case it is found that $\mathcal{M} \approx 1.2-1.5$. The shock width [see, e.g., *Mellott & Greenstadt, 1984*] can be estimated from

$$\Delta_{\text{sh}} \approx \frac{2\pi\lambda_i}{(\mathcal{M}_A^2 - 1)^{\frac{1}{2}}} \cos \Theta_{Bn} \quad (4.2)$$

Assuming protons and a solar wind density $N \sim 10^7 \text{ m}^{-3}$, this yields a shock ramp width of roughly $\sim(60-80) \text{ km}$ the order of or less than an ion inertial length [*Balikhin et al, 2008*]. Observations reviewed by *Mellott [1984]* usually gave a thickness of the order of a few times the ion inertial length. But there is an inconsistency between theory and observation. The observed inbound shock exhibits an indication of weak phase-locked upstream whistler oscillations (which in the case of the outbound crossing are absent) in addition to weakly damped long wavelength downstream oscillations. The latter are also seen in the outbound crossing where they have somewhat shorter wavelengths. Such damped downstream oscillations in quasi-perpendicular shocks should be present predominantly in the case of supercritical conditions. In strictly perpendicular subcritical shocks we have seen in Chapter 3 that such oscillations occur as a downstream wave trail as solutions of the dissipative Korteweg-deVries equation behind the steepened wave front under particular conditions on wave dispersion.

In this case, however, both the presence of precursor whistlers and the oblique shock normal angle should rule out this kind of wave trail generation. *Balikhin et al [2008]* therefore suggest that they discovered a particular kind of cold subcritical shock the formation of which is not due to straightforward dissipation but is the cause of the cold gyrating solar wind beam. It had been shown in simple ballistic model calculations by *Zilbersher [1998]* that in subcritical shocks such conditions can indeed be realised when the inflowing beam is cold and when the shock ramp electric potential is not large enough to reflect it. This is probably the case for very cold subcritical beams which generate a moderate shock potential only. The ion beam then becomes deflected in the shock potential and leaves the shock to downstream as a beam that gyrates about coherently as a whole. Even in the case of upstream phase locked whistlers this gyration causes weakly damped long wavelength downstream oscillations similar to those which have been observed in the Venus bow shock crossings.

It should be noted at this occasion that the Venus bow shock is a particular case as Venus does not possess an own magnetic field but has a dense atmosphere which mass loads the solar wind. Hence, the solar wind becomes heavily retarded before it can evolve into a shock, and the resulting shock becomes subcritical and has properties that differ from shocks generated in unloaded collisionless streams. More properties of the Venusian bow shock will be discussed in Chapter 10 in comparison to the bow shocks of the other planets as well as to cometary bow shocks.

4.3 Theory

The foundations of the theory of subcritical shocks have been laid in the first three sections of Chapter 3. There we discussed different models for shocks evolving out of a nonlinear wave that becomes locally damped. Such a damped oscillator model had been proposed by *Sagdeev* [1966]. Early summaries of the theory of subcritical shocks can be found in the classical text of *Zeldovich & Raizer* [1966], the monograph of *Tidman & Krall* [1971] and the reviews by *Biskamp* [1973] and *Mellott & Greenstadt* [1984]; *Mellott* [1985]. According to the theory the evolution of subcritical shocks is due to the competition between the nonlinear steepening of a large amplitude low frequency plasma wave and the dispersive properties of the plasma. Dispersion can lead to the formation of localised waves of the soliton type. In a medium that contains a small amount of dissipation the waves become weakly damped, and these localised structures evolve into a ramp which mimics a shock.

It is important to note that the dissipation is in many cases quite unimportant as long as the amplitude and the steepness of the wave packet remain small. However when both increase, the gradient scale enters into the scale of the local dissipative interactions, and dissipation starts becoming important. Before constructing a model of the subcritical shock we therefore discuss the shock potential in a subcritical shock and the relevant scale of dissipation.

4.3.1 Subcritical Shock Potential

Subcritical shocks do not reflect ions. This means that all upstream ions pass the shock. However, the shock ramp, being of the order of the ion inertial length $\Delta_{\text{sh}} \sim \lambda_i$ allows for the ions to be non-magnetised while the electrons are tied to the magnetic field. As a consequence, the shock ramp contains a finite electric potential U which retards the ions, however is not large enough to stop their motion and to reflect them. Neglecting any upstream thermal spread of the ions which move at upstream bulk velocity V_1 , the ion speed v_x across the shock is therefore given simply by

$$v_x(x) = [V_1^2 - 2eU(x)/m_i]^{1/2} \quad (4.3)$$

independent on how complicated the real motion of the ions would be in crossing the ramp. For a cold upstream flow with $\beta \ll 1$ this is also the bulk flow velocity inside the ramp. Flux conservation (or the first Rankine-Hugoniot relation applied locally) then yields that the local density becomes

$$N(x) = N_1 \frac{V_1}{v_x(x)} = N_1 \left[1 - \frac{2eU(x)}{m_i V_1^2} \right]^{-1/2} \quad (4.4)$$

The transmitted ion density is trivially related to the shock potential which, thus, can be determined from a single measurement of the transmitted ion density. On the other hand, pressure balance requires that

$$m_i N v_x^2 + P + \frac{B^2}{2\mu_0} = \text{const}, \quad P \equiv P_i + P_e \quad (4.5)$$

Here the (scalar) pressure is retained as in the ramp the particles become heated and thus the pressure raises. Substituting for v_x , this relation can be used to express the shock ramp density and ion pressure as [Gedalin, 1997]

$$\frac{N(x)}{N_1} = \frac{1}{[1 - \bar{U}(x)]^{\frac{1}{2}}} \left\{ 1 + \frac{3}{4} \frac{\beta_i}{\mathcal{M}_A^2} \frac{\bar{U}(x)}{[1 - \bar{U}(x)]^2} \right\}, \quad \frac{P_i(x)}{P_{1\text{ram}}} = \frac{\beta_i / \mathcal{M}_A^2}{[1 - \bar{U}(x)]^{\frac{3}{2}}} \quad (4.6)$$

where we introduced the normalised potential $\bar{U} \equiv 2eU/m_i V_1^2$ and the upstream flow ram pressure $P_{1\text{ram}} = \frac{1}{2} m_i N_1 V_1^2$. This is an implicit expression for the shock potential $\bar{U}(x)$ at position x in the shock ramp as function of the local plasma density $N(x)$ and Mach number $\mathcal{M}_A$. Given a measured shock density profile the shock potential is completely determined under the conditions of validity of the assumptions made. The main restriction is that there should be only negligibly few ions in the upstream flow of energy less than the shock potential as those would be reflected and would thus violate the condition on the normal ion velocity. Hence, the theory applies to a cold upstream ion plasma only.

The ion velocity at the top of the ramp is given by

$$V_R(x_R)/V_1 = (1 - \bar{U}_{\text{tot}})^{\frac{1}{2}} \quad (4.7)$$

where $\bar{U}_{\text{tot}}$ is the normalised total ramp potential drop. This is less than the upstream speed but does not coincide with the downstream velocity V_2 determined from the Rankine-Hugoniot relations. The difference $V_R - V_2$ determines the downstream gyration speed and the downstream ion temperature $T_2 \sim m_i (V_R - V_2)^2 / 2$. This implies that the potential drop obeys the condition $\bar{U} < 1 - V_2^2 / V_1^2$. Ion deceleration in the ramp is solely due to the shock potential, but the flow deceleration is related to the heating of the ions across the shock.

It is noteworthy here to add that the shock ramp sees the ions unmagnetised, a conditions which we have made use of here. The electrons on the other hand are magnetised. They gyrate and experience the shock potential which for them is a transverse electric field. Thus, inside the ramp the electrons undergo an $\mathbf{E} \times \mathbf{B}$ -drift along the shock which generates an electron drift current $\mathbf{j}_{ed}(x) = eN(x)\nabla_x U(x) \times \mathbf{B}(x)/B^2(x)$ in the direction perpendicular to the shock ramp electric and magnetic fields. The magnitude of the tangential velocity $v_{ed} \equiv v_{ey}$ of this electron drift current can easily be expressed in terms of the local shock potential as

$$\frac{v_{ey}}{V_{A1}} = \frac{\mathcal{M}_A^2}{2} \frac{B_1}{B(\xi)} \left| \frac{\partial \bar{U}(\xi)}{\partial \xi} \right| \quad (4.8)$$

where $\xi = x/\lambda_i$ is the normalised coordinate in shock normal direction, and the direction of electron flow is along the shock. The magnetic field of this current causes a magnetic overshoot in the subcritical shock. In addition, if the electron current becomes strong enough it serves as source of instability and generates dissipation.

4.3.2 Dissipation Length

For a subcritical shock with no primary particle reflection the shock transition must be capable of generating sufficient dissipation to come up for the entropy increase in the flow

when it crosses the shock. The shock has thickness Δ_{sh} , and the flow has velocity V_1 . Hence the dissipation must be active during the crossing time $\tau_{\text{sh}} = \Delta_{\text{sh}}/V_1$. It is thus clear that the dissipation time $\tau_d < \tau_{\text{sh}}$ must be shorter than the crossing time. Alternatively, the dissipation scale $L_d < \Delta_{\text{sh}}$ must be shorter than shock width. This simple conclusion contains some important physics. First, the dissipation scale is not the same as the mean free path of the particle which is $\lambda_{\text{mfp}} = 1/(\sigma_c N)$, with σ_c the collisional cross section. In collisionless plasma this mean free path is infinite. Second, the idea is that shock formation generates a transition region that is so narrow that it becomes the order of the internal dissipative processes. Since the shock layer is a current layer this implies that the shock current must flow in such a narrow region that it can excite current driven instabilities that generate anomalous resistance or anomalous viscosity, and the scale on that this anomalous resistivity η_{an} is large enough to sustain the dissipation in the shock must be shorter than the shock width.

In the case when the dissipation is dominated by Joule heating of the plasma in the shock (which is believed to be the case in subcritical shocks which are not subject to ion reflection and therefore do not generate large ion viscosities) one can give this condition a quantitative expression from observing that from the induction equation

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times \mathbf{V} \times \mathbf{B} - \frac{\eta_{\text{an}}}{\mu_0} \nabla^2 \mathbf{B} \quad (4.9)$$

follows that dissipation can only dominate the flow if the resistive second term on the right is larger than the convection term. Dimensionally this implies for the dissipation length

$$L_d^2 \lesssim \frac{\eta_{\text{an}}}{\mu_0} \frac{\Delta_{\text{sh}}}{V_1} \quad (4.10)$$

Expressing the anomalous resistivity in terms of the anomalous collision frequency ν_{an} and the electron inertial length $\lambda_e = c/\omega_{pe}$ yields

$$L_d \lesssim \left(\frac{\nu_{\text{an}}}{V_1/\Delta_{\text{sh}}} \right)^{\frac{1}{2}} \lambda_e \quad (4.11)$$

Thus, on requiring that the dissipation scale $L_d = \sqrt{\alpha} \Delta_{\text{sh}}$ is of the order of the shock width, this expression becomes a condition on the anomalous collision frequency

$$\nu_{\text{an}} \gtrsim \alpha \left(\frac{V_1}{\lambda_e} \right) \left(\frac{\Delta_{\text{sh}}}{\lambda_e} \right) \quad \text{or equivalently} \quad \nu_{\text{an}} \tau_{\text{sh}} \gtrsim \alpha \left(\frac{\Delta_{\text{sh}}}{\lambda_e} \right)^2 \quad (4.12)$$

where α is a numerical factor. For anomalous dissipation to be sufficiently large to sustain the shock the ratio of the flow transition time to the anomalous collision time must thus be larger than the square of the shock width measured in electron inertial lengths.

As for an example let us assume that the shock width is $\Delta_{\text{sh}} = 1000$ km. Then, for a plasma density of $N \sim 5 \times 10^6 \text{ m}^{-3}$ and a subcritical flow velocity of $V_1 = 100$ km/s the anomalous collision frequency should be larger than $\nu_{\text{an}} > 10^4 \alpha \text{ Hz}$, which is of similar order as the electron plasma frequency $f_{pe} \sim 20 \text{ kHz}$. Since such high anomalous collision

frequencies are unrealistic, one must require $\alpha \sim 0.1$ corresponding to a substantially narrower dissipation scale $L_d \sim 0.3\Delta_{sh}$ or, correspondingly, narrower current sheets inside the shock transition.

If the subcritical shock evolves out of a magnetosonic solitary wave, the requirement on the anomalous collision frequency can also be expressed in terms of the upstream shock Mach number. For magnetosonic shocks there is a distinct relation between the shock width and the Mach number [see, e.g., *Sagdeev, 1966*] given approximately by $\Delta_{sh} \simeq \lambda_e / \sqrt{\mathcal{M}_A - 1}$. Inserting into the first of the above conditions and noting that $\lambda_e^2 / V_A^2 = (m_e / m_i) \omega_{ci}^{-2}$, one obtains

$$\frac{v_{an}}{\omega_{ci}} \gtrsim \alpha \left(\frac{m_i}{m_e} \right)^{\frac{1}{2}} \frac{\mathcal{M}_A^2}{\sqrt{\mathcal{M}_A - 1}} \quad (4.13)$$

Similar relations also hold for other types of solitons. Assuming a subcritical Mach number of $\mathcal{M}_A = 1.5$ this expression yields $v_{an} > 140 \alpha \omega_{ci} \sim 3 \alpha \omega_{lh}$, a value still quite high for this kind of soliton as anomalous collision frequencies will barely exceed the lower-hybrid frequency ω_{lh} . Thus one again requires that the current sheets are narrower than the shock transition rather being of the order of the electron skin depth. In view of our introductory remarks it is the steepness of the localised wave packets which is responsible for the generation of the dissipation. Its scale is shorter than the full shock transition.

4.3.3 Subcritical Cold Plasma Shock Model

Following *Sagdeev [1966]* we refer to a one-dimensional cold ($\beta_1 \ll 1$) two fluid damped oscillator model assuming that the dissipation is entirely resistive in order to describe the structure of a stationary subcritical shock. This procedure leads to the Sagdeev potential method for the fast magnetosonic wave which in the *Sagdeev [1966]* case is a perpendicular wave. The equations are the continuity and momentum conservation equations, the latter reducing to pressure balance for the ions, and the stationary electromagnetic equations. They have been given in Chapter 3. Here, on identifying the only spatial direction of interest as x , we write them as

$$\begin{aligned} \nabla_x B_z &= \mu_0 e N v_y, & \nabla_x E_y &= 0 \\ \nabla_x N v_x &= 0, & \nabla_x \left(\frac{1}{2} m_i N v_x^2 + \frac{B_z^2}{2\mu_0} \right) &= 0 \\ m_e N v_x \nabla_x v_y &= -e N E_y + e N v_x B_z - v_{an} m_e N v_y \end{aligned} \quad (4.14)$$

where we have introduced the anomalous collision frequency v_{an} which is assumed to account for the dissipation. We repeat in passing that the plasma is collisionless and, therefore, binary collisions are excluded from contribution to dissipation.

Looked at from the fluid picture, the writing is formally so as if the electrons would interact with ions via anomalous collisions. Transforming to the shock velocity V_{sh} these equations can be reduced in the usual way to an equation for $b \equiv B_z / B_1$, the ratio of the

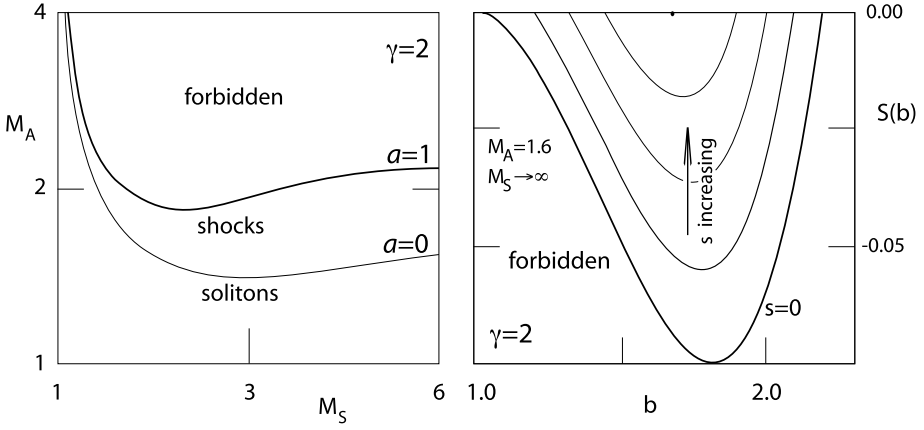


Figure 4.7: *Left:* The allowed regions in $(\mathcal{M}_A, \mathcal{M}_s)$ -space for solitons and shocks depending on the resistive dissipation coefficient a for a given adiabatic index $\gamma = 2$. Subcritical shocks exist only in the narrow domain between the two curves for $0 < a < 1$. *Right:* The shape of the Sagdeev pseudo-potential $S(b)$ for $\mathcal{M}_A = 1.6$, $\mathcal{M}_s \rightarrow \infty$ and $\gamma = 2$ as function of the entropy density s . For $s = 0$, corresponding to $a = 0$, the Sagdeev pseudo-potential has its largest excursion into the negative domain. For this case solitons exist throughout the entire region inside the curve. With increasing s the domain shrinks, and no soliton solutions exist anymore being replaced by shock solutions. A shock starts at one of those curves and wanders upward in the diagram until it reaches the maximum entropy point on the $S = 0$ axis (black dot) [data taken from Ziegler & Schindler, 1988].

magnetic field to the upstream magnetic field. With the help of the electron inertial length λ_e we define a new dimensionless variable $\xi = x/\lambda_e$. Then this equation assumes the form

$$\frac{d^2 b}{d\xi^2} = b - 1 + Ab(1 - b^2) - a \frac{db}{d\xi} \quad (4.15)$$

where $A \equiv B_1^2/2\mu_0 N m_i V_{sh}^2$, and $a \equiv \nu \lambda_e / V_{sh}$. This is the equation of motion of a hypothetical particle with coordinate b and time ξ including friction. In the absence of dissipation one defines the Sagdeev pseudo-potential

$$S(b) = \frac{1}{2}(b-1)^2 \left[\frac{A}{2}(b+1)^2 - 1 \right], \quad b < b_{\max} = \frac{2V_{sh}}{V_{A1}} - 1 \quad (4.16)$$

Its negative values constitute a potential trough for the hypothetical particle for field amplitudes $b < b_{\max}$. For a given value of $S(b)$ the hypothetical particle performs a stationary oscillation in this potential trough between its walls with amplitude equal to the distance of the walls. The shape of this potential trough is shown on the right in Figure 4.7 for the special case of Mach number $\mathcal{M}_A = 1.6$ as the heavy line marked $s = 0$. $b_{\max}$ is the maximum possible amplitude of the soliton which is reached at $S(b) = 0$. The ratio $V_{sh}/V_{A1} = \mathcal{M}_{sh}$ is the shock-Mach number, i.e. the Mach number of the stationary soliton solutions that are possible in this dissipationless case. However, including the resistive damping, this amplitude decreases during the oscillation, as we have discussed in Chapter 3, and the pseudo-

particle will ultimately settle in the final state at the minimum of the pseudo-potential $S(b)$. This minimum is at the stationary downstream value of the normalised magnetic field

$$b(S_{\min}) \equiv b_2 = \frac{1}{2} \left[(1 + 8\mathcal{M}_{\text{sh}}^2)^{\frac{1}{2}} - 1 \right] \quad (4.17)$$

Most interestingly, this final stationary shock state of a subcritical fast magnetosonic shock of this kind is independent of the dissipation even though it has been reached only due to the action of the anomalous dissipation a . A shock of this kind is a weak shock since the plasma pressure contribution has been neglected and the plasma has been assumed to be cold such that the heating of the plasma by the shock itself is also small and does not appear anywhere. However, the shock profile can be determined from the solution of the above equation for b . It is found to be a spatial oscillation reaching maximum at the shock and decreasing exponentially with damped amplitude $b(\xi) \sim \exp(-a\xi) \sin(\xi \sqrt{\mathcal{M}_{\text{sh}}^2 - 1})$ behind the shock [see, e.g., *Tidman & Krall, 1971*]. Note, however, that because of the simplifying assumptions this solution holds only for small shock amplitudes.

4.3.4 Extension to Warm Plasma

The zero temperature small amplitude case of subcritical perpendicular shocks can be analytically extended to the consideration of subcritical shocks of arbitrary amplitudes in a warm two component plasma consisting of electrons and ions [*Ziegler & Schindler, 1988*]. We again normalise all quantities to their upstream values, use quasi-neutrality $N_i = N_e$ as we have done above, and measure the length along the coordinate x in $\tilde{\lambda} = \sqrt{\lambda_e \lambda_i}$. In this case the equation for the normalised magnetic field component b reduces to

$$bv_x - 1 - \frac{a}{v_x} b' = b'' \quad (4.18)$$

The prime $'$ denotes the derivative with respect to the integrated coordinate $\bar{x} = \int_{-\infty}^x dz/v_x(z)$. This equation still contains the flow velocity v_x which is determined from the momentum conservation equations. Not neglecting the thermal pressure one must include the equation of energy conservation, with entropy density s given by the integral

$$s(\bar{x}) = (\gamma - 1) \int_{-\infty}^{\bar{x}} dz a v_x^{\gamma-2} [b'(z)]^2 \quad (4.19)$$

This leads to the wanted equation for v_x which is more complicated [*Biskamp, 1973; Ziegler & Schindler, 1988*]. It includes the upstream Alfvén and sound wave Mach numbers $\mathcal{M}_A$, $\mathcal{M}_s$, respectively. The latter appears because of the finite thermal pressure. The resistive Joule dissipation is again contained in the coefficient $a = v_{an} m N / \tilde{\lambda}$, with $m = m_i + m_e$. In the form given by *Biskamp [1973]* it reads

$$\begin{aligned} \frac{1.6v_x}{V_1} = & \left(1 - \frac{b^2 - 1}{2\mathcal{M}_A^2} \right) + \frac{0.6}{\mathcal{M}_s^2} \pm \left\{ 0.36 \left(1 - \frac{1}{\mathcal{M}_s^2} \right)^2 \right. \\ & \left. - \frac{(b^2 - 1)}{\mathcal{M}_A^2} \left(1 - \frac{1.28}{b+1} + \frac{0.6}{\mathcal{M}_s^2} - \frac{b^2 - 1}{4\mathcal{M}_A^2} \right) \right\}^{\frac{1}{2}} \end{aligned} \quad (4.20)$$

Here the upper sign of the root must be chosen because $v_x/V_1 \rightarrow 1$ for $b \rightarrow 1$. At large Mach numbers $(\mathcal{M}_A, \mathcal{M}_s) \rightarrow \infty$ the lower sign yields $v_x/V_1 = \frac{1}{4}$ for the downstream state. Thus at high Mach numbers there are two solutions for the velocity indicating the absence of a continuous solution independent of the resistivity. This is in agreement with the earlier claim that high Mach number shocks are supercritical and require additional dissipation which is found by reflecting part of the upstream flow back upstream.

The small amplitude Sagdeev solution follows for $\mathcal{M}_s \rightarrow \infty$, negligible plasma pressure $P/B_1^2(b^2 - 1) \ll 1$, and $v_x \rightarrow V_{sh}$. Then $s(\bar{x})$ can be neglected and the problem reduces to the cold plasma soliton case which was described above. In the more general case, real stationary solutions that connect the homogeneous states at $x = \pm\infty$ exist only in a limited Mach number regime for flows that are supersonic with respect to the phase velocity of linear waves propagating perpendicular to the magnetic field $\mathbf{B}$. The regions of existence of these solutions are given on the left in Figure 4.7. In the $(\mathcal{M}_A, \mathcal{M}_s)$ -plane the forbidden and allowed regimes are separated by the curves in this figure. Solutions exist only below these curves depending on dissipation a . In the absence of dissipation, solitons exist below the curve with $a = 0$. With increasing dissipation up to its maximum value $a = 1$ there are no soliton solutions; here the solutions are shock waves. These are confined to the region between the curves $a = 0$ and $a = 1$. On the right-hand side of this figure the parts of the Sagdeev pseudo-potential curves $S(b) < 0$ that correspond to real solutions are given as function of the entropy density s that is generated by the anomalous dissipation. It is seen that with increasing dissipation the domain of shock solutions shrinks gradually.

Figure 4.8 shows the results obtained by numerically integrating the full dissipative two-fluid warm plasma equations. The left part of the figure is the dissipationless case $a = 0$ for different combinations of Mach numbers $\mathcal{M}_A, \mathcal{M}_s$. Solitons of different form and amplitude are obtained. The equations permit for multi-soliton solutions. Here only one soliton of each kind is shown. The black curve on the right part of the figure is a typical shock solution with weak dissipation $a = 0.01$ and for negligible thermal pressure. The solution in this special case is a steep shock front corresponding to the steep leading edge of the first soliton on the left. However, the dissipation does not allow for the soliton amplitude to return to its initial upstream level because the steep gradient scale in the shock front approaches the dissipation length scale such that dissipation becomes important and generates entropy. This causes the oscillation downstream of the shock to be spatially damped. The oscillation amplitude exponentially decreases with distance from the shock as predicted by the simplified small amplitude treatment [Sagdeev, 1966] until the field settles on the value of $b(S_{min})$ far behind the shock transition.

Also shown in this figure is the case of maximum dissipation $a = 1$ when the shock transition scale is larger. Under these maximum dissipation conditions the shock does not exhibit a downstream oscillation but becomes a steep and smooth shock ramp with small overshoot and laminar downstream state. Again the amplitude of the latter corresponds to $b(S_{min})$ for the chosen combination of Mach numbers, yielding a smaller final downstream amplitude. It is interesting to see that for these marginal conditions the downstream state is indeed completely laminar with no oscillations evolving. Of course, the solutions are strictly time independent such that no turbulence can evolve. Under real conditions allowing for time variations stationarity will be violated.

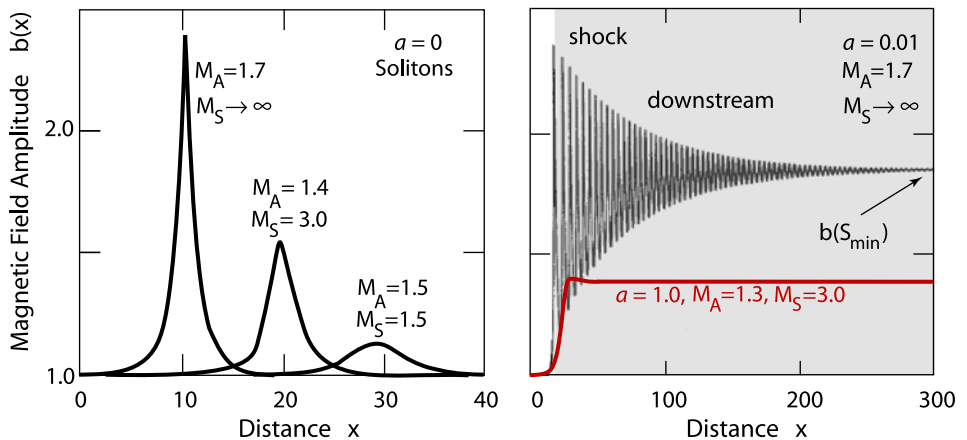


Figure 4.8: *Left:* Numerical calculation of soliton forms for $a = 0$ and different Mach number combinations. *Right:* The shape of the subcritical shock in the two-fluid model for finite dissipation $a = 0.01$. The solution has a steep ramp and a short wavelength downstream oscillation which damps slowly out with normalised distance from the shock ramp until it reaches the final amplitude $b(S_{\min})$ corresponding to the minimum of the Sagdeev pseudo-potential for the given combination of Mach numbers. The red line shows the asymptotic shock profile for maximum dissipation $a = 1$. In this case the shock does not exhibit oscillations but is a smooth ramp [data taken from Ziegler & Schindler, 1988].

A subcritical shock model that is a combination of the steeping of an ion-acoustic wave and the excitation of an Alfvén wave which dispersively evolves into a magnetic shock ramp has been developed early by Kennel & Sagdeev [1967]. The idea is that in a warm plasma stream conditions are given under that the ion-acoustic mode can be excited along the magnetic field to large amplitude, thereby producing a pressure anisotropy which can excite the firehose instability. Both effects working in tandem would generate a subcritical shock with a magnetic Alfvén ramp and anomalous resistivity provided by the ion-acoustic instability. Kennel & Sagdeev [1967] also proposed another variant of this model in which the coupling of the ion-sound wave and the Alfvén wave was not needed as the entire process of steeping was attributed to a magneto-acoustic wave while the generation of anomalous resistivity was still left to the action of ion acoustic waves excited in the shock. However, both models have been replaced by the realisation that whistlers are more important in shock generation. Still, the problem of generation of anomalous dissipation remains.

4.3.5 Dissipation in Subcritical Shocks

So far nothing has been said about the origin of the dissipation. Biskamp [1973], following Sagdeev [1966], reviewed the then known theoretical foundations of subcritical shocks suggesting that a subcritical shock must necessarily be time-dependent. The argument is simply the shock transition in a subcritical shock is due to the presence of dissipation and

causes an increase in entropy which, assuming pressure isotropy and a constant specific heat C_V across the shock, is given by

$$s_2 - s_1 = C_V \ln \left\{ \frac{P_2}{P_1} \left(\frac{N_1}{N_2} \right)^7 \right\} \quad (4.21)$$

On the other hand, the stationary collisionless Vlasov equation conserves the particle number and entropy flux such that $\nabla \cdot \int \mathbf{v} d\mathbf{v} F(\mathbf{v}) \ln F(\mathbf{v}) = 0$. Hence, a stationary shock state is possible only in the average over a rapid small scale variation taking place in the shock ramp in the sense we have discussed in Chapter 3. In this case the entropy flux can indeed increase and is given by

$$\nabla \cdot \int \mathbf{v} d\mathbf{v} \langle F(\mathbf{v}) \rangle \ln \langle F(\mathbf{v}) \rangle = \frac{e}{m} \int d\mathbf{v} \langle (\mathbf{e} + \mathbf{v} \times \mathbf{b}) f(\mathbf{v}) \rangle \cdot \frac{\partial \ln \langle F(\mathbf{v}) \rangle}{\partial \mathbf{v}} \quad (4.22)$$

where $f = F - \langle F \rangle$. From here it is clear that the entropy generation is due to the fluctuations f in the distribution function F and the related electric $\mathbf{e}$ and magnetic $\mathbf{b}$ field fluctuations. The averages are taken over the width of the shock transition and the time it takes an upstream fluid element to cross the shock. In a subcritical shock one can still speak of such fluid elements as no groups of particles are singled out by shock reflection. If this would be the case – as it happens in supercritical shocks – these particles would separate from the bulk flow and have their own life, gyrating around the field until ultimately as well crossing the shock causing the fluid element to become strongly deformed in phase and real space. It is thus the “turbulent” fluctuations that are responsible for the generation of entropy and cause the subcritical shock transition. One may speak here of turbulence that is confined to the shock transition in the sense that macroscopically it can be considered to cause a quasi-stationary state in which the fluctuations are confined to the shock transition while being of scales shorter than the shock width. Macroscopically speaking they cause the anomalous dissipation. The latter should in principle be possible to determine from knowledge about the nature, generation and saturation level of the turbulent fluctuations.

The fluctuations leading to turbulence in the shock are caused by instabilities that are excited in the shock transition. In Chapter 3 we have discussed a variety of such instabilities that can grow in the shock transition and are potentially responsible for the anomalous dissipation. Such instabilities are divided into macroscopic and microscopic instabilities and among them into resonant and non-resonant instabilities. Since the shock transition width is of the order of the ion-inertial length $\lambda_i = c/\omega_{pi}$, macroscopic instabilities mainly affect the gross two- or three-dimensional structure of the shock front. In the generation of dissipation it are the microscopic (kinetic) instabilities which are of main interest. Moreover, pure wave-wave interactions of the kind $k_1 = k_2 + k_3$, $\omega_1(k_1) = \omega_2(k_2) + \omega_3(k_3)$, so-called decay instabilities where one large amplitude wave ($k_1 \omega_1$) decays into two smaller amplitude waves (k_2, ω_2), (k_3, ω_3) may transfer turbulent energy from one region of (k, ω) -space into another region where anomalous dissipation can be caused. In order to satisfy the momentum conservation (wave number) condition in this three-wave process of the same kind the dispersion relation $\omega(k)$ must be concave from below. Thus decay takes preferentially place between waves of different kinds.

In Chapter 3 we provided a review of instabilities which may be responsible for generation of anomalous resistivity in collisionless shocks. Other reviews can be found in *Papadopoulos* [1985] and *Sagdeev* [1979]. It is not yet definitely known which of the instabilities and which of the mechanisms is the most viable source of anomalous dissipation. According to the above discussion of the dissipation length scale the favoured instability is the one which generates a collision frequency of the order of the lower hybrid frequency ω_{lh} , i.e. it must involve non-magnetic ions and magnetised electrons. There are basically two types of instabilities fitting this condition. The first is the gradient drift instability, in particular the lower-hybrid drift instability which, however, is should not be strong enough in shocks as it depends heavily on β_i and is damped for not too small β_i . Moreover, it requires density gradient scales shorter than the ion inertial length λ_i , implying a rather though not completely unrealistic steep shock profile. If the shock evolves out of a solitary structure, which is believed to be the case for subcritical shocks, it requires formation of a steep and narrow magnetosonic soliton. Though this might be possible it seems unlikely that it is the general case in subcritical shocks.

The other candidate is the modified two-stream instability. It has a very low threshold but requires the presence of a shock-normal electric field across the shock ramp transition on a scale $\lesssim \lambda_i$ in order to keep the ions non-magnetic. Fortunately, in contrast to generating a very steep density gradient, production of a cross-shock electric potential drop is quite natural in any magnetised shock formation. Even though this field is too weak to enable ion reflection, it serves accelerating electrons across the shock in an ambipolar manner thereby reducing the scale below the ion inertial length.

More importantly, however, the magnetised electrons perform an $\mathbf{E} \times \mathbf{B}$ -drift in the shock magnetic and cross-shock electric fields. This drift is directed tangential to and along the shock and produces an electron drift current $\mathbf{j}_d = -eN\mathbf{E} \times \mathbf{B}/B^2$ that is flowing along the shock, most probably in the shock ramp, and whose magnetic field contributes to the steepening of the magnetic shock ramp and the first shock magnetic field maximum. Its magnitude in terms of the upstream Mach number $\mathcal{M}_A$ and the cross-shock potential $\bar{U}(x)$ has been given in Eq. (4.8). This drift current stores free energy, providing it to the excitation of the modified two-stream instability. The modified two-stream instability does not exist for strictly parallel and strictly perpendicular shocks, however, but this is not a severe restriction, for these two marginal cases will rarely be realised in practice. It has other consequences which we will briefly mention below. However, before coming to this we discuss the possible direct role of the modified two-stream instability in anomalous dissipation.

Referring to Figure 3.16 we note that the modified two-stream instability has wavelength $\lambda \gg \lambda_D$, considerably larger than the Debye length, and growth rate of the order of roughly ten times the ion cyclotron frequency. This is of the right order of magnitude which is needed in subcritical shocks. Its frequency is in the range of $\omega \sim \omega_{lh}$. It is excited in the shock transition on transverse scales shorter than the ion-inertial length and on tangential scales of many wavelengths along the shock surface where it produces a high anomalous collision frequency $\nu_{an}^{\text{MTSI}} \propto \omega_{lh}$. By Eq. (4.13) this is large enough to account for the required dissipation in subcritical shock formation.

Based on these considerations we may complete the scenario of evolution of a subcritical shock. The subcritical shock is the result of the nonlinear steepening of a magnetosonic whistler wave. Dispersive effects may or may not lead to the presence of phase locked whistlers. During wave steepening a solitary wave forms with ramp steepness on the scale or larger than the ion inertial length. The difference in the ion and electron gyroradii produces a cross shock potential when the upstream particles hit the soliton ramp. This potential is not strong enough for reflecting the ions but causes a cross-field drift of the electrons which corresponds to a current flowing along the shock in the shock ramp. This current, even when it is weak, is strong enough to excite the modified two-stream instability because of the very low threshold of this instability. When the modified two-stream waves grow and quasilinear effects set on, anomalous collisions are generated which scatter the electrons and produce the required dissipation, Joule heating of the electrons and ions, and entropy for the transformation of the soliton into a shock. At the same time the presence of the anomalous collisions causes the electrons to diffuse ambipolarly across the shock to shorten the length scale below the ion inertial length thus keeping the ions non-magnetic and the electron drift scale short enough for the modified two-stream instability to be maintained.

This scenario is to some degree self-consistent. Whether it is realised remains uncertain. The modified two-stream instability is an oblique instability. Consequently, it has an electric field component $E_{\parallel}$ parallel to the shock magnetic field $\mathbf{B}$. The electric field accelerates electrons along the magnetic field thereby generating field aligned electron beams. These, on their hand, excite plasma waves which have the tendency to readily disintegrate into electron holes and BGK modes. A description of this process can be found in Chapter 3. Electron trapping in the electron holes respectively BGK modes causes strong and irreversible electron heating. Such a kind of heating is the result of the electron vortex formation by electron trapping in the holes and the related broadening of the electron distribution function, and is manifestly nonlinear and also non-resistive.

Since hole formation is a fast process it competes with the quasilinear generation of collisionless resistivity and may be more important than the resistive Joule heating which is caused by the anomalous resistance that is generated in the shock by quasilinear effects on the modified two-stream instability. In the fluid equations in Chapter 3 this process is implicitly contained in the ponderomotive force term in the generalised Ohm's law but has so far not been investigated in collisionless shock physics even though more recently electron holes on scales much shorter than the ion inertial length have been detected in the shock transition region, both in the shock feet and ramps [Bale *et al*, 1998; Bale & Mozer, 2007; Balikhin *et al*, 2005]. Observations of this kind suggest that the shock dissipation mechanism takes place on the microscopic scale and probably in a way that is different from the conventional thinking in terms of dissipation.

Still, it is a collective plasma process in which several effects are involved, wave steepening, cross shock potential generation, electron cross-field drift current production, modified two-stream instability, electron beam acceleration, BGK modes, electron trapping, vortex formation, broadening of the electron distribution function yielding collisionless non-resistive heating and the required non-resistive dissipation, and finally ion heating by thermal transport of energy from electrons to ions. The chain of these processes, however,

cannot be treated by an expansion procedure with respect to a small parameter since some of the processes involved are of the same order.

We note in passing that, in addition, other processes of electron acceleration and transport have been proposed to be also at work in the subcritical shock transition [Balikhin *et al*, 1993, 1998; Gedalin *et al*, 1995; Gedalin, 1998; Zilbersher, 1998]. These processes are based on the chaotisation of the electron orbits in the combined shock magnetic and electric fields and might as well contribute to collisionless dissipation, electron heating, thermalisation and the necessary entropy production for sustaining the shock.

4.4 Conclusions

Briefly summarising the content of this chapter, which is the first in this volume that deals with an application of the methods and ideas in Chapter 3 to shocks, more precisely to shocks the Mach number of which is so low that they can be described by the combined action of dispersion and dissipation present in dispersive waves in collisionless plasmas. Naturally, such shocks have been, historically, the first which had been considered when collisionless shocks were discovered first in the laboratory and then in space. The laboratory came first, simply because it was difficult to produce the high Mach numbers needed by super-critical shocks. The discovery in space of collisionless shocks with their particular shock properties (like shock-particle reflection) came as a surprise. In space, high Mach numbers are the rule when observing the large violent processes related to the shocks that we today are familiar with, the planetary Bow Shocks, CME shocks, the Termination shock and others. Subcritical shocks have been believed to be rare there; they were mostly restrictedly associated to heavy mass loading of the solar wind as is the case in the vicinity of comets and Venus and Mars as the unmagnetised planets, in particular at Venus with its dense atmosphere. However, they might be much more frequent simply due to the properties of dispersive waves which nonlinearly are capable of steepening and evolving into shocks.

Evolution of subcritical shocks in the latter case is now quite well understood, even though the generation of anomalous resistance and anomalous dissipation below the critical Mach number still poses many unresolved problems. It is well established that the subcritical shock evolves through the various phases of steepening of a low frequency magnetosonic wave the character of which has been identified of being on the whistler mode branch. This steepening process is completely non-collisional. The modes propagate against the upstream flow, forming a train of localised wave modes where the steepening is produced by sideband generation of higher spatial harmonics all propagating (approximately) at the same phase (group) velocity such that their amplitudes are in phase and superimpose on the mother wave. When the gradient length of the leading wave packet becomes comparable to the dissipation scale L_d , dissipation sets on. At this time the smaller scale higher harmonic sidebands either outrun the leading wave packet ending up as standing, spatially damped precursor wave modes in front of the shock, or forming a spatially damped trailing wake of the packet. This depends on whether the dispersion is convex or concave as we have described in Chapter 3. This dispersive effect limits the amplitude of the shock.

At the same time the ramp is formed out of the wave packet by the dissipation generated inside the shock.

Generation of dissipation is most likely due to electron current instabilities of the shock ramp current on a scale that is shorter than the ion inertial scale. So far the instability has not yet been identified, but we have given strong arguments that it is the modified two stream instability which signs responsible. The anomalous collision rate is at the lower hybrid frequency in the shock ramp, quite high in this case and sufficient for providing the necessary dissipation for entropy generation, shock heating and compression. In addition, other small scale effects might occur which we have only given a hint on but not discussed in depth. Some processes have been discussed in the literature [see, e.g., *Biskamp, 1973; Tidman & Krall, 1971; Galeev, 1976*, and others]. In the following chapters we turn to the investigation of the substantially more interesting shocks which form the extended class of supercritical shocks.

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Quasi-perpendicular Supercritical Shocks

Abstract. Quasi-perpendicular shocks are the first and important family of collisionless magnetised shocks which reflect particles back upstream in order to satisfy the shock conditions. Discussion of the particle dynamics gives clear definition for distinguishing them from quasi-parallel shocks by defining a shock normal angle with respect to the upstream magnetic field. They exist for shock normal angles $<45^\circ$. Reflected particles at quasi-perpendicular shocks cannot escape far upstream along the magnetic field. They form a foot in front of the shock ramp. We discuss the reflecting shock potential and the explicit shock structure. Most theoretical insight is provided by numerical simulations which confirm reflection, foot formation and reformation of the shock. The latter being caused by steeping of the foot disturbance until the foot itself becomes the shock transition, reflecting particles upstream. Reformation modulates the shock temporarily but on the long terms guarantees its stationarity. Ion and electron dynamics are explicitly discussed in view of the various instabilities involved as well as particle acceleration and shock heating. Finally, a sketchy model of a typical quasi-perpendicular shock transition is provided.

5.1 Setting the Frame

As long as the shocks are subcritical with Mach numbers $\mathcal{M} < \mathcal{M}_c$ the distinction between quasi-perpendicular and quasi-parallel shocks is not overwhelmingly important, at least as long as the shock normal angle is far from zero. The mechanisms of dissipation in such sub-critical (or laminar) shocks have been discussed in the previous chapter. However, when the Mach number increases and finally exceeds the critical Mach number, $\mathcal{M} > \mathcal{M}_c$, the distinction becomes very important.

We speak of quasi-perpendicular super-critical shocks when the shock-normal angles $\Theta_{Bn} < 45^\circ$, and this because of good reasons. First, super-critical shocks cannot be maintained by dissipation alone. This has been clarified in Chapters 2 and 3. The inflow of matter into a supercritical shock is so fast that the time scales on which dissipation would take place are too long for dissipating the excess energy and lowering the inflow velocity below the downstream magnetosonic velocity. Hence, the condition for criticality, as we have shown in Chapter 2, is that the downstream flow velocity becomes equal to the downstream magnetosonic speed, which yielded the critical Mach number, $\mathcal{M}_c \lesssim 2.76$. We have also shown that $\mathcal{M}_c(\Theta_{Bn})$ is a function of the shock normal angle and can become quite small, even though $\mathcal{M}_c(\Theta_{Bn}) \gtrsim 1$ for existence of a shock, of course.

In order to help maintain a shock in the supercritical case the shock must forbid an increasing number of ions to pass across its ramp, which is done by reflecting some particles back upstream. This is not a direct dissipation process, rather it is an emergency act of the shock. It throws a fraction of the incoming ions back upstream and by this reduces

both the inflow momentum and energy densities. Clearly, this reflection process slows the shock down by attributing a negative momentum to the shock itself. The shock slips back and thus in the shock frame also reduces the difference velocity to the inflow, i.e. it reduces the Mach number. In addition, however, the reflected ions form an unexpected obstacle for the inflow and in this way reduce the Mach number a second time.

These processes are very difficult to understand, and we will go into more detail of them in this chapter. However, we must ask first, what the reason is for this rigid limit in Θ_{Bn} for calling a shock a quasi-perpendicular supercritical shock. The answer is that a shock as long belongs to the class of quasi-perpendicular shocks as reflected particles cannot escape from it upstream along the upstream magnetic field. After having performed half a gyro-circle back upstream they return to the shock ramp and ultimately traverse it to become members of the downstream plasma population.

5.1.1 Particle Dynamics

To see this we must return to the orbit a particle performs in interaction with a supercritical shock when it becomes reflected from the shock. In the simplest possible model one assumes the shock to be a plane surface, and the reflection being specular turning the component v_n of the instantaneous particle velocity $\mathbf{v}$ normal to the shock by 180° , i.e. simply inflecting it. In a very simplified version we have already considered this problem in Chapter 3. Here we follow the explicit calculation for these idealised conditions as given by *Schwartz et al* [1983] who treated this problem in the most general way. One should, however, keep in mind that the assumption of ideal specular reflection is the extreme limit of what happens in reality. In fact, reflection must by no means be specular because of many reasons. One of the reasons is that the shock ramp is not a rigid wall; the particles penetrate into it at least over a distance of a fraction of their gyroradius. In addition, they interact with waves and even excite waves during this interaction and during their approach of the shock. Altogether, it must be stressed again that the very mechanisms by which they become reflected are poorly known, indeed. Specular reflection is no more than a convenient assumption. Nevertheless, observations suggest that assuming specular reflection seems to be quite a useful approximation to reality as long as nothing more precise is known about the inelastic reflection processes.

Figure 5.1 shows the coordinate frame used at the planar (stationary) shock, with shock normal $\mathbf{n}$, magnetic $\hat{\mathbf{b}}$ and velocity $\hat{\mathbf{v}}$ unit vectors, respectively. Shown are the angles Θ_{Bn} , θ_{Vn} , θ_{BV} . The velocity vector $\mathbf{V}_{HT}$ is the de Hoffmann-Teller velocity which lies in the shock plane and is defined in such a way that in the coordinate system moving along the shock plane with velocity $\mathbf{V}_{HT}$ the plasma flow is along the magnetic field. $\mathbf{V} - \mathbf{V}_{HT} = -v_{\parallel} \hat{\mathbf{b}}$. Because of the latter reason it is convenient to consider the motion of particles in the de Hoffmann-Teller frame. The guiding centres of the particles in this frame move all along the magnetic field. Hence, using $\mathbf{V} = -V\hat{\mathbf{v}}$, $\mathbf{n} \cdot \hat{\mathbf{v}} = \cos \theta_{Vn}$, $\mathbf{n} \cdot (\hat{\mathbf{b}}, \hat{\mathbf{x}}, \hat{\mathbf{y}}) = (\cos \Theta_{Bn}, \sin \Theta_{Bn}, 0)$,

$$v_{\parallel} = V \frac{\cos \theta_{Vn}}{\cos \Theta_{Bn}}, \quad \mathbf{V}_{HT} = V \left(-\hat{\mathbf{v}} + \frac{\cos \theta_{Vn}}{\cos \Theta_{Bn}} \hat{\mathbf{b}} \right) \equiv \frac{\mathbf{n} \times \mathbf{V} \times \mathbf{B}}{\mathbf{n} \cdot \mathbf{B}}, \quad V_{HT,n} \equiv 0 \quad (5.1)$$

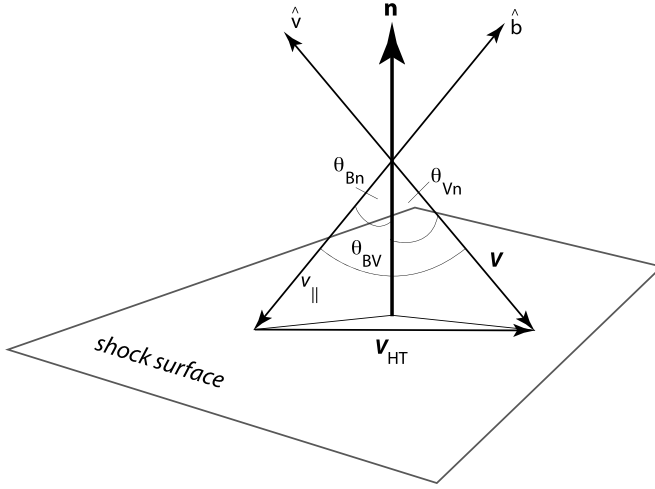


Figure 5.1: The shock coordinate system showing the shock normal $\mathbf{n}$, velocity and magnetic field directions $\hat{\mathbf{v}}$, $\hat{\mathbf{b}}$, the three angles Θ_{Bn} , θ_{Vn} , θ_{BV} between $\hat{\mathbf{b}}$ and $\mathbf{n}$, velocity $\mathbf{V}$ and $\mathbf{n}$, and velocity $\mathbf{V}$ and $\hat{\mathbf{b}}$, respectively. The velocity $\mathbf{V}_{HT}$ in the shock plane is the de Hoffmann-Teller velocity [after Schwartz *et al*, 1983, courtesy American Geophysical Union].

The de Hoffmann-Teller velocity is the same to both sides of the shock ramp, because of the continuity of normal component B_n and tangential electric field $\mathbf{E}_t$. Thus, in the de Hoffmann-Teller frame there is no induction electric field $\mathbf{E} = -\mathbf{n} \times \mathbf{V} \times \mathbf{B}$. The remaining problem is two-dimensional (because trivially $\mathbf{n}$, $\hat{\mathbf{b}}$ and $-\nu_{\parallel} \hat{\mathbf{b}}$ are coplanar, which is nothing else but the coplanarity theorem holding under these undisturbed idealised conditions).

In the de Hoffmann-Teller (primed) frame the particle velocity is described by the motion along the magnetic field $\hat{\mathbf{b}}$ plus the gyromotion of the particle in the plane perpendicular to $\hat{\mathbf{b}}$:

$$\mathbf{v}'(t) = v'_{\parallel} \hat{\mathbf{b}} + v_{\perp} [\hat{\mathbf{x}} \cos(\omega_{ci} t + \phi_0) \mp \hat{\mathbf{y}} \sin(\omega_{ci} t + \phi_0)] \quad (5.2)$$

The unit vectors $\hat{\mathbf{x}}$, $\hat{\mathbf{y}}$ are along the orthogonal coordinates in the gyration plane of the ion, the phase ϕ_0 accounts for the initial gyro-phase of the ion, and $\pm$ accounts for the direction of the upstream magnetic field being parallel (+) or antiparallel to $\hat{\mathbf{b}}$.

In specular reflection (from a stationary shock) the upstream velocity component along $\mathbf{n}$ is reversed, and hence (for cold ions) the velocity becomes

$$\mathbf{v}' = -v_{\parallel} \hat{\mathbf{b}} + 2v_{\parallel} \cos \Theta_{Bn} \hat{\mathbf{n}}$$

which (with $\phi_0 = 0$) yields for the components of the velocity

$$\frac{v'_{\parallel}}{V} = \frac{\cos \theta_{Vn}}{\cos \Theta_{Bn}} (2 \cos^2 \Theta_{Bn} - 1), \quad \frac{v_{\perp}}{V} = 2 \sin \Theta_{Bn} \cos \theta_{Vn} \quad (5.3)$$

These expressions can be transformed back into the observer's frame by using $\mathbf{V}_{HT}$. It is, however, of greater interest to see under which conditions a reflected particle turns around

in its upstream motion towards the shock. This happens when the upstream component of the velocity $v_x = 0$ of the reflected ion vanishes. For this we need to integrate Eq. (5.2) which for $\phi_0 = 0$ yields

$$\mathbf{x}'(t) = v'_{\parallel} t \hat{b} + \frac{v_{\perp}}{\omega_{ci}} \{ \hat{x} \sin \omega_{ci} t \pm \hat{y} (\cos \omega_{ci} t - 1) \} \quad (5.4)$$

Scalar multiplication with $\mathbf{n}$ yields the ion displacement normal to the shock in upstream direction. The resulting expression

$$\mathbf{x}'_n(t^*) = v'_{\parallel} t^* \cos \Theta_{Bn} + \frac{v_{\perp}}{\omega_{ci}} \sin \Theta_{Bn} \sin \omega_{ci} t^* = 0 \quad (5.5)$$

vanishes at time t^* when the ion re-encounters the shock with normal velocity $v_n(t^*) = v'_{\parallel} \cos \Theta_{Bn} + v_{\perp} \sin \Theta_{Bn} \cos \omega_{ci} t^*$. The maximum displacement away from the shock in normal direction is obtained when setting this velocity to zero, obtaining for the time t_m at maximum displacement (again including the initial phase here)

$$\omega_{ci} t_m + \phi_0 = \cos^{-1} \left(\frac{1 - 2 \cos^2 \Theta_{Bn}}{2 \sin^2 \Theta_{Bn}} \right) \quad (5.6)$$

This expression must be inserted in $\mathbf{x}_n$ yielding for the distance a reflected ion with gyro-radius $r_{ci} = V/\omega_{ci}$ can achieve in upstream direction

$$\Delta x_n = r_{ci} \cos \theta_{Vn} \{ (\omega_{ci} t_m + \phi_0) (2 \cos^2 \Theta_{Bn} - 1) + 2 \sin^2 \Theta_{Bn} \sin(\omega_{ci} t_m + \phi_0) \} \quad (5.7)$$

For a perpendicular shock $\Theta_{Bn} = 90^\circ$ and $\phi_0 = 0$ this distance is $\Delta x_n \simeq 0.7 r_{ci} \cos \theta_{Vn}$, less than an ion gyro radius. The distance depends on the shock normal angle, decreasing for non-planar shocks. Note that the argument of $\cos^{-1}$ in Eq. (5.6) changes sign for $\Theta_{Bn} \leq 45^\circ$. Equation (5.5) has solutions for positive upstream turning distances only for shock normal angles $\Theta_{Bn} > 45^\circ$, for an initial particle phase $\phi_0 = 0$. (A finite initial phase $\phi_0 \neq 0$ may, however, modify this conclusion shifting the boundary between quasi-parallel and quasi-perpendicular shock to angles larger or smaller than 45° , depending on the sign of the initial phase.) Reflected ions can return to the shock in one gyration time only when the magnetic field makes an angle with the shock normal that is larger than this value. For less inclined shock normal angles the reflected ions escape along the magnetic field upstream of the shock and do not return within one gyration. This sharp distinction between shock normal angles $\Theta_{Bn} < 45^\circ$ and $\Theta_{Bn} > 45^\circ$ thus provides the natural (kinematic specular) discrimination between quasi-perpendicular and quasi-parallel (planar) shocks we were looking for.

The theory of shock particle reflection holds, in this form, only for cold ions, which implies complete neglect of any velocity dispersion and proper gyration of the ions. The ions are considered of just moving all with one and the same oblique flow velocity $\mathbf{V}$. In a warm plasma each particle has a different speed, and it is only the group of zero velocity ions which are described by the above theory. Fortunately, these are the particles which experience the reflecting shock potential strongest and are most vulnerable to specular reflection. When temperature effects will be included, the theory is more involved in

a number of ways. Firstly, the de-Hoffmann-Teller velocity must be redefined to include the microscopic particle motion. Secondly, the assumption of ideal specular reflection becomes questionable, as the particles themselves become involved into the generation of the shock potential. This problem is still open to investigation. Observations in space suggest that, for high flow velocities and supercritical Mach numbers, the simple kinematic reflection is a sufficiently well justified mechanism, however.

We may, formally, extend the above approach for warm ions to include the proper particle motion. In order to distinguish the different velocity components in this case, we indicate the bulk flow components by a tilde, $\tilde{V}'$, and write the proper microscopic velocity $\mathbf{v}_i$. Warm particles obey a phase space velocity distribution with different velocity components parallel, $v_{\parallel i}$, and perpendicular, $v_{\perp i}$, to the magnetic field, as well as gyration phases $\omega_{ci}t + \phi_0$. These add to the gyro-centre velocity. (One should note that in principle for high Mach numbers $V \gg v_i$ the velocity addition theorem should be applied in its relativistic version even if the proper particle motions particles are considered non-relativistic.) In the completely non-relativistic case where the velocities simply add linearly, the specularly reflected normal component of the particle velocity after reflection becomes sufficiently complicate:

$$v_n = (v_{\parallel i} + \tilde{V}'_{\parallel}) \cos \Theta_{Bn} + [v_{\perp i} \cos(\omega_{ci}t + \phi_0) + \tilde{V}'_{\perp} \cos \omega_{ci}t] \sin \Theta_{Bn} \quad (5.8)$$

The gyration phase ϕ_0 of the particle must be retained in this case as it is different from that of the bulk flow phase. The inclusion of many gyrating particles inhibits to identify it with that flow phase. Setting this normal component of velocity after reflection to zero does not lead to a simple expression for the turning point of an upstream reflected particle orbit, and the distinction between parallel and perpendicular shocks becomes blurred as it can at most be defined only approximately and in the average over the particle distributions and phases. On defining a new phase ψ through

$$\tan \psi = \frac{v_{\perp i} \sin \phi_0}{\tilde{V}'_{\perp} + v_{\perp i} \cos \phi_0} \quad (5.9)$$

the above expression set to zero can be rewritten into the form

$$\cos(\omega_{ci}t + \psi) = -\frac{v_{\parallel i} + \tilde{V}'_{\parallel}}{\tilde{V}'_{\perp}} \frac{\cos \Theta_{Bn}}{\sin \Theta_{Bn}} \quad (5.10)$$

again with the modulus of the expression on the right-hand side required to be <1 as the new condition to discriminate between parallel and perpendicular shocks. Clearly this condition makes sense only when averaged over the thermal particle distribution and over the phases of all particles. This average should, however, be taken already in Eq. (5.8).

In order to do so, we assume a simple thermal distribution of the upstream particles with anisotropic temperature T_i . In carrying out the averaging one must observe that the phase of the reflected particles can vary only in the forward half-space interval $0 \leq \phi_0 \leq \pi$, while the limits on the reflected parallel and perpendicular velocities are $-\tilde{V}'_{\parallel} \leq v_{\parallel} \leq \infty$, and $0 \leq v_{\perp} \leq \infty$, respectively.

Now carrying out the integration and defining the new mock phase through $\tan \psi = v_{\perp i} / \sqrt{\pi} V'_{\perp}$, we find that for a hot plasma (for generality including a temperature anisotropy $T_{i\perp} / T_{i\parallel}$) the condition for a turning upstream gyration orbit after specular reflection can formally be written as

$$\frac{\cos(\omega_{ci}t + \psi)}{\sqrt{\pi}} = \frac{1 - 2 \cos^2 \Theta_{Bn}}{2 \sin^2 \Theta_{Bn}} + \frac{v_{\parallel i}}{V} \frac{\cos \Theta_{Bn}}{2 \sin^2 \Theta_{Bn} \cos \theta_{Vn}} \frac{\exp(-\tilde{V}'^2_{\parallel} / v_{\parallel i}^2)}{\sqrt{\pi} [1 + \operatorname{erf}(\tilde{V}'_{\parallel} / v_{\parallel i})]} \quad (5.11)$$

which in this case replaces the earlier simple condition (5.6) of the cold plasma. This is contained here in the first term on the right to which the second term adds a rather complicated thermal correction which becomes important at large parallel temperatures only, however, as indicated by the appearance of the ratio $v_{\parallel i} / V$. The presence of this term modifies the angle at which the sign of the right-hand side changes and thus also modifies the angle of transition from a quasi-parallel to a quasi-perpendicular shock. This modification will, of course, be more substantial the larger the upstream parallel temperature is and will further be modified if the plasma becomes relativistic, in which case it should depend on Γ . One, on the other hand, expects that highly relativistic plasmas with $\Gamma \gg 1$ will behave about like cold plasmas and therefore yield the ordinary distinction between quasi-parallel and quasi-perpendicular shocks. However, the thermal modification, though possibly not of vital importance for this distinction, will substantially affect the distance up to that a specularly reflected particle at the quasi-perpendicular shock can penetrate the upstream flow, i.e. it affects the width of the quasi-perpendicular shock-foot, even in the case when the reflection process is genuinely specular.

5.1.2 Foot Formation and Acceleration

Shock reflected ions in a quasi-perpendicular shock cannot escape far upstream (see Figure 5.2). Their penetration into the upstream plasma is severely restricted by formula (5.7). Within this distance the ions perform a gyration orbit before returning to the shock.

Since the reflected ions are about at rest with respect to the inflowing plasma they are sensitive to the inductive convection electric field $\mathbf{E} = -\mathbf{V}_1 \times \mathbf{B}_1$ behaving very similar to pick-up ions and becoming accelerated in the direction of this field to achieve a higher energy [Schwartz *et al.*, 1983]. When returning to the shock their maximum (minimum) achievable energy is

$$\mathcal{E}_{\max} = \frac{m_i}{2} [(v'_{\parallel} + V_{\text{HT}\parallel})^2 + (V_{\text{HT}\perp} \pm v_{\perp})^2] \quad (5.12)$$

This energy is larger than their initial energy with that they have initially met the shock ramp and, under favourable conditions, they now might overcome the shock ramp potential and escape downstream. Otherwise, when becoming reflected again, they gain energy in a second round until having picked up sufficient energy for passing the shock ramp.

In addition to this energisation of reflected ions which in the first place have not made it across the shock, the reflected ions when gyrating and being accelerated in the convection electric field constitute a current layer just in front of the shock ramp of current density

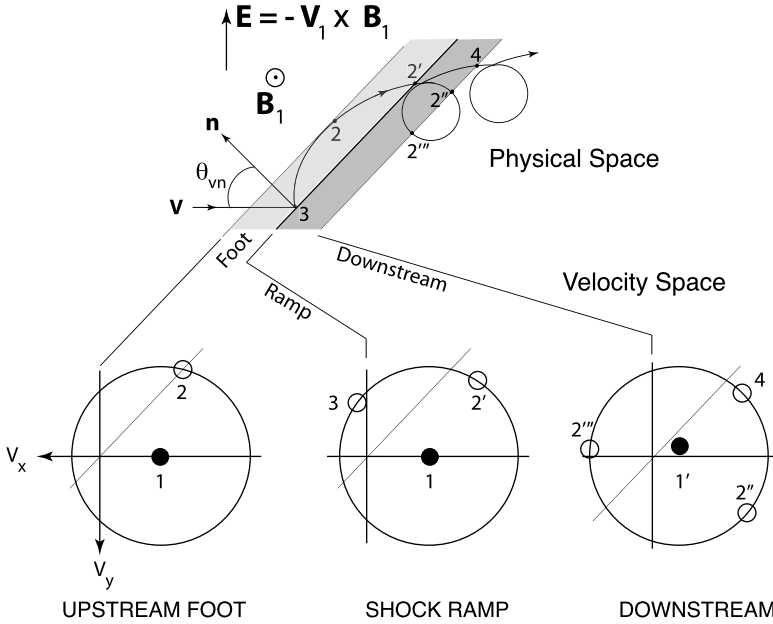


Figure 5.2: *Top:* Reflected ion orbits in the foot of a quasi-perpendicular shock in real space. The ion impacts under an instantaneous angle θ_{vn} , is reflected from the infinitely thin shock, performs a further partial gyration in the upstream field $\mathbf{B}_1$ where it is exposed to the upstream convection electric field $\mathbf{E} = -\mathbf{V}_1 \times \mathbf{B}_1$ in which it is accelerated as is seen from the non-circular section of its orbit in the shock foot. It hits the shock ramp a second time now at energy high enough to overcome the shock potential, passing the ramp and arriving in the compressed downstream magnetic field behind the shock where it performs gyrations of reduced gyro-radius. *Bottom:* The ion distribution function mapped into velocity space v_x, v_y for the indicated regions in real space, upstream in the foot, at the ramp, and downstream of the shock ramp. Upstream the distribution consists of the incoming dense plasma flow (population 1, dark circle at $v_y = 0$) and the reflected distribution 2 at large negative v_y . At the ramp in addition to the incoming flow 1 and the accelerated distribution 2' there is the newly reflected distribution 3. Behind the ramp in the downstream region the inflow is decelerated 1' and slightly deflected toward non-zero v_y , and the energised passing ions exhibit gyration motions in different instantaneous phases, two of them (2'', 4) directed downstream, one of them (2''') directed upstream. [redrawn after *Skopke et al*, 1983, courtesy American Geophysical Union].

$j_y \sim eN_{i,\text{refl}}v_{y,\text{refl}}$ which gives rise to a foot magnetic field of magnitude $B_{z,\text{foot}} \sim \mu_0 j_y \Delta x_n$. It is clear that this foot ion current, which is essentially a drift current in which only the reflected newly energised ion component participates, constitutes a source of free energy as it violates the energetic minimum state of the inflowing plasma in its frame. Being the source of free energy it can serve as a source for excitation of waves via which it will contribute to filling the lack of dissipation. However, in a quasi-perpendicular shock there are other sources of free energy as well which are not restricted to the foot region.

Figure 5.3 shows a sketch of some of the different free-energy sources and processes across the quasi-perpendicular shock. In addition to the shock-foot current and the presence of the fast cross-magnetic field ion beam there, the shock ramp is of finite thickness. It contains a charge separation electric field E_x which in the supercritical shock is strong

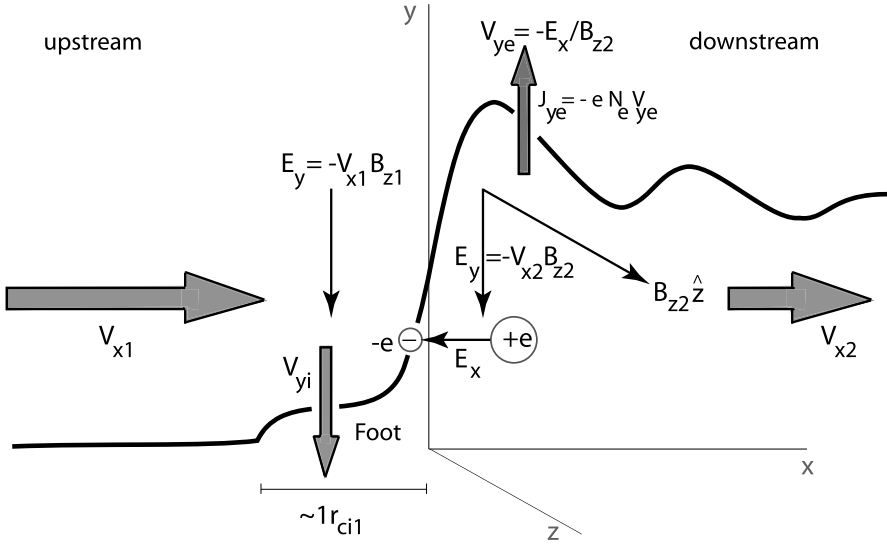


Figure 5.3: Geometry of an ideally perpendicular supercritical shock showing the field structure and sources of free energy. The shock is a compressive structure. The profile of the shock thus stands for the compressed profile of the magnetic field strength $|\mathbf{B}|$, the density N , temperature T , and pressure NT of the various components of the plasma. The inflow of velocity V_1 and outflow of velocity V_2 is in x direction, and the magnetic field is in z direction. Charge separation over an ion gyroradius r_{ci} in the shock ramp magnetic field generates a charge separation electric field E_x along the shock normal which reflects the low-energy ions back upstream. These ions see the convection electric field E_y of the inflow, which is along the shock front, and become accelerated. The magnetic field of the current carried by the accelerated back-streaming ions causes the magnetic foot in front of the shock ramp. The shock electrons are accelerated antiparallel to E_x perpendicular to the magnetic field. The shock electrons also perform an electric field drift in y -direction in the crossed E_x and compressed B_z fields which leads to an electron current j_y along the shock. These different currents are sources of free energy which drives various instabilities in different regions of the perpendicular shock.

enough to reflect the lower energy ions. In addition it accelerates electrons downstream thereby deforming the electron distribution function.

The presence of this field, which has a substantial component perpendicular to the magnetic field, implies that the magnetised electrons with their gyro radii being smaller than the shock-ramp width experience an electric drift $V_{ye} = -E_x/B_{z2}$ along the shock in the ramp which can be quite substantial giving rise to an electron drift current $j_{ye} = -eN_{e,ramp}V_{ye} = eN_{e,ramp}E_x/B_{z2}$ in y -direction. This current has again its own contribution to the magnetic field, which at maximum is roughly given by $B_z \sim \mu_0 j_{ye} \Delta x_n$. Here we use the width of the shock ramp. The electron current region might be narrower, of the order of the electron skin depth c/ω_{pe} . However, as long as we do not know the number of magnetised electrons which are involved into this current nor the width of the electric field region (which must be less than an ion gyro-radius because of ambipolar effects) the above estimate is good enough.

The magnetic field of the electron drift current causes an overshoot in the magnetic field in the shock ramp on the downstream side and a depletion of the field on the upstream

side contributing to the steepness of the ramp. When this current becomes strong it contributes to current-driven cross-field instabilities like the modified two-stream instability.

Finally, the mutual interaction of the different particle populations present in the shock at its ramp and behind provide other sources of free energy. A wealth of instabilities and waves is thus expected to be generated inside the shock. To these micro-instabilities add the longer wavelength instabilities which are caused by the plasma and field gradients in this region. These are usually believed to be less important as the crossing time of the shock is shorter than their growth time. However, some of them propagate along the shock and have therefore substantial time to grow and modify the shock profile. In the following we will turn to the discussion of numerical investigations of some of these processes reviewing their current state and provide comparison with observations.

5.1.3 Shock Potential Drop

One of the important shock parameters is the electric potential drop across the shock ramp – or if it exists also across the shock foot. This potential drop is not necessarily a constant but changes with location along the shock normal. We have already noted that it is due to the different dynamical responses of the inflowing ions and electrons over the scale of the foot and ramp regions. Its theoretical determination is difficult, however when going to the de Hoffmann-Teller frame the bulk motion of the particles is only along the magnetic field, and in the stationary electron equation of motion the $\mathbf{V}_e \times \mathbf{B}$ -term drops out and, to first approximation, the cross shock potential is given by the pressure gradient (when neglecting any contributions from wave fields). The expression is then simply

$$\Delta\Phi(x) = \int_0^x \frac{1}{eN_e(n)} [\nabla \cdot \mathbf{P}_e(n)] \cdot d\mathbf{n} \quad (5.13)$$

Integration is over n along the shock normal $\mathbf{n}$. For a gyrotropic electron pressure, valid for length scales longer than an electron gyroradius, $\mathbf{P}_e = P_{e\perp} \mathbf{I} + (P_{e\parallel} - P_{e\perp}) \mathbf{BB}/B^2$ one obtains [Goodrich & Scudder, 1984], taking into account that $\mathbf{E} \cdot \mathbf{B}$ is invariant,

$$\frac{d}{dn} \Phi(n) = -\frac{E_{\parallel}}{\cos \Theta_{Bn}} = \frac{1}{eN_e} \left[\frac{d}{dn} P_{e\parallel} - (P_{e\parallel} - P_{e\perp}) \frac{d}{dn} (\ln B) \right] \quad (5.14)$$

which, when used in the above expression, yields

$$e\Delta\Phi(x) = \int_0^x dn \left\{ \frac{dT_{e\parallel}}{dn} + T_{e\parallel} \frac{d}{dn} \ln \left[\frac{N(n)}{N_1} \frac{B_1}{B(n)} \right] + T_{e\perp} \frac{d}{dn} \ln \left[\frac{B(n)}{B_1} \right] \right\} \quad (5.15)$$

This expression can approximately be written in terms of the gradient in the electron magnetic moment $\mu_e = T_{e\perp}/B$ as follows:

$$e\Delta\Phi(x) \simeq \Delta(T_{e\parallel} + T_{e\perp}) - \int_0^x dn \frac{d\mu_e(n)}{dn} B(n) \quad (5.16)$$

with T_e in energy units. When the electron magnetic moment is conserved, the last term disappears, yielding a simple relation for the potential drop $e\Delta\Phi(x) \simeq \Delta(T_{e\parallel} + T_{e\perp})$ as the

sum of the changes in electron temperature. The perpendicular temperature change can be expressed as $\Delta T_{e\perp} = T_{e\perp,1} \Delta B / B_1$ which is in terms of the compression of the magnetic field. Non-adiabatic effects contribute via the dropped integral term.

The parallel change in temperature is more difficult to express. One could express it in terms of the temperature anisotropy $A_e = T_{e\parallel} / T_{e\perp}$ as has been done by *Kuncic et al* [2002], and then vary A_e . But this depends on the particular model. It is more important to note that this adiabatic estimate of the potential drop does not account for any dynamical process which generates waves and substructures in the shock. It thus gives only a hint on the order of magnitude of the potential drop across the foot-ramp region in quasi-perpendicular shocks.

5.2 Shock Structure

Figure 5.4 shows observations from one of the first unambiguous satellite crossings of a quasi-perpendicular supercritical (magnetosonic Mach number $\mathcal{M}_{ms} \sim 4.2$) shock in near Earth space. The crossing occurred at the Earth's bow shock, the best investigated shock in the entire cosmos! A complete discussion of its properties will be given in Chapter 10. Here it should mainly serve for visualisation of the properties of a real collisionless shock how it appears in the data. The shock crossing shown in the figure is indeed a textbook example.

5.2.1 Observational Evidence

The crossing occurred on an inbound path of the two spacecraft ISEE 1 (upper block of the figure) and ISEE 2 (lower block of the figure) from upstream to downstream in short sequence only minutes apart. In spite of some differences occurring on the short time scale the two shock crossings are about identical, identifying the main shock transition as a spatial and not as a temporal structure. Temporal variations are nevertheless visible on the scale of a fraction of a minute.

From top to bottom the figure shows the electron density (N_E), energetic ion density (N_I), proton and electron temperatures (T_P, T_E), bulk flow velocity (V_P), electron pressure (P_E), magnetic field (B), and Θ_{Bn} . The latter is close to 90° prior to shock crossing (in the average $\Theta_{Bn} \sim 85^\circ$), and fluctuates afterwards around 90° identifying the shock as quasi-perpendicular. Accordingly, the shock develops a foot in front of the shock ramp as can be seen from the slightly enhanced magnetic field after 22:51 UT in ISEE 1 and similar in ISEE 2, and most interestingly also in the electron pressure. At the same time the bulk flow velocity starts decreasing already, as the result of interaction and retardation in the shock foot region. The foot is also visible in the electron density which increases throughout the foot region, indicating the presence of electrons which, as is suggested by the increase in pressure, must have been heated or accelerated.

The best indication of the presence of the foot is, however, the measurement of energetic ions (second panel from top). These ions are observed first some distance away from the shock but increase drastically in intensity when entering the foot. These are the shock-

reflected ions which have been accelerated in the convection electric field in front of the shock ramp. Their occurrence before entrance into the foot is understood when realising that the shock is not perfectly perpendicular. Rather it is quasi-perpendicular such that part of the reflected ions having sufficiently large parallel upstream velocities can escape along the magnetic field a distance larger than the average upstream extension of the foot. For nearly perpendicular shocks, this percentage is small.

The shock ramp in Figure 5.4 is a steep wall in B and P_E , respectively. The electron temperature T_E increases only moderately across the shock while the ion temperature T_P jumps up by more than one magnitude, exceeding T_E downstream behind the shock. This behaviour is due to the accelerated returning foot-ions which pass the shock. P_E , B , and N_E exhibit overshoots behind the shock ramp proper. Farther away from the shock they merge into the highly fluctuating state of lesser density, pressure, and magnetic field that can be described as some kind of turbulence. Clearly, this region is strongly affected by the presence of the shock which forms one of its boundaries, the other boundary being the obstacle which is the main responsible for the formation of the shock.

The evidence provided by the described measurements suggests that the quasi-perpendicular shock is a quasi-stationary entity. This should, however, not been taken as apodictive. Stationarity depends on the spatial scales as well as the time scales. A shock is a very inhomogeneous subject containing all kinds of spatial scales. Being stationary on one scale does not imply that it is stationary on another scale. For a shock like the Earth's bow shock considered over times of days, weeks or years the shock is of course a stationary subject. However on shorter time scales of the order of flow transition times this may not be the case. A subcritical shock of the kind discussed in Chapter 4 may well be stationary on long and short time scales. However, for a supercritical shock the conditions for forming a stationary state are quite subtle. From a single spacecraft passage like that described above it cannot be concluded to what extent, i.e. on which time scale and on which spatial scale and under which external conditions (Mach number, angle, shock potential, plasma- β , ...) the observed shock can be considered to be stationary [a discussion of the various scales has been given, e.g. by *Galeev et al*, 1988]. Comparison between the two ISEE spacecraft already shows that the small-scale details as have been detected by both spacecraft are very different. This suggests that – in this case – on time scales less than a minute variations in the shock structure must be expected.

Generally spoken, one must be prepared to consider the shock locally (on the ion gyro-scale) and temporarily (on the ion-cyclotron frequency scale) as a non-stationary phenomenon [this has been realised first by *Morse et al*, 1972] which depends on many competing processes and, most important though only secondarily related to non-stationarity, a shock as a whole is not in thermal equilibrium. It needs to be driven by some energy source external to the shock in order to be maintained. It will thus be very sensitive to small changes in the external parameters and will permanently try to escape the non-equilibrium state and to approach equilibrium. Since its non-equilibrium is maintained by the conditions in the flow, it is these conditions which determine the time scales over which a shock evolves, re-evolves and changes its state. In the following we will refrain from analytical theory and, forced by the complexity of the problem, mainly discuss numerical experiments on shocks. However, at a later stage we will return to the problem of non-

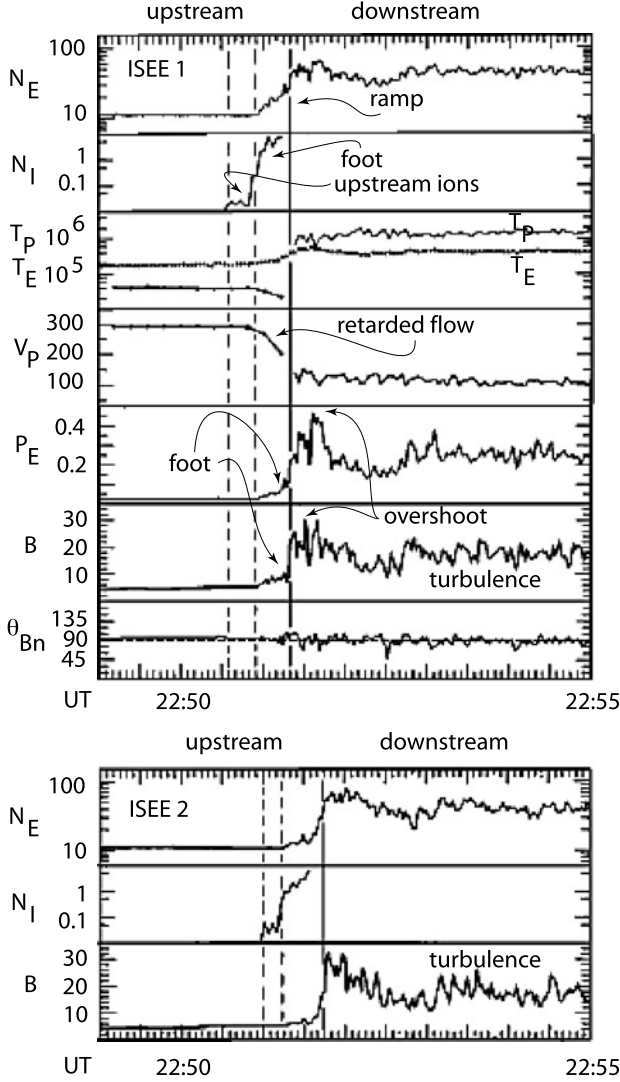


Figure 5.4: Time profiles of plasma and magnetic field parameters across a real quasi-perpendicular shock that had been crossed by the ISEE 1 and 2 spacecraft on November 7, 1977 in near-Earth space [after *Scopke et al, 1983*, courtesy American Geophysical Union]. The shock in question is the Earth's bow shock wave which will be described in detail Chapter 10. Here the measurement serve as typical for a quasi-perpendicular shock. N_E is the electron density, N_I the reflected ion density, both in cm^{-3} , T_P , T_E are proton and electron in K. V_P is the proton (plasma) bulk velocity in km s^{-1} , P_E electron pressure in 10^{-9} N m^{-2} , B the magnitude of the magnetic field in nT, and θ_{Bn} . The vertical lines mark the first appearance of reflected ion, the outer edge of the foot in the magnetic profile, and the ramp in the field magnitude, respectively. The *abscissa* is the Universal Time UT referring to the measurements. The *upper block* are observations from ISEE 1, the *lower block* observations from ISEE 2.

stationarity again. Real supercritical shocks, whether quasi-perpendicular or quasi-parallel are in a permanently evolving state and thus are *intrinsically* nonstationary.

5.2.2 Simulation Studies of Quasi-perpendicular Shock Structure

Perpendicular or quasi-perpendicular collisionless shocks are relatively easy to treat in numerical simulations. Simulations of this kind have been mostly one-dimensional. Only more recently they have begun to become treated in two dimensions.

Already from the first one-dimensional numerical experiments on collisionless shocks [for an early review cf., e.g., *Biskamp, 1973*] it became clear that such shocks have a very particular structure. This structure, which we have describe in simplified version in Figure 5.3 and which could to some extent also be inferred from the observations of Figure 5.4, becomes ever more pronounced the more refined the resolution becomes and the better the shorter scales can be resolved.

As already mentioned, collisionless shocks are in thermodynamic non-equilibrium and therefore can only evolve if a free energy source exists and if the processes are violent enough to build up and maintain a shock. Usually in a freely evolving system the free energy causes fluctuations which serve dissipating and redistributing the free energy towards thermodynamic and thermal equilibria. (Thermal equilibria are characterised by equal temperatures among the different components, e.g. $T_e = T_i$ which is clearly not given in the vicinity of a shock as seen from Figure 5.4. Thermodynamic equilibria are characterised by Gaussian distributions for all components of the plasma. To check this requires information about the phase space distribution of particles. Shocks contain many differing particle distributions, heated, top-flat, beam distributions, long energetic tails, and truncated as well as gyrating distributions which we will encounter later. Consequently, they are far from thermodynamic equilibrium.)

For a shock to evolve the amount of free energy needed to dissipate is so large that fluctuations are unable to exercise their duty. This happens at large Mach numbers. The shock itself takes over the duty of providing dissipation. It does it in providing all kinds of scales such short that a multitude of dissipative processes can set on.

Scales

For a quasi-perpendicular shock propagating and evolving in a high- β plasma there is a hierarchy of such scales available (we recall that $\beta = 2\mu_0 nT/B^2$ refers to the thermal energy of the flow. The kinetic $\beta_{kin\perp} = 2\mu_0 Nm_i V_n^2 / 2B^2 \equiv \mathcal{M}^2 > 1$ implies that the kinetic energy in the flow exceeds the magnetic energy. Hence the flow dominates the magnetic field, which is transported by the flow. In plasmas with $\beta_{kin\perp} = \mathcal{M}^2 < 1$ the magnetic field dominates the dynamics, and shock waves perpendicular to the magnetic field cannot evolve. Parallel shocks are basically electrostatic in the $\beta_{kin\perp} \ll 1$ -case and can evolve when the flow is sufficiently fast along the field, as is observed in the auroral magnetospheres of the magnetised planets in the heliosphere. On the other hand, for large Mach numbers and $\beta \gtrsim 1$ conditions shocks do exist, as the example of the solar wind shows).

The different scales can be organised with respect to the different regions of the shock.

1. The macroscopic scale of the foot region, which determines the width of the foot, is the ion gyroradius based on the inflow velocity $r_{ci,1} = V_1/\omega_{ci,1}$. With the slight modification of replacing the upstream magnetic field with the (inhomogeneous) ramp magnetic field $B_r(x)$ this also becomes approximately the scale of the macroscopic electric potential drop in the ramp, $\Delta\phi_{,r} \sim r_{ci,r} \sim V_1/\omega_{ci,r}$.

Other scales are

2. the ion inertial length c/ω_{pi} , which is also a function of space inside the ramp because of the steep density increase $N(x)$. It determines the dispersive properties of the fast magnetosonic wave which is locally responsible for steeping and shock ramp formation;

3. the thermal ion gyroradius $r_{ci} = v_i/\omega_{ci}$. It determines the transition from unmagnetised to magnetised ions and from non-adiabatic to adiabatic heating of the ions;

4. the density gradient scale $L_P = (\nabla_x \ln P)^{-1}$. It determines the importance of drift waves along the shock which, when excited, structure the shock in the third dimension perpendicular to the shock normal and the magnetic field;

5. the electron inertial length c/ω_{pe} . It is the scale length of whistlers which are excited in front of the shock and are generally believed to play an essential role in shock dynamics;

6. the thermal electron gyroradius $r_{ce} = v_e/\omega_{ce}$. It determines whether electrons behave magnetised or non-magnetised. In the shock they are usually magnetised under all conditions of interest. However, when non-adiabatic heating becomes important for electrons it takes place on scales comparable to r_{ce} ;

7. the Debye length λ_D . It determines the dispersive properties of ion acoustic waves which are responsible for anomalous resistivity and for smaller scale density substructures in the shock like the phase space holes mentioned earlier which evolve on scales of several Debye lengths. It also determines the scales of the Buneman two-stream (BTS) and modified two-stream (MTS) instabilities which are the two most important instabilities in the shock foot.

The importance of some of these scales has been discussed by [Kennel *et al.*, 1985] assuming that some mostly anomalous resistance has been generated in the plasma. In this case the speed of the fast magnetosonic wave, which is responsible for fast shock formation, is written as

$$c_{ms}^2 = c_{ia}^2 + \frac{V_A^2}{1 + k^2 R^2}, \quad R = \begin{cases} R_\eta = (\eta/\mu_0)(k/\omega), & \eta \neq 0 \\ \lambda_e = c/\omega_{pe}, & \eta \rightarrow 0 \end{cases} \quad (5.17)$$

taking explicitly care of the dispersion of the wave which leads to wave steeping. The macroscopic scale of shock formation enters here through the definition of R which in the collisionless case becomes the electron skin depth. Starting from infinity far away from

the shock one seeks for growing solutions of the linear magnetic disturbance $b_z \sim \exp \lambda x$ in the stationary point equation

$$R_e^2 b_z'' + R_\eta b_z' = D b_z, \quad D \equiv \frac{1 - \mathcal{M}^{-2}}{1 - c_{ia}^2/V^2}$$

where the prime $' \equiv \partial/\partial x$ indicates derivation with respect to x . With $b_z \rightarrow 0$ for $x \rightarrow -\infty$ this yields for the spatial growth rate

$$\lambda_{>} = -\frac{R_\eta}{2\lambda_e^2} + \left[\frac{D}{\lambda_e^2} + \left(\frac{R_\eta}{2\lambda_e^2} \right)^2 \right]^{\frac{1}{2}} \rightarrow \frac{D^{\frac{1}{2}}}{\lambda_e} \quad \text{for} \quad R_\eta \ll \lambda_e \quad (5.18)$$

which identifies the approximate shock transition scale as proportional to the electron skin depth, $\Delta \simeq c/\omega_{pe} D^{\frac{1}{2}}$, just what one intuitively would believe to happen for freely moving electrons and ions. Since the upstream sound speed $c_{ia} \ll V$ is small compared with the fast flow V , we have $D \approx 1 - \mathcal{M}^{-2}$, and the shock ramp width becomes slightly larger than the electron skin depth $\lambda_e = c/\omega_{pe}$, viz.

$$\Delta \simeq \mathcal{M} (\mathcal{M}^2 - 1)^{-\frac{1}{2}} \lambda_e \quad (5.19)$$

For large Mach numbers this width approaches λ_e . However, we have already seen that at large Mach numbers the competition between dispersion and dissipation does not hold anymore in this simple way.

With increasing wave number k the fast magnetosonic mode merges into the whistler branch with its convex dispersion curve. This implies that dispersive whistler waves will outrun the shock becoming precursors of the shock, a problem we have discussed in Chapter 3. Whistlers propagate only outside their resonance cone. The limiting angle between $\mathbf{k}$ and the magnetic field $\mathbf{B}$ for which the whistler outruns the shock is given by $\theta_{\text{wh,lim}} \lesssim \cos^{-1}[\mathcal{M}_A(m_e/m_i)^{\frac{1}{2}}]$, artificially limiting the Alfvénic Mach number $\mathcal{M}_A = V/V_A < 43$.

In one-dimensional simulations with all quantities changing only along the shock normal $\mathbf{n}$ and the $\mathbf{k}$ -vectors of waves along $\mathbf{n}$ as well, one choses angles between $(\mathbf{k}, \mathbf{n})$ and $\mathbf{B}$ larger than this in order to have clean effects which are not polluted by those whistlers. However, the maximum phase speed of whistlers does not exceed the Alfvén speed by much (see Figure 3.10). Hence, as long as the upstream velocity is less than this maximum whistler speed, a standing whistler precursor will be attached to the shock in front. When the upstream velocity exceeds this velocity, phase standing whistlers become impossible. This happens at the critical whistler Mach number given in Chapter 3. The shock structure becomes more complicated then by forming shock substructures [Galeev et al, 1988] on scales of c/ω_{ce} , and the shock might become non-stationary [Krasnoselskikh et al, 2002].

One-Dimensional Structure

One-dimensional observations as those presented in Figure 5.4 confirm the theoretical prediction of the gross structure of a quasi-perpendicular shock. They can, however, when

taken by themselves, not resolve the spatial structure of the shock on smaller scales, nor do they allow to infer about the evolution of the shock. To achieve a clearer picture of both, the structure and the evolution, the observations must be supported by numerical simulations.

Such simulations have been performed in the past in various forms either as hybrid simulations or as full particle simulations. In hybrid simulations the electrons form a neutralising, massless background with no dynamics, i.e. the electrons react instantaneously while maintaining and merely adjusting their equilibrium Boltzmannian distribution to the locally changing conditions. Such simulations overestimate the role of the ions and neglect the dynamical contribution of the electrons. They nevertheless give a hint on the evolution and gross structure of a shock on the ion scales. Hybrid simulations have the natural advantage that they can be extended over relatively long times $\omega_{ci}t \gg 1$.

On the other hand full particle simulations are usually done for unrealistically small mass ratios $m_i/m_e \ll 1836$ much less than the real mass ratio. The electrons in these simulations are therefore heavy even under non-relativistic conditions. Their reaction is therefore unnaturally slow, the electron plasma and cyclotron frequencies are low, and the electron gyroradius, inertial length, and Debye length are unnaturally large. Under these conditions electrons readily become unmagnetised, non-adiabatic electron heating is prominent, and dispersive effects on ion-acoustic waves are overestimated.

Moreover, because of the large electron mass the electron thermal speed is reduced, and the Buneman two-stream instability sets on earlier and grows faster than under realistic conditions. This again should affect electron heating and structuring of the shock. On the other hand, the reduced electron gyroradius also reduces the shock potential, because the differences in ion and electron penetration-depths into the shock are smaller than in reality. This reduces the reflection capability of the shock, reduces the direct electric field heating of the impacting electrons, reduces the electron drift current in the shock ramp and shock transition and thus underestimates the dynamic processes in the shock, its structure, time dependence, formation and reformation and the strength of the foot effect and density of the foot population.

It is very difficult to separate all these effects, and comparison of different simulations is needed.

Low-Mass Ratio Simulations

Figure 5.5 shows an early one-dimensional low-mass-ratio perpendicular shock simulation with $m_i/m_e = 128$. Simulation times are short, not more than four ion-gyration times when energy conservation starts breaking down. Moreover, only a very small number of macro-particles (see Chapter 2) per simulation grid cell could be carried along in these simulation. Thus the noise in the simulations is large, not allowing for long simulation times, readily introducing diverging fake modes and fake dissipation/heating.

Nevertheless, the left-hand side of the figure shows the evolution of the magnetic field from the homogeneous state into a shock ramp and further the destruction and, what is known by now from much longer and better resolved simulation studies, the reformation of the shock profile. It should be noted that the shock in this case forms by reflection of the

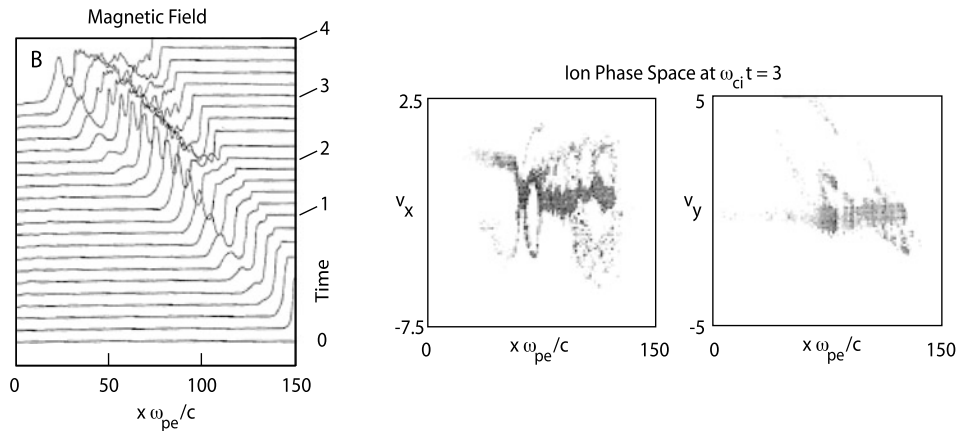


Figure 5.5: One-dimensional full particle-in-cell (PIC) simulations [after *Biskamp & Welter, 1972*] of shock formation assuming a mass ratio of $m_i/m_e = 128$. *Left:* Time-evolution of the magnetic field in stack-plot representation. Time is measured in units of ω_{ci}^{-1} , space in units of (heavy) electron inertial lengths c/ω_{pe} . The simulations are for a supercritical shock with $\mathcal{M} = 2.3$. Note the evolution of the magnetic field and the formation of a ramp, a foot and an overshoot. *Right:* Ion phase space plots (v_x, x) and (v_y, x) at time $\omega_{ci}t = 3$. Velocities are measured in units of the Alfvén velocity $V_A = B/\sqrt{\mu_0 N m_i}$.

fast initial supercritical flow with Mach number $\mathcal{M} = 2.3$ (which is above critical for the conditions of the simulation), entering the one-dimensional simulation ‘box’ from the left, from a ‘magnetic piston’ located at the right end of the box. This reflection causes a back-streaming ion-beam that interacts with the inflowing ions and drives an electromagnetic ion-ion instability which grows to large amplitude. The system is not current-free. In the interaction region of the two ion components the magnetic field forms a shock ramp. But after a short time of a fraction of an ion gyro-period a new ramp starts growing in the foot of the ramp, which itself evolves into a new ramp while the old ramp becomes eroded. This new ramp has not sufficient time to evolve to a full ramp as another new ramp starts growing in its foot. This causes the shock ramp to jump forward in space in steps from one ramp to the next, leaving behind a downstream compressed but fluctuating magnetic field region. The jump length is about the width of the foot region. It will become clear later why this is so.

Hybrid simulations [*Leroy et al, 1981, 1982; Leroy, 1984*] with fluid electrons and an artificially introduced anomalous resistivity show similar behaviour even though a number of differences have been found which are related to shock reformation. In particular shock reformation is slow or absent in hybrid simulations if not care is taken on the reaction of the electrons. The responsible instability in the foot region cannot evolve fast enough even though the hybrid simulation which take care of the ion dynamics also find reflection of ions and the evolution of a foot in front of the ramp. These differences must be attributed to the above mentioned lesser reliability of hybrid simulations than full particle codes.

Extended low-mass ratio full particle simulations in one space dimensions over a wide range of shock-normal angles $\Theta_{Bn} < 45^\circ$ have been performed by *Lembège & Dawson*

[1987a, b] with the purpose to study plasma heating. These simulations used mass ratios of $m_i/m_e = 100$ and a magnetic-piston generated shock. The simulations were completely collisionless, relatively small Mach number but nevertheless supercritical when taking into account the decrease in critical Mach number with Θ_{Bn} [Edmiston & Kennel, 1984]. They showed the formation of a foot and overshoot, the generation of an electric charge separation field in the shock transition from the foot across the shock with a highly structured electric field which was present already in the shock foot. Moreover, indications for a periodicity of the electric field structure in the foot region were given which we now understand as standing whistler wave precursors in the shock foot for oblique shock angles and supercritical but moderate Mach numbers [Kennel *et al*, 1985; Balikhin *et al*, 1995]. In addition to the field variations these simulations already demonstrated much of the supercritical particle dynamics related to shock reflection and foot formation which we will discuss separately below.

Before discussing the ion phase space plots in Figure 5.5 on the right we are going to describe recent investigations on the effects of the mass ratio dependence of the one-dimensional full particle simulations on the shock structure. Of course, in the end only such simulations can be believed which not only take into account the full mass ratio but which are long enough for following the evolution of the shock from a small disturbance up to a stage where the shock on some time scale has approached kind of a state that in a certain sense does not further evolve. This state is either stationary or it repeats and restores itself such that it is possible to speak at all of a quasi-perpendicular supercritical shock.

The Shock Transition Scale

Determination of the shock foot scale is relatively easy both in simulations as also from observation. From observations, as already mentioned, it has been first determined by Sckopke *et al* [1983] who found that the foot scale is slightly less but close ($\sim 0.7 r_{ci, \text{refl}}$) to the reflected ion gyroradius in quasi-perpendicular shocks. The reasons for this number have been given by Schwartz *et al* [1983] and are related to the reflected ions coupling to the upstream convection electric field in which they are accelerated. This can also be checked in simulations. Of more interest is the determination of the shock transition, i.e. the width of the shock ramp which from theory is not well determined since it depends on several factors which can hardly be taken into account at once.

The width of the shock transition is particularly important in its relation to the width of the electrostatic potential drop across the shock. There are essentially three transition scales: the magnetic scale Δ_B , the density scale Δ_N , and the electric potential scale Δ_E . Since the shock is not in pressure equilibrium, the first two scales must not necessarily be proportional to each other. However, the electric field and density gradient might be related, so one expects that $\Delta_N \sim \Delta_E$ even though this is not necessarily so, in particular not when instabilities arise which cause very small scale electric field gradients. In principle one can distinguish three different cases [Lembège *et al*, 1999] which describe different physics:

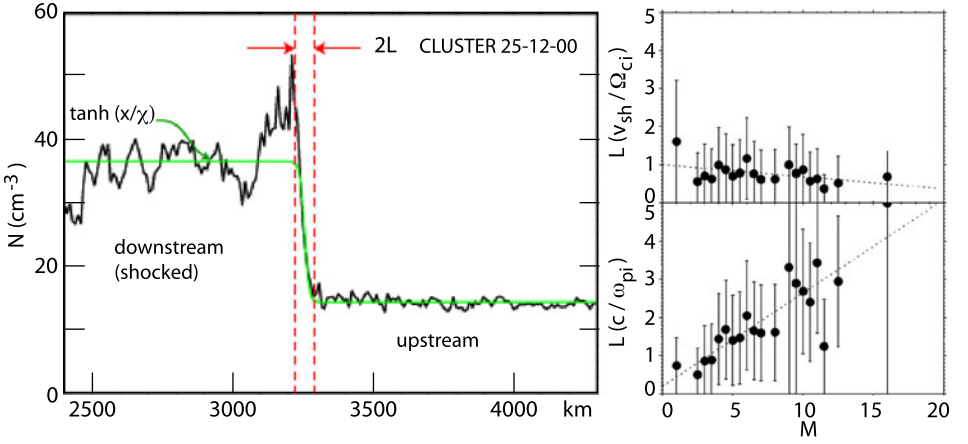


Figure 5.6: *Left:* Shock density transition-fit by a tanh-function in order to determine the shock-ramp transition scale [after *Bale et al, 2003*]. 98 of those shock transitions have been used in order to find a dependence of the shock ramp width from some physical parameter. *Right:* The dependence inferred by *Bale et al [2003]*. The *upper part* of the figure scales the dependence of the gyroradius with Mach number, the *lower part* the dependence of the ion inertial scale. Apparently there is no dependence of gyroradius on Mach number, while there is a clear linear dependence of the inertial scale on Mach number.

1. $\Delta_E \gg \Delta_B$. This is a case that has been reported to have been observed in Bow shock crossings [*Scudder et al, 1986; Scudder, 1995*]. The magnetic ramp is much steeper in this case than the structure of the electric field. The latter smears out over the foot and ramp regions. In this case the electrons will behave adiabatically, while the ions may be only partially or even non-magnetised.
2. $\Delta_E \sim \Delta_B$. In this case there will be a significant deviation from adiabatic behaviour of the electrons in the shock transition. Electron heating and motion will not be adiabatic anymore, and the electron distribution will significantly be disturbed [see, cf., *Balikhin et al, 1995*]. Observations of such cases have been reported [*Formisano & Torbert, 1982*].
3. $\Delta_E \ll \Delta_B$. This case which is also called the ‘isomagnetic’ transition [*Eselevich, 1982; Kennel et al, 1985*] corresponds to shock transitions with electrostatic substructuring which are sometimes also called subshocks.

The most recent *experimental* determination of the density transition scale has been provided by *Bale et al [2003]* using data from 98 Bow Shock crossings by the Cluster spacecraft quartet. The result is shown in Figure 5.6 for an example of this fit-determination by fitting a tanh-profile to the shock density transition. The point is that these authors found a dependence of the shock ramp transition on Mach number when the transition is scaled in ion inertial units, while there is no dependence when scaled in ion gyroradii. This scaling suggests that the shock scales with the gyroradius, since $(V/c)(\omega_{pi}/\omega_{ci}) \sim \mathcal{M}_A$.

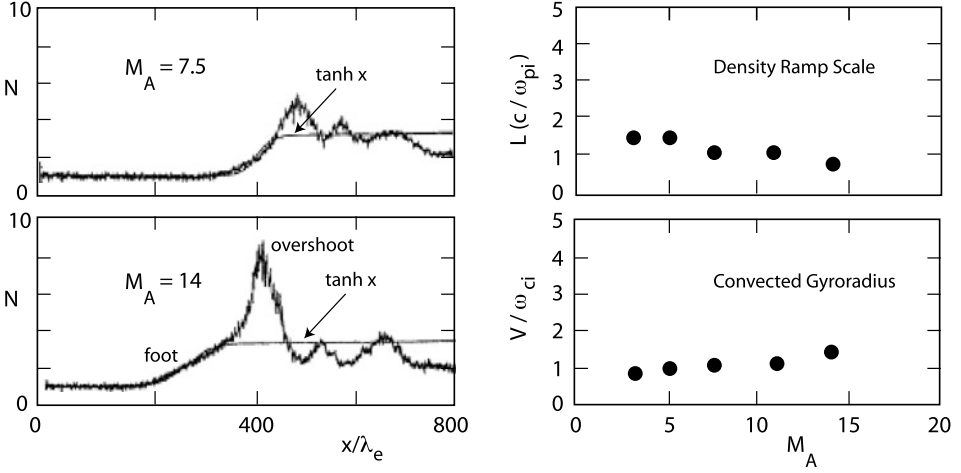


Figure 5.7: *Left:* Density profile in two full PIC simulations of large Mach numbers. Indicated is the pronounced overshoot and the long extended foot. The *straight lines* are $\tanh x$ -fits to the simulations showing the neglect of the overshoot and ramp during such fits which only account for the foot region. From fitting the ramp with the curves on the right are obtained. *Right:* Density ramp scales and convected ion gyroradii (in units of upstream inertial length) obtained in one-dimensional full particle PIC simulations of quasi-perpendicular shocks [after Scholer & Burgess, 2006] as function of Alfvénic Mach number. Use has been made of the full particle mass ratio 1838, $\Theta_{Bn} = 87^\circ$, and $\omega_{pe}/\omega_{ce} = 4$. The magnetic field used is that of the overshoot. One observes that the ratio of ion gyroradius to ion inertial length is constant. Also the scale of the ramp is about $\sim 1c/\omega_{pi}$, supporting a narrow ramp. The simulations also show that the scale of the ramp sharpens with increasing Mach number.

In order to check this behaviour numerically, Scholer & Burgess [2006] performed a series of one-dimensional full-particle PIC simulations with the correct mass ratio $m_i/m_e = 1838$ and for the Alfvénic Mach number range $3.2 \leq M_A \leq 14$ and a shock normal angle $\Theta_{Bn} = 87^\circ$ in order to have a component of $k_{\parallel}$ parallel to $\mathbf{B}$, but with small ratio $\omega_{pe}/\omega_{ce} = 4$ to compromise computing requirements. Figure 5.7 shows the results of these simulations. A $\tanh x$ -fit neglects in fact the entire ramp and takes account only of the foot region. Correcting the above described measurements it is thus found that the ramp thickness is just of the order of $\sim 1\lambda_i = c/\omega_{pi}$ and decreases slightly with increasing Mach number. However, from the form of the density profile it seems clear that the shock ramp is basically determined by the overshoot, and one must take the overshoot magnetic field value in calculating the gyroradius. The convected gyroradius based on the overshoot magnetic field B_{ov} and measured in λ_i is about constant very close to unity. Thus the shock ramp scale is given by the convective ion gyroradius based on the overshoot magnetic field. One should, however, note that the computing power in the simulations does not yet allow for larger ratios ω_{pe}/ω_{ce} which may affect the result. Moreover, higher dimensional simulations would be required to confirm the general validity of those calculation and conclusions.

Hence, combining the observations of Bale *et al* [2003] and the results of the simulation studies of Scholer & Burgess [2006] we may conclude that the scale of the shock foot

is given by the upstream-convected ion gyroradius $r_{ci} = V/\omega_{ci,1}$ based on the upstream field B_1 , while the scale of the shock ramp is given by the ramp-convected ion gyroradius $r_{ci,ov} = V/\omega_{ci,ov}$ based on the value of the magnetic field B_{ov} overshoot. This is an important difference which can be taken as a *golden rule* for estimates of the structure of quasi-perpendicular shocks even though, of course, these values are dynamical values which change from position to position across the foot and ramp. The scale differences are the reasons for the large upstream extension of the foot and the relative steepness of the shock ramp.

The observed constancy of the overshoot magnetic field-based convective ion gyroradius $r_{ci} \propto V/B_{ov}$ with Mach number $\mathcal{M}_A \propto V$ can be understood when considering the about linear increase of the overshoot magnetic field $B_{ov} \propto \mathcal{M}$ with Mach number (or with upstream velocity V) which holds for supercritical Mach numbers $\mathcal{M} > \mathcal{M}_{crit}$ as long as $\mathcal{M}$ is not too large. At very large – but still non-relativistic – Mach numbers $\mathcal{M} < \mathcal{M}_{max}$ the increasing steepness of the shock ramp and the increasing extension of the foot ultimately lead to the excitation of smaller scale structures in the ramp and the foot, which smear out any further increase in the overshoot.

The generation of these structures by a variety of instabilities might even turn the shock foot and ramp regions into regions where large anomalous collisions and thus resistances are generated as the result of wave-particle interactions. In this case the shock returns to become resistive again due to preventing large numbers of reflected ions from passing across the steep shock ramp and large shock potential, using the kinetic energy of the reflected particle population for the generation of a broad wave spectrum which acts to scatter the particles around in the foot and ramp regions and, possibly, also up to some distance in the transition region behind the ramp. This kind of confinement of reflected particles over long times will then be sufficiently long for providing the heating and dissipation which is required for sustaining a resistive shock which, then, is the result of the combined action of ion viscosity and anomalous resistivity, i.e. anomalous collisions. In addition, the scattering of the trapped reflected particle population necessarily results in plasma heating, and some particles will become accelerated to high velocities in these interactions as well. It is then possible that these particles provide the seed population for energetic particles which have been accelerated to high energies in the well-known shock-Fermi-one and shock-Fermi-two acceleration mechanisms.

So far the range of Mach numbers $\mathcal{M}_{max} < \mathcal{M} < \mathcal{M}_{rel}$ where this will happen is unknown, as it is hardly accessible to numerical simulations. However, the available simulations seem to point in this direction as long as the Mach numbers remain non-relativistic. In relativistic shocks with $\mathcal{M} = \mathcal{M}_{rel}$ different effects arise which are not subject to our discussion at this place.

5.2.3 Shock Reformation

It has already been mentioned several times that supercritical shocks do under certain conditions reform themselves periodically – or quasi-periodically –, which is kind of a non-stationarity of the shock that does not destroy the shock but, at the contrary, keeps it intact in a temporarily changing way. We will come later to the problem of real non-stationarity.

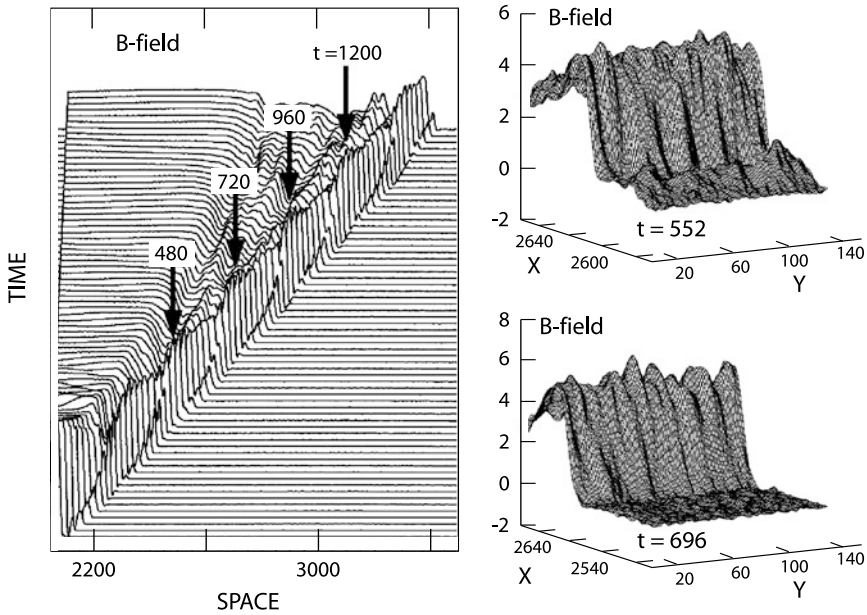


Figure 5.8: Magnetic field from full particle PIC simulations of shock reformation [after *Lembège & Savoini, 2002*, courtesy American Geophysical Union]. *Left:* Reformation cycles of the magnetic field in the shock. Time is measured in inverse electron plasma frequencies ω_{pe}^{-1} . The reformation times are indicated by the *arrows* in the plot with time given when the cycle is complete. *Right:* Two snapshots in time of the view of the shock front in the magnetic field at reformation. The interesting finding is that the front in this two-dimensional view is not a smooth plane but is quite distinctly structured in space and at the same time evolving.

Reformation in One Dimension: Mass Ratio Dependence

Reformation of quasi-perpendicular shocks is thus an important shock property which is closely related to highly super-critical shocks and the formation of a foot region, i.e. to the reflection of ions from the shock ramp.

In fact, reformation was already observed by *Biskamp & Welter [1972]* in the early short-simulation time PIC simulations shown in Figure 5.5, where we have noted it explicitly. Reformation of quasi-perpendicular shocks has also been reported, for instance, by *Lembège & Dawson [1987a]*, *Lembège & Savoini [1992, 2002]*, *Hellinger et al [2002]* and others who all used small ion-to-electron mass ratios. For illustration, Figure 5.8, on its left, shows a low mass-ratio example of the temporary evolution of a shock during shock reformation in a magnetic field stack plot. On its right the structure of the shock ramp at two different two reformation times is plotted along the shock ramp. There is a distinct reformation cycle in this simulation and also a distinct structure of the ramp/shock front in the tangential direction which is far from being smooth, a fact to which we will return during discussion of non-stationarity of shocks. The shock not only reforms cyclically in time, it also develops ripples along its surface which travel like waves along the shock.

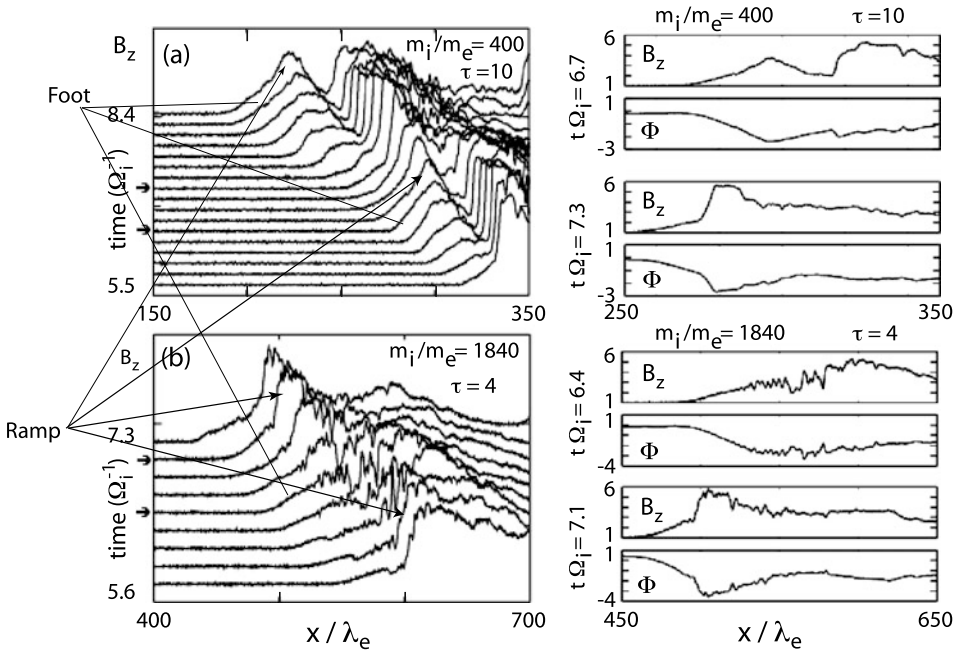


Figure 5.9: *Left:* One-dimensional PIC simulations [after Scholer *et al.*, 2003, courtesy American Geophysical Union] of quasi-perpendicular $\Theta_{Bn} = 87^\circ$ shock reformation for mass ratios $m_i/m_e = 400$ and 1840. Time is in ω_{ci}^{-1} (here denoted as Ω_i^{-1}), space in $\lambda_e = c/\omega_{pe}$. The parameter $\tau = \omega_{pe}^2/\omega_{ce}^2$ is taken small in both cases. The higher mass ratio shows a violent time evolution because of the high electron mobility. Reformation of the shock is due to evolution of the shock feet. The original foot region builds up until becoming itself the shock assuming the role of the ramp. Afterwards a secondary foot evolves in front of this new ramp. *Right:* Spatial shock profiles at two time sections (see arrows on the left). The higher mass ratio run shows a more subtle structure in B_z and shock potential Φ , but the gross features are similar. The potential drop exists already in the foot but the main drop occurs in the ramp. The lower mass ratio has a more concentrated foot region.

Full particle electromagnetic PIC simulations with realistically large mass ratios have been performed only very recently [Matsukiyo & Scholer, 2003; Scholer *et al.*, 2003; Scholer & Matsukiyo, 2004] and only in one spatial dimension, showing that reformation at least occurs at small ion- $\beta_i \sim 0.2$. In these simulations the shock is produced by injecting a uniform plasma from $-x$ and letting it reflect from a stationary wall at the right end of the simulation box. The plasma carries a uniform magnetic field in the (x, z) -plane, and the plasma is continuously injected in the $+x$ -direction. Since the right-hand reflecting boundary is stationary the shock, which is generated via the ion-ion beam instability in the interaction of the incoming and reflected ion beams, moves to the left at velocity given by the supercritical shock Mach number $\mathcal{M}_A \sim 4.5$. The upstream plasma has $\beta_i = \beta_e = 0.05$, and the shock normal angle is $\Theta_{Bn} = 87^\circ$.

Two runs of these simulations are shown in Figure 5.9, one is for a mass ratio of 400, the other for a mass ratio of 1840. The left-hand side of the figure shows stack plots of time profiles of the nearly perpendicular magnetic field B_z with time running in equidistant

units upward on the ordinate. Since the plasma is injected from the left and reflected at the right boundary the shock is seen to move from the right to the left in this pseudo-three-dimensional representation. Time is measured in ion cyclotron periods ω_{ci}^{-1} , while space on the abscissa is in units of the electron inertial length c/ω_{pe} . The magnetic profiles are strikingly similar for both mass ratios. In both cases a relatively flat foot develops in front of the steeper shock ramp caused by the shock reflected ions. The magnetic field of this foot starts itself increasing with time with growth being strongest close to the upstream edge of the foot until the foot field becomes so strong that it replaces the former shock ramp and itself becomes the new and displaced shock ramp.

This is seen most clearly in the upper low mass-ratio part of the figure. The foot takes over, steepens and becomes itself the shock. One can recognise in addition that, even earlier, the intense foot already had started reflecting ions by himself and developing its own flat pre-foot region. This pre-foot evolves readily to become the next foot, while the old ramps become part of the downstream turbulence.

During this reformation process the shock progresses upstream from right to left. This progression is not a continuous motion at constant speed. Both the foot and the ramp jump forward in steps. One such step ahead is seen, for instance, at time $t\omega_{ci} = 7.6$. Sitting in the shock frame one would experience some forward acceleration at this time, seeing the ramp moving downstream as a magnetic wave front the apparent source of which is the instantaneous shock ramp, while it is just the old shock foot. Hence the shock ramp and shock overshoot act as a source of a pulsating magnetic waves that are injected downstream from the shock with periodicity of roughly $\Delta t \sim 1.8\omega_{ci}^{-1}$ (for $m_i/m_e = 400$) and add to the downstream turbulence.

The realistic mass-ratio run in the lower part on the left also shows reformation of the shock. However there are some differences. First, the magnetic profiles are much stronger disturbed exhibiting much more structuring. Second, the foot region is considerably more extended in upstream direction. Third, the ramp is much steeper, and reformation is faster, happening on a time scale of $\Delta t \sim 1.3\omega_{ci}^{-1}$, roughly 30% faster than in the above case. Reformation is, however, more irregular at the realistic mass ratio with the property of reforming the shock ramp out of a long extended relatively smooth shock foot which exhibits pronounced oscillations.

The right-hand side of the figure shows two shock profiles at constant times for the two different mass-ratio simulations. The first profile at $t\omega_{ci} = 6.7$ has been taken when a well developed foot and ramp had been formed on the shock, the second profile at $t\omega_{ci} = 7.3$ is at the start of the new foot towards the end of the simulations. At the low mass ratio the foot profile is quite smooth showing that the foot is produced by the accumulation of reflected ions at the upstream edge of the foot where the ions have the largest velocity in direction y along the shock. This is where, during their upstream gyration in the upstream magnetic field, they orbit about parallel to the upstream convection electric field and gain most energy.

Hence, here, the current density is largest due to the accumulation of the reflected ions, due to the retardation of some ions from the inflow already at this place, and due to the speeding up of the reflected and retarded ions in y -direction by the convection electric field

E_y . All this leads to a maximum in the current density j_y and thus causes a maximum in the magnetic field B_z close to the upstream edge of the foot.

Most interestingly, the electric potential exhibits its strongest drop right here in the foot region with a second but smaller drop in the ramp itself. It is the electric field that belongs to this potential drop that retards the inflow already before it reaches the shock ramp. At the contrary, when the shock ramp is well developed, the main potential drop is for short time right at the ramp and extends even relatively far into the downstream region.

For the realistic mass ratio the foot- and ramp-transitions are both highly structured at $t\omega_{ci} = 6.4$ exhibiting fluctuations in both the magnetic field and electric potential, but the electric potential drop extends all over the foot region with nearly no drop in the ramp. When the ramp has been reformed at $t\omega_{ci} = 7.1$, the foot region still maintains a substantial potential drop, but 50% of the total drop is now found in the ramp with the downstream potential recovering. This is interesting as it implies that lower energy electrons will become trapped in the overshoot region, an effect which is much stronger for the large mass-ratio than for small mass-ratios and thus closer to reality.

Some recent one-dimensional full particle PIC simulations by *Umeda & Yamazaki [2006]* at Mach number $\mathcal{M}_A = 10$ and medium mass ratio $m_i/m_e = 100$ throw additional light on the reformation process when keeping in mind that reformation is not as strongly dependent on the mass ratio as originally believed. Figure 5.10 shows a collection of their results which this time are represented in the shock frame of reference.

The simulations have been performed by assuming an initial Rankine-Hugoniot equilibrium in the PIC code. The non-physicality of this initialisation is manifested in the initial evolution over the first few ion cyclotron periods. During this time the simulation adjusts itself to the correct physics, and the non-physical disturbance decays. The shock frame has shifted by this to a new position, which in the shock frame is located farther downstream (which takes into account of the momentum transferred to the shock by the reflection of the upstream ions who lower the shock speed).

The further evolution of the shock shows the quasi-periodic reformation and the play between the foot and the ramp formation. The periodicity is roughly $\sim 10\omega_{ci,2}^{-1}$. When the foot takes over to become the ramp, the ramp jumps ahead in a fraction of this time. Afterwards the formation of the foot retards the ramp motion, and the ramp softens and displaces itself downstream to become a downstream moving spectrum of magnetic oscillations which is injected into the downstream region in the form of wave packets. The various plasma parameters in the left part of the figure show in addition the compression of plasma and field, and the dominance of perpendicular ion heating which is, of course, due to the accelerated foot ions which pass into downstream.

Two-Dimensional Reformation: Whistlers and Mach Number Dependence

First two-dimensional simulations of a strictly perpendicular $\Theta_{Bn} = 90^\circ$ shock formation have recently been performed by *Hellinger et al [2007]*. These simulations were intended to study the reformation process in two dimensions when the perpendicular shock is supercritical. Since PIC simulations are very computer-time consuming, most of the simulation runs by *Hellinger et al [2007]* used a two-dimensional hybrid code with the shock being

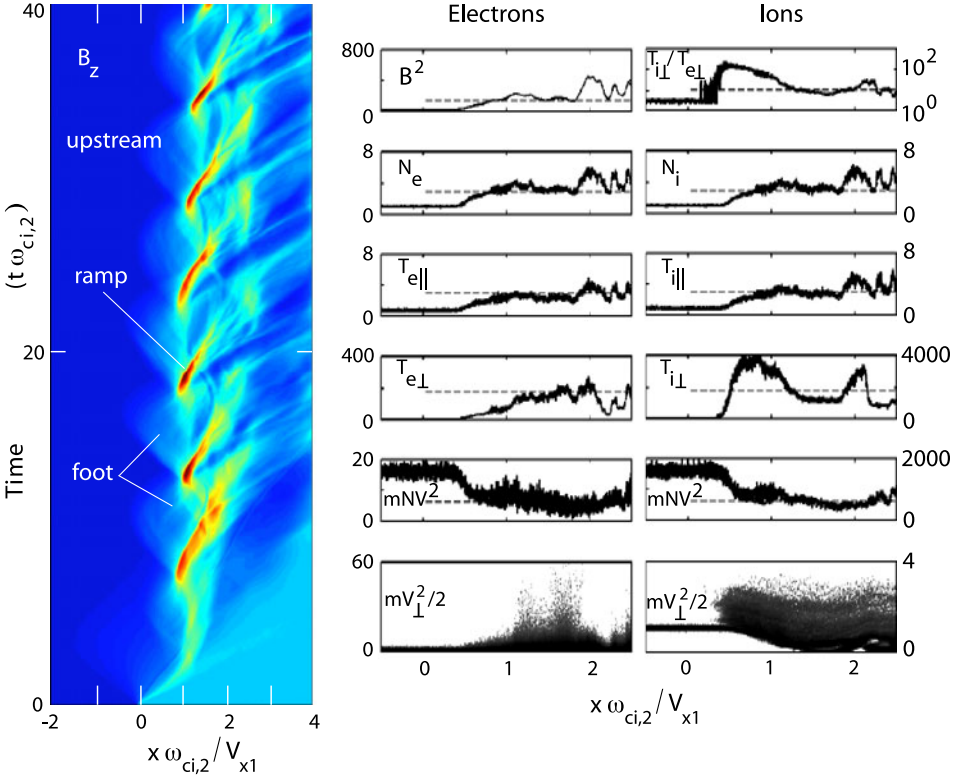


Figure 5.10: *Left:* Evolution of the magnetic field in a quasi-perpendicular high Mach number $\mathcal{M}_A = 10$ PIC simulation [after Umeda & Yamazaki, 2006]. Here the presentation is in the shock frame of reference, and the shock has been initialised by assuming Rankine-Hugoniot initial jump conditions. The non-physical nature of this assumption is visible in the initial evolution and fast displacement of the shock to the right. After the initial unphysical disturbance has disappeared a self-consistent physical state is reached in which the shock quasi-periodically reforms itself. The competition between the shock foot and ramp formation is nicely seen in the colour plot of the magnetic field B_z . *Right:* Electron and ion plasma parameters in computational units. Of interest is only their relative behaviour, not the absolute values. The profiles are taken at time $t \omega_{ci,2} = 38.1$. They show the compression of the plasma and heating of electrons and ions. Parallel electron and ion heating is comparable, but ions are heating much stronger than electrons in perpendicular direction causing a large perpendicular temperature anisotropy downstream of the shock.

generated by a magnetic piston as in the case of the simulations by Lembège & Dawson [1987a]. The interesting result of this simulation study was that no shock reformation was found while phase locked whistlers were detected which formed a characteristic interference pattern in the shock foot regions. This result is surprising as for strictly perpendicular shocks no whistlers should be generated according to the one-dimensional theory [see the above discussion on whistlers and, e.g., Kennel et al, 1985; Balikhin et al, 1995].

In order to cross check their hybrid simulation results Hellinger et al [2007] also performed a two-dimensional PIC [Lembège & Savoini, 1992] simulation choosing a mass

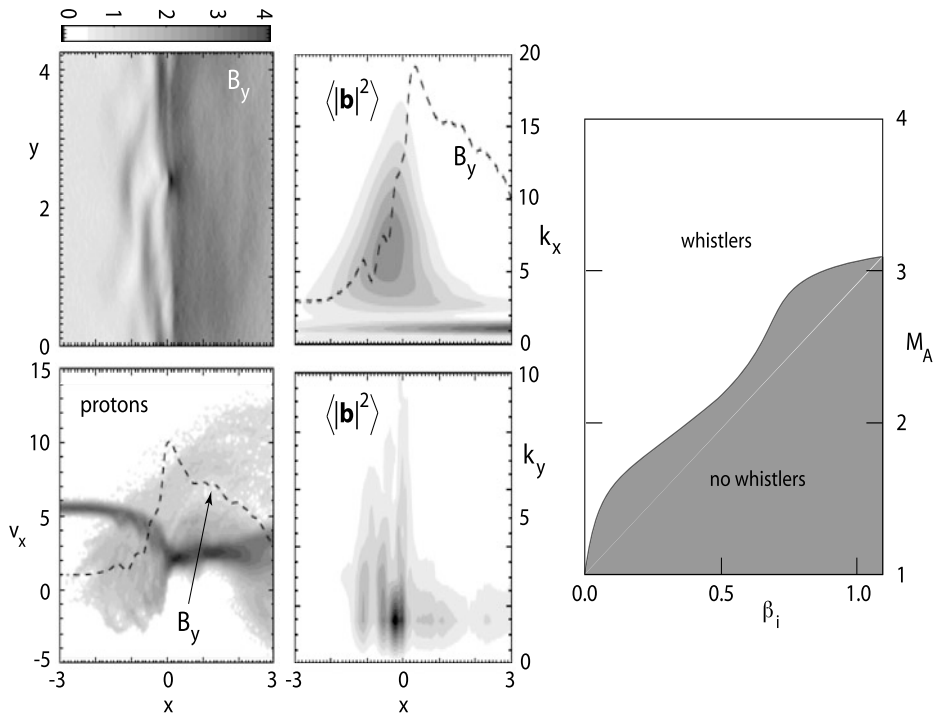


Figure 5.11: *Left:* Two-dimensional PIC simulations [after Hellinger et al., 2007, courtesy American Geophysical Union] of the end time $\omega_{cit} = 28$ of the evolution of a *strictly perpendicular* shock using $m_i/m_e = 400$. Shown is the magnetic structure in the (x, y) -plane, the proton phase space (x, v_x) and the power of magnetic fluctuations in dependence on space x and wave numbers (k_y, k_x) . Lengths are measured in ion inertial lengths c/ω_{pi} , velocities in Alfvén speeds V_A , wave numbers in inverse ion inertial lengths ω_{pi}/c . Magnetic fields and powers are in relative units (see grey scale bar). No shock reformation is seen in the *upper panel* of B_y on the left. A periodic foot evolves periodically causing a higher and steeper ramp overshoot when its cycle ends, but the shock ramp does not become exchanged with a new ramp. Note also that the next foot cycle begins before the end of the former cycle, i.e. the shock foot itself reflects ions. The power spectra show a periodic spatial spectrum of whistlers standing in and restricted to the shock foot. Periodicity in k_y is caused by interference between outward and inward moving whistlers. The proton phase space shows the retardation of the incoming flow in the shock foot, the occurrence of reflected ions in the foot and the heating of foot ions. Forward heating is also seen in the overshoot. *Right:* A parametric (2D-hybrid simulation) investigation of the evolution of phase locked whistlers in the shock foot in dependence on Mach number M_A and β_i . Large Mach numbers and small β_i support the excitation of standing whistlers.

ratio $m_i/m_e = 400$, Mach number $M_A = 5.5$, electron plasma-to-cyclotron frequency ratio $\omega_{pe}/\omega_{ce} = 2$, upstream $\beta_e = 0.24$, $\beta_i = 0.15$ and 4 particles per cell. The results of this simulation have been compiled in Figure 5.11, showing only the PIC simulations and no hybrid simulations. Deviating from the former one-dimensional simulations by Matsukiyo & Scholer [2003] the magnetic field is in y . The block consisting of the four panels on the left in the figure are the simulation results at the end of the simulation run, showing the compression of the magnetic field B_y in the (x, y) -plane, proton velocity space (v_x, x) – only

the normal velocity component is shown –, and the average magnetic fluctuation spectra $\langle b^2 \rangle$ as functions of wave number components k_x, k_y . As the authors describe, after a short initial time when the shock foot forms and the shock reforms, reformation stops and does not recover again in these two dimensional run. Instead, the shock foot starts exhibiting large magnetic fluctuations. These are seen in the low-frequency magnetic power spectra being confined solely to the shock foot (as recognised from the B_y -profile) as seen from the k_x dependence of the magnetic power spectrum and forming an interference pattern in k_y .

These fluctuations are identified as whistler waves propagating obliquely (in k_x and k_y) across the foot and the magnetic field. Since $k_x \sim 3k_y$ their perpendicular wave numbers are large, they are quite oblique, and their parallel wavelengths are long. They are excited in the foot and because of their obliqueness probably propagate close to the resonance cone. Their main effect is to resonantly suppress shock reformation by inhibiting the ions to accumulate in the foot. Hence, under the conditions of these simulations the shock turns out to be stable and does not reform. It maintains its structure thanks to the generation of oblique whistlers in the shock foot which dissipate so much energy that the shock becomes about resistive. In one-dimensional simulations this regime has not been seen and is probably inhibited for strictly perpendicular shocks. In two-dimensional simulations, on the other hand, the additional degree of freedom provided by the introduction of the second spatial dimension allows for the generation of the whistlers which are suppressed in the one-dimensional case (where $\mathbf{k}$ has only the component k_x).

Guided by these simulations *Hellinger et al* [2007] have undertaken a parametric study of the regime where whistler excitation and thus presumably stationary shock structures lacking reformation should exist. Their results are given on the right in Figure 5.11 in $(\beta_i, \mathcal{M}_A)$ -space. According to this figure, whistlers will not be excited at low Mach numbers. Here the two-dimensional perpendicular shocks will reform. At higher Mach numbers whistlers should be excited, and the shock should become stabilised in two dimensions.

This surprising result suggests that sufficiently high Mach numbers are needed in order to excite whistlers; on the other hand, when the Mach number will become large (*Hellinger et al* [2007] investigated only the range of Mach numbers < 5) then other effects should set on, and the shock should become non-stationary with whistlers becoming unimportant and reformation becoming possible again.

5.3 Ion Dynamics

Until now we have avoided discussing the behaviour of particles in the simulations. The mere idealised reflection process we have already discussed, as far as this could be done analytically. The complicated geometry and dynamics of particle motion in shocks necessitates to return to simulations.

We have mentioned that reflected particles are forming a foot on the shock and may contribute to the reformation of the shock. We have, however, not yet gone into detail and into the investigation of the relation of the particle distributions in phase and real space

observed in the simulations and their relation to the shock dynamics. This will be done in the present section for the ions on the basis of simulation studies.

5.3.1 Ion Dynamics in Shock Reformation

The two plots on the right-hand side of Figure 5.5 show phase space representations of ions in the early simulations of shock formation by *Biskamp & Welter* [1972] at $t\omega_{ci} = 3$ close to the end of the simulation. The first box is the dependence of the ion velocity v_x on position x (direction of inflow), the second plot the transverse ion velocity component v_y . Each macroparticle present in the box is identified by a dot. Hence high particle phase space density is reflected by accumulation of many such dots. The plasma enters the central region of the box from left at high positive speed v_x and $v_y = 0$ and leaves the central region to the right at high density and very low speed v_x . The inflowing plasma was extended relatively narrow in v_x indicating low temperature. The shocked plasma distribution is broader thus having higher temperature. Moreover, some of the ions are seen to be smeared out from large positive up to large negative velocities in v_x indicating that strong particle scattering has occurred and that there are particles of high velocity in both directions being equivalent to heated plasma. In the foot region the closed circle in the particle distribution signifies the presence of the reflected trapped ions which are accelerated in both direction. their presence decelerates the inflow as can be seen by the drop in the velocity v_x just before entering the closed loop trapped ion region. The reflected particles are seen on the left of this closed loop having negative speeds in v_x therefore being directed upstream. The (v_y, x) -box shows that the reflected and trapped ions gyrate and are accelerated into y -direction. Also many of the ions downstream of the shock possess non-zero v_y velocity components, some of the large, indicating acceleration along the shock front.

Even though it seems clear from this figure that reflected ions are involved into shock reformation and the dynamics of the shock profile, the process remained to be unclear until simulations at higher resolution and much higher mass ratio had been performed. Figure 5.9 gave a clearer idea of the ion dynamics. The corresponding ion phase space plots are shown in Figure 5.12 for the two mass ratios $m_i/m_e = 400$ (top) and 1840 (bottom) respectively. Both plots show just the enlarged shock foot transition region over the same scale of $100c/\omega_{pe}$. The electron $\beta_e = 0.2$ has been kept constant in both simulations, while the ion β_i has been changed. Only the normal component of the ion velocity is shown for the nearly perpendicular supercritical shock. In both plots the magnetic field B_z has been drawn as a thin continuous line showing the magnetic shock profile over the spatial distance Δx .

The upper (low-mass-ratio) low- β_i panel shows the cold dense ion inflow at velocity $v_{ix} \sim 5$ (in units of the upstream Alfvén velocity) being retarded to nearly Mach number 1 already when entering the foot. This retardation is due to its interaction with the intense but cold (narrow in velocity space) reflected ion beam which is seen as the narrow negative v_{xi} -velocity beam originating from the shock ramp. This reflected ion beam needs a certain distance to interact with the upstream plasma inflow. This distance is the length the beam-beam excited waves need to grow. But once the interaction becomes strong enough, the

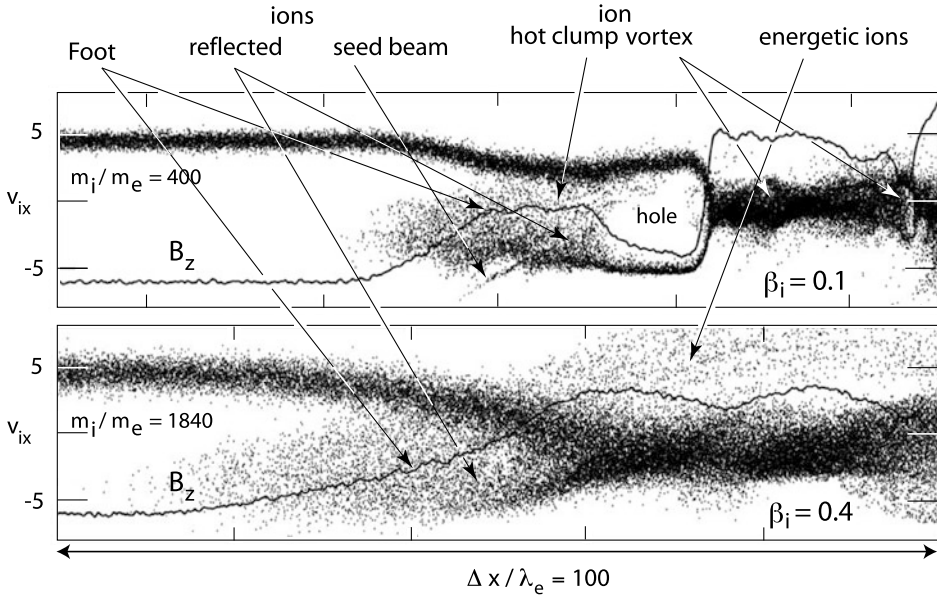


Figure 5.12: Ion phase space for two simulations of supercritical quasi-perpendicular shock formations for same Mach number and $\beta_e = 0.2$ but different mass ratios and β_i [after Scholer *et al.*, 2003, courtesy American Geophysical Union]. *Top:* Foot formation and shock reformation in a low- β_i shock simulation. The reflected ion beam accumulates in the shock foot forming a hot clump-vortex ion distribution which causes scattering of the reflected and upstream ions and generates a high foot magnetic field. The hole of the vortex corresponds to low magnetic field, the ring to large magnetic fields. When the foot field further increases it takes over and becomes the new ramp. At this time the ramp position jumps from the current position to that of the foot edge thereby reforming the shock. Such former reformation cycles can be recognised in the downstream distribution as remains of ion vortices (holes) in the otherwise hot downstream distribution. *Bottom:* The same simulation with realistic mass ratio but large $\beta_i = 0.4$. No reformation and no ion vortex is observed. Instead, the foot ions because of their high temperature smear out the entire gap between the original inflow and the reflected ion beams. Note also the occurrence of a large number of downstream diffuse energetic ions in this case.

reflected ions are scattered by the waves into a hot ion clump in addition to being turned around by gyration. Both effects cause a reduction in velocity v_x of the reflected ions which, being accelerated by the convection electric field, turn to flow in y direction and cause the magnetic bump that develops in this region of the foot. It is interesting to remark that in the (v_x, x) -plane the reflected ions close with the upstream flow into an hot ion ring distribution (vortex) just in front of the ramp of which the hot ion clump that brakes the inflow is the upstream boundary. Behind the ramp, which is the point of bifurcation of the ion distribution, i.e. the location where the reflection is at work, a broad hot ion distribution arises which at some locations shows rudimentary remains of ion vortices from former reformation cycles. Their magnetic signatures are the dips seen in the magnetic field. The next reformation cycle can be expected to completely close the ion vortex in the foot and to transform the ramp from its current position to the position of the foot. The first sign of this process is already seen in the foot ion distribution, which shows the birth of a faint new

reflected ion beam at high negative speeds. This beam is not participating in the formation of the ring but serves as the seed of the newly reflected population.

The same behaviour is found in large mass ratio simulations as long as β_i is small. This is obvious from the large mass-ratio magnetic field shown in Figure 5.9. As long as β_i remains to be small, the shock undergoes reformations also for realistic mass ratios. In other words, as long as the plasma is relatively cool the real shocks found in nature should develop feet which at a later time quasi-periodically become the shock ramp.

This changes completely, when β_i increases as is suggested from consideration of the lower panel in Figure 5.12. There a realistic mass ratio has been assumed, but $\beta_i = 0.4$, and no reformation is observed, at least not during the realisable simulation times. Instead, the shock develops a very long foot region that is extended twice as far into the upstream region as in the low- β_i case. Clearly, the high ion temperature smears out the reflected ion population over the entire gap region between the upstream and reflected beam regions, and no vortex can develop. This implies that the foot remains smooth and does not evolve into a ramp.

However, inspecting the panels of Figure 5.12 it becomes immediately clear that suppression of reformation is a relative process. Reformation will be suppressed only when the thermal speed v_i of the ions is large enough to bridge the gap between the reflected and incoming ion beams, i.e. large enough to fill the hole. Semi-empirically one can establish a condition for shock reformation as $v_i < \alpha V_{n1}$ when taking into account that the normal speed of incoming ions is simply specularly turned negative. Since this is never exactly the case, the coefficient will roughly be in the interval $1.5 < \alpha < 2$. This condition for reformation to occur can be written as

$$\mathcal{M}_A > \frac{\beta_i^{\frac{1}{2}}}{\alpha} \quad (5.20)$$

where the Alfvénic Mach number is defined on V_{1n} . The larger the Mach number becomes the less suppression of reformation will play a role, and at high Mach numbers one expects that either reformation becomes a normal process or that other time-dependent processes set on which lead to a non-stationary state where the shock reformation becomes a chaotic and unpredictable process. As we have argued earlier this is quite normal as the shock is thermodynamically and thermally in a non-equilibrium state: it is a region where electrons and ions have violently different temperatures; it is not in pressure equilibrium; upstream and downstream temperatures are different; and it hosts a number of non-Boltzmannian phase space distributions all concentrated in a small volume of real space. Under such conditions stationary states suppressing reformation [advocated e.g. by *Hellinger et al, 2007*] will occur only very exceptionally.

Hada et al [2003] recently attempted to semi-empirically determine the fraction of reflected ions needed for reformation to occur. They performed a large parametric search based on hybrid simulations by changing the Mach number and determined the fraction of reflected ions in the foot when reformation occurred. Their result is shown in Figure 5.13 for cold electrons $\beta_e = 0$ (assuming that electrons do not contribute to reformation and reformation being exclusively due to ion-viscosity, and for fixed thermal velocity of upstream and reflected ions $v_i/\omega_{pe}d = 0.2$, where d is the spatial simulation-grid spacing). The latter

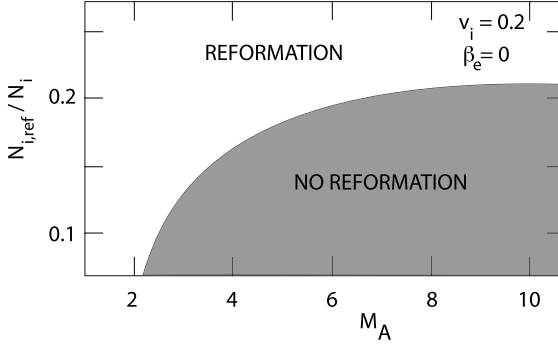


Figure 5.13: Parametric determination of the fraction of shock-reflected ions in the foot of a quasi-perpendicular shock as function of Alfvénic Mach number for the special case of $\beta_e = 0$ and ion thermal velocities $v_i = 0.2$ measured in $\omega_{pe}d$, where d is the numerical grid spacing [after *Hada et al, 2003*, courtesy American Geophysical Union].

restriction is certainly not satisfied as the realistic mass ratio simulations show. Therefore these results must be taken with caution.

5.3.2 Ion Instabilities and Ion Waves

So far we have described the field and particle properties as have been observed in simulations of quasi-perpendicular supercritical shocks. It has become obvious that in the different regions of the shock transition the particle distributions carry free energy. This is true for the foot region, the shock transition and overshoot as well as the downstream region. And it is true for both species, electrons and ions. This free energy is the source of a number of instabilities which excite waves of different kinds in the various shock transition regions which can be measured. Figure 5.14, taken from *Gurnett et al [1979]*, shows an example of such measurements when the ISEE-1 spacecraft crossed an interplanetary shock travelling outward in the solar wind. The passage of the shock over the spacecraft is seen in the wave instrument in the various channels as a steep increase of the power spectral density of the electric field which is highest in the crossing of the ramp and at medium frequencies of a few 100 Hz under the conditions of the crossing. After the crossing took place the wave power in the downstream region of the shock still remained high but was lower than during the shock crossing. In the right part of the figure the different electric wave power spectra are shown in their time sequence as a stack plot. From this plot the dramatic increase of the power in the medium frequencies during the crossing is nicely visible. Most of the waves excited in this frequency range are caused by electron-ion instabilities which we will discuss in the next section.

The free energy source of the instabilities is less the temperature anisotropy than the direct differences in bulk flow properties of the different species components. We therefore ignore the temperature anisotropy differences even though such instabilities may arise, in particular when ion-whistlers are excited of which we know that they can be driven

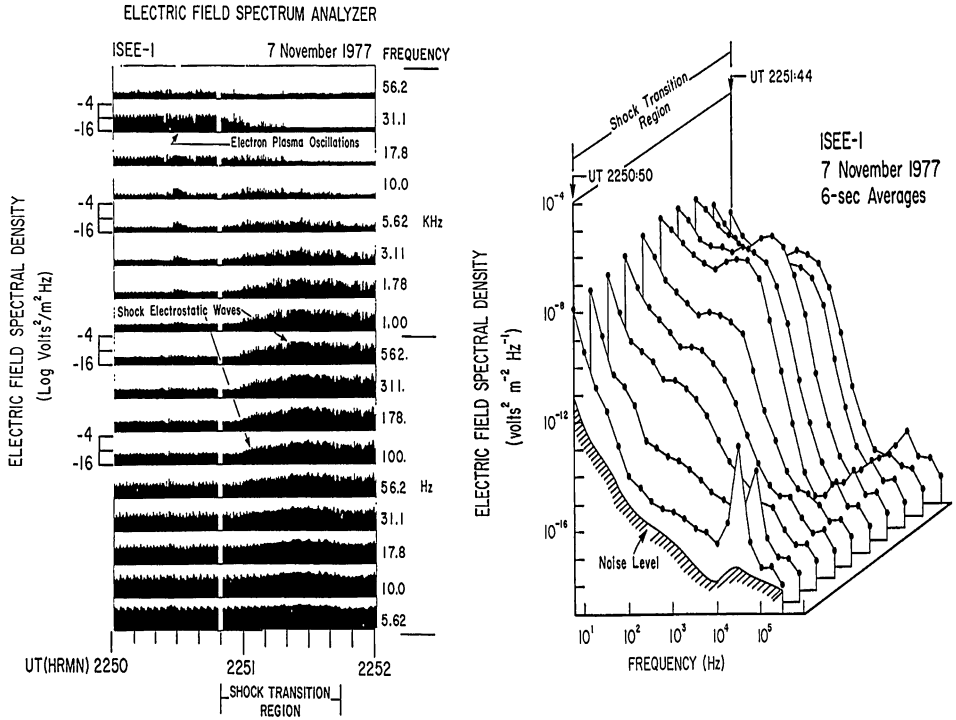


Figure 5.14: Electric wave spectra measured during the spacecraft crossing of an interplanetary shock [after Gurnett *et al*, 1979, courtesy American Geophysical Union]. *Left:* Power spectra (in $\text{V}^2/\text{m}^2 \text{Hz}$) with respect to time in a number of frequency channels. The spacecraft approaches the shock from left and crosses over it. The increase in power is well documented from low to high frequencies when coming into the shock transition region. *Right:* A sequence of shock electric spectra during this crossing given as power spectral density with respect to frequency. The dramatic increase of the low frequency wave power is seen when the spacecraft approaches and crosses over the shock. Behind the shock the power remains high but lower than in the transition region. The Bump around a few 100 Hz is the most interesting from the point of view of instability. These waves are excited by electron-ion instabilities discussed in the next section.

by a temperature anisotropy. To some extent the occurrence of two (counter-streaming in direction x of the shock normal) ion beams already fakes a bulk temperature on the ion component thus generating some relationship between a two-beam situation and a temperature anisotropy. Similarly transversely heated reflected ions superposed on a low perpendicular temperature inflowing ion background fakes a perpendicular temperature anisotropy. For our purposes, however, the bulk flow differences are more interesting and have, in fact, been more closely investigated right from the beginning [Forslund & Shonk, 1970; Forslund *et al*, 1970; Papadopoulos *et al*, 1971; Wu *et al*, 1984].

Foot Region Waves

Let us first consider the foot region. The free energy sources here are the relative drifts between the incoming electrons and reflected ions and the incoming electrons and incom-

Instability	Excitation by	Source of free energy	Direction of propagation
Ion-ion streaming instability	Reflected ions and transmitted ions	Relative streaming between the ion species	$(\mathbf{k} \cdot \mathbf{B}_0) = 90^\circ$
Kinetic cross-field streaming instability	Reflected ions	Relative streaming between the reflected ions and the solar wind electrons	$0 < (\mathbf{k} \cdot \mathbf{B}_0) < 90^\circ$ (In the coplanarity plane)
	Transmitted ions	Relative streaming between the transmitted ions and the electrons	$(\mathbf{k} \cdot \mathbf{B}_0) \leq 90^\circ$ (In the coplanarity plane)
Lower-hybrid-drift instability	Reflected ions	(i) Relative cross-field drift between the reflected ions and the electrons (ii) Density gradient	$(\mathbf{k} \cdot \mathbf{B}_0) < 90^\circ$ (Out of the coplanarity plane)
	Drifting electrons	(i) Relative cross-field drift between the electrons and the transmitted ions (ii) Density gradient	$(\mathbf{k} \cdot \mathbf{B}_0) < 90^\circ$ (Out of the coplanarity plane)
Ion-acoustic instability	Transmitted ions	Relative streaming between the ion species and the electrons	$(\mathbf{k} \cdot \mathbf{B}_0) < 90^\circ$ (Out of the coplanarity plane)
Electron-cyclotron drift instability	Drifting electrons	Electron drift relative to the solar wind ions	$(\mathbf{k} \cdot \mathbf{B}_0) \simeq 90^\circ$ (Out of the coplanarity plane)
Whistler instability	Electrons	Electron thermal anisotropy $T_{e\perp} > T_{e\parallel}$	$(\mathbf{k} \cdot \mathbf{B}_0) \simeq 0^\circ$

Figure 5.15: Instabilities in the Foot Region [copied from *Wu et al*, 1984].

ing ions. The presence of the reflected ions causes a decrease of the ion bulk velocity in the foot region. This implies that the incoming electrons are decelerated so that the current in shock normal direction is zero, i.e. the flow is current-free in normal direction. However, this has the consequence that a relative bulk velocity between electrons and reflected ions or electrons and incoming ions arises. These differences will contribute to the excitation of instabilities. Electrons are not resolved in hybrid simulations, however. In this section we will restrict to ion instabilities leaving the essentially more interesting ion-electron instabilities for the next section. A list of the most important instabilities in the foot region is given in Figure 5.15. This list has been compiled by *Wu et al* [1984]. It is interesting to note that only a few of these instabilities have ever been identified in actual observations and in the simulations even though theoretically they should be present. This can have several reasons, too small growth rates, too strong Landau damping, for instance, in the presence

of hot ions, convective losses or very quick saturation due to heating effects, competition with other waves or wave-wave interaction and so on.

The ion-ion instability [Papadopoulos *et al*, 1971; Wu *et al*, 1984] generates waves in the whistler/lower hybrid frequency range $\omega_{ci} \ll \omega\omega_{ce}$ as we have discussed in Chapter 3. Its energy source is the beam-beam free energy of two counter streaming ion beams, one the reflected ion beam, the other the inflow. As long as the wavelengths are shorter than the reflected ion beam gyroradius the instability is relatively high frequency and electrostatic close to perpendicular and at the lower hybrid frequency. However, at slightly longer wavelength the magnetisation of the ions comes into play and the instability generates electromagnetic whistlers. These are the whistlers which are observed at angles larger than the critical whistler angle mentioned earlier and probably also in the two-dimensional simulation case by Hellinger *et al* [2007] for the parameters used there. An important ion-driven instability in the foot region of quasi-perpendicular supercritical shocks is the whistler instability which we have already mentioned several times. It is related either to reflected ion beams, to assumed temperature anisotropies [as assumed to exist – though never been confirmed by observations – by Wu *et al*, 1984], or to result from diamagnetic density-gradient drifts in the lower-hybrid band as the electromagnetic branch of the lower-hybrid/modified two stream instability, the electrostatic part of which we will discuss in the next section on electronic instabilities and electron dynamics.

Hellinger *et al* [2007], in their two-dimensional hybrid simulations discussed above, have seen the evolution of whistlers without identifying their sources. Recently Scholer & Burgess [2007] performed an extensive parametric search for whistler waves in the foot region of oblique shocks between $60^\circ \leq \Theta_{Bn} \leq 80^\circ$, the region where in one-dimensional PIC simulations intense whistlers should become excited theoretically when reflected ions are present. This is the case for the more oblique but still quasi-perpendicular supercritical shocks. A wide range of Alfvénic Mach numbers was used, and strong excitation of whistlers in the parametric range was found indeed. We will discuss these observations here in more detail as they are the currently existing best available simulation results representing the current state of the art in the field of whistler excitation in connection with the formation, stability and time dependence of supercritical shocks at the time of writing this review.

Before discussing their results we briefly review the physics involved in the importance of whistlers in shock foot stability as had already been anticipated by Biskamp & Welter [1972] following a suggestion by Sagdeev [1966]. As we have mentioned earlier, a linear whistler wave precursor can stand in front of the quasi-perpendicular shock as long as the Mach number $\mathcal{M} < \mathcal{M}_{wh}$ is smaller than the critical whistler Mach number

$$\mathcal{M}_{wh} = \frac{1}{2} \left(\frac{m_i}{m_e} \right)^{\frac{1}{2}} |\cos \Theta_{Bn}| \quad \text{or nonlinearly} \quad \frac{\mathcal{M}_{wh, nl}}{\mathcal{M}_{wh}} = \sqrt{2} \left[1 - \left(\frac{27\beta}{128\mathcal{M}_{wh}} \right)^{\frac{1}{3}} \right] \quad (5.21)$$

The second expression results when including the nonlinear growth of the whistler amplitude during the steepening process [for the derivation of this expression see, e.g., Kazantsev, 1961; Krasnoselskikh *et al*, 2002]. The nonlinear critical whistler Mach number $\mathcal{M}_{wh, nl}$ is larger by a factor of $\sqrt{2}$ than the whistler Mach number $\mathcal{M}_{wh}$. It depends weakly on the plasma- β which has a decreasing effect on it, slightly reducing the whistler range.

For instance, with realistic mass ratio $m_i/m_e = 1840$ and $\Theta_{Bn} = 87^\circ$ the whistler critical Mach number is quite small, $\mathcal{M}_{wh} \simeq 1.14$. Otherwise no standing whistlers can thus be expected. It has also been speculated that above the above critical whistler Mach number the shock ramp is replaced by a nonlinear whistler wave train with wavelength of the order of λ_e . Approximating such a train as a train of whistler solitons one realises that the amplitude of the solitons increases with Mach number. Hence, the critical whistler Mach number in this case must become dependent on the whistler amplitude [Krasnosel'skikh et al, 2002]. This leads to the slightly larger nonlinear whistler critical Mach number $\mathcal{M}_{wh,nl}$ on the right in the above equation. At such realistic low Mach numbers the shocks are subcritical, however, and we are back to the problems discussed in Chapter 4.

Between the two Mach numbers $\mathcal{M}_{wh} \leq \mathcal{M} \leq \mathcal{M}_{wh,nl}$ the nonlinear whistler soliton train can exist attached to the ramp. However, when the nonlinear whistler Mach number is exceeded this is not possible anymore, and the whistler wave should turn over due to a so-called gradient catastrophe leading to non-stationarity of the shock front, which we will discuss later. In the simulations of Biskamp & Welter [1972] the simulation range was in favour of the excitation of whistlers which have also been seen and by these authors had been attributed to a nonlinear instability between the two ion beams and the electric field of a standing whistler wave an interpretation which is overturned by the new simulation results. Scholer & Burgess [2007] performed PIC simulations with physically realistic mass ratio. For all shocks with $\Theta_{Bn} \leq 83^\circ$ the whistler critical Mach number is well above the critical Mach number such that the shock is supercritical. This is in order to check the excitation of whistlers in the different regimes of $\mathcal{M}$. In the left part of Figure 5.16 the Mach number is below the critical whistler Mach number, and in the shock foot region a group of phase locked whistler waves is excited with increasing amplitudes towards the shock ramp. This is nicely seen on the left in both the magnetic stack plot as also in the time profile in the second panel on the left from top. The whistlers slow the incoming flow V_{ix} down before it reaches the ramp. In the phase space plot the incoming and reflected beams are seen as is the scattering and trapping of the resonant ions in the whistler waves. The right part of the figure shows an identical simulation with Mach number above the critical whistler Mach number but below the nonlinear whistler Mach number. The stack plot shows two well developed reformation cycles with all signs of normal reformation. The magnetic field profile chosen in the second panel is at time $t\omega_{ci} = 7.6$ in the second reformation cycle when the foot loop is well developed. The magnetic field signature shows the distortion due to the foot which is caused by a large amplitude non-phase-standing whistler wave that evolves nonlinearly. However, the reformation is not due to this whistler but due to the accumulation of gyrating ions at the foot edge as known from previous simulations.

Near the ramp the ions become trapped in a large-amplitude whistler loop, as is seen in the phase space plot. The loop coincides with a minimum in the B_z -component of the magnetic field. The whistler, on the other hand, does practically not affect reformation, even though it structures the overshoot region. Reformation time is defined thus, as we know already, by the gyration time of ions in the foot, being of the order of a few ion-cyclotron periods. When the shock becomes more oblique, the whistler effect increases again at fixed Mach number as the Mach number enters the domain below critical whistler Mach num-

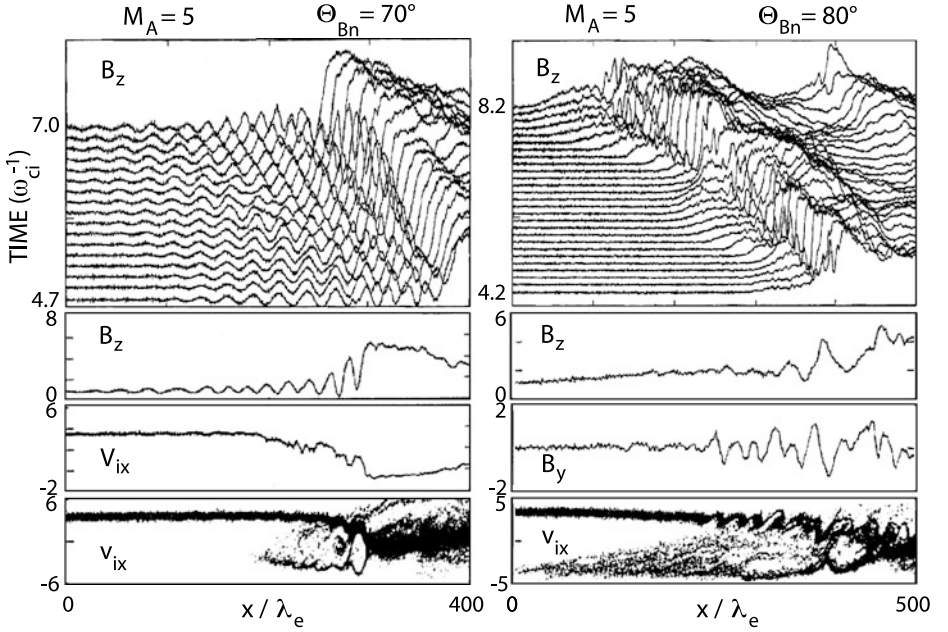


Figure 5.16: One-dimensional full particle PIC simulations with realistic mass ratio for the same Mach number but different angle $\Theta_{Bn} = 70^\circ$ and 80° [after Scholer & Burgess, 2007]. *Left top:* Mach number $\mathcal{M} < \mathcal{M}_{wh}$. The shock foot region is filled with waves of two polarisations, one of the expected standing whistler waves which interfere with the other kind of waves. No substantial shock reformation is observed in this case. *Right top:* Here the Mach number is in the range Mach number $\mathcal{M}_{wh} < \mathcal{M} < \mathcal{M}_{wh, nl}$. Two reformation cycles are visible during the simulation run in the magnetic field, and no whistler waves occur because the Mach number exceeds the first whistler Mach number. However there is also no nonlinearly steepened whistler in the shock front which is simply taken over by the foot after one reformation cycle. *Bottom three panels:* On the *left* shown the standing whistler oscillations in the magnetic field on the *left*, the decrease in the flow velocity when entering the foot due to the whistler scattering of the ions, and the particle phase space in v_{ix} with the reflected beam and the scattered foot ions with little vortex formation such that the reformation is inhibited. On the *right* seen is the more irregular structure of the transition, the component B_y in the magnetic field and the accelerated foot ions in the absence of whistlers.

ber $\mathcal{M}_{wh}$, and the shock transition becomes much more structured and reformation less important.

In order to see what kind of waves are excited during the whistler cycles, a separation of the magnetic wave spectrum B_y into positive B_y^+ and negative B_y^- helicity components has been performed for a $\mathcal{M}_A = 9$, $\Theta_{Bn} = 70^\circ$ simulation run. Figure 5.17 shows the result. The negative helicity waves B_y^- propagate toward the shock, i.e. to the right. After correcting for the convection velocity which is also to the right, these waves turn out to be left-hand polarised. The lower panel shows positive helicity B_y^+ waves propagating to the left, so they are upstream propagating waves and are also left-hand polarised. The positive helicity waves have longer wavelength than negative helicity waves. They propagate close to the shock speed upstream. They are thus almost standing in the shock frame. These are the

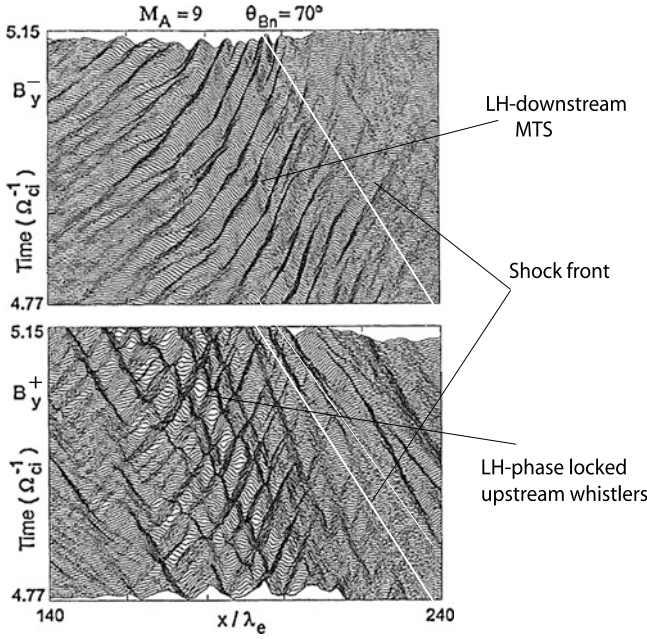


Figure 5.17: One-dimensional full particle PIC simulations with realistic mass ratio for Mach number $M_A = 9$ and $\Theta_{Bn} = 70^\circ$ in the non-reformation whistler regime [after Scholer & Burgess, 2007]. *Top:* Negative helicity waves B_y^- propagating to the right are left-hand polarised short wavelength waves moving downstream toward the shock and being mostly absorbed in the shock transition. *Bottom:* Positive helicity waves B_y^+ . These waves move upstream and thus are also left-hand circularly polarised waves. They move at shock velocity which identifies them as the phase-locked standing whistler precursors in the shock-upstream region with decaying upstream amplitude and long wavelength. Some interference is seen on these waves. Their left-hand polarisation identifies them as ion-beam excited whistlers and not as electron temperature anisotropy excited whistlers.

upstream left-hand polarised (ion beam and not electron temperature anisotropy driven) whistlers.

The downstream propagating negative helicity waves are no whistlers. They are caused in quite a different way which is related to the electromagnetic modified two-stream instability which we will discuss in the next section on electron waves.

Ramp Transition Waves

Stability of the ramp is a question that is not independent of the stability of the foot as both are closely connected by the reformation process of the quasi-perpendicular shock front. Since the suggestion of Sagdeev [1966] it is widely believed that the whistler waves excited in the foot are the main responsible for the stability of the foot and formation of the ramp by steeping. In fact, they might accumulate their, store energy in both magnetic and electric field, trap particles and excite different waves.

Instability	Nature of wave mode	Typical wavelength	Frequency and growth rate	Remarks
Ion-ion streaming instability	Magnetosonic waves	$k \sim \frac{\omega_e}{c}$	$\gamma \sim \Omega_i$	Stabilized when the streaming velocity exceeds the Alfvén speed.
Kinetic cross-field streaming instability	Whistler mode waves with oblique propagation	$k \gtrsim \frac{\omega_{LH}}{V_0}$	$\omega \simeq \omega_{LH}$ $\gamma > \Omega_i$ $\omega \simeq \omega_{LH}$ $\gamma > \Omega_i$	The instability persists even if $V_0 \gg v_A$
Lower-hybrid-drift instability	Lower hybrid waves and drift waves	$k \sim \frac{\omega_{LH}}{V_0}$	$\omega \simeq \omega_{LH}$ $\gamma \gg \Omega_i$	Instability enhanced by ∇T_e
	Doppler-shifted whistler mode	$k > \frac{\omega_{LH}}{V_0}$	$\omega \simeq \omega_{LH}$ $\gamma \gg \Omega_i$	
Ion-acoustic instability	Ion waves	$k\lambda_D \lesssim 1$	$\omega \lesssim \omega_i$ $\gamma > \Omega_i$	Instability enhanced by ∇T_e
Electron-cyclotron drift instability	Doppler-shifted Bernstein waves and ion waves	$k\lambda_D \lesssim 1$	$\omega \simeq n\Omega_e$ $\gamma > \Omega_i$	Instability suppressed by ∇B
Whistler instability	Whistler-mode waves with parallel propagation	$k \lesssim \frac{\omega_e}{c}$	$\omega \ll \Omega_e$ $\gamma \gg \Omega_i$	

Figure 5.18: Instabilities in the Ramp Transition [copied from Wu *et al.*, 1984].

A very early and to large extent out of time list is given in Figure 5.18 of the theoretical expectations for possible instabilities in the shock ramp region [Wu *et al.*, 1984]. As in the case of the foot region, the instabilities in the ramp which might be of real importance have turned out to not fit very well into the scheme of this listing. In both cases, the case of the foot and the case of the ramp, this disagreement reflects the weakness of theoretical speculation which is not supported by direct observations on the one hand and clear parametric searches in numerical experiments on the other.

Waves excited or existing in the ramp cannot be considered separate from the stability of the shock ramp. They are mostly related to electron instabilities and will to some extent been considered there in the next section. On the other hand they are also related to the non-stationarity of a shock ramp. We will therefore return to them also in the respective section on the time dependence of evolution of shock ramps and their stability. It is, however, worth mentioning that recently very large electric fields have been detected during

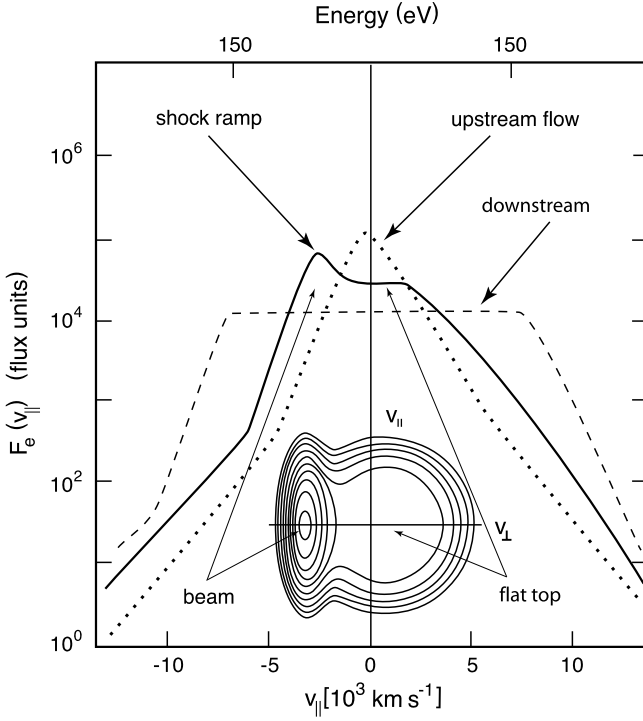


Figure 5.19: Several successive reduced parallel electron distribution functions $F_e(v_{\parallel})$ during the crossing of the supercritical bow shock of the Earth by ISEE 2 on December 13, 1977. The cuts through the distribution show the transition from the Maxwellian-plus-halo upstream flow distribution through the shock ramp distribution to the close to the shock downstream distribution. The shock ramp distribution is intermediate in evolving into a flat-top distribution of the kind of the downstream distribution but contains in its upstream directed part a well expressed shock-reflected electron beam of velocity of a few 1000 km s^{-1} which is sufficiently fast to destabilise the shock front and excite electron plasma waves [after Gurnett, 1985, courtesy American Geophysical Union].

passages of the Polar satellite across the quasi-perpendicular bow shock when the spacecraft was traversing the shock ramp.

These observations showed that in the shock ramp electric fields on scales $\lesssim \lambda_e = c/\omega_{pe}$ exist with amplitudes of the order of several 100 mV m^{-1} . These are amongst the strongest localised electric fields measured in space [Bale & Mozer, 2007]. Clearly, these localised fields are related to the electron dynamics in the shock and in particular in the shock overshoot/ramp region. Excitation of intense electron waves in the shock ramp has been expected for long time already since the observation of the (reduced) electron distribution across the shock from a drifting to a flat-topped distribution [Feldman *et al*, 1983]. Figure 5.19, taken from Gurnett [1985], shows this transition. The interesting point is that right in the ramp/overshoot region the reduced electron distribution function shows the presence of an electron beam in addition to an already quite well developed flat top

on the distribution. Such a beam will almost inevitably serve as the cause of instability.

5.3.3 The Quasi-perpendicular Shock Downstream Region

We have already mentioned earlier that a shock without a downstream region does not exist. Detached shocks that evolve in front of blunt obstacles like a magnetosphere, in particular, possess extended downstream regions which separate the obstacle from the undisturbed upstream region. These regions are the domains of compressional waves – of the family of magnetosonic waves – with upstream directed phase and group velocities larger than the downstream flow velocity that in a characteristic flow time across the downstream distance from the shock to the obstacle can make it upstream to reach the shock, which forms the spatial envelope of these waves thereby confining them to downstream of the shock transition.

This is the fluid picture of the evolution of the region downstream of the shock. In this subsection we are not going into an extended discussion of the downstream region of a quasi-perpendicular shock but restrict only to a few remarks. The reason is that for a curved collisionless supercritical shock it is very difficult to distinguish between the processes triggered by the quasi-perpendicular and quasi-parallel parts separately. In order to do this one needs to have plane shocks which are found possibly at the giant distant planets (cf. Chapter 10 on Planetary Bow Shocks) and at interplanetary travelling shocks. The latter, unfortunately, are mostly subcritical as they are convected by the stream and therefore their downstream regions differ from those of supercritical shocks. Because of this reason we delay a more substantial discussion of downstream regions to the next chapter after having presented the main properties of quasi-parallel shocks.

Collisionless shocks cannot be described solely in the fluid picture, however. Therefore the downstream region has quite a different behaviour from that inferred from the hydrodynamic description. This is particularly true for quasi-parallel shocks which (see the next chapter). Downstream of a quasi-perpendicular or perpendicular super-critical shock the conditions are as well very different from those which the fluid picture prescribes which of the plasma properties basically predicts an increase in the fluid pressure anisotropy due to compression of the magnetic field and plasma under conservation of the fluid magnetic moment (the ratio $T_{\perp}/B$), the deviation of the flow around the obstacle, draping of the magnetic field, and a pile-up region of the magnetic field close to the obstacle where pressure balance requires dilution of the plasma.

The main reasons for a deviation from this average and laminar behaviour of the downstream region are that in the supercritical case the fluid picture does not contain the ion (and electron) reflection processes and their consequences for shock foot formation, ramp physics and formation of the shock overshoot. Since collisional dissipation is out of question, the reflection mechanisms replace the necessary dissipation. It is easy to estimate, for instance for the quasi-perpendicular Earth's bow shock, how much energy must be dissipated in order to adjust for the differences in the flow properties between upstream and downstream. For an upstream solar wind of nominal density $N = 5 \times 10^6 \text{ m}^{-3}$ and velocity $V_1 = 500 \text{ km/s}$ at 1 AU, Mach number $\mathcal{M} = 5$ and moderate shock strength (compres-

sion ratio) of $B_2/B_1 = 3$, the energy density in the flow behind the shock (Mach number $\mathcal{M} = 1$) is just half the energy density in the solar wind. Thus 50% of inflow energy must be dissipated by the shock. Since at most 10% of the ions are reflected (corresponding to 10% of incoming energy only), the remaining 40% are absorbed and dissipated in other ways.

A first simple effect that affects the region behind the shock ramp is that the reflected ions are accelerated by the motional electric field in the shock foot along the shock surface in the direction perpendicular to the upstream magnetic field to energies that are much higher than the kinetic energy of the incoming ions, thereby absorbing a substantial part of inflow energy. Acceleration of the 10% reflected ions by a factor 2 gives already 40% of energy dissipation. Hence this is the main dissipation mechanism.

We have already described how this acceleration affects the nearly periodic reformation of a quasi-perpendicular shock when these ions carry a drift current and gyrate at high speeds and large ion gyroradii in the upstream magnetic field. However, after a few such gyrations they have gained sufficient energy to break through the reflecting shock potential and overcome the shock barrier to enter the downstream flow region. Downstream of the quasi-perpendicular shock they appear as an energetic gyrating ion component with much larger perpendicular than parallel energy.

In fact these gyrating energetic ions have been observed [Sckopke *et al*, 1983]. They are a source of free energy and excite a number of instabilities like electromagnetic ion-cyclotron waves which can grow to large amplitudes downstream of the shock and contribute to the magnetic fluctuation that are observed downstream. Even mirror modes just behind the shock have been reported [Czaykowska *et al*, 1998]. The rms amplitudes $b_{\text{rms}} = \sqrt{\langle |\mathbf{b}|^2 \rangle} \sim B_2$ of these waves downstream of the shock are comparable to the average downstream field B_2 . Therefore these waves contribute to the energy balance. Moreover, since they are convected downstream and damp away, they contribute to heating of the downstream plasma and shock energy dissipation.

The large amplitude waves do not damp but decay into other shorter amplitude waves or by ponderomotive force interaction may stabilise and form large amplitude localised magnetic structures, density pulses or also narrow current sheets. The physics of these structures is still barely understood. Sometimes one speaks about downstream turbulence which, however, is not qualified and also unjustified as turbulence can hardly develop in the downstream region because it has not sufficient time for reaching a well mixed state. This would require a very large extended region and relatively quiet boundary conditions with continuous inflow at an approximately constant energy level. It would also require stationarity of the shock which is definitely not given, as we will see later when discussing the problem of stationarity. Shock reformation is already kind of a rather non-stationary process that makes the development of downstream turbulence rather doubtful.

The narrow current sheets which the presence of such wave structures implies have transverse scales comparable to the ion inertial length. They must close in themselves. They thus form current vortices on scales of several ion gyroradii. Even though they are very thin, their dynamics is different from the dynamics of narrow current sheets that undergo reconnection, drift kink modes and tearing instabilities like the current sheets

in the tail of the Earth's magnetosphere. The reason is that these vortices are convected with the flow and there is very little velocity difference between them to trigger reconnection.

In addition, adjacent current sheet vortices have antiparallel currents and experience repulsing Lorentz forces such that they avoid close contact. However, if the currents become strong enough, i.e. if the amplitudes of the magnetic field fluctuations become large enough, they can decay due to generation of current instabilities which again are mostly in the drift modes and lower hybrid modes, accelerate electrons along the magnetic field and heat the plasma by dissipating their energy. Electron heating and the appearance of lower hybrid modes in connection to the observation of these current vortices thus rather indicates their decay than reconnection.

To these downstream excited waves a couple of other waves can be added. These are waves that are generated in the upstream shock foot having phase and group velocities less than the flow velocity and thus are swept down into the shock, where they accumulate and contribute to local dissipation in case their group velocities compensate for the reduced downstream flow across the shock. Otherwise they pass the shock and contribute to the downstream fluctuations.

Among those waves are the various low frequency electrostatic waves that are excited in the foot and will be described below. They propagate in the ion acoustic and lower hybrid mode frequency ranges and appear downstream as broadband low frequency noise where they have been observed for long time [Rodriguez & Gurnett, 1975]. The electrostatic modes in the lower-hybrid range have electric field components along the magnetic field and thus accelerate electrons into beams, a process that serves the transport of energy from the ion flow to the electron component and is responsible for electron heating, formation of electron beamlets and generation of the recently detected downstream solitary structures in the electron plasma wave component [Pickett *et al.*, 2004].

In the next chapter on quasi-parallel shocks we compare the downstream wave spectra for both quasi-perpendicular and quasi-parallel shocks in Figure 6.50. However, even this short reconciliation of the fluctuating state of the downstream region behind a quasi-perpendicular shock shows that the downstream medium is in a rather complex plasma state investigation of which is interesting in itself though poorly understood yet.

5.4 Electron Dynamics

When talking about the dynamics of electrons hybrid simulations cannot be used anymore. Instead, one must return to the more involved full particle PIC simulation codes or to Vlasov codes, which directly solve the Vlasov equation in the same way as a fluid equation, this time, however, for the “phase space fluids” of ions and electrons. In both cases short time scales of the order of the electron gyro-period ω_{ce}^{-1} or even the electron plasma period ω_{pe}^{-1} must be resolved, and resolution of spatial scales of the order of the electron inertial scale λ_e and Debye scale λ_D is required. It is thus no surprise that reliable simulations of this kind became available only within the last decade with the improved computing capacities.

5.4.1 Shock Foot Electron Instabilities

Papadopoulos [1988] proposed that in the foot region of a perpendicular highly supercritical shock the velocity differences between reflected ions and electrons from the upstream plasma inflow should be responsible for the excitation of the Buneman two-stream instability thus heating the electrons, generating anomalous conductivity and causing dissipation of flow energy which contributes to shock formation.

Buneman Two-Stream Heating in Strictly Perpendicular Shocks

Shimada & Hoshino [2000] and *Schmitz et al* [2002] building on this idea performed full particle PIC simulations in strictly perpendicular shocks discovering that the Buneman two-stream instability can indeed work in the foot region of the shock and can heat and accelerate the electrons. *Shimada & Hoshino* [2000] initiated their one-dimensional simulations for a small mass ratio of $m_i/m_e = 20$, $\beta_i = \beta_e = 0.15$, and Alfvénic Mach numbers $3.4 \leq \mathcal{M}_A \leq 10.5$.

Figure 5.20 shows some of their simulation results. It is interesting to inspect the right part of the figure which shows the (shaded) ramp and foot regions on the left in expanded view. The electron phase space shows the development of electron holes which are generated by the Buneman two stream instability in this strictly perpendicular shock simulation. The signature of the electrostatic field E_y in the lowest panel shows the bipolar electric field structure the holes cause. The average field is zero, but in the hole it switches to large negative values, returns to large positive values and damps back to zero when passing along the direction normal to the hole. This is exactly the theoretical behaviour expected for both, solitons and electron holes of the form of BGK modes. As known from simulations (see Chapter 3) such BGK-hole structures will trap electrons and heat them, they do, on the other hand, also accelerate passing electrons to large velocities. Both is seen here also in the simulations in the vicinity of the shock: Three such holes are completely resolved in the right high resolution part of the figure, with decreasing amplitude when located closer to the shock ramp. All of them contain a small number of trapped electrons over a wide range of speeds which on the gross scale on the left fakes the high temperature of the electrons they achieve in the hole. This is just the effect of heating by the two-stream instability. In addition the electron velocity shows two accelerated populations, one with positive velocity about 2–3 times the initial electron speed, the other reflected component with velocity almost as large as the positive component but in the opposite direction suggesting that the electron current in the holes is almost compensated by the electron distribution itself.

Obviously the further strong heating of electrons in the ramp is caused by many smaller amplitude overlapping holes as is suggested by the structures in the inflowing and reflected ion distributions which do also strongly interact with the electric field of the holes. This is seen in the incoming ion component in the first hole as a dip in the velocity. The hole retards the incoming flow. It is also seen in the reflected ion component as strong distortion of their backward directed velocity when encountering a hole. The smaller speed ions are obviously retarded in their backward flow and are partially trapped in the negative electric

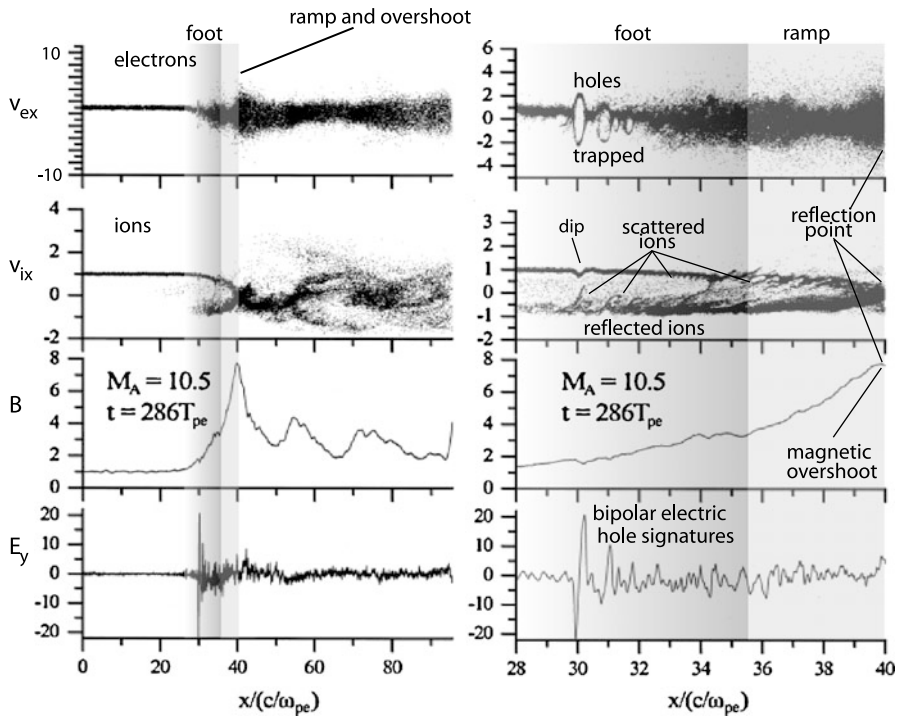


Figure 5.20: One-dimensional full particle PIC simulations with mass ratio 20 for Mach number $\mathcal{M}_A = 10.5$ and $\Theta_{Bn} = 90^\circ$ resolving the electron scales [after Shimada & Hoshino, 2000]. *Left:* Simulation overview for electron and ion phase spaces, magnetic field and electric field. Ion reflection from the ramp and foot formation is seen in the *second panel from top*. The electrons are heated in the foot region. the heating coincides with large amplitudes in the electric field in the *lowest panel*. *Right:* Expanded view of the shaded foot and ramp regions on the *left*. The electron heating location turns out to be a site of electron hole formation. Three Buneman two-stream instability holes are nicely formed on this scale with trapped electrons. The broadening of the distribution and thus heating is due to the holes. Retardation of the ions in interaction with the holes is seen in the *second panel* which is due to the retarding electric potential in the large amplitude electric field oscillations (*lowest panel*). Interestingly enough, ion reflection takes place in the very overshoot! The ion distribution is highly structured in the entire region which is obviously due to interaction with many smaller scale electron holes.

field part of the hole. Very similar strong scattering of the incoming ion component is seen in the ramp region. This suggests that a large number of electric field structures are located in the ramp which scatter the incoming ions. These must be related to the highly fluctuating electric field component in the ramp seen on the right in the lowest panel.

Two further observations which are related to the ion component are of considerable interest: The first is that the retardation of the incoming ion flow and the scattering of the reflected ions in the foot region cause a signal on the magnetic field component. The second is that the reflection of the main incoming ion beam, i.e. the incoming plasma takes place at the location of the magnetic overshoot and not in the shock ramp. Therefore, physically spoken, the shock ramp is also part of the foot, while it is the narrow overshoot

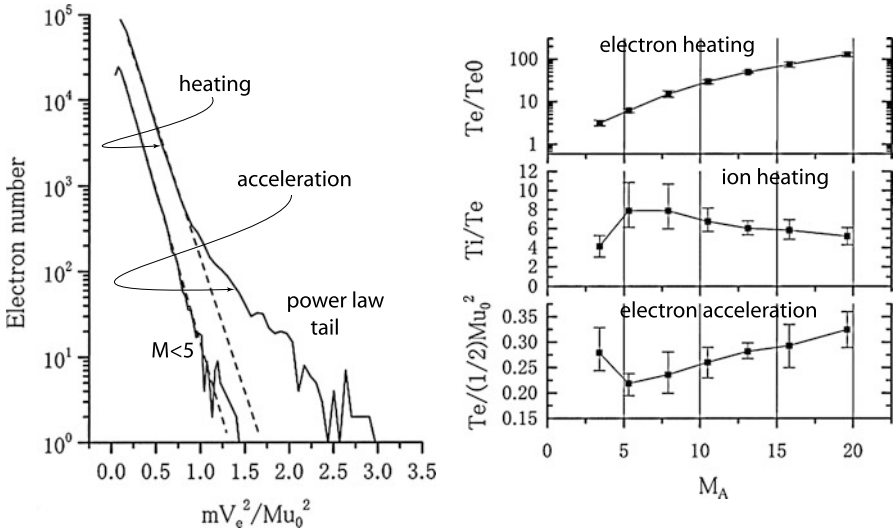


Figure 5.21: *Left:* The electron distribution in the shock arising from the action of the Buneman electron hole interaction at large Mach numbers $\mathcal{M}_A > 5$. The interaction not only causes heating but also an energetic tail on the electron distribution function. This tail has the shape of a power law $F\varepsilon \propto \varepsilon^{-\alpha}$, with power $\alpha \approx 1.7$. Note that this power gives a very flat distribution close to marginal flatness $\alpha = \frac{3}{2}$ below that an infinitely extended distribution function has no energy moment. *Right:* Evolution of the average electron temperature, ion temperature and ratio of electron temperature to initial kinetic energy in the simulations as function of Alfvén Mach number [after Shimada & Hoshino, 2000]. All quantities are in relative units of computation.

region where the reflection occurs in a strictly perpendicular supercritical shock with cold ion inflow. The actual ramp region is much narrower than for instance shown in the figure. Its actual width is only of the order of $\Delta \sim (1-2)c/\omega_{pe}$.

Electron Heating and Acceleration

Shimada & Hoshino [2000] followed the evolution of the electron vortices (holes) and showed that a hole once evolved distorts the ion and electron velocities in such a way that nonlinearly the velocity difference can increase and cause the generation of secondary vortices, which leads to excessive electron heating [see also Shimada & Hoshino, 2005]. The result is the generation of an extended electron tail on the electron distribution. This is seen from the left part of Figure 5.21 in a log-lin representation of the electron number versus normalised electron energy. When plotting the data on a log-log scale (not shown) one realises that the newly produced tail of the electron distribution has a power law slope $F(\varepsilon) \propto \varepsilon^{-\alpha}$, notably with power $\alpha \approx 1.7$. (Note that this power is close to the marginally flattest power $\alpha = \frac{3}{2}$ below that an infinitely extended power law energy distribution has no energy moment more and thus ceases to be a distribution. In fact, any real nonrelativistic power law will not be infinitely extended but will be truncated due to the finite extent of the volume and loss of energetic particles.)

The dependence of electron heating and ion cooling on Mach number for the investigated range of Alfvénic Mach numbers is plotted on the right. The effect does not occur for small Mach numbers, too small for the Buneman two-stream instability to be excited. However, once excited, the heating increases strongly with $\mathcal{M}_A$. Over the range $5 < \mathcal{M}_A < 20$ the increase in electron temperature (electron energy stored mainly in the tail of the distribution) is a factor of 40–50, which demonstrates the strong non-collisional but anomalous transfer of kinetic flow energy into electron energy via the two-stream instability. However, one should keep in mind that this result holds merely for a one-dimensional simulation of strictly perpendicular shocks.

At this place we should look again at observations during crossings of real collisionless shocks in space. Recently, during passages of the Polar satellite of the quasi-perpendicular Bow Shock of the Earth, very strong localised electric fields have been detected. These fields exist on scales $\lesssim \lambda_e = c/\omega_{pe}$, less than the electron skin-depth, and reach enormous values of $\lesssim 100 \text{ mV m}^{-1}$ parallel and $\lesssim 600 \text{ mV m}^{-1}$ perpendicular to the magnetic field. They must naturally be related to the electron dynamics and should play a substantial role in the formation and dissipation processes of the quasi-perpendicular shock. They should also be of utmost importance in accelerating electrons (and possibly also ions) at shocks. Their nature still remains unclear, however, it is reasonable to assume that they are generated by some electron-current instability via either the Buneman-two-stream instability, the modified two-stream instability which we discuss below, or the ion-acoustic instability, depending on the current strength. In any case they will turn out to belong to the family of Bernstein-Green-Kruskal modes which are encountered frequently in collisionless plasmas.

5.4.2 Modified-Two Stream Instability

The Buneman two-stream instability works on scales $\leq \lambda_e = c/\omega_{pe}$. This condition is less easily satisfied in quasi-perpendicular shocks. However, here other instabilities can evolve which are relatives of the Buneman two-stream instability.

The condition that there is no current flowing in the shock normal direction during foot formation and reflection of ions at the shock requires that the electron inflow from upstream is decelerated when entering the foot region. This causes a difference in the ion and electron inflow velocities. In a quasi-perpendicular shock the wave vector $\mathbf{k} = (k_{\parallel}, \mathbf{k}_{\perp})$ is allowed to have a component $k_{\parallel}$ along the magnetic field. The velocity difference between ions and electrons can then excite the modified two-stream (MTS) instability, a modification of the Buneman instability acting in the direction perpendicular to the magnetic field. This instability is electromagnetic coupling the Buneman two-stream instability to the whistler mode. The waves generated propagate on the whistler mode branch with frequency $\omega_{\text{mtsi}} \sim \omega_{lh} \ll \omega_{ce}, \omega_{pe}$, close to the lower-hybrid frequency, but far below both the electron cyclotron and electron plasma frequencies, respectively. These waves are capable to modify the shock profile when being swept downstream towards the shock ramp. Their obliqueness generates a magnetic field aligned wave electric field component which accelerates, traps and eventually pre-heats the electrons in the shock foot along the magnetic field.

Relation to the Buneman Instability

Scholer & Matsukiyo [2004] investigated the transition from Buneman to modified two-stream (MTS) instabilities as function of mass ratio m_i/m_e and for various β_i, β_e in the regime where no upstream standing whistlers exist, i.e. above the critical whistler Mach number $\mathcal{M}_A > \mathcal{M}_{wh}$. This investigation is restricted to shocks, however, with $\mathbf{k}$ -vectors being strictly perpendicular to the shock along the shock normal and for one-dimension only. This excludes any waves which could propagate along the inclined magnetic field. Nevertheless, this investigation is interesting in several respects. First it showed that for mass ratios $m_i/m_e \lesssim 400$ no modified two-stream instabilities occur since their growth rates are small. The electron dynamics and the shock behaviour in this range are determined by the Buneman two-stream instability unless the electron temperature is large enough to inhibit its growth in which case ion-acoustic instability should (or could) set on (but has not been observed or has not been searched for). For larger mass ratios (and particularly for the realistic mass ratio) the Buneman two stream instability ceases to be excited. Instead, the modified two-stream instability (MTSI) takes over which is strong enough to completely determine the behaviour of the electrons. A summary of their results is schematically given in Figure 5.22.

The evolution of the MTS-waves for realistic mass ratio simulations is shown in Figure 5.23 for three instants of time. The wave spectrum has been determined in the shaded area. Large amplitude waves of left-hand polarisation propagate toward the shock during this reformation cycle. These waves are related to the electron dynamics. They are excited by the modified two-stream instability in the foot (top panel) in interaction between the retarded electrons and the fast ions. The simulations also show the evolution of large amplitude electron holes and ion holes (right lower panel). Such structures have been observed in the electric field in the quasi-perpendicular shock region [*Bale et al, 2002; Pickett et al, 2004; Balikhin et al, 2005; Hull et al, 2006; Hobara et al, 2008*] with differing interpretations. From the simulations it is concluded that both kinds of holes/solitary structures are excited near a quasi-perpendicular shock on similar scales while being related to the combined electron and ion dynamics.

Modified Two-Stream Instability and Quasi-perpendicular Shock Reformation

They cause reformation of the shock, but in a different way than it is caused for low mass ratio by the Buneman-instability. There the reformation was the result of accumulation of ions at the upstream edge of the foot, while here it is caused by participation of the foot ions in the MTSI all over the foot and particularly close to the shock ramp and presumably also at the ramp itself. Phase mixing of the ions leads to bulk thermalisation and formation of a hot retarded ion component in the foot region which has similar properties like the downstream population and, when sufficiently compressed takes over the role of the shock ramp. This can be seen from the lower right part of Figure 5.23 which is a snapshot at time $t\omega_{ci} = 3.7$ showing the magnetic profile, the density profile with its strong distortions, and the evolution of the ion distribution which evolves into large thermalised vortices towards

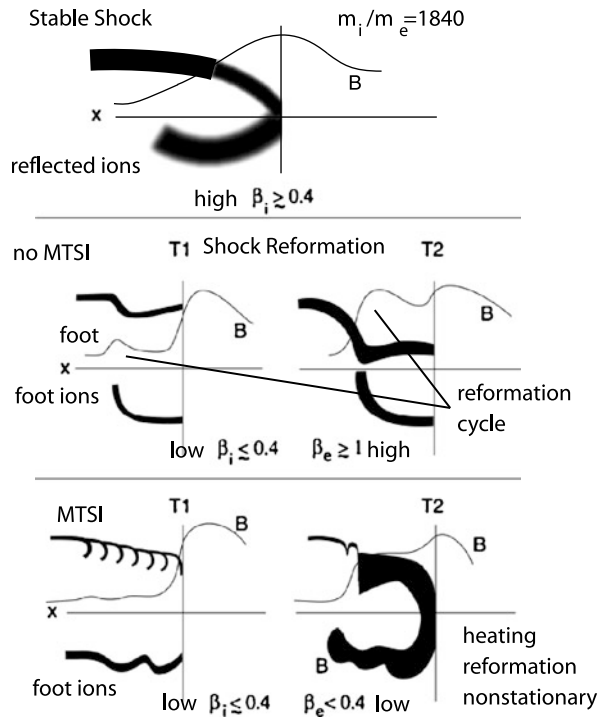


Figure 5.22: Schematic of the dependence of the shock structure on the combinations of β_i, β_e for quasi-perpendicular supercritical but non-whistler shocks. For large β_i the shock is stable even though ions are reflected. At small β_i , large β_e the shock reforms due to accumulation of ions at the edge of the foot forming a reformation cycle. For small β the MTSI evolves in the foot, strong heating and complicated dynamics evolves due to nonlinear interaction, heating and hole-vortex formation [after Scholer & Matsukiyo, 2004].

the front of the shock (note that here the shading indicates also the spatial domain where the wave spectra have been taken).

The generation of MTS-waves by the modified two-stream instability has been investigated in depth theoretically and with the help of specially tailored one-dimensional numerical simulation studies by Matsukiyo & Scholer [2003], and in two-dimensional simulations by Matsukiyo & Scholer [2006a, b] which we are going to discuss in detail.

Modified Two-Stream Generation Mechanism: Tailored Simulations. Figure 5.24 in its left-hand parts shows the set-up of the two-dimensional simulation and the resulting time histories of fields and particles. The incoming and reflected ion velocities are shown for time zero in the (x, y) -plane where the co-ordinate y is about parallel to the magnetic field. The phase space at time zero contains the three distributions of inflow and reflected ions and hot incoming electrons. The slight displacement between the latter and the incoming ions accounts for zero normal current flow in presence of reflected ions. Clearly this configuration is unstable causing instabilities between the ion beams and electrons (in

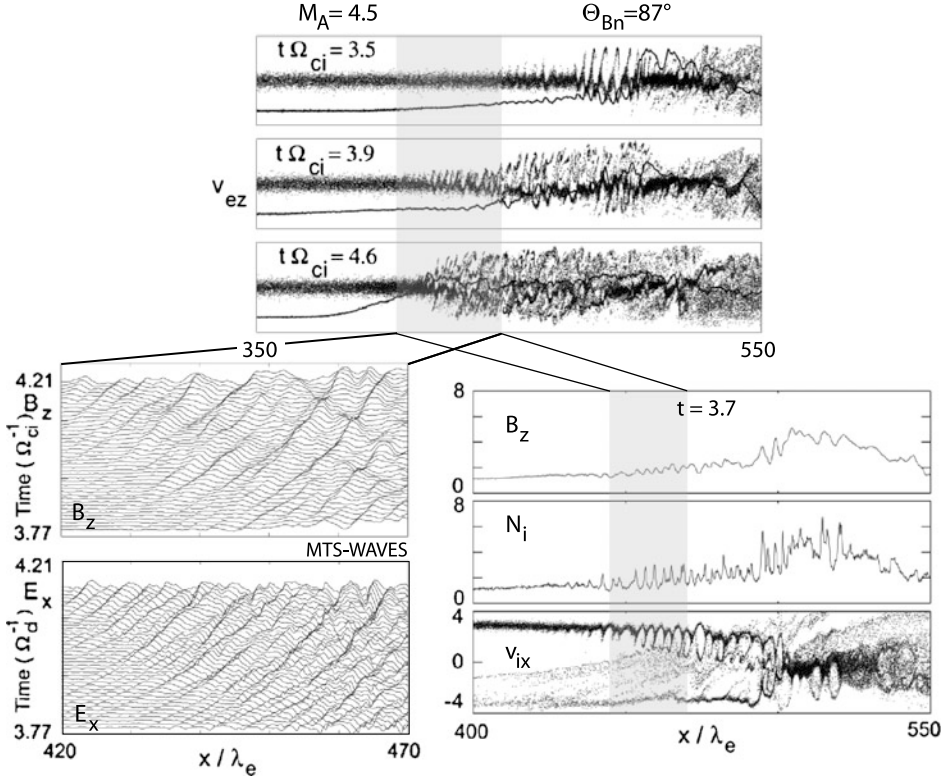


Figure 5.23: *Top:* Electron phase space evolution showing the distortion of the electrons until thermalisation during the modified two-stream instability. The evolution of narrow electron holes can be seen of increasing amplitude (phase space width) the closer one comes to the shock ramp. *Lower left:* Magnetic and electric wave components of MTSI waves present in the *grey shaded area* in space during part of the time shown and are moving toward the shock ramp in the left-hand mode as discussed earlier. These waves steepen when reaching the shock front. Distortion of the ion distribution is the result as shown in *Lower right*. Large amplitude ion holes are formed as well [after Scholer & Matsukiyo, 2004].

addition to the slowly growing ion-ion instabilities discussed earlier). The basic physics of the instability can be readily identified from the time histories of the fields and particles in the middle of Figure 5.25. The first exponential growth phase of the E_x -component for times $\omega_{pet} < 500$ is due to the Electron-Cyclotron-Drift instability (ECDI) which we have omitted in our theoretical analysis in Chapter 3 [cf., also, Muschietti & Lembège, 2006]. This instability is driven by the ion beam when it interacts with obliquely propagating electron-Bernstein waves (electron-cyclotron waves). In fact, this instability, in the present case is nothing else but the Buneman instability (BI) which for the given set-up is initially unstable (as is seen from the bulk velocity difference between the ion and electron phase space distributions on the left of the figure) due to the interaction of the ion beam mode with the lowest order electron-cyclotron mode. Initially there is some growth also in the

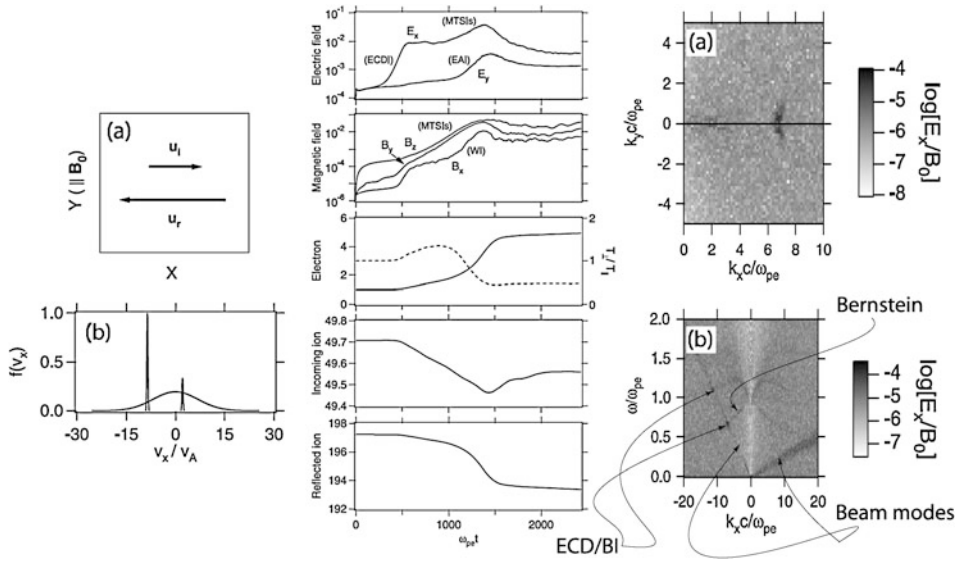


Figure 5.24: Left: The phase space distribution set-up for the simulation. The original magnetic field is in z -direction. The upper panel shows the incoming and reflected ion beams. The lower panel shows the two cold ion distributions, incoming and reflected, and the hot electron distribution, shifted slightly in order to satisfy the zero-current condition in shock-normal direction. Centre: Time histories of the energy densities of the simulation quantities: electric and magnetic wave fields, electrons and the two ion components. Right: Wave power spectra in k -space at early times $t\omega_{pe} < 404.8$ showing the excited power in the Buneman mode in the upper panel. The lower panel shows the dispersion relation. The two straight lines correspond to the damped beam modes of the reflected (negative slope) and direct (positive slope) ion beams. The enhanced power in the two dark spots is due to the ECD-instability, which is the Buneman mode which excited under these early conditions in the simulation as the interaction between the reflected ion beam mode and the first and second Bernstein harmonic waves [after Matsukiyo & Scholer, 2006a, courtesy American Geophysical Union].

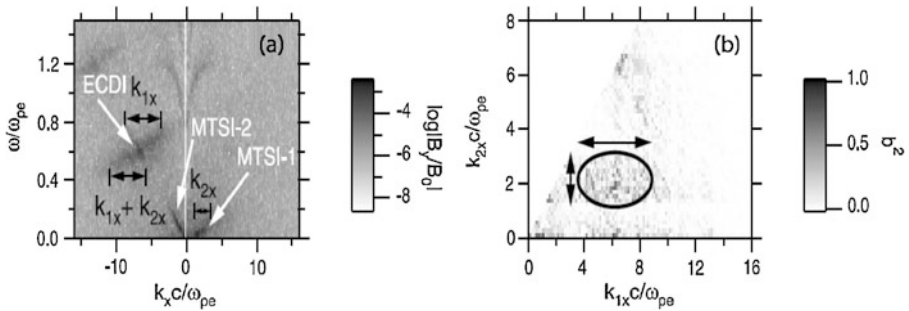


Figure 5.25: Top: The dispersion relation for the time interval $607.2 < t\omega_{pe} < 1011.9$ showing the ECDI (Buneman mode), the original MTSI-1 and the secondary MTSI-2 which is generated by wave-wave interaction. The corresponding reaction in k_x numbers is indicated for the waves which participate in the three-wave process. Bottom: The power spectral density in the (k_x, k_y) -plane. The ellipse indicates the wave numbers that contribute to the wave-wave interaction of the MTSI-1 and ECDI [after Matsukiyo & Scholer, 2006a, courtesy American Geophysical Union].

magnetic field which is strongest in B_z and much weaker in B_y and B_x . However, until the MTSI sets on the magnetic field energy does not grow substantially. This changes with onset of the MTSI when all components increase with B_y , B_z dominating and being of equal intensity, showing that due to the magnetic wave field of the MTSI the instantaneous magnetic field develops a transverse component.

The MTSI does not grow in this initial state because its growth rate is small for these conditions. During the saturation phase the ECDI still dominates in the flat regime until the MTSI takes over causing further growth of the already large amplitude electric field fluctuations. This stage after $\omega_{pet} > 10^3$ is characterised by a growth phase also in E_y (which is due to the electron acoustic instability EAI which can be excited in presence of both a cold and a hot electron component) and, surprisingly, the normal component B_x .

This latter component might be caused by the Weibel instability (WI) when a substantial anisotropy is generated. Such an anisotropy exists for the ions, in fact, in our case as they propagate solely in $\pm x$ -direction at grossly different speeds. The growth rate of the instability, neglecting the magnetic field, i.e. setting $\omega_{ce} = 0$, is $\gamma_{WI} = (V_i/\lambda_i)(1 + 1/k^2\lambda_e^2)^{-\frac{1}{2}}$ [Weibel, 1959], where $\lambda_{e,i} = c/\omega_{pe,i}$ are the ion and electron inertial lengths, respectively. When the magnetic field is not neglected but the ions are taken as non-magnetised, as is the case in the shock foot, then

$$\gamma_{WI} = \frac{V_i}{\lambda_i} \left(1 + \frac{1}{k^2\lambda_e^2}\right)^{-\frac{1}{2}} \quad \text{for} \quad \left(1 + \frac{1}{k^2\lambda_e^2}\right) \frac{\omega_{ce}^2}{\omega_{pi}^2} \gg \frac{V_i^2}{c^2} \left(1 + \frac{1}{k^2\lambda_e^2}\right) \quad (5.22)$$

At short wavelengths the growth rate of this instability can be quite large. Its maximum is assumed for $k^2 \rightarrow \infty$ when it becomes the order of $(\gamma_{WI}/\omega_{ci})_{\max} \sim V_i/V_A \simeq \mathcal{M}_A$. At the expected wave-number $k\lambda_i \sim 1$ is just a factor $\sqrt{2}$ smaller than its maximum value and decreases rapidly towards longer wavelengths.

One may thus expect that large Mach number shocks generate magnetic fields by the Weibel instability, in which case the field becomes non-coplanar, and small-scale stationary magnetic structures appear in the shock foot and ramp. Still, this is a little speculative. However, if the Weibel instability exists it will generate many small-scale magnetic structures in the shock. This is, in itself, sufficiently interesting to be noted. The simulations show the presence of $B_x \neq 0$, suggesting that the magnetic field becomes three-dimensional since the Weibel instability has zero frequency and thus produces a steady normal field component. At very high Mach numbers the Weibel instability will help reflecting a much larger fraction of ions thereby contributing to sustaining the shock and increasing the (viscous) dissipation rate. This is of substantial interest in astrophysical applications.

The right outermost part of the figure shows two power spectra of the electric field in (k_x, k_y) -space at times $t\omega_{pe} = 253$ (top) and $t\omega_{pe} < 404.8$ (bottom). In the top panel the power of the waves concentrates at $(k_x\lambda_e, k_y\lambda_e) = (6.8, 0)$. These waves propagate nearly perpendicular to the ambient magnetic field.

The lower panel shows the dispersion relation $\omega(k_x)$ for these waves in a grey scale representation. The two straight dark lines with negative and positive slopes belong to the damped ion beam modes for the reflected (negative slope) and incoming (positive slope) ion beams.

There are two dark specks on the reflected beam mode where the intensity of the electric field (which is shown here only) is enhanced. These specks are separated by about the electron cyclotron frequency in frequency. They belong to the crossings of the reflected ion beam mode and the two lowest harmonics of the electron Bernstein modes which is the ECD-instability which in this case is also the Buneman instability (BI). This mode has been investigated by *Muschiatti & Lembège* [2006] in one-dimensional PIC simulations and has been shown to be present in the foot region. Since we now know that it is the Buneman mode, it is no surprise to find it in the early stage here in two dimensions, when the conditions are favourable for the Buneman mode and the initial situation is still close to one-dimensional. It is, however, important to note that the Buneman two-stream instability is excited by the large difference in bulk speeds between electrons and reflected ions. In the later stages, as the existence of the electromagnetic left-hand polarised negative helicity waves in Figure 5.23 confirms, the ECDI/Buneman mode is replaced by the MTS-instability which generates oblique, nearly perpendicularly propagating large amplitude electromagnetic waves which also form hole structures and heat the plasma.

Figure 5.25 shows the next time slot in the presentation of the dispersion relation (left). At this time the waves have reached large amplitudes, large enough to cause various interactions among the waves which react on the wave and particle distributions and, in addition, cause nonlinearity of the plasma state at wave saturation. The ECDI forms as a broad spot in the (ω, k) -domain. The MTSI is the short nearly straight line at low frequencies and small positive k_x (indicated as MTSI-1 in the figure). In the wave spectrogram these waves move towards the shock ramp. This means their slope is positive in the dispersiogram!

Secondary Modified Two-Stream Instability: Wave-Wave Interaction. In addition to these modes another negatively moving low frequency wave appears. This is also an MTSI, but it is a secondary one, which *Matsukiyo & Scholer* [2006a] have shown to arise in a three wave process when the ECDI-BI and the MTSI-1 interact causing a wave with $k_x = k_{\text{BI}} + k_{\text{MTSI-1}} = k_{\text{MTSI-2}}$. The right part of the figure shows the enhanced wave power for this process extracted from the data on the way of a bi-spectral analysis and represented in the k -plane. The ellipse encircles the wavenumbers which are involved into the three-wave interaction, the original ECD-wave, and the resulting MTS-2-wave. Clearly, a whole range of waves participates in the interaction because the ECD-spectrum has broadened when saturating, and many combinations of ECD and MTS-1 waves satisfy the nonlinear three-wave interaction condition.

The top-left panel of Figure 5.26 shows the evolution of the electron velocity v_y during this interval and averaged over a range of x -values along the normal. This velocity is about perpendicular to the magnetic field; its dynamic range of variation is impressive. The panel at the lower left shows the electron phase space distribution. Two electron beams are seen to propagate at counter streaming velocities. These beams can already be identified in the upper panel. Due to the interaction with the unstable waves the region between the beams is partially filled. These distributions have been taken in the interval $17.6 < y/\lambda_e < 19$ as indicated in the top panel.

Another distribution a little further in the interval $39.1 < y/\lambda_e < 40.5$, at the rear end of the top panel, is shown in the lower right panel. Here the distribution has evolved

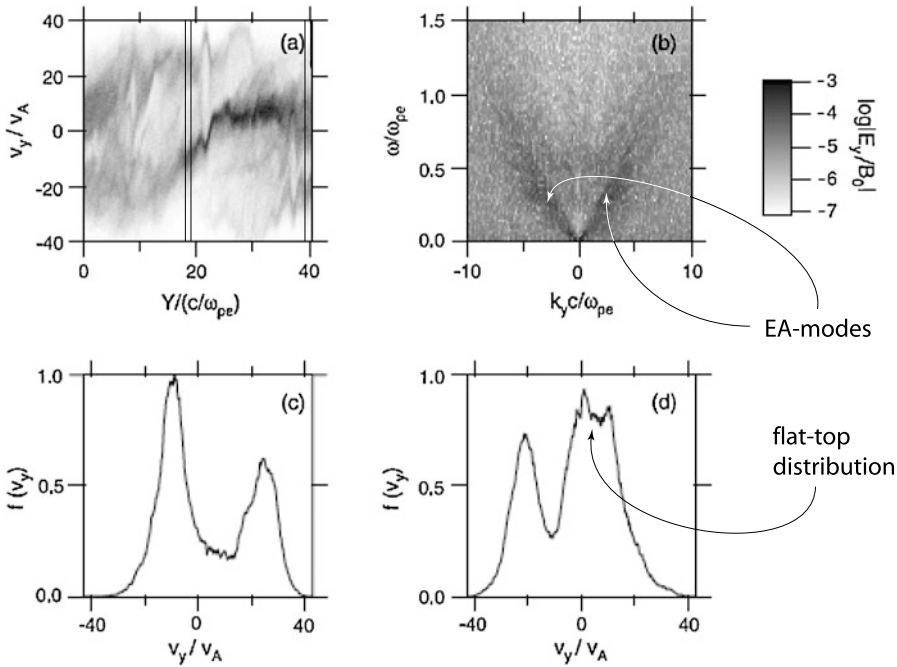


Figure 5.26: *Top:* Electron phase space plot (*left*) at $19.4 < x/\lambda_e < 21$, and (*right*) dispersion relation $\omega(k_y)$ for the period $910.7 < t\omega_{pe} < 1315.5$ as obtained from E_y . This dispersion relation shows the occurrence of EA-waves with strictly linear dispersion and frequency below $\omega_{ea} \lesssim \omega_{pe}$ while propagating in both directions. These are generated in the presence of the two-electron component structure seen in the distribution function below. They are responsible for the subtle fine-structuring of the electron distribution in the phase space as is seen in the *upper left panel* representation of v_y versus y which exhibits trapping and scattering of electrons on very small scales. *Bottom:* Two electron distribution functions at $17.6 < y/\lambda_e < 19$ and $39.1 < y/\lambda_e < 40.5$ at $t = 1000$, as indicated in the *upper left panel* by the vertical lines, showing the large electron hole distribution that is generated by the MTSI and some smaller substructures [after Matsukiyo & Scholer, 2006a, courtesy American Geophysical Union].

into a totally different combination of two electron populations, one top-flat and hot, the other one narrow, i.e. cold, but of same height indicating the retardation of one beam and heating of the other. Altogether the electron plasma has been heated to high temperature. Returning to the upper left panel the complicated structure of the distribution is nicely seen with several sub-beams evolving and also with electron trapping in some vortices being visible for instance in the upper part around $y \simeq 22\lambda_e$.

Coming now to the upper right panel, which shows the dispersion relation in the y -direction, one recognises two low-frequency linear wave modes propagating in positive and negative y -directions. These waves are electron-acoustic (EA) modes which are excited in the presence of the two electron distributions, the hot top-flat distribution and the cold beam distribution. They have strictly linear dispersion and frequencies below the electron plasma frequency $\omega_{ea} \lesssim \omega_{pe}$. Because they interact strongly with the electron distribution, they are responsible for the fine-structuring in the electron distribution function

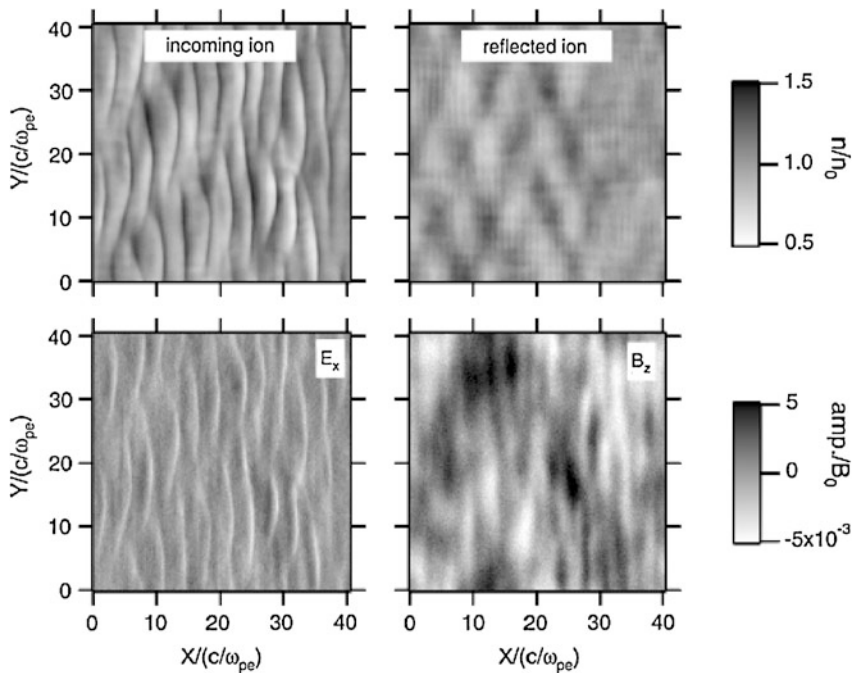


Figure 5.27: *Top:* Incoming (*left*) and reflected (*right*) ion densities at the late time $t\omega_{pe} = 2023.9$ shown in the (x, y) -plane. This time, in terms of the ion cyclotron period, corresponds to $t\omega_{ci} = 0.55$, i.e. about half an ion-cyclotron period. At later times the ion magnetisation would come into play as well. *Bottom:* The corresponding electric field E_x and magnetic field B_z profiles. One observes that the ECD-waves (Bernstein modes) have decayed by feeding their energy into electron heating. The two MTS-modes are still visible as the wavy variations in the incoming and reflected ion beams. The original MTS-1 wave modulates the incoming beam, which is seen in its downward propagation towards the shock. The secondary MTS-2 mode modulates, in addition, the reflected ion beam causing the interference pattern seen in the reflected beam density. The electric field is modulated by the MTS-2 wave, while the magnetic field contains signatures of both, MTS-1 and MTS-2 [after *Matsukiyo & Scholer, 2006a*, courtesy American Geophysical Union].

seen in the top-left panel of Figure 5.26 where they cause electron trapping and scattering which results in electron heating and electron acceleration.

We close this section by presenting ion densities of the reflected and incoming beams and the corresponding modulations of the electric E_x and magnetic variation fields B_z , respectively, in Figure 5.27 in two-dimensional grey-scale representation in the (x, y) -plane. One observes that the ECD-waves (Bernstein and Buneman-modes) have decayed. They have been feeding their energy into electron heating, creating electron holes, trapping electrons and shaking them, as we have discussed above. The two MTS-modes are still visible. They dominate the ion density structure being visible as the wavy variations in the incoming and reflected ion beams. The original MTS-1 wave only modulates the incoming beam. This is recognised from its long-wavelength downward propagation towards the shock. The secondary shorter scale MTS-2 mode modulates, in addition, the reflected ion

beam causing the interference pattern seen in the reflected ion-beam density. The electric field is merely modulated by the short wavelength MTS-2 wave.

Weibel Instability Caused Effects. On the other hand, the magnetic field contains signatures of both, MTS-1 and MTS-2 thus exhibiting a more irregular structure than the electric field. Here, probably, also the small-scale structures of the Weibel instability do contribute. We have noted its effect already above, but it would be rather difficult to extract them from the figure as they should appear as stationary vortices, which are convected downstream towards the shock front with the speed of the average bulk flow. Their dynamics remains to be unresolved, i.e. it is not clear what will happen to them when encountering the shock front. One possibility would be that they accumulate there and generate a non-coplanar magnetic component. Nevertheless, the possibility for the Weibel instability to evolve in supercritical quasi-perpendicular shocks is of interest as Weibel vortices could, if confirmed by observations, cause an irregular fine structuring of the magnetic field in the shock ramp transition, which would have consequences for the particle dynamics, trapping, scattering, reflecting and acceleration of particles from the shock front. It could, moreover, also lead to small scale reconnection in the shock front, which so far has not been believed to exist in the shock, including the various side-effects of reconnection. Weibel vortices could also pass into the downstream region where they might contribute to the downstream magnetic turbulence. There they would occur as magnetic nulls or holes for which the shock would be the source.

5.5 The Problem of Stationarity

In this last section of the present chapter we will be dealing with the time-dependence of quasi-perpendicular shocks. Since in the previous sections we have frequently dealt with time variations, there is little new about time-dependencies of shocks. Nevertheless, in the past few years there has been much ado about the so-called “problem” of shock non-stationarity which has grown into an own field of shock research. We have already spoken on this before.

What does it mean that shocks can be or even are non-stationary? In principle, stationarity means time independence and, hence, is a question of the time scale under consideration. For instance, any cascade looked at from far is stationary. Such are shocks. As for an example take Earth’s bow shock which stands in front of Earth for eons, or astrophysical shocks which for the human eye are practically invariable. On the time scale of their existence they are stationary, while on much shorter time scales they undergo global and local variations. Non-stationarity, at its best, thus just means nothing else but dealing with such variations. However, at a physically much deeper level would be the question for the *equilibrium state* of shocks. Clearly, as already explicated, *shocks are not in thermodynamic equilibrium*. They need to be driven. Thus the time scale of the global “stationarity” of shocks is the life time of the driver.

One kind of non-stationarity is shock reformation. This is a periodic or better quasi-periodic process in which the shock ramp for the time of foot formation remains about

stationary, i.e. the shock ramp is well defined and moves ahead only very slowly. In fact, during the following reformation the shock is everything else but stationary: from one reformation cycle to the next the ramp flaps, it flattens and broadens while the new shock foot grows and steepens. And towards the end of the reformation cycle the shock ramp suddenly jumps ahead from its old position to the edge position of the shock foot. Could one follow this evolution over very many reformation cycles, one would find that there is no real periodicity but that the process of reformation, i.e. the time sequence of final forward jumps of the shock ramp would form a quasi-periodic or even chaotic time series. Unfortunately, computer capacities do not yet allow to simulate more than a few reformation cycles such that this conjecture cannot be proved yet. But it is simple logic that reformation cannot be strictly cyclic; there are too many processes involved into it, too many instabilities cooperate, and the particle dynamics is too complicated for a strictly periodic process to be maintained over longer times than one or two ion-gyroperiods. In addition, once the shock is considered two-dimensional – or even three-dimensional – the additional degrees of freedom introduced by the higher dimensions multiply and the probability for the shock of becoming a stationary or even cyclic entity decreases rapidly. This is particularly true for high-Mach number shocks even under non-relativistic conditions. Hence, we may expect that a realistic high-Mach number, i.e. supercritical shock will necessarily be locally non-stationary.

Principally, stationarity is a question of scales. On the macroscopic scale, the scale of the macroscopic obstacle and the macroscopic flow, a shock will be stationary as long as the flow and the obstacle are stationary. For instance, such stationary shocks are the planetary bow shocks that stand in front of the planetary magnetospheres or ionospheres. On the scales of the magnetospheric diameters their variation is of the same order as the variation of the solar wind or – if the magnetospheres themselves behave dynamically – the time and spatial scales of their variations are of the same order as the time and spatial scales of the magnetospheric variations, for instance the diurnal precession of the Earth's magnetic axis which causes a strictly periodic flapping of the magnetosphere and thus a strictly periodic variation of the position and shape of the Earth's bow shock.

On spatial scales of the order of the ion gyroradius and temporal scales of the order of the ion gyroperiod there is little reason to believe in stationarity of a collisionless supercritical shock wave. The whole problem of stationarity reduces to the investigation of instabilities and their different spatial and temporal scales and ranges, their evolution, saturation, being the sources of wave-wave interactions and nonlinear wave-particle interactions and so on. These we have already discussed in the former sections as far as the current state of the investigations do allow. What, thus, remains is to ask how a shock surface can become modulated in higher dimensions and what reasons can be given for such modulations.

5.5.1 Theoretical Reasons for Shocks Being “Non-stationary”

That collisionless shock waves might exhibit non-stationary behaviour was suggested early on from the first laboratory experiments on collisionless shocks [Auer *et al*, 1962; Paul *et al*, 1967]. Morse *et al* [1972] were the first to definitively conclude from their one-dimensional shock full particle PIC simulations that collisionless shocks seem to be non-

stationary on the scale of the ion gyroperiod. Afterwards, time variations in the behaviour and evolution of collisionless shocks have been recovered permanently in shock simulations [e.g., *Lembège & Dawson, 1987a, b*].

This is no surprise as we have mentioned several times already. The *principal reason* is that shocks, and in particular supercritical shocks which are not balanced by collisional dissipation, are in thermal non-equilibrium and are thermodynamically not in balance. Hence, locally they are longing for any opportunity to escape this physically unpleasant situation in order to achieve balance and thermal equilibration. However, as simple as this reason might look, as difficult is it to find out what under certain given conditions will actually happen and which way a shock will locally direct itself for a try to escape non-equilibrium and to achieve equilibrium. Even though when it is permanently driven by an unchangeable flow and a stationary obstacle it will chose any kind of irregularity, fluctuation or detuning to drive some kind of instability, cause dissipation and, when driving will become too hard in any sense, it will overturn and break; and in this way it will maximise dissipation if it is not possible to achieve it in any smoother way.

Non-stationary behaviour of quasi-perpendicular shocks has been anticipated theoretically, following *Sagdeev [1966]*, by *Kennel et al [1985]* who noted the existence of the critical whistler Mach number $\mathcal{M}_{wh}$, which we have discussed above in comparison to numerical simulations. *Galeev et al [1988]* tried to give a theoretical account for reasons of the anticipated non-stationary character of supercritical shocks. They investigated the role of whistlers in the nonlinear domain at the ramp, finding that whistlers for flow speeds sufficiently above the Alfvén speed do not possess soliton solutions and thus do not sustain the steady state of a shock. This means very simply that neither dissipation nor dispersion can sustain the nonlinear steeping of the waves, and therefore the waves should cause breaking of the flow and lead to non-stationary behaviour of the ramp and crest, a process called by them ‘gradient catastrophe’. These authors also dealt with quasi-electrostatic waves of frequencies close to the lower-hybrid frequency ω_{lh} to which they attributed responsibility for wave breaking.

Simulations by *Quest [1986]*, *Lembège & Savoini [1992]*, *Savoini & Lembège [1994]* and *Hada et al [2003]* for low mass-ratios have attempted to illuminate some aspects of this non-stationary behaviour. *Lowe & Burgess [2003]* and *Burgess [2006a, b]* have investigated two-dimensional rippling of the shock surface in hybrid simulations and its consequences. Full particle simulations up to realistic mass ratios have been performed by *Scholer & Matsukiyo [2004]*, *Matsukiyo & Scholer [2006a, b]* and *Scholer & Burgess [2007]*. We will return to these attempts. Here we first follow the analytical and simulation attempts of *Krasnoselskikh et al [2002]* to advertise the general non-stationarity of quasi-perpendicular shocks. We should, however, note that there is no principle reason for a shock to behave like we wish, i.e. to behave stationary. It might, if necessary, break and overturn or mike not; the only requirement being that it follows the laws of physics.

Nonlinear Whistler Mediated Non-stationarity

Krasnoselskikh et al [2002] rely on a method developed by *Whitham [1974]* to describe the nonlinear breaking of simple waves by adding to the simple wave evolution equation

a nonlocal term that takes care of accumulating short wavelength waves. The Whitham equation reads

$$\frac{\partial v}{\partial t} + v \nabla_x v + \int_{-\infty}^{\infty} d\xi K(x - \xi) \nabla_{\xi} v(t, \xi) = 0, \quad K(x) = \frac{1}{2\pi} \int_{-\infty}^{\infty} dk \frac{\omega(k)}{k} e^{ikx} \quad (5.23)$$

If this additional term is purely dispersive, it reproduces the Korteweg-de Vries equation, if it is dissipative it reproduces Burgers' equation. In general, stationary solutions $\partial v / \partial t \rightarrow 0$ peak for $K(x) \sim |x|^{-\alpha}$ for $x \rightarrow 0$, and $\alpha > 0$. *Krasnoselskikh et al [2002]* use the whistler dispersion relation

$$\frac{\omega}{kV_A} = \left(\frac{m_i}{m_e} \right)^{\frac{1}{2}} \frac{|\cos \theta| k \lambda_e}{1 + k^2 \lambda_e^2}, \quad \text{for } k \lambda_i \gg 1, \cos^2 \theta \gg \frac{m_e}{m_i} \quad (5.24)$$

describing low frequency whistlers at oblique propagation which, when inserted into the above integral for $K(x)$ asymptotically for $|x| \rightarrow 0$ yields $K(x) \sim \pi^{-1} \sqrt{m_i/m_e} |\cos \theta| [C + \ln|x| + \dots]$. Here $C = 0.577\dots$ is Euler's constant. Since $|x|^\alpha \ln|x| \rightarrow 0$ for all positive $\alpha > 0$ and $|x| \rightarrow 0$, nonlinear low-frequency whistler waves will necessarily break by the above condition. Thus, when whistlers are involved into shock steepening, and when $\alpha > 0$, they will necessarily break as their dispersion does not balance the nonlinear steepening. This happens when the Mach number exceeds the nonlinear whistler Mach number $\mathcal{M} > \mathcal{M}_{\text{wh,nl}}$. Both, whistler dispersion and dissipation by reflected ions cannot stop the whistlers from growing and steepening anymore, then.

In order to prevent breaking, another mechanism of dissipation is required. Still based on the whistler assumption, *Krasnoselskikh et al [2002]* argue that the shock ramp would radiate small wavelength whistler trains upstream as a new dissipation mechanism. This works, however, only as long as the Mach number remains to be smaller than another critical Mach number $\mathcal{M} < \mathcal{M}_{\text{wh,g}}$ that is based on the whistler group velocity, $\partial \omega / \partial k$,

$$\mathcal{M}_{\text{wh,g}} = \left(\frac{27 m_i}{64 m_e} \right)^{\frac{1}{2}} |\cos \Theta_{Bn}| \lesssim 19.8, \quad (5.25)$$

since for larger $\mathcal{M}$ the whistler-wave energy will be confined to the shock and cannot propagate upstream. The right-hand estimate holds for an electron-proton plasma and $\cos \Theta_{Bn} \sim 0.707$, i.e. at the largest shock-normal angle $\Theta_{Bn} = 45^\circ$ of quasi-perpendicular shocks.

In other words, since the nonlinear whistler Mach number is always larger than the whistler-group Mach number, whistler energy will leave the shock upstream only in a narrow Mach number range $\mathcal{M}_{\text{wh}} < \mathcal{M} < \mathcal{M}_{\text{wh,g}} < \mathcal{M}_{\text{wh,nl}}$. This range corresponds to $15 < \mathcal{M} / |\cos \Theta_{Bn}| < 19.8 < 21.3$. One-dimensional full particle PIC simulations with realistic mass ratio $m_i/m_e = 1840$ have been performed by *Scholer & Matsukeyo [2004]* and *Scholer & Burgess [2007]* and confirm that whistlers affect the stationary or non-stationary behaviour of nearly perpendicular shocks.

At larger $\mathcal{M} > 19.8 |\cos \Theta_{Bn}|$ the whistler energy is again confined to the shock and will be swept downstream towards the shock when transported by the passing though continuously retarded flow. In the region of the foot and ramp where the energy accumulates it

will cause different instabilities some of which propagate downstream. Such processes can be wave-wave interactions driven by the high whistler energy, as had already been envisaged by *Sagdeev* [1966], or nonlinear wave-particle interactions. In addition to causing anomalous resistivity and anomalous dissipation, these processes should lead to emission of plasma waves from the shock, preferably into the direction downstream of the shock, as only there $\mathcal{M} \lesssim 1$, and the wave group and phase velocities can exceed the speed of the flow.

This whole discussion refers only to whistler waves and follows the traditional route. We have, however, seen in the previous sections that the foot of the shock is capable of generating waves of another kind, electromagnetic Buneman modes, modified-two-stream waves, and possibly even Weibel modes. These waves are highly productive in generation of electron heating; they cause magnetic disturbances that move towards the shock or also upstream. The role in dissipation and dispersion these waves play has not yet been clarified and is subject to further investigation. It is, however, clear that their excitation and presence in the shock foot produces electron heating, retarding of ions and ion heating as well and will thus provide an efficient dissipation mechanism. Whether this can prevent shock breaking and overturning at very large Mach numbers is not known yet.

Shock Variability Due to Two-Stream and Modified Two-Stream Waves

Variability of the quasi-perpendicular shock has been demonstrated from numerical full particle PIC simulations in one and two dimensions to come about quite naturally for a wide range of – sufficiently large – Mach numbers. While in the low Mach number range whistlers are involved in the variability, reformation and non-stationarity, the simulations have clearly demonstrated that at higher Mach numbers the responsible waves are the Buneman two-stream mode and the modified two-stream instability. This has been checked [cf., *Matsukiyo & Scholer, 2006a, b*] by shock-independent simulations where the typical electron and ion phase space distributions have been used which occur in the vicinity of supercritical shocks during particle reflection events. So far the importance of these waves over whistlers has been investigated only for the shock-foot region. The shock ramp and overshoot are more difficult to model because of the presence of density and field gradients, their electrical non-neutrality, and the fuzziness of the particle phase-space distribution functions.

Differences were also found between strictly perpendicular and quasi-perpendicular shocks. The former are much stronger subject to the Buneman two-stream instability that completely rules the reformation process in one and two dimensions in this case, causing phase-space holes to evolve and being responsible for quasi-periodic changes in the positions, heights and widths of the shock front and foot regions, respectively. We have already put forward arguments that an investigation of the long-term behaviour of this quasi-periodic variation should reveal that this process is irregular in a statistical sense. Even under apparently periodic reformation conditions the shock will presumably not behave stationary on the short time and spatial scales. This, however, can be checked only with the help of long-term simulations which so far are inhibited if done with sufficiently many macro-particles, realistic mass ratios and in more than one-dimension.

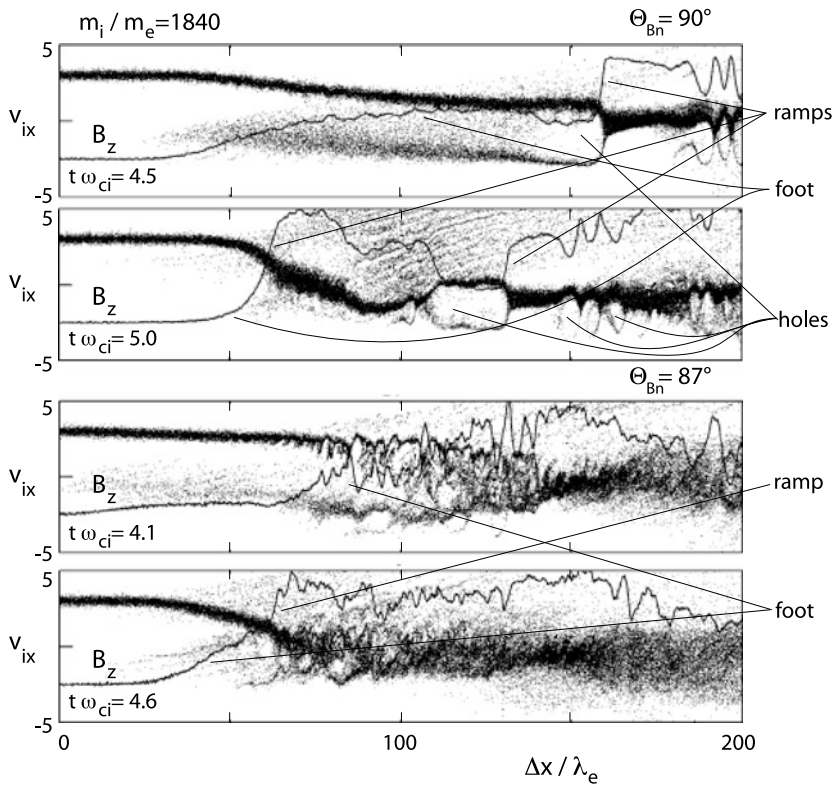


Figure 5.28: The completely different reformation behaviour of shocks in one-dimensional PIC simulations with realistic mass ratio of 1840 for strictly perpendicular and oblique quasi-perpendicular shocks at exactly same parameter settings and scales. Shown is the magnetic field B_z and the ion phase space at two subsequent simulation times for each of the respective simulations. Since the evolution is different in both cases the x -coordinate is given as a relative scale not in x but for the same interval lengths in Δx for the instance when reformation takes place in both cases [compiled from *Matsukiyo & Scholer, 2006a*, courtesy American Geophysical Union]. *Top:* Reformation at $\Theta_{Bn} = 90^\circ$ at two times showing the evolution of the foot in the magnetic field and the taking-over of the ramp by the foot while a new foot evolves. This process is governed by the Buneman two-stream instability. Large holes evolve on the ion distribution. Note the correlation of the ion holes with depressions in the magnetic field. In the *second panel* the old ramp is still visible as the boundary of the large ion hole. Farther downstream many holes are seen, each of them corresponding to a magnetic depression, and the regions between characterised by magnetic overshoots. *Bottom:* The corresponding evolution at $\Theta_{Bn} = 87^\circ$. High variability of the shock profile is observed which is identified as being due to the large amplitude MTS-waves travelling into the shock. The foot region is extended and very noisy both in the magnetic field and ion distributions, the latter being highly structured. The foot is extended much longer than in the perpendicular case. The *two bottom panels* might also show signatures of wave breaking in the ion velocities when groups of ions appear which overturn the main flow in forward downstream direction.

Figure 5.28 provides an impression of the variability of shock reformation in the two cases of a strictly perpendicular, highly supercritical shock, and the case of an oblique supercritical shock at $\Theta_{Bn} = 87^\circ$, when whistler excitation is absent. The settings of the

simulations are otherwise identical, but the evolution of the two simulations is completely unrelated. This is because the perpendicular shock does not allow, in these one-dimensional simulations, for the modified two-stream instability to grow. So only the Buneman two-stream instability grows. It reforms the foot in the way we have already described, forming large holes and letting the shock ramp jump ahead in time-steps of the order of roughly an ion gyro-period. The shock foot acts decelerating on the flow, and already during reformation begins to reflect ions and to form a new foot. Most interestingly is that the holes survive quite a while downstream while being all the time related to magnetic depressions. At their boundaries large magnetic walls form which can be interpreted as magnetic compressions (or otherwise signatures of current vortices).

The oblique case looks different. It is highly variable both in time and in space. The magnetic profile is more irregular, and the ion-phase space exhibits much more structure than in the perpendicular case. This has been identified to be due to the combined action of the Buneman two-stream and the modified two-stream instabilities with the two-stream instability being important only during the initial state of the reformation process, while the modified two-stream instability dominates the later nonlinear evolution. Both, foot and ramp, are extended and vary strongly. It is quite obvious, that in this case one can speak of a stationary shock front only when referring to the long-term behaviour of the shock, much longer than the irregular reformation cycle lasts. At the scale of reformation and below there is no stationarity but variability and evolution, which can be attributed to the growth and interaction times of the MTSI and the various secondary processes caused by it.

To complete this section, we note in passing that the low-mass ratio two-dimensional full particle PIC simulations with small particle numbers performed by *Lembège & Savoini* [1992, 2002] and *Savoini & Lembège* [1994, 2001] also showed non-stationary behaviour of the quasi-perpendicular – or perpendicular – shock leading to so-called “rippling” of shocks, which we will briefly describe in the next paragraph.

5.5.2 Formation of Ripples

One-dimensional theory and one-dimensional simulations implicitly treat the shock as an infinitely extended plane surface. In addition they allow only for instabilities to evolve in the direction of the shock normal at an angle relatively close to 90° such that any waves along the shock surface are completely excluded and waves parallel to the magnetic field have very small wave numbers $k_{\parallel} = k_x \cos \Theta_{Bn} \ll k_x$ corresponding to very long parallel wavelengths. To be more realistic, two-dimensional PIC simulations have been performed to investigate the effect of the additional freedom given by the second spatial dimension which allows instabilities to evolve in other than the shock normal direction. The cost of these simulations is being restricted to low mass ratios only. In the simulation of *Lembège & Savoini* [1992, 2002] and *Savoini & Lembège* [1994, 2001] the mass ratio has been taken as $m_i/m_e = 42$ which implies from comparison to the high-mass ratios in one-dimensional simulations that the modified two-stream instability will presumably be excluded. The structure of the shock front in these simulations has been shown in Figure 5.8, the left-hand side of which shows the cyclic reformation of the shock – which at these low mass ratios is clearly expected to occur – at a period comparable to the ion-gyro period of the reflected

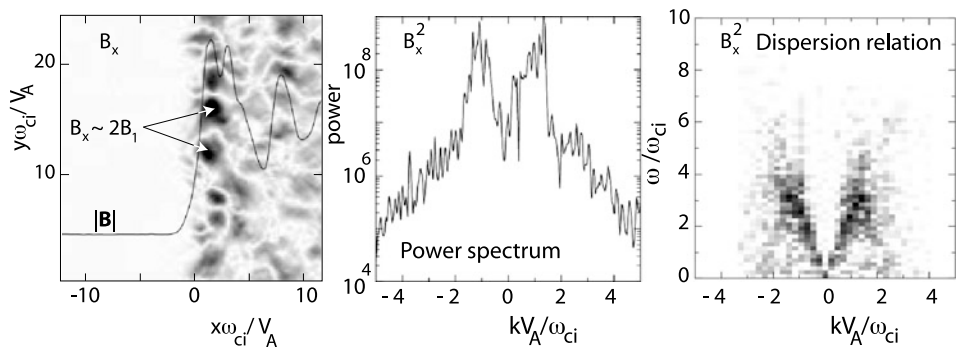


Figure 5.29: *Left:* Spatial distribution of the $B_n = B_x$ -component of the magnetic field in the hybrid simulations of *Lowe & Burgess [2003]* with Mach number $\mathcal{M}_A = 5.7$, and for $\Theta_{Bn} = 88^\circ$. The B_x -component is not zero; it reaches values twice the upstream magnetic field B_1 and shows quite structured behaviour along the shock surface which indicates that the surface is oscillating back and forth and that waves are running along the surface. These waves are interpreted as surface waves. *Centre:* The power in the presumable surface waves as determined from the simulations. Obviously the power concentrates around the ramp. *Right:* Apparent dispersion relation $\omega(k)$ of the fluctuations [after *Lowe & Burgess, 2003*].

ions in the foot of the shock. The right-hand side shows a pseudo-three-dimensional profile of the shock in the two spacial dimensions, in the top part at the time when the foot is fully developed, in the bottom part when the ramp has just reformed, i.e. when the foot has taken over to become the ramp. What interests us here is that the shock ramp surface is by no means a smooth plane in the direction tangential to the shock. It exhibits large variations both in space and time which are correlated but not directly in phase with the presence of reflected ions in the foot. The overshoot, steepness and width of the ramp and ramp position oscillate at a not strictly periodical time-scale. In addition, the structure of the ramp also exhibits shorter scale fluctuations.

Further hybrid simulations [*Lowe & Burgess, 2003*; *Burgess & Scholer, 2007*] in two dimensions at $\Theta_{Bn} = 88^\circ$ for and $\mathcal{M}_A = 5.7$ satisfying Rankine-Hugoniot conditions with $\mathbf{n} \cdot \mathbf{B} = B_n$ constant, attribute fluctuations in B_n to these so-called ripples [*Lembège & Savoini, 2002*; *Lowe & Burgess, 2003*; *Burgess, 2006a, b*] as given in Figure 5.29 for the $B_n = B_x$ (left), power B_x^2 (centre), and the ‘dispersion relation’ $\omega(k)$ of the fluctuations in B_x . These fluctuations are of the same value as the main component of the magnetic field B_z reaching maximum values of twice the upstream magnetic field B_1 . They are concentrated in the ramp, foot and overshoot. The dispersion relation is about linear and low frequency but exceeds the ion-cyclotron frequency for shorter wavelengths. There is no mode known which corresponds to these waves, so they are attributed to surface waves flowing in the shock transition. Maximum wavelengths are a few ion inertial lengths.

The lesson learned is, however, quite simple: the shock exhibits structure along its surface which can presumably be attributed to waves running along the shock front and modulating it temporarily and spatially. The caveat of these simulations is however, their hybrid character which does not account for the full dynamics of the particles and therefore it cannot be concluded about the nature of the waves. *Burgess & Scholer [2007]* have

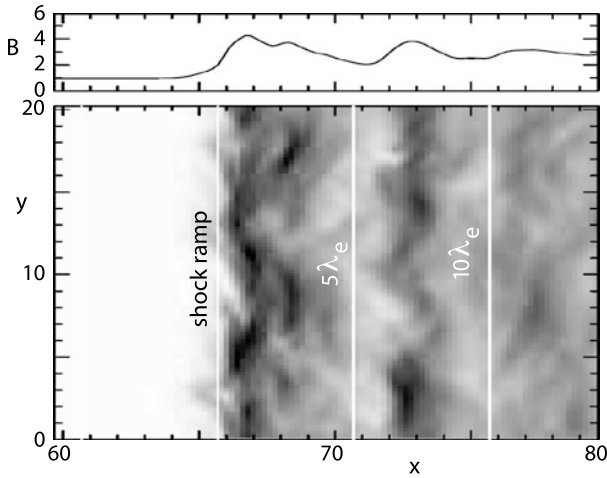


Figure 5.30: Two-dimensional structure of the surface waves [Burgess & Scholer, 2007]. *Top:* Magnetic field average across the shock. *Bottom:* Grey scale plot of the surface waves. The three white lines show the presumable location of the nominal shock and two distances from it downstream [after Burgess, 2006b, courtesy American Geophysical Union].

extended these investigations to infer about the driver of these waves. They find that it is the reflected ion component in the shock foot which flows along the shock surface and at large Mach numbers becomes unstable. Figure 5.30 shows a grey scale plot of the two-dimensional structure of the surface wave oscillation. Its growth is proportional to the Mach number, i.e. it must therefore be proportional to the number of reflected ions, their velocity and to the upstream convection electric field that accelerates the ions. Presumably it is some variant of Kelvin-Helmholtz instability along the shock surface, which is driven by the velocity shear introduced by the reflected ion flow along the shock surface. It causes undulations (or flow vortices) at the shock which, in the magnetic field, appear as ripples.

It should be clear, however, that long-term full particle simulations must be performed at real mass ratio before any reliable conclusion can be drawn about the existence of surface waves. We have seen that part of these waves is nothing but the exchange between the foot and the ramp during the reformation process. This applies to the long wave part. In addition it is indeed possible that the strong and fast flow of reflected ions along the shock surface could excite a Kelvin-Helmholtz-like instability. This depends on the fraction of ions which the shock is able to reflect, and this is a question of Mach number dependence. If the fraction of reflected ions is small the velocity shear will be too small to drive a Kelvin-Helmholtz instability. Giving a quantitative argument is, unfortunately, impossible. However, it is not clear whether these oscillations are the sole action of the modulation of the shock surface in two or three dimensions. The only conclusion we can safely draw is that the shock surface, even under ideal non-curved and quiet upstream conditions, will at high Mach numbers not remain to be a quiet smooth and stable shock surface but will

exhibit fluctuations in position, structure, overshoot amplitude and width on the scales of the ion inertial length and the ion cyclotron period.

5.6 Summary and Conclusions

Among the collisionless shocks, quasi-perpendicular supercritical shocks are the best investigated. They are also the shocks which for other purposes like particle acceleration to very high energies have naturally been favoured and find wide application in astrophysics where one of the central problems consists in the explanation of the presence of energetic particles (Cosmic Rays). Theory has predicted that quasi-perpendicular shocks reflect ions by some basically unspecified mechanism that for simplicity is assumed to be specular, implying elastic scattering from the shock front. This might be provided for the low energy part of the upstream distribution by the transverse shock potential. These reflected ions form feet which are located upstream but adjacent to the shock ramps. Quasi-perpendicular shocks possess either whistler precursors or trails. Theory also predicts that whistlers could be phase-locked and stand in front of the shock ramp only for a limited range of Alfvénic Mach numbers.

We have reviewed here the theoretically expected shock structure, the relevant scales, the most relevant particle simulations for perpendicular and quasi-perpendicular supercritical shocks, the shock-reformation process and its physics as far as it could be elucidated from one-dimensional and to a certain part also from two-dimensional simulations. The most relevant instabilities generated in the shock foot have been identified as the whistler instability for nearly perpendicular supercritical but low-Mach number shocks, leading to foot formation but not being decisive for feet, as it has turned out that feet in this Mach number and shock-normal angle ranges are produced by accumulation of gyrating ions at the upward edge of the foot.

More important than whistlers have turned out the Buneman and modified two-stream modes, the former dominating shocks at perpendicular angles, the latter growing slowly but dominating at more oblique angles and at later simulation times with the effect of completely restructuring both the shock feet and ramps. Both instabilities generate phase space holes which during reformation survive and are added to the downstream plasma and, in addition, being responsible for low magnetic field values.

Most interestingly, the plasma state just downstream of the shock is, at least to large extent, nothing else but the collection of the old shock ramps which have been left over from former reformation cycles and move relative to the shock frame in the direction downstream of the shock. This is best seen in the simulations when analysing the ion phase space where each of the old ramps can still be identified as an ion clump, the remainder of the former ramp reflection position.

The modified-two-stream instability in addition generates waves which flow into the shock ramp where they contribute to the dynamics of the ramp. Wave-wave interaction and wave particle interaction lead to the generation of secondary waves and to particle heating. Finally, simulations show that the shock front in more than one dimension is not a plane surface but exhibits a strong variability in time and space. This can be explained

as surface waves on the shock front which might be driven by the reflected ion current flow along the surface similar to a Kelvin-Helmholtz instability. This question is still open to investigation. So far the evolution of the shock ramp, its stability and time variation as well as the physics of the region just downstream of the ramp is not yet well explored. It is, however, clear from the available intelligence that any serious investigation must be based on full particle simulations and experimenting with appropriate sets of distribution functions suggested by the simulations in order to investigate the instabilities and interactions between the waves and particles as well as between waves and waves in order to understand the physics. This has to a certain degree already been achieved for the foot region. In the shock ramp and in the strongly disturbed region behind the ramp it is more difficult as the conditions there are less clean and the definition of the responsible distributions is more difficult. Moreover, plasma and field gradients must be taken into account in this region, and the electric charge separation field that is partially responsible for ion reflection cannot be neglected as well. With the further increase of computing capacity and the refinement of the models one expects that within the next decade also the physics of the shock ramp will become more transparent.

An interesting problem is the stability of shocks. Above the critical Mach number they reflect ions and generate ion viscosity that helps dissipating the excess energy. This dissipation goes via the above mentioned instabilities and less on the way of ion viscosity in the classical sense of fluid theory. For even larger Mach numbers these processes will not suffice to stabilise a quasi-perpendicular shock. What will then happen? It has been suggested that strongly nonlinear processes driven by whistlers will set on and lead to non-stationarity of the shock ramp. This might be the case. However, only simulations at high Mach numbers and full mass ratios in large enough systems can answer this question.

We state that the problem of stationarity or cyclic behaviour of the shock is not the problem of the shock; rather it is the problem of our understanding. For the shock nothing else counts than dissipating the excess inflow energy and momentum. If this turns out to be impossible by either anomalous dissipation, shock reflection, foot formation, precursors and early flow retarding, then the shock will not care but will break and turn over as this will be the only way for reducing the scale to microscopic dimensions producing violent heating and energy dissipation.

At the time of writing it remains unclear whether breaking takes place. Magnetic field lines cannot break-off; they kink but remain to be simply connected. Hence, any breaking that is going on takes place in the particle component requiring non-adiabaticity. It will be connected with vortex formation. The appearance of phase space vortices at high Mach numbers resembles a tendency towards shock breaking. In the light of this discussion the lower two panels in Figure 5.28 can also be interpreted as breaking and overturning of the quasi-perpendicular ($\Theta_{Bn} = 87^\circ$) realistic mass-ratio supercritical shock. In particular during the phase before reformation (third panel from top) the magnetic field behaves irregular, and both the incoming and reflected beams form many partial vortices prior to the reflection point (at $\Delta x \sim 140 \lambda_e$). Behind the reflection point the ion velocity shows formation of bursts of ions which run away in forward direction, which is just what is expected in breaking. A sketch of the dynamical processes is given in Figure 5.31.

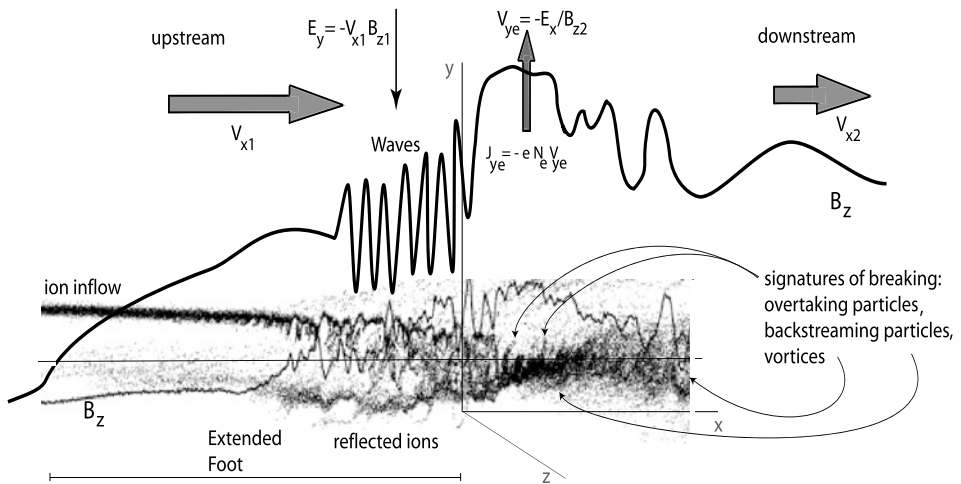


Figure 5.31: Schematic of the profile of a highly supercritical shock with waves just before shock reformation and signatures of beginning wave breaking. The sketch has been completed with a copy of the ion phase space from the simulations of *Matsukiyo & Scholer* [2006b] showing the structure of the ions in the ramp with the signatures of overtaking ions and backstreaming ions as well as ion vortices, all an indication of onset of breaking.

5.7 Update – 2012

Nonstationarity viz. shock reformation has been investigated at the example of Earth's bow shock wave with the help of CLUSTER spacecraft measurements by *Mazelle et al* [2010] showing that the ramp-gradient scale changes with time and can become as narrow as an electron inertial length, which is a most interesting observation as this is the shortest possible transverse length scale for a collisionless magnetised shock.

Related to this observation may be another observation by *Bale et al* [2008] who used CLUSTER measurement across the quasi-perpendicular super-critical bow shock wave at low plasma β to determine the cross-shock electric potential finding that the potential is typically in the range of $500 < \Phi < 2500$ V which amounts to large variations in the ion energy change in the range $20 < \Delta \mathcal{E}_{ion} / \mathcal{E}_{ion} < 240\%$. This would be indication of the high variability of the shock.

The cross-shock potential is also related to electron heating as described in this chapter. However, direct multi-spacecraft CLUSTER measurements analysed by *Schwartz et al* [2011] that this electron heating is a more complex and obviously multi-scale process which is related to the spatial and temporal structure of the shock layer electromagnetic field. These authors determined the electron temperature gradient across the quasi-perpendicular Earth's bow shock from the electron distribution measured *in situ*. According to these observations, roughly $\sim 50\%$ of the electron heating must be attributed to the narrow few λ_e thick shock ramp layer. This heating comes from an inflation the electron phase space and indicates irreversibility which is attributed to wave particle interaction. In view of this the conclusion that wave turbulence in a narrow transition region practically

independent on ion dynamics is responsible for the electron heating. This is, however, consistent with the assertion that violent instabilities work in this narrow region like the Buneman mode, the modified two-stream instability caused by the shock-electric potential. Possibly, as the authors suggest, oblique electron whistlers are also involved. *Sundkvist et al* [2009]; *Sundkvist et al* [2012], based on CLUSTER multi-spacecraft observations from crossings of Earth's bow shock, calculated the Poynting flux of such oblique whistler waves in two quasi-perpendicular ($\Theta_{Bn} \sim 85^\circ$) relatively high-Mach ($\mathcal{M}_A = 5.5$ and 11) number shocks finding – in the shock-normal frames – that the Poynting flux is directed obliquely upstream, away from the shock, as predicted by quasi-perpendicular supercritical shock theory. Their assumption is that the parallel electric field vanishes. Thus the whistlers generated at the shock are oblique and dispersive with k -vector approaching 90° with respect to the ambient magnetic field at the shock ramp. Downstream of the shock, the wave spectrum becomes incoherent, indicating that the waves in the whistler frequency range are turbulent. We note that recently [*Wilson et al*, 2012] electromagnetic whistler precursor waves have also been inferred from Wind data around some supercritical interplanetary shocks.

Shinohara & Fujimoto [2010] have performed first three-dimensional PIC simulations of quasi-perpendicular shocks at a realistic mass ratio $\mu = 1840$ finding very strong wave activity in the shock foot region which is permitted by the inclusion of the third dimension. These waves cause a stronger than known upstream electron acceleration and the generation of non-thermal electrons. PIC simulations at smaller mass ratio and in two dimensions [*Savoini et al*, 2010] had already shown indications of reflected electrons before and attributed them to a magnetic mirror effect in the foot-shock ramp field.

Also of interest is the observation of a variability of electron counts at 500 eV electron energy observed downstream of the super-critical quasi-perpendicular bow shock by CLUSTER [*Matsui et al*, 2011]. It suggests a quasi-periodic modulation of electron flux at this energy with period ~ 3 s, suggested to be related to mirror mode waves excited by the presumably large downstream temperature anisotropy that is generated by the quasi-perpendicular shock transition.

Near the Saturnian bow shock, *Masters et al* [2011] found substantial solar wind electron heating from Cassini data during 2005 and 2007, attributing the heating to the action of the comparably strong bow shock at this outer strongly magnetised planet. The heating observed is correlated to the ram flow energy of the incident solar wind, e.e. to the upstream Mach number. It amounts to between 3% and 7% of the incident ram energy.

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Quasi-parallel Supercritical Shocks

Abstract. Quasi-parallel shocks are probably more common in nature than quasi-perpendicular. They exist for shock normal angles $>45^\circ$. Their main distinctions from the latter are, first, their turbulent nature implying that the shock transition is less sharp and thus less well defined; second, the existence of an extended turbulent foreshock instead of a shock foot. This foreshock has distinct properties which are discussed. It consists of an electron and an ion foreshock. The main population is a diffuse ion component. The turbulence in the foreshock is generated by the (reflected and accelerated) foreshock particle populations. We discuss the various processes of wave generation referring to observations and sophisticated numerical simulations. An important point in quasi-parallel shock physics is the reformation of the shock which works completely differently from quasi-perpendicular shocks. Here it is provided by upstream low-frequency electromagnetic waves excited by the diffuse ion component. Steeping of these waves during shockward propagation and addition of the large amplitude waves at the shock transition reforms the shock front. The old shock front is expelled downstream where it causes downstream turbulence. During the reformation process the shock normal angle locally turns to low values. Thus the shock becomes locally quasi-perpendicular for the ions supporting particle reflection. Finally hot flow anomalies are briefly discussed.

6.1 Introduction

At a first glance it is surprising that a change in the shock-normal angle Θ_{Bn} by just a few degrees from, say, $\Theta_{Bn} = 50^\circ$ to $\Theta_{Bn} = 40^\circ$ should completely change the character of the supercritical shock. We have seen in the last chapter, when having discussed the conditions for reflection of ions from the shock ramp that such a change in the shock properties is theoretically predicted. The critical shock-normal angle $\Theta_{Bn} = 45^\circ$ does indeed separate two completely different phases of a supercritical shock. At this angle the shock experiences a ‘phase transition’ from the quasi-perpendicular to the quasi-parallel shock state, with the shock-normal angle Θ_{Bn} having the property of a ‘critical control parameter’.

We have already learned that the reason for the different behaviour and different physics involved of the two shock phases is that in a quasi-perpendicular shock all reflected particles from the foot region of the shock return to the shock after no more than a few gyrations when they have picked up sufficient energy in the upstream convection electric field $\mathbf{E}_1 = -\mathbf{V}_1 \times \mathbf{B}_1$ to ultimately overcome the shock ramp potential, pass the shock ramp and merge into the downstream flow. We have not discussed what happens to the accelerated ions in the downstream region as this is not of primary importance in the shock formation mechanism which to good approximation depends only on the upstream conditions. This question will be treated in a separate section on shock-particle acceleration.

In contrast, in a quasi-parallel shock the combined geometries of the upstream magnetic field and generally curved shock surface prevent the shock-reflected particles from immediate return to the shock. The reason is that their gyro-orbits, after having suffered reflection from the shock ramp, lie completely upstream, outside the shock ramp, such that they do not touch the shock ramp again after reflection. Since, in addition, their upstream velocities have a large component parallel to the upstream magnetic field, which increases the more the shock-normal turns parallel to the upstream magnetic field, the reflected particles are enabled to escape upstream from the shock along the magnetic field thereby escaping backward convection and forming fast upstream particle beams.

A quasi-parallel supercritical collisionless shock thus populates the upstream space with a reflected particle component. This population moves a long distance away from the shock along the magnetic field. Was the upstream flow, in the case of the quasi-perpendicular shock, completely uninformed about the presence of the shock up to a distance of the mere width of the shock foot so, at the quasi-parallel shock, it receives a first signal of the presence of a shock already at quite a large upstream distance when the first and fastest reflected particles arrive on the magnetic field lines that connect the flow to the shock. It is, in fact, only these particles that can inform the flow about the presence of a supercritical shock, because any low-frequency plasma wave cannot propagate far upstream for supercritical Mach numbers $\mathcal{M} > M_c$, while any electromagnetic radiation that is generated at the shock has frequency $\omega_{rad} > \omega_{pe}$. For it the plasma flow presents a vacuum. The upstream flow recognises the reflected particles in its own frame of reference as a high-speed magnetic-field aligned beam. Thereby a beam-beam configuration is created which leads to a number of beam-driven instabilities. These excite various plasma waves that fill the space in front of the shock and modify its properties.

Figure 6.1 shows a sketch of the magnetic profile of a supercritical quasi-parallel shock which contrasts the profile of a quasi-perpendicular shock that has been given in the previous chapter. The quasi-parallel magnetic shock profile is much stronger distorted than that of a quasi-perpendicular shock, such that it becomes difficult to identify the location of the genuine shock ramp on the profile.

The main difference between quasi-perpendicular and quasi-parallel shocks is that quasi-perpendicular shocks possess a narrow $\sim 1r_{ci}$ wide foot region that is tangential to the shock surface, while quasi-parallel shocks possess an extended *foreshock* region. Interestingly, in curved shocks which arise, for instance, in front of spatially confined obstacles, both phases of a supercritical shock can coexist at the same time, being spatially adjacent to each other. An example is shown in Figure 2.6 in the sketch of the curved Earth's bow shock. Dealing with quasi-parallel shocks means to a large extent dealing with the processes that are going on in the foreshock. It will thus be quite natural to start with a discussion of the properties of the foreshock. This discussion will occupy a substantial part of this chapter.

However, before continuing we point out that in spite of the strict distinction between the quasi-perpendicular and quasi-parallel shocks there is also a close relation between the two. Both, being supercritical, can exist only because they reflect ions; and both possess an upstream region in front of the shock transition that is populated by the reflected ions. That this region is narrow in the case of quasi-perpendicular shocks is a question of the ions

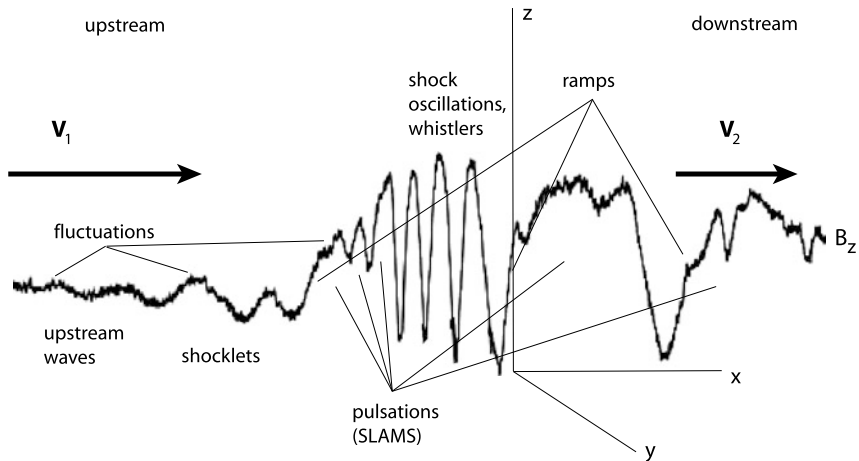


Figure 6.1: Schematic one-dimensional profile taken along the nominal instantaneous shock normal of a supercritical quasi-parallel shock as seen in the magnetic field component B_z . This is the analogue to the quasi-perpendicular shock profile. (This figure is essentially a compilation of different magnetic features observed in space and seen in full particle simulations.) It shows the main features in the vicinity of the quasi-parallel shock transition: the large amplitude upstream waves with the turbulent fluctuations on top of the waves, the formation of shocklets, i.e. steep flank formation on the waves exhibiting small-scale fluctuations on top of the wave, which act already like small shocks, very-large amplitude pulsations (magnetic pulsations or SLAMS) which turn out to be the building blocks of the shock, multiple shock-ramps at the leading edges of the pulsations belonging to diverse ramp-like steep transitions from upstream to downstream lacking a clear localisation of the shock transition and any attached phase-locked whistlers. (Note that the entire figure is, in fact, the shock transition, as on this scale no clear decision can be made where the shock ramp is located.) Not shown here are the out of plane oscillations of the magnetic field that accompany the waves. Also not shown is the particle phase space.

being tied to the magnetic field, which also holds in the case of the quasi-parallel shock. At quasi-perpendicular shocks the reflected ions do readily return to the shock. At the quasi-parallel shock they ultimately do also return to the shock, but only after having been processed far away from the shock in the foreshock, having coupled to the flow, having passed several stages in this processing, and having become the energetic component of the main flow. As such they finally arrive at the shock together with the stream. In between their main duty was to dissipate the excess energy, which they possessed when arriving for the first time at the shock and which could not be dissipated in the narrow shock-ramp transition region. This could be achieved only in the broad extended foreshock which, seen from this point of view, *is already the shock*. It belongs inextricably to the quasi-parallel shock transition. Here, a substantial fraction of the energy of the incident flow is dissipated in a way which is completely different from the flow being shocked.

It is these dissipation processes that cause the main difference between the quasi-perpendicular and quasi-parallel states of a collisionless supercritical shock. To stress the analogy with phase transitions a little further, we may say that quasi-perpendicular shocks are in the solid – or ordered – shock state, while quasi-parallel shocks are in the fluid – or partially disordered – shock state.

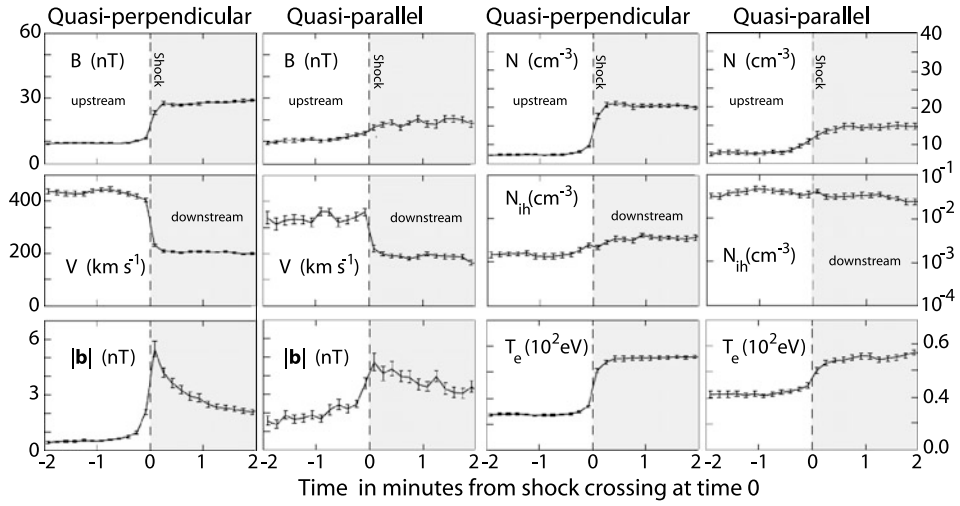


Figure 6.2: Comparison of average bulk plasma parameters in quasi-perpendicular and quasi-parallel shocks [data from Czaykowska *et al.*, 2000, courtesy American Geophysical Union]. The figure shows (*on the left*) the mean magnetic field B , bulk flow velocity V , average magnetic fluctuation amplitude at ultra-low frequencies, and (*on the right*) plasma density N , density of high energy ions N_{ih} , and electron temperature T_e . The shaded regions are the downstream parts of the AMPTE IRM crossings of the bow shock. The data have been obtained by normalising the time with respect to crossing the nominal shock ramp by using the measured normal component (not shown) of the bulk flow velocity. This for the many observations included in this figure implies a stretching (or squeezing) of each individual shock crossing, causing some uncertainty, in particular for quasi-parallel shock crossings as there the shock ramp is not well defined. However, this figure serves for an immediate overview of the differences in both types of shocks.

6.2 The (Quasi-parallel Shock) Foreshock

Quasi-parallel shocks are abundant in space because in most cases when shocks develop the flow direction is independent of the direction of the magnetic field. Moreover, as pointed out earlier, when a bow shock forms around an obstacle (planet, magnetosphere, moon, ...) this bow shock is curved around the obstacle, and the shock is quasi-perpendicular only in a certain region on the shock surface that is centred around the point where the upstream magnetic field touches the shock tangentially. Farther away the shock turns gradually to become quasi-parallel. On the other hand, when a shock survives over a long distance in space as for instance in supernova remnants then it sweeps the upstream magnetic field and pushes it to become more tangential to the shock surface. In this case the shock is about quasi-perpendicular. We will later provide arguments that any quasi-parallel supercritical shock on the small scale close to the shock surface, i.e. on the electron scale, behaves quasi-perpendicularly while on the larger ion scale it remains to be quasi-parallel. This has consequences for the differences in the dynamics of electrons and ions during their interaction with quasi-parallel shocks.

Figure 6.2 shows at one glance the main differences in the (average) bulk plasma parameters between quasi-perpendicular and quasi-parallel shocks as measured with the

AMPTE IRM spacecraft at many crossings of the bow shock. The data used in this figure have been stapled, averaged and plotted with respect to the time normalised to the shock ramp crossing. For such a normalisation one uses the shock-normal upstream velocity to recalculate the time. This procedure is not very certain for quasi-parallel shocks since – as we will see later – the shock ramp is ill defined in a quasi-parallel shock. However, for a simple comparison of the main differences this uncertainty is less severe.

The shaded area in the figure corresponds to the downstream region. Shown are – in pairs of quasi-perpendicular/quasi-parallel values – the magnetic field B , bulk velocity V , average fluctuation amplitude in the ultra-low frequency waves $|\mathbf{b}|$, plasma density N , high-energy ion density N_{ih} of energy >15 keV, and electron temperature T_e . The general conclusion from this figure is that all quantities in the quasi-perpendicular case exhibit a much sharper transition than in the quasi-parallel case. Moreover, the quasi-perpendicular averages are quieter than those of the quasi-parallel case. Also, in general, the quasi-parallel levels are higher than the quasi-perpendicular. In almost all cases the pre-shock levels are enhanced in the quasi-parallel shock case with over the pre-shock levels of quasi-perpendicular shocks. This is seen most impressively in the energetic ion density, which is nearly constant over this distance/time scale at quasi-parallel shocks and much higher than that in quasi-perpendicular shocks, signifying on the one hand the importance of energetic particles in quasi-parallel shock dynamics, on the other hand the capability of quasi-parallel shocks to accelerate particles to substantial energies. The presence of energetic ions (particles) far in front of the quasi-parallel shocks and the enhanced pre-shock levels indicate the importance of foreshocks in quasi-parallel shock dynamics. In the following we will therefore first concentrate on the foreshock.

The physics of quasi-parallel shocks cannot be understood without reference to the foreshock. The foreshock is that part of the upstream shock region that is occupied with reflected particles. At a curved shock, like the Earth's bow shock, the foreshock starts on the shock surface at the location where the upstream magnetic field shock-normal angle exceeds $\Theta_{Bn} \gtrsim 45^\circ$. From that point on electrons and ions escape along the magnetic field in upstream direction. Since electrons generally move at a larger parallel velocity than ions they are less vulnerable to the convective motion of the upstream magnetic field line to which they are tied, and so there is generally a region closer to the foreshock-boundary magnetic field line where only upstream electrons are found. This region is confined approximately between the line that marks the electron foreshock boundary and the more inclined line that marks the ion foreshock boundary. An example of this geometry was depicted in Figure 2.6 for the bow shock of the Earth.

More schematically this is shown in a simplified version in Figure 6.3 for the particular case that the upstream magnetic field forms an angle of 45° with the symmetry axis of the shock. In this case half of the shock is quasi-perpendicular and the other half is quasi-parallel. The figure also shows the directions of three shock normals, the narrow foot region in front of the quasi-perpendicular shock, and the two (electron and ion) foreshocks. Particles escape from the quasi-parallel shock along the upstream magnetic field. The magnetic field is convected toward the shock by the perpendicular upstream velocity component $\mathbf{V}_\perp$ as shown in the figure. This component adds to the velocity of the upstream particles leading to an inclined foreshock boundary. Since the ions have much

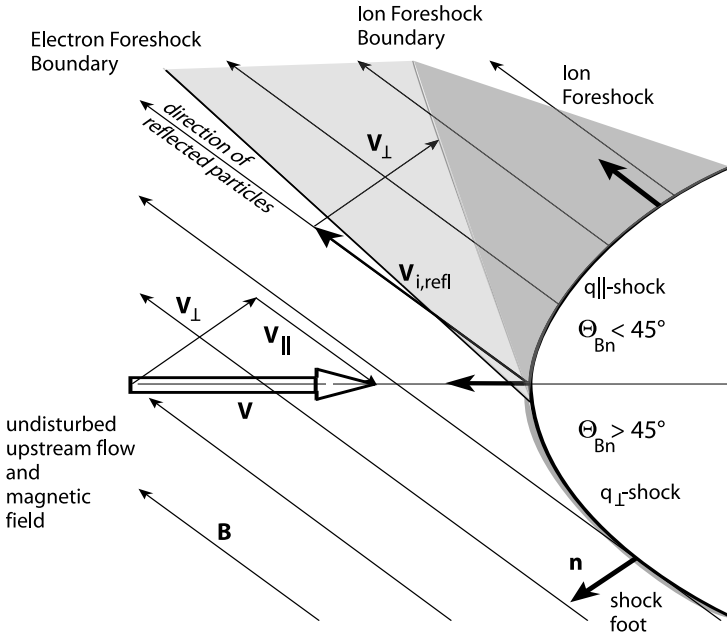


Figure 6.3: Schematic of the relation between a curved shock and its foreshock in dependence on the direction of the upstream magnetic field $\mathbf{B}$, shock-normal $\mathbf{n}$, and shock-normal angle Θ_{Bn} for the special case when the magnetic field is inclined at 45° with respect to the symmetry axis of the shock. In this case the upper half of the shock becomes quasi-parallel ($\Theta_{Bn} < 45^\circ$), the lower half is quasi-perpendicular ($\Theta_{Bn} > 45^\circ$). The velocity of reflected particles is along the magnetic field. However, seeing the flow the field-line to which they are attached displaces with perpendicular velocity. This velocity shifts the foreshock boundary toward the shock as shown for electrons (light shading) and ions (darker shading). The ion foreshock is closer to the shock because of the lower velocity of the ions than the electrons. For the electrons the displacement of the electron foreshock boundary is felt only at large distances from the shock.

smaller speed than the electrons, the ion foreshock boundary is more inclined than the electron foreshock boundary.

In discussing the properties of the foreshock one thus has to distinguish of which foreshock is the talk. However, the properties of the electron foreshock are not as decisive for the formation of a quasi-parallel shock as are the properties of the ion foreshock. Because of this reason we will, in the following, refer to the ion foreshock as the foreshock. The electron foreshock properties we will mention only later.

6.2.1 Ion Foreshock

The ion foreshock is not a homogeneous and uniform region. The reflected ion component evolves across the ion foreshock from the ion foreshock boundary to the centre of the ion foreshock and from there towards the shock. Speaking of a reflected ion component that can unambiguously be identified as being reflected, i.e. streaming into the upward

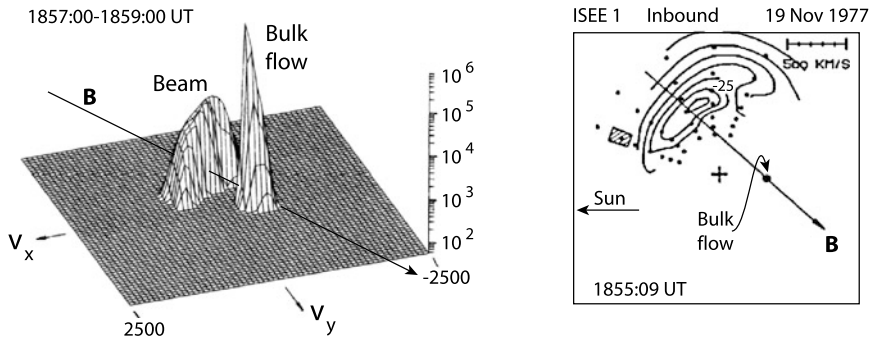


Figure 6.4: ISEE 1 observation of a reflected ion beam on November 19, 1977 propagating along a magnetic field line that was connected to the quasi-parallel Earth's bow shock. *Left:* A two-minute average pseudo-three-dimensional ion velocity-space profile in the (v_x, v_y) -plane showing the undisturbed and cold (narrow) plasma inflow in negative v_x -direction, and the fast and warm (broad) beam of reflected ions propagating in positive v_x -direction and spreading in v_y . This beam is quite anisotropic in temperature. Velocities are in km s^{-1} . The scale on the right is count rates, and background count rates were suppressed by choosing only values above 50 s^{-1} . *Right:* Contour plot of a similar beam a little earlier showing that the beam is centred on the magnetic field that connects to the shock, is quite narrow along the field and about 2–3 times as broad perpendicular to the field [data from Paschmann *et al*, 1981, courtesy American Geophysical Union]. The cross indicates the origin (zero velocity), the dot the bulk flow centre. The $10^{-25} \text{ s}^3 \text{ cm}^{-6}$ level flux contour has been marked.

direction, makes sense only in the immediate vicinity of the ion foreshock boundary. Here the reflected ions appear as a fast ion beam the source of which can be traced back to the shock. Deeper in the foreshock the beam component cannot be identified anymore.

The Ion Foreshock Boundary Region

First identifications of reflected beam protons in space in the magnetic flux tube connected to the Earth's bow shock wave were reported by Gosling *et al* [1978] and Paschmann *et al* [1981] who distinguished those beams by their poorly resolved distribution functions from more diffuse protons deeper in the foreshock. Interestingly, observations in the foreshock of interplanetary travelling shocks did not show any indication of such beams but only the diffuse ion component. Figure 6.4 gives an observational example of such a reflected ion beam that propagates very close to the foreshock boundary upstream away from the shock. The bulk flow is the narrow cold beam in the left part of the figure which is displaced in negative v_x -direction (note that in this figure the positive direction points away from the shock). The reflected beam is less dense (lower count rates) but much more energetic. It is displaced in $+v_x$ -direction, i.e. streaming away from the shock, and has also a $-v_y$ -component, i.e. it constitutes a gyrating bunch of ions moving away from the shock. In the right part of the figure it is seen that the beam is moving away along the magnetic field line that is connected to the shock, while the bulk of the plasma flows in positive direction.

These beams along the foreshock boundary play some role in the foreshock dynamics as they seem to represent a source population for the entire ion foreshock. Whether and

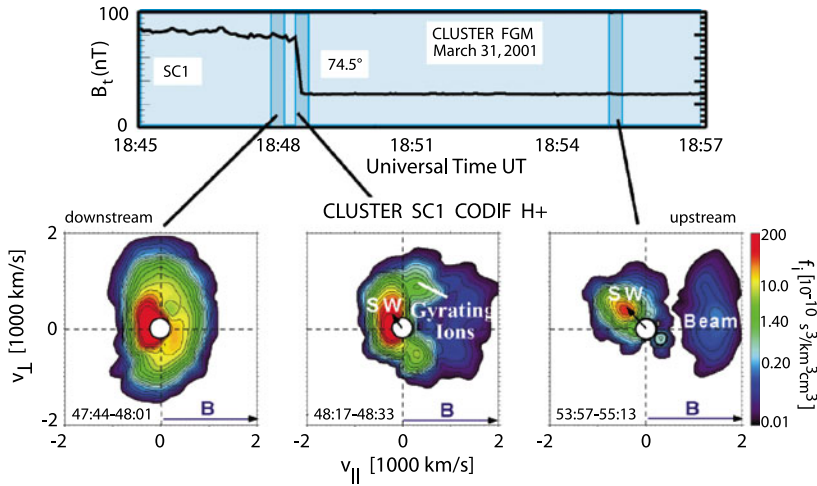


Figure 6.5: Ion phase space at three locations along the shock-connected magnetic field line in a high-Mach number supercritical quasi-perpendicular shock the magnetic field of which is shown in the top panel. The shock-normal angle is $\Theta_{Bn} = 74.5^\circ$. The three lower plots show downstream, shock ramp/shock foot, and distant upstream phase-space plots. The ramp/foot plot shows the presence of the incoming flow (SW) and gyrating ions in the foot. The upstream plot shows the upstream field-aligned relatively hot (large velocity spread) beam well separated from the inflow (SW) [data from Kucharek *et al.*, 2004].

why this is really so is not yet fully understood as the shock should reflect ions at almost every place in its quasi-parallel state. However, it seems that only the group of ions that escape from the shock along the foreshock boundary can form such beams. This points on a further interesting relation between quasi-perpendicular and quasi-parallel shocks at a curved shock surface with a smooth transition from quasi-perpendicular to quasi-parallel as sketched in Figure 6.3 and realised in space, for instance, at planetary bow-shocks. It seems as so these beams escape from the quasi-perpendicular region of the shock along the nearly tangential field lines. This would also be in agreement with the observation [Gosling *et al.*, 1984] that the foreshocks of extended interplanetary shocks do not show any signs of reflected ion beams. They are only very weakly curved being nearly planar, and do not possess a recognisable quasi-perpendicular area on the surface.

Kucharek *et al.* [2004] analysed ion distributions along magnetic field lines that were connected to the quasi-perpendicular area of Earth's bow shock. Investigating the origin of those beams they found that, indeed, the observed ion beams at the foreshock boundary result from reflection at the quasi-perpendicular shock and that, without the presence of a quasi-perpendicular region, there would presumably be no distinct foreshock boundary and no ion beams escaping into the foreshock. If this is the case, then the population in the quasi-parallel ion foreshock is indeed provided by two different ion sources, the beams from the quasi-perpendicular shock region and the genuine foreshock ion population. The latter has no beam character but is rather a diffuse ion component [see also Meziane *et al.*, 2004]. Interplanetary shocks possess only the latter one.

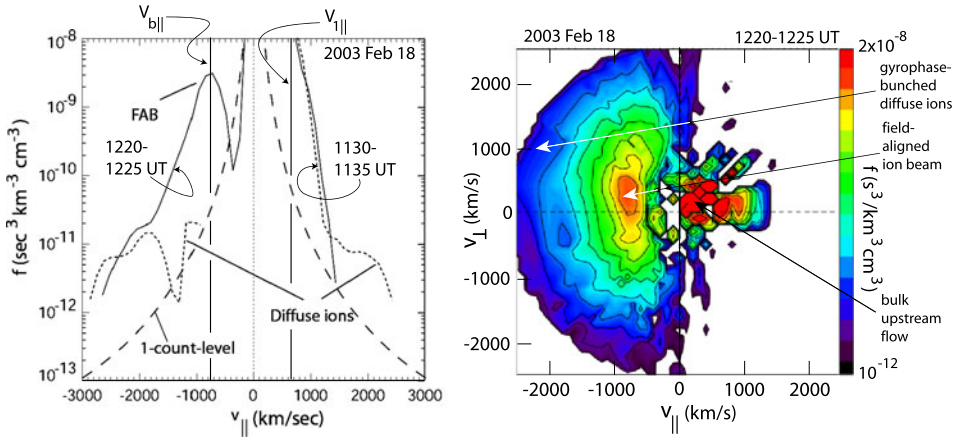


Figure 6.6: *Left:* The reduced field-aligned ion distribution functions in the foreshock region showing the field-aligned reflected ion beam (FAB) at 1220–1225 UT on February 18, 2003 which arrives at the spacecraft location along the magnetic field line that is connected to the quasi-perpendicular region of the Mach number $\mathcal{M}_A \sim 8$ supercritical shock (smooth solid line). The upstream bulk velocity was $V_1 \simeq 660 \text{ km s}^{-1}$. Also shown is the reduced parallel distribution function (dashed) for the diffuse ion distribution observed at 1130–1135 UT the same day (for the geometry see Figure 6.7). It is clearly seen that the latter is about symmetric to the magnetic field direction indicating the symmetry of the ring of diffuse ions, while the foreshock-boundary field-aligned beam is flowing in the direction away from the shock into the upstream medium (negative velocities). Note also that the diffuse distribution appears as a smooth tail on the full ion distribution. The small gap on the left is uncertain as it dips into the 1-count level. *Right:* The two-dimensional phase space plot for the time interval 1220–1225 UT when the field-aligned beam was observed. Indicated are the bulk upstream flow, field-aligned ion beam, and the gyro-phase bunched residue of the diffuse upstream ions [data from *Kis et al., 2007*].

Roughly 2% of the ion inflow leaves the shock ramp upstream in the form of a beam along the magnetic field. *Kucharek et al [2004]* argue that the ions which escape along the magnetic field, are reflected from the very ramp/overshoot region where they have been in resonance with low-frequency plasma waves, which they assume to be large amplitude Alfvén-whistler waves. These ions experience pitch-angle scattering and pitch-angle diffusion towards small pitch angles, and subsequently can escape along the magnetic field in the upstream direction. Since the conditions for escape depend in the first place on the pitch-angle scattering process, the beams should be highly variable in time and location. The mechanism might still sound a bit speculative as long as no simulation proves its reality, but any mechanism which is able to pitch-angle scatter ions along the magnetic field in a quasi-perpendicular shock will naturally cause ion beams to escape from the ramp both in the upstream and in the downstream directions. Such simulations require a three-dimensional treatment which is not in reach yet.

A measured example of the ion foreshock-boundary field-aligned ion beam distribution [*Kis et al., 2007*] is shown in Figure 6.6 for an upstream flow velocity $V_1 \simeq 660 \text{ km s}^{-1}$, Mach number $\mathcal{M}_A = 8$, and average shock-normal angle $\Theta_{Bn} \simeq 15^\circ$ [*Archer et al., 2005*], at an upstream distance from the shock (in this case again the Earth’s bow shock). This distribution is a so-called reduced distribution; it is the integrated over pitch-angle ϕ and

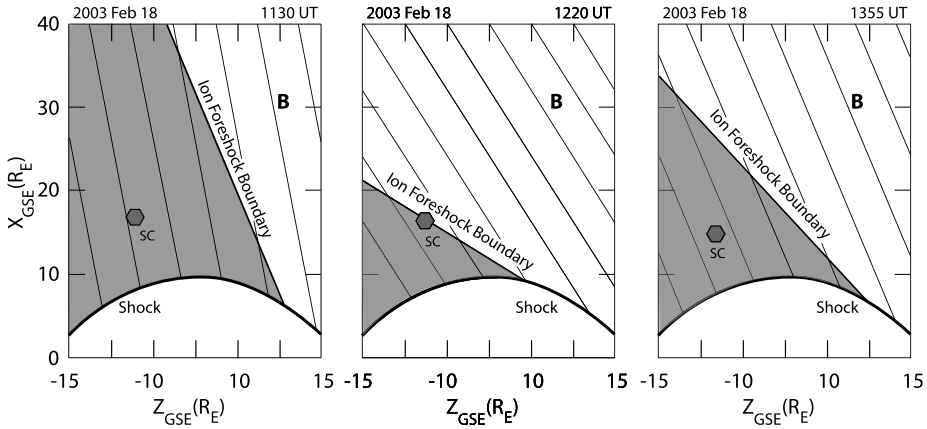


Figure 6.7: Reconstruction of the shape and location of the ion-foreshock boundary from the measured upstream plasma properties (direction of magnetic field, speed and density) for three successive times in order to explain the ion-phase space observations in Figure 6.6 [data from *Kis et al., 2007*]. The cold magnetic field-aligned ion beam is observed when the spacecraft is located in the vicinity of the ion-foreshock boundary at 1220 UT, while at 1130 UT and 1355 UT no beam was detected. It is assumed that the beam is generated along the foreshock boundary at the position where at the quasi-perpendicular shock surface the shock-normal angle is roughly about $\Theta_{Bn} \sim 60^\circ$.

perpendicular velocity $v_\perp$ magnetic field-aligned phase space distribution function $f(v_\parallel) \sim 2\pi \int v_\perp^2 dv_\perp f(v_\perp, v_\parallel)$ that has been appropriately binned and smoothed. The information to be taken out of this figure is that the reduced ion-beam distribution (solid line) is narrow in velocity, maximising at a speed (in absolute terms $|v_{b\parallel}| \simeq 800$ km/s) that is only slightly larger than the parallel flow velocity ($V_{\parallel 1} \simeq 640$ km/s, when taking into account the shock-normal angle), i.e. $|v_{b\parallel}| \simeq 1.25 V_{\parallel 1}$. It is directed opposite to the flow. For comparison, a reduced parallel diffuse ion distribution is shown in the same plot taken deeper in the foreshock. This distribution is about symmetric to the magnetic field direction, indicating the about circular phase-space distribution of the diffuse ion component which appears as an energetic tail on the main ion distribution (note that the small gap on the left of the dashed diffuse-ion distribution curve is questionable as it dips below the 1-count level). *Meziane et al [2004]* reported CLUSTER observations of foreshock-boundary ion beams simultaneously with diffuse ions. They found that the nominal ion-beam velocity $V_b \simeq 1.7 V_1$ had no relation to any known shock-reflection mechanism like specular reflection, a conclusion which supports pitch-angle scattering as the beam-injection process as this is independent of the speed of the inflow.

Kis et al [2007] have carefully analysed the relation between the observed ion distribution at the upstream spacecraft position and the calculated shock-normal angle Θ_{Bn} , determined from the local upstream magnetic field direction and the predicted shape of the bow shock. The result is shown in Figure 6.7 for three successive times on February 18, 2003 when the CLUSTER spacecraft was outside the bow shock. At 1130 UT and 1355 UT the shock normal angles $\Theta_{Bn} \sim 15^\circ$ and shock-spacecraft distances ~ 7.5 respectively $\sim 6.7 R_E$

were similar. Despite this similarity the observed ion phase-space distributions were completely different. At 1130 UT no low-energy gyrating ions were observed, while they were present at the later time 1355 UT. Hence, at 1130 UT the spacecraft must have been closer to the ion foreshock boundary, i.e. the foreshock boundary was more inclined than at the later time such that beam particles scattered from the foreshock boundary have not arrived at the location of the spacecraft. Due to the velocity filter effect they have been separated out at spacecraft distance. The reconstruction of the position of the ion-foreshock boundary using the measured upstream conditions and shape of the shock for this period is shown in Figure 6.7. Indeed, at 1220 UT the spacecraft was close to the foreshock boundary and, as expected, detected the ion beam (as seen in Figure 6.6).

This observation supports the above advocated view that the ion-foreshock beam is generated in the transition region from quasi-perpendicular to quasi-parallel shock. The remaining questions to be answered are: what pitch-angle scattering mechanism is responsible for the generation of such a beam, and what is the fate of the foreshock-boundary beam-ions during their propagation along the foreshock boundary? Do they contribute to the foreshock ion population and if, in what way? Currently we are not able to answer either of these questions definitely. In particular, the pitch-angle scattering mechanism is unknown or at least uncertain. When discussing wave generation, we will touch on the problem of the fate of the beam. Below we present evidence for the scattering of the ion-beam ions and merging into the upstream foreshock diffuse-ion population.

Diffuse Ions

Reflected upstream magnetic field-aligned ion beams are observed at the foreshock boundary only. The second (and main) ion component encountered in the foreshock is the diffuse ion population which is detected there as the energetic extension of the inflow plasma. It is widely assumed that the origin of this component is also at the shock as there is no other energetic particle source available. However, close investigation of the diffuse ion component in the foreshock of the Earth's bow shock wave has demonstrated that these ions are not produced in a specular reflection process at the shock. Rather their origin is of diffusive nature.

So far it has not been possible to identify the source of these ions though some models have been proposed which we will briefly discuss below in relation to the appearance of foreshock waves. We will also return to these particles in the next chapter on shock particle acceleration. Here, we merely discuss some of their properties and provide evidence that they are indeed coming from the shock, proving that the shock is an energetic ion source.

The foreshock-boundary ion beams also merge into the foreshock particle distribution by scattering from the foreshock boundary off their self-generated wave spectrum and subsequently being convected downstream by the bulk flow. *Paschmann et al* [1981] have mapped this merging process by following the evolution of the foreshock-boundary beam distribution shown in Figure 6.4 during the convection. In that figure the beam was detected close to the foreshock boundary. Figure 6.8 show its form deeper in the foreshock when it has spread substantially in angle evolving into half of a ring distribution already. Even

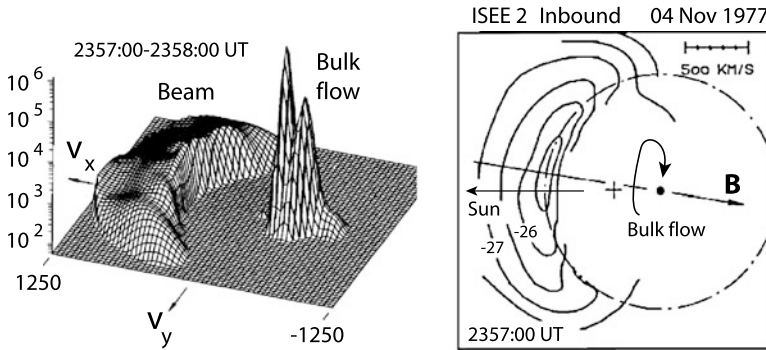


Figure 6.8: The evolution of the ion foreshock-boundary ion beam during its convection downstream into the foreshock as has been seen by ISEE 2 on 04 November 1977. *Left:* A one-minute average pseudo-three-dimensional ion velocity-space profile in the (v_x, v_y) -plane showing the spreading of the beam in angle around the bulk flow without merging into the bulk flow. Velocities are in km s^{-1} . The scale on the right is count rates, and background count rates are suppressed by choosing only values above 50 s^{-1} . *Right:* Contour plot of the partial ring distribution. The direction of the magnetic field is also shown [data from *Paschmann et al.*, 1981, courtesy American Geophysical Union]. The 10^{-26} and $10^{-27} \text{ s}^3 \text{ cm}^{-6}$ level flux contours have been marked.

deeper inside the foreshock the reflected ion distribution assumes the shape of a full ring around the bulk distribution as is shown in Figure 6.9.

The three observations depicted in Figures 6.4, 6.8 and 6.9 are from different times; it has, however, been checked that they are at distances corresponding to increasing distance from the foreshock boundary such that the assumption of the convectively processed beam evolution is well founded (or at least reasonable) even though it has not been directly proven. It is interesting to note that a gap remains between the original ion-phase space beam distribution and the bulk-flow distribution, which is another indication that the evolving distribution is part of the evolution of the ion-foreshock beam. Of a distribution that does not evolve out of a beam one expects a less regular behaviour and, generally, no such well expressed gap between bulk flow and beam in velocity space. In fact, the main energetic ion component in the foreshock is irregular and lacks a well expressed gap.

The discrepancy between the smooth no-gap foreshock distributions [*Sentman et al.*, 1981a] and the gap-observations of *Paschmann et al.* [1981] has, in fact, been noted much earlier [*Sanderson et al.*, 1981; *Scholer*, 1985] without giving an explanation but suggesting a continuous ion source at the parallel shock. Below we provide further arguments for the two-source, foreshock-boundary beam and continuous extended shock-surface source hypothesis. Nonetheless, this conclusion must be taken with care because Figure 6.9 shows the gap progressively closing. Being sufficiently far, i.e. even farther away from the ion foreshock boundary than in this figure, it will not anymore be possible to distinguish between beam-evolved and genuine diffusive-ion distributions. Ultimately, both distributions will have merged indistinguishably. Still there is no agreement whether and where, i.e. at what distance from the ion foreshock boundary this merging of the two populations occurs.

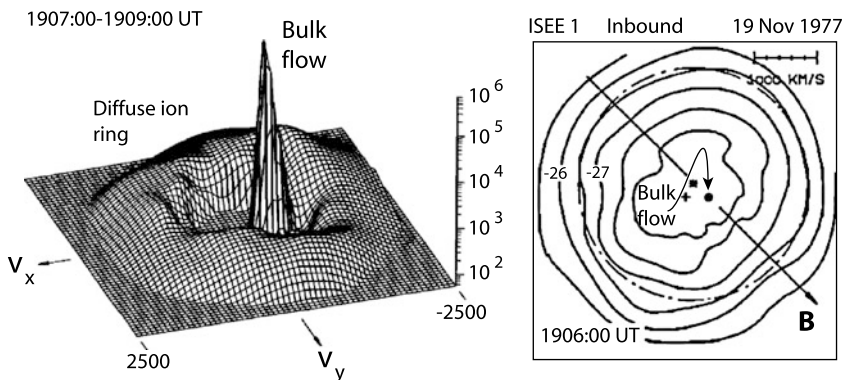


Figure 6.9: ISEE 1 observation of diffuse ion beam on November 19, 1977 propagating deep inside the foreshock away from the ion foreshock boundary. *Left:* A two-minute average pseudo-three-dimensional ion velocity-space profile in the (v_x, v_y) -plane showing the undisturbed and cold (narrow) plasma inflow in negative v_x -direction about centred and surrounded by a ring distribution of fast and warm (broad) diffuse ions that have been reflected from the shock but have been processed in the foreshock region when propagating from the foreshock boundary into the foreshock. These ions are hollow in the sense that they separate from the bulk distribution but have a nearly isotropic distribution function. Velocities are in km s^{-1} . The scale on the right is count rates, and background count rates were suppressed by choosing only values above 50 s^{-1} . *Right:* Contour plot of a similar ring a little earlier showing that the ring centre (star) is slightly displaced from the bulk flow (dot) on the magnetic field that connects to the shock and from the centre (cross) of the phase space frame [data from Paschmann *et al.*, 1981, courtesy American Geophysical Union]. The 10^{-26} and $10^{-27} \text{ s}^3 \text{ cm}^{-6}$ level flux contours have been marked. Note the near isotropy of the ring distribution.

Trattner *et al.* [1994], using AMPTE IRM measurements of diffuse ion densities upstream of the quasi-parallel (bow) shock found that the diffuse ion density decreases exponentially with shock distance. This investigation was substantially improved by Kis *et al.* [2004] in an attempt to infer about the source of the upstream diffuse energetic ions. These authors determined the partial-density gradient of diffuse ions in the energy range from 10 to 32 keV as a function of distance from the (bow) shock. This investigation was made possible due to the availability of the CLUSTER spacecraft, a four identical-spacecraft mission which during this measuring period had an inter-spacecraft separation distance between ~ 1 and $\sim 1.5 R_E$. They used a nominal bow shock model [Peredo *et al.*, 1995], based on the measured upstream flow parameters (basically the dynamic pressure of the flow), providing the shock-spacecraft distances along the magnetic field flux tube for the individual CLUSTER spacecraft and the local shock-normal angles Θ_{Bn} . The average Θ_{Bn} over the whole 10 hours of observation time was $20^\circ \pm 8^\circ$ proving that CLUSTER was in front of the quasi-parallel shock far away from the ion-foreshock boundary, and the Mach number was $\mathcal{M}_A \sim 8$.

Diffuse ion partial densities were determined as function of distance from the shock and in four consecutive energy bands every 32 s at two spacecraft. This allowed to determine the partial diffuse ion-density gradients along the magnetic field as function of energy from the density differences between two spacecraft and the differences of the spacecraft distances along the magnetic field to the shock intersection point. Perpendicular density

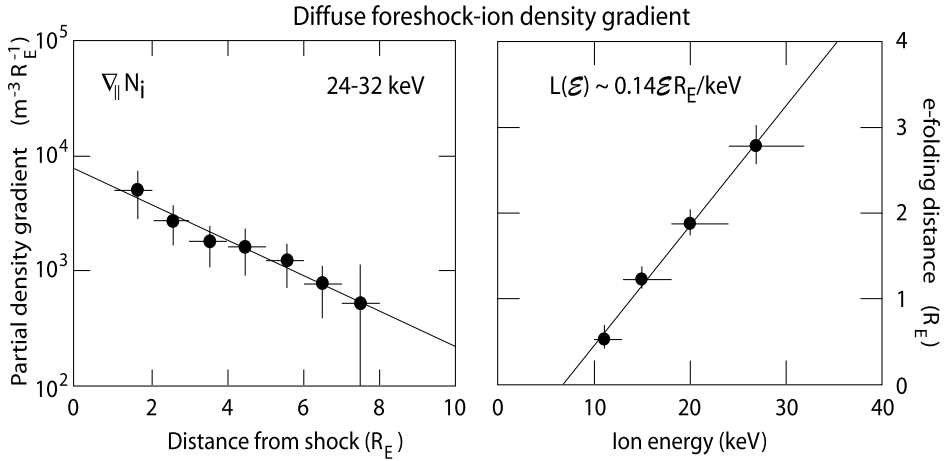


Figure 6.10: Partial-density gradient of the diffuse foreshock-ion component along the shock-connected magnetic field, and density e-folding lengths, determined from CLUSTER observations in the Earth's bow shock foreshock [data from *Kis et al., 2004*, courtesy American Geophysical Union]. *Left:* The parallel partial diffuse-ion density as function of distance from the shock along the magnetic field flux tube connecting the spacecraft to the shock. The diffuse ion density falls off exponentially with distance suggesting an ion-energy dependent diffusive process being responsible for transport of the ions upstream of the shock. *Right:* The e-folding distance of diffuse ions along the magnetic field as determined as function of energy from the exponential decay of the ion density. The e-folding distance increases linearly with ion energy.

gradients were neglected. The obtained parallel gradients were attributed to the average CLUSTER location. These gradients were then used to find the e-folding distance of the density variation.

The results of this investigation are shown in Figure 6.10. It is learned from this figure that in the deep foreshock, i.e. at distances far away from the ion-foreshock boundary, the diffuse ion component is densest close to the shock with density decaying exponentially with increasing upstream distance from the shock along the magnetic field. This spatial decay of the diffuse partial ion density $N_i(\mathcal{E}, z) \sim \exp[-z/L(\mathcal{E})]$ is different for particles of different energy $\mathcal{E}$. The e-folding distance $L(\mathcal{E}) \sim \mathcal{E}$ turns out to increase linearly with energy, i.e. low energy particles are confined to the shock. The higher the ion energy the deeper can the ions penetrate into the upstream plasma. The proportionality constant determined from these data under the special conditions of the Earth's bow shock is $\sim 0.14 R_E / \text{keV}$. This behaviour of the energetic foreshock ions provides indisputable evidence for the extended parallel shock-surface origin of the diffuse ion component. The source of the diffuse ions lies at the quasi-parallel shock. In order to be found at a distance upstream of the shock the ions undergo a diffusion process along the magnetic field. These ions are thus completely different from the beam ions found at the ion-foreshock boundary.

The e-folding distance for the diffuse ions is given by $L(\mathcal{E}) = \kappa_{\parallel}(\mathcal{E})/V_1$, with spatial diffusion coefficient $\kappa_{\parallel}(\mathcal{E}) = \frac{1}{3} v \ell_{\parallel}(\mathcal{E})$, where $\ell_{\parallel}$ is the diffusion length (parallel ion mean free path) and v the particle velocity (note that the diffusion coefficient has the correct

dimension $[\kappa_{\parallel}] = \text{m}^2\text{s}^{-1}$; justification of the diffusion assumption will be given in the Chapter 7 on particle acceleration). From balance between convective inflow and diffusion into upstream direction, one can write $\ell_{\parallel}(\mathcal{E}) = 3L(\mathcal{E})\sqrt{\mathcal{E}_1/\mathcal{E}} \sim \sqrt{\mathcal{E}} \cdot E_1$. The diffusion length increases as the root of the product of particle energy $\mathcal{E}$ and upstream flow energy $\mathcal{E}_1$. In the solar wind the flow energy is a few keV, and a 20-keV diffuse ion will have a typical parallel diffusion length (or mean free path) of $\ell_{\parallel} \sim (1-2)R_E$. This is a rather short distance, orders of magnitude shorter than the collisional mean free path of an ion. Hence, the diffusion estimate suggests that strong wave-particle interactions can be held responsible for the scattering and acceleration of the diffuse particle component, which enables it to diffuse and escape upstream from the shock and populate the foreshock. The diffusion process is energy dependent with the most energetic ions diffusing fastest.

These interactions should take place in the quasi-parallel shock transition because, as we have shown above, the diffuse upstream-ion density maximises closest to the shock. It is interesting to estimate the corresponding upstream-ion collision frequency $\nu_{c,ui} \simeq v/\ell_{\parallel}$. For the 20 keV-upstream ions this yields $\nu_{c,ui} \sim 0.2$ Hz. This value is comparable to the ion cyclotron frequency $\omega_{ci}/2\pi = (0.1-0.3)$ Hz in the $B \simeq 8$ nT upstream to $B \simeq 30$ nT shock ramp magnetic field [Kis *et al*, 2004] during the time of observation.

It seems that waves, electromagnetic and/or electrostatic, related to the ion-cyclotron frequency are involved into the process of upstream ion diffusion. Since this diffusion is energy dependent, this process is not a simple pitch-angle diffusion as in the case of the generation of the ion beam that propagates along the foreshock boundary. The diffuse ion component must have experienced a substantial acceleration in this process, and this acceleration is located at or around the shock transition and contrasts with the ion-beam acceleration which is a scattering process followed up by pick-up acceleration when the upstream-propagating beam ions are subject to the effect of the main-bulk-stream convection-electric field in which they become accelerated in the direction perpendicular to the magnetic field to roughly four times the energy of the bulk flow, thereby evolving into the ring distribution that characterises their phase space distribution.

6.2.2 Low-Frequency Upstream Waves

In the frame of the upstream bulk flow the two ion components that populate the ion foreshock carry a substantial amount of free energy which is subject to dissipation. Since this dissipation is collisionless it can proceed only through the excitation of waves and wave turbulence through instability upstream of the quasi-parallel supercritical shock. On the other hand it is obvious that the presence of neither of the components can be understood without complete knowledge of the waves in the foreshock and their interaction with the particles.

Since their first detection by Olson *et al* [1969], Russell *et al* [1971], Fairfield [1974] and Fairfield & Feldman [1975], observation of shock-upstream waves has been a long-standing issue. Their existence was predicted by Tidman & Northrop [1968], followed by hydromagnetic [Barnes, 1970] and kinetic [Hasegawa, 1972] theories of electromagnetic wave excitation and propagation upstream of a collisionless quasi-parallel supercritical shock. Wu [1972] suggested that they might develop into discrete wave-packets as had

been inferred from observation by *Russell et al* [1971]. The ISEE 1-3 spacecraft allowed for a more elaborate investigation of the properties of upstream waves [cf., e.g., *Hoppe & Russell*, 1981, 1983; *Russell & Hoppe*, 1981, 1983; *Sentman et al*, 1981a, b, 1983; *Hoppe et al*, 1982; *Thomsen*, 1985; *Mellott*, 1986; *Russell*, 1988; *Le & Russell*, 1992, among others]. More recently, CLUSTER measurements have been used to investigate the temporal and spatial structure of upstream waves and wave turbulence [cf., e.g., *Mazelle et al*, 2000, 2003; *Eastwood et al*, 2002, 2003; *Meziane et al*, 2004; *Narita et al*, 2004, 2006; *Narita*, 2007, among others]. We will briefly review the properties of the upstream waves in the ion-beam and diffuse ion region in view of the observations and mechanisms of their generation. Wave generation is coupled to particle-energy loss and to particle scattering both playing a substantial role in particle acceleration. It will therefore be quite natural that in the next Chapter 7 on particle acceleration at shocks we will return to the upstream-wave problem.

Burgess [1997] gave a comprehensive review of the various types of waves encountered upstream of quasi-parallel shocks in the ion foreshock. In his words, “upstream particles cause upstream waves. . . . Once a wave is created. . . , it then propagates, and so its continued existence relies on it remaining in a region where it is undamped. Its properties might even change as it propagates. Observationally, the wave propagation is superimposed on the convection of the plasma frame, which introduces Doppler shifts in frequency, and possible reversal of polarisation sense. . . . And that is not the end of the story, since one must take account of the feedback of the waves on the particle distribution function. . . , and even the possibility that the. . . shock injection of particles into the foreshock is modulated either by intrinsic processes or even by the foreshock waves themselves.”

Burgess’ review was organised by the observed wave frequencies. He distinguishes between Low Frequency Waves (5 mHz- few 100 Hz) and High Frequency Waves (>1 kHz), the latter covering the electrostatic waves from ion-sound to electron plasma waves, as well as radiation. Radiation generation will be discussed in detail later. Here, we only note the almost continuous presence of waves in the ion-acoustic band which have been known since *Rodriguez & Gurnett* [1975] to populate the complete foreshock region. These spectra might be composed of several different modes, ion-sound, electron-acoustic, Buneman modes, electron-cyclotron harmonics, and others. Their generation mechanism is not clear yet. They might, via a number of different instabilities, be the result of the presence of the hot foreshock-electron component, which also invades the ion foreshock, or they are excited by unresolved narrow electron beamlets that emanate from the quasi-parallel shock. They might also be excited by plasma inhomogeneities, spatial inhomogeneities in the electron distribution, or they are the result of nonlinear wave-wave interaction which is expected to take place in the foreshock. Currently these questions are difficult to answer and await further observation, simulation and theory. Little has changed so far since Burgess’ remarks concerning high frequency waves.

However, there has been substantial progress in the understanding of the low frequency waves and their role in quasi-parallel shock dynamics. The waves that are most important in shock formation propagate in the ultra-low frequency range <0.1 Hz. Usually they have large (magnetic) wave amplitudes, around $|\mathbf{b}|/B \sim 0.2-1.0$, which identifies them as highly nonlinear. These large amplitude waves had already been observed by *Russell et al* [1971] to have wave forms from monochromatic to solitary waves, frequently with

steep edges resembling shocklets and suggesting that the waves have experienced non-linear steepening during their evolution and propagation. They sometimes show the typical signs of fluctuations that are connected to these edges and obviously propagate in the whistler mode. Thus these forms are indeed little shock-like structures.

In the same frequency window, large-amplitude pulsations have been identified. These are very typical for quasi-parallel shocks. *Schwartz & Burgess [1991]* and *Schwartz et al [1992]*, identifying them in the AMPTE magnetic field measurements, coined the (some-what ugly as in German it means mud) name SLAMS for them, which stands for ‘Short, Large Amplitude Magnetic Structures’. We prefer to call them upstream pulsations, here. Their duration is 10–20 s; they have very large amplitudes $|\mathbf{b}|/B \sim 5$, indeed, but appear as a more coherent structure that is embedded in the ultra-low frequency wave turbulence. Like the ultra-low frequency waves, they propagate in upstream direction in the plasma rest frame while being swept toward the shock by the convective flow. Their polarisation is mixed with – possibly – left-hand polarisation (in the plasma rest frame) slightly dominating, suggesting their ultra-low-frequency wave origin. Sometimes the polarisation is different on both sides of the upstream pulsation, indicating that they have evolved by some process which produces both kinds of polarisation, which is similar to a solitary wave.

The propagation velocity of the upstream pulsations exhibits an interesting amplitude dependence. The upstream directed pulsation speed increases with amplitude, which also is a solitary wave-like property. Moreover, they grow when approaching the shock and entering the increasing density gradient of diffuse ions, and they play an important role in shock reformation. *Schwartz et al [1992]* suspect that these pulsations are the ‘building blocks’ [an expression used by *Schwartz & Burgess, 1991*] of quasi-parallel shocks.

Another interesting property is that upstream pulsations contain thermal plasma with properties of the upstream flow, while being surrounded by the hot foreshock plasma. One would therefore believe that their source region is located at the ion-foreshock boundary. Being created there by an ion-ion beam plasma interaction they might grow nonlinearly until reaching quasi-equilibrium like solitary structures, having captured the upstream plasma, and afterwards being convected toward the shock into the heart of the ion foreshock. We will return to these interesting structures when discussing simulations below.

Ion-Beam Waves

Each of the two different upstream-ion populations is responsible for the excitation of its own instabilities. In this section we deal only with those waves which are excited by the foreshock-boundary ion beam. From Figure 6.6 suggests that the situation is that of an ion-ion beam interaction (for the moment neglecting the gyrophase-bunched diffuse-ion component). In the frame of the upstream flow the reflected foreshock-boundary ion beam propagates upstream along the magnetic field at parallel speed $v_{b\parallel} \simeq -(2-3)V_{1\parallel}$, where $V_{1\parallel} \simeq 600 \text{ km s}^{-1}$. The upstream ion temperature (in energy units) was $T_i \simeq (1-2) \text{ keV}$, yielding a thermal ion velocity of $v_i \simeq 200 \text{ km s}^{-1}$. The beam can be taken as warm with thermal speed $v_{b,th} \lesssim v_{b\parallel}$. Hence, $v_{b\parallel} > v_i$ with magnetised background and beam ions. On the other hand, the electrons are hot with $T_e \simeq 100 \text{ eV}$.

The Expected Wave Modes. Since the large parallel speed of the beam corresponds to a large parallel temperature anisotropy, it is clear that the beam can excite long-wavelength negative-helicity Alfvén waves via the firehose instability. In addition, because the beam is mildly warm, it can excite the resonant left-hand ion-ion beam instability, which is possible for $v_{b\parallel} > V_A \sim 100 \text{ km s}^{-1}$. Under conditions, when the beam thermal speed can be considered to be small, it excites right-hand resonant ion-ion beam modes. Both modes propagate upstream along the magnetic field with the beam on the background of the upstream flow. They are not as fast as the beam, however, and are thus subject to downstream convection with the flow towards the shock. The firehose mode, at the contrary, moves against the beam and thus by itself moves downstream in the direction of the shock position, when excited.

In all three cases the foreshock-boundary beam will lead to the excitation of low-frequency Alfvén and ion cyclotron waves which in the shock frame approach the shock while having their source on the foreshock-boundary field line. During this shock-directed convection and/or propagation they populate the foreshock region with low frequency electromagnetic fluctuations, which might further interact with the diffuse foreshock-ion component. On the other hand, since the phase and group velocities $\omega/k_{\parallel} \sim (\partial\omega/\partial k_{\parallel}) \sim V_A \ll v_{b\parallel}$ of these waves are of the order of the Alfvén velocity and are, thus, much less than the beam and flow velocities, downstream convection will quickly remove them from the foreshock-boundary source region. Hence, their further evolution in the foreshock is determined by the competition between nonlinear wave steeping and – most important as will become clear below – interaction with the diffuse foreshock-ion component.

Including the electrons (while so far neglecting the electron foreshock component) leaves us with an ion-acoustic unstable phase-space configuration. In the upstream ion-plasma frame a relatively dense foreshock-boundary ion-beam is propagating upstream on a cold ion-hot electron plasma. In order to keep the plasma current-free the electron component is slightly retarded, creating conditions under which ion-acoustic waves can be excited. On the other hand, the configuration is not able to excite neither the Buneman-two stream instability nor – because the beam propagates solely parallel to the magnetic field – the modified-two stream instability.

This is shown in Figure 6.11. The instability is excited by the velocity difference between the ion foreshock-boundary ion beam $F_{ib}(v_{\parallel})$ and the slightly shifted to the left hot ($T_e \sim 100 \text{ eV}$) electron distribution because the velocity difference ΔV between ion beam and electron component $v_e > \Delta V \sim 10^3 \text{ km s}^{-1} > c_{ia} \sim 100 \text{ km s}^{-1}$ exceeds the ion acoustic speed while being less than the electron thermal velocity v_e . These ion acoustic waves occupy a relatively broad spectrum with downstream parallel phase velocities $< v_{ib}$ and become swept towards the shock. Since they quickly leave the foreshock boundary and since the beam is not hot, ion-Landau damping in the source region plays no role. However, when entering the diffuse-ion foreshock region these waves encounter the hot diffuse ion component, will interact with it, and will thereby experience (positive and negative) Landau damping.

Observations. Observationally, it is difficult to distinguish between foreshock-boundary waves and diffuse-ion generated waves. Two types of waves that can be related to the

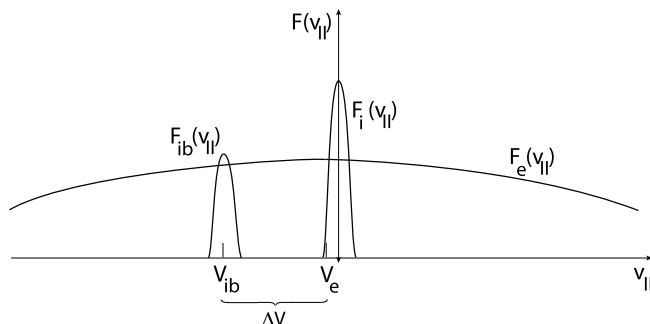


Figure 6.11: Reduced parallel distribution functions along the ion-foreshock boundary showing the cold main flow ion distribution $F_i(v_{\parallel})$, the hot main flow electron distribution $F_e(v_{\parallel})$ shifted to the left into the direction of the foreshock ion beam for keeping zero-current conditions, and the ion-foreshock ion beam distribution $F_{ib}(v_{\parallel})$. This configuration is unstable with respect to ion-ion beam instabilities and the ion-beam driven ion-acoustic instability.

presence of the upstream beams have been reported in the vicinity of the bow shock. *Fairfield* [1974] and *Fairfield & Feldman* [1975] found monochromatic large-amplitude discrete wave packets in the frequency range $\omega/2\pi \sim 0.4$ Hz which *Hoppe & Russell* [1980] could show to be right-hand polarised whistlers at $\omega \simeq 10\omega_{ci}$ propagating on the plasma rest frame and being unable to escape far upstream. The other class of waves is of smaller amplitude $|b|/B \sim 0.1$ and frequency of the order of $\omega/2\pi \sim 1$ Hz. These waves form trains which are directly tied to the upstream foreshock-boundary ion beams [*Hoppe et al*, 1982]. In fact, as *Hoppe et al* [1982] demonstrated they occur only in the presence of the observed reflected 2–5 keV foreshock-boundary ion beams [*Eastman et al*, 1981] which *Kis et al* [2007] have shown to evolve from a beam into a gyrating particle component that in the deep foreshock is superimposed on the diffuse ion component before it merges into it. However, as *Feldman et al* [1983] have shown from measuring the electron distribution function, it is not the ion beam who excites these whistlers. The probability that they are driven by the particular electron distribution in the foreshock is rather higher.

The waves observed propagated obliquely ($<60^\circ$) with respect to the magnetic field, at an average angle of $\sim 45^\circ$. The observed spectrum and the dispersion relation determined from the measurements are shown in the high-frequency part of Figure 6.12. This figure has been combined from the high-frequency observations of *Hoppe et al* [1982] and the low-frequency observations of *Le & Russell* [1992], both obtained from the ISEE spacecraft measurements. Tentative wave numbers have been determined for the high-frequency waves from measuring the time delay of the wave front arrivals at the two spacecraft ISEE 1 and ISEE 2 spacecraft yielding (surprisingly) short wavelengths <100 km. Knowing the wave number, the frequency has been back-Doppler shifted $\omega = \omega_{\text{ISEE}} - \mathbf{k} \cdot \mathbf{V}_1$ into the plasma frame yielding unusually high frequencies $\omega \simeq (20\text{--}100)\omega_{ci}$. This procedure determines the dispersion relation (naturally with large errors as indicated by the bars), which is found to be about linear in the narrow high frequency range of the waves in

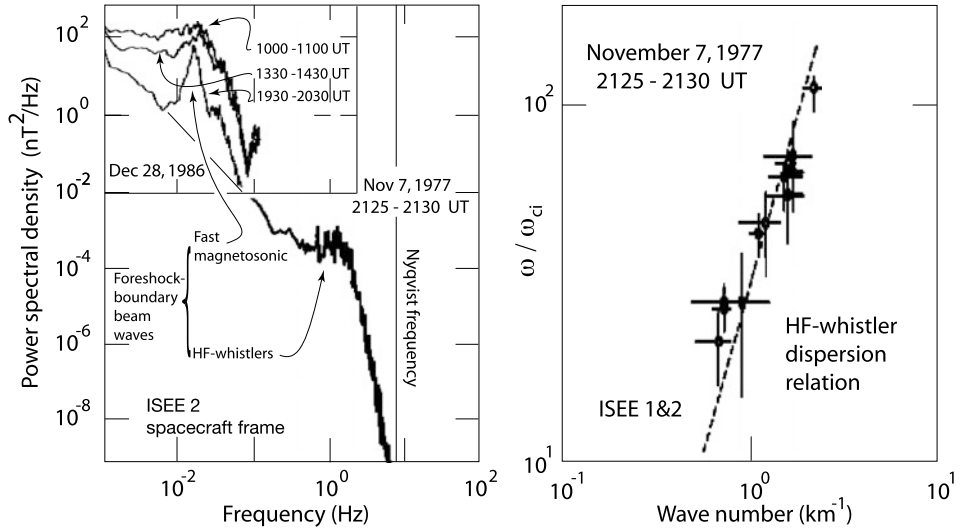


Figure 6.12: *Left:* The combined power spectral density of magnetic fluctuations excited by the ion foreshock-boundary ion beam with central frequency around ~ 1 Hz [lower part, data from *Hoppe et al., 1982*, courtesy American Geophysical Union] and at < 0.1 Hz [left upper part, data from *Le & Russell, 1992*, courtesy American Geophysical Union] as measured by the ISEE spacecraft. This spectrum is obtained in the ISEE spacecraft frame. The low frequency spectrum is shown at three different times corresponding to (from below upward) increasing distance from inside the foreshock to the foreshock boundary. It is about five orders of magnitude more intense than the high-frequency waves. It is worth noting that the high frequency spectrum of *Hoppe et al. [1982]* toward lower frequencies in the overlap with the *Le & Russell [1992]* spectrum does not show any indication of the lower frequency peak. As indicated by the *thin straight line*, it would smoothly continue into the 1930 UT branch at 0.01 Hz. This lack of the low frequency peak is probably accidental and due to the particular conditions at the time of measurement. *Right:* The wave dispersion relation for the high frequency ~ 1 Hz whistler-mode waves determined from a tentative estimate of the relevant wave numbers and transformed into the plasma rest frame. It is seen that these waves in the observation range have about linear dispersion and frequencies $\omega \sim (20-100) \omega_{ci}$, far above the ion-cyclotron frequency ω_{ci} [data from *Hoppe et al., 1982*, courtesy American Geophysical Union].

the rest frame of the plasma. The waves that have been left-hand polarised in the spacecraft frame turn out to become right-hand polarised in the plasma rest frame.

From these properties it was initially concluded that the waves propagate in the whistler branch and have been excited by the cool ion-ion beam instability, which would be consistent with our initial discussion. This conclusion is, however, questionable. The high rest-frame frequencies are not in agreement with model calculation [*Sentman et al., 1981a; Gary, 1993*] using the realistic observed beam properties. These yield wave frequencies of the order of $\omega \sim 0.1 \omega_{ci}$ for both the ion-ion beam and firehose modes. In fact, *Eastwood et al [2005b]* report CLUSTER observations of similar waves but with much lower frequency $\omega/\omega_{ci} \sim 0.1$ and wavelength the order of $\sim 1 R_E \sim 6000$ km, which is in excellent agreement with the theoretical predictions. These waves are cold ion beam-excited fast-kinetic (magnetosonic) whistlers, in the terminology of *Gary [1993]* and *Krauss-Varban et al [1994]*. The high-frequency waves ~ 1 Hz whistlers are instead most probably driven by the particular electron distribution in the foreshock [*Feldman et al., 1983*] rather than

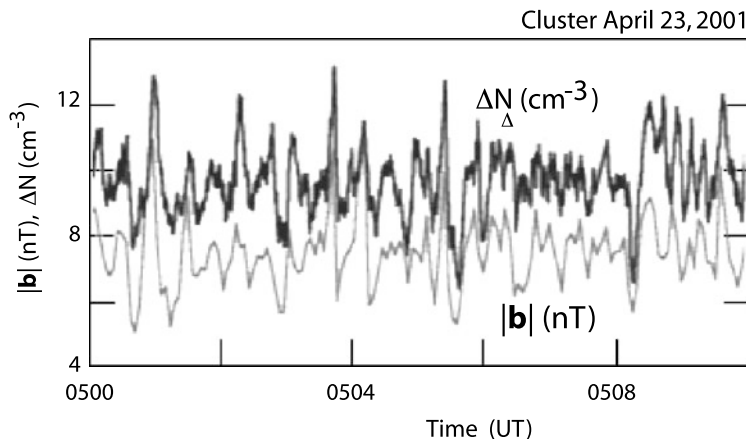


Figure 6.13: *Left:* The excellent correlation between the density and magnetic field variations in the ion-beam excited foreshock-boundary low-frequency waves as measured by CLUSTER [data from *Eastwood et al., 2002*, courtesy American Geophysical Union]. The variations are practically in phase thus identifying the fluctuations as fast magnetosonic.

by the cold-ion beam instability. This claim is also supported by the fact that deeper in the ion foreshock, where the electron distribution becomes more isotropic, these waves do not occur separately from smaller magnetic field structures (large-amplitude magnetic pulsations of the SLAMS type or shocklets).

In quite good agreement with the CLUSTER measurements are the ISEE foreshock waves analysed by *Le & Russell [1992]* who found the spectral peak at frequencies $\omega \sim 10^{-2}$ Hz with the spectrum broadening with increasing distance from the ion-foreshock boundary. This is shown in the upper left part of the spectrum in Figure 6.12 where the measurements of *Hoppe et al [1982]* and *Le & Russell [1992]* have been combined. Even though the conditions on the two observation times were not identical, one finds that the high-frequency whistlers identified by *Hoppe et al [1982]* occur on the approximate high-frequency extension of the *Le & Russell [1992]* 1930–2030 UT spectrum which is closest to the ion foreshock boundary. Note, however, that the low-frequency part of the *Hoppe et al [1982]*-spectrum (not shown here other but indicated by a thin straight line in the figure) was flatter and in the (relatively short) overlap region with the *Le & Russell [1992]*-spectrum did not show any indication of the magnetosonic foreshock-boundary waves, which must be due to the particular conditions prevailing during this observation period. Thus the *Hoppe et al [1982]* waves are a (high frequency) wave species that is different from these more common low-frequency/long-wavelength fast-magnetosonic waves. These latter waves are also of much higher (~ 5 orders of magnitude) spectral density as seen from Figure 6.12 and has been confirmed by the later CLUSTER measurements. Moreover, the magnetosonic waves are left-hand polarised in the plasma frame, and their compressive fast-magnetosonic character is proved from the in-phase variation of the density and magnetic field fluctuations shown in Figure 6.13.

The presence of these low-frequency ion-beam generated waves implies that the ion beams interact with the fluctuating electromagnetic field. In this process they become scattered and diffuse in phase-space, which is the reason for the spreading of the ion beam in phase space and the final merging into the diffuse background distribution. *Kis et al* [2007] have followed this evolution of the beam as we have described above. It is responsible also for the gradual spreading of the spectrum in the upper left part of Figure 6.12. A spectrum similar to that given there can be found in the paper of *Kis et al* [2007]. *Archer et al* [2005], from CLUSTER observations, inferred the nature of these ultra-low frequency (ULF) waves and showed that their correlation lengths along the wave vector direction $\mathbf{k}$ is of the order of $1-3R_E$, while it can be a factor of three larger in the direction perpendicular to $\mathbf{k}$, rendering these waves oblate though nearly planar.

During times of hot beams *Eastwood et al* [2003] observed upstream propagating left-hand low-frequency waves which have been excited on the Alfvén-ion-cyclotron branch of the kinetic dispersion relation. These waves are kinetic Alfvén waves which have been excited by the hot upstream propagating beam similar to the one shown on the right in Figure 6.6. In the spacecraft frame these waves because they are swept downstream by the flow, appear as right-handed waves. *Eastwood et al* [2003] report that with onset of the waves the beam gets more diffuse. This can either be interpreted as the reaction of the waves on the beam or that the spacecraft enters the region where the initially cold beam enters the foreshock, spreads in velocity and after becoming hot enough generates the observed waves.

The question arises whether there are simulation studies available of the evolution of the reflected ion beam and the generation of upstream waves. This question cannot be definitely answered at present. It is clear that a simulation of this kind should reconstruct the region of the foreshock boundary which requires that a curved shock surface must be assumed from the beginning. In other word, this problem can be investigated only in a two-dimensional simulation. Two-dimensional simulations are available in hybrid form but have not been applied to curved shocks. Hence, this problem remains an open simulation problem. On the other hand, in the following we will extensively discuss the role diffuse particles play in shock formation and upstream wave generation for planar shocks in one and two dimensions and hybrid as well as PIC simulations.

Diffuse Ion Waves

Long-period ~ 30 s waves are the rule in the foreshock. They occur together with the diffuse foreshock-ion component [cf., e.g., *Sanderson et al*, 1983; *Mellott*, 1986] which is located deeper inside the foreshock. Because of this reason, any waves that are excited by the diffuse ion component are restricted to the interior of the foreshock. One even has defined some fuzzy boundary of these waves called ULF-wave boundary [*Russell & Hoppe*, 1983; *Greenstadt & Baum*, 1986], which is even more inclined against the upstream magnetic field than the ion-foreshock boundary and outside of which the ULF-wave activity should be weak. The reason is that ULF waves, if propagating upstream, can move at most with fast-magnetosonic velocity which for a supercritical shock is smaller than the stream and also less than the reflected ion beam velocities. The advection by the

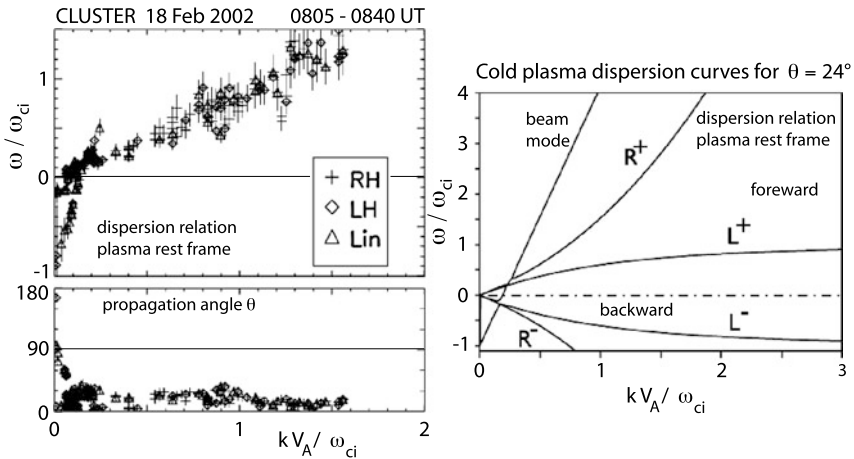


Figure 6.14: *Left:* The dispersion relation of low-frequency waves as measured by CLUSTER in the foreshock away from the foreshock boundary [data taken from [Narita & Glassmeier, 2005](#), courtesy American Geophysical Union], determined from the “wave-telescope” analysis method. The *upper panel* shows the dispersion relation $\omega(k)$ consisting of several parts and distinguished by the sense of polarisation as right-hand, left-hand and linear. The different branches cluster. There are forward (positive frequency parallel to the magnetic field) and backward (negative frequency antiparallel) branches. The *lower panel* shows the wave propagation angle for the shorter wavelengths being mostly oblique $\theta \approx (30-25)^\circ$, for longer wavelengths being more perpendicular to the magnetic field. *Right:* The five low-frequency modes in a cold beam plasma system of similar conditions as the foreshock plasma and for the propagation angle $\theta = 24^\circ$ showing similarity to the observed dispersion relation [data from [Narita et al, 2003](#), courtesy American Geophysical Union].

flow will blow them downstream towards the shock and confine them to a region relatively close to the shock [[Meziane & d’Uston, 1998](#)] bounded to upstream by the ULF-boundary.

[Narita & Glassmeier \[2005\]](#) determined the dispersion relation of the low-frequency waves in the foreshock from the three-dimensional observations of CLUSTER. Their result is shown in Figure 6.14. The left part of the figure shows the dispersion relation in the plasma rest frame and the angle of wave propagation with respect to the local average upstream magnetic field. The scatter is quite large. Nevertheless it is surprising to find a well expressed nearly linear part on the dispersion relation even though the dispersion relation is composed of the contributions of several distinct modes with different polarisations. The representation is linear, which, on the k -axis, emphasises the short wavelengths (larger k values). From linear theory one expects that the long wavelengths (small k) should be more pronounced. Unfortunately, this region could not be resolved sufficiently at the available CLUSTER separation distances at that time [[Narita et al, 2003](#)]. These low frequency waves, in the plasma frame, seem to propagate close to perpendicular to the magnetic field and, therefore, are probably in the fast magnetosonic mode. Moreover, it seems that a straight beam dispersion relation contributes to the perpendicular propagating waves. It also seems that in the low-frequency (magnetosonic) waves have negative frequency (i.e. propagating in opposite direction). Linear dispersion theory for the parameters during the observation time yields the curves on the right in the figure showing the cou-

pling between the presumable beam mode and the four plasma modes at wave numbers $kV_A/\omega_{ci} < 0.3$. The experimentally determined dispersion relation resembles this clustering of couplings at small k and very low frequencies ω . However, application of linear theory is dangerous if not questionable because the foreshock plasma is highly disturbed, the low frequency waves have rather large amplitudes and can barely be considered as linear waves. In addition the plasma is highly inhomogeneous exhibiting steep gradients in density and magnetic field. For an Alfvén speed $V_A \simeq 30 \text{ km s}^{-1}$ and an ion-cyclotron frequency $\omega/2\pi \simeq 1 \text{ Hz}$, wavelengths around the linear wave coupling are of the order of $k/2\pi \simeq 100 \text{ km}$ and should thus be affected by the plasma inhomogeneities (note that decreasing the reference cyclotron frequency decreases the wavelength even further). Nevertheless, at the very low frequencies and very long waves the waves are probably in the fast magnetosonic wave band.

More mysterious are the higher frequency-short wavelength waves. From comparison with the linear dispersion relation they seem to fit on the right-handed whistler branch (high frequency magnetosonic whistler R^+). However, inspection of the dispersion relation indicates that all kinds of polarisation are scattered along the dispersion curve. The curve itself is very irregular even though the propagation angle of the waves seems to be about constant at weakly oblique angles. This wave composition suggests that we are not dealing here with one single wave mode but rather with the short wavelength part of a turbulent spectrum which has generated all kinds of short wavelength fluctuations with different almost randomly distributed polarisation in a forward cascading process from long to short wavelengths. This idea has been elaborated in more detail in *Narita et al* [2006] who found that the shorter wavelength spectrum is indeed about featureless and power-law, close to a $-\frac{5}{3}$ -Kolmogorov-spectrum of stationary turbulence with most of the power in the perpendicular (non-compressive) magnetic component providing another argument for Alfvénic and magnetosonic turbulence. This is shown on the left in Figure 6.15. The figure also indicates the respective gyro- and inertial wave numbers, $k_{ci} = v_i/\omega_{ci}$ and $k_{in} = c/\omega_{pi}$. There is indication that the spectra change slope at around these numbers due to onset of anomalous ion viscosity and ion inertia. The large maximum on the perpendicular power curve near $k \sim 10^{-3} \text{ km}^{-1}$ corresponds to a wavelength of $\lambda \sim 6000 \text{ km}$ reported earlier [e.g., *Eastwood et al*, 2003] and is at the right position for energy injection by low frequency instability.

However, the conclusion that the deep foreshock is subject to fully developed fast magnetosonic/Alfvénic turbulence is by no means fully justified. This becomes obvious from considering the probability distribution of the field fluctuations as given on the right in Figure 6.15. Would the fluctuations be normally distributed then their distribution functions would have the shape of the dashed curves. Instead the distributions are skewed into tails. In the case of the parallel component the distribution lacks symmetry. Such tails may indicate that the turbulence is non-stationary, intermittent or inhomogeneous. All three cases might hold in the foreshock, in particular as the foreshock is a rather limited spatial region which is bounded from two sides and is subject to plasma injections and plasma losses. Nevertheless, the spectra determined contain signs of strongly nonlinear and turbulent interactions, and the waves, particularly the low-frequency waves, are of large amplitude and interact with the plasma ion component as also with other waves. These waves affect

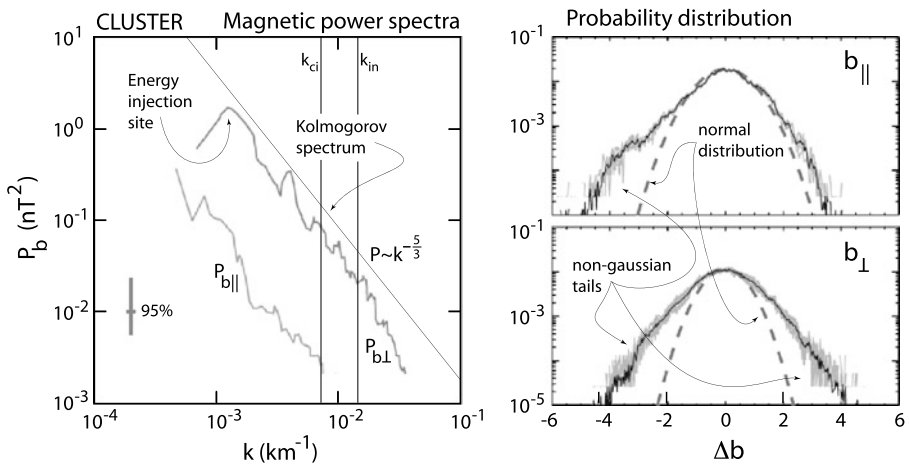


Figure 6.15: *Left:* Short wavelength spectra for the magnetic fluctuations measured by CLUSTER in the foreshock away from the foreshock boundary. These spectra are obtained by application of the “wave-telescope” analysis method. The perpendicular fluctuations are more than one magnitude more intense than the parallel fluctuations while the spectra decay about according to the Kolmogorov $-\frac{5}{3}$ -law of stationary inertial turbulence. Indication of a cut-off is seen at wave numbers larger than the inertial wave number k_{in} where strong dissipation sets on. From the maximum at $k \sim 10^{-3}$ one concludes that in this wavelength range the spectral energy is injected by instability of fast and Alfvénic magnetosonic waves which cascade nonlinearly forward towards shorter wavelengths. *Right:* The probability distributions of the parallel and perpendicular magnetic fluctuations. The dashed distributions are log-normal of same maximum. These distributions exhibit extended tails and thus indicate non-stationary and probably intermittent not fully developed magnetic turbulence [after Narita *et al*, 2006].

the upstream backstreaming ion component scattering, heating and accelerating it. On the other hand the upstream ion component is responsible for the existence of the waves as it is the ultimate energy source of the waves. And, to close the cycle, the waves cannot escape upstream very far from the foreshock since in a supercritical shock the Mach number of the flow is higher than any Mach number based on the wave speed and is thus high enough for the flow to advect the wave spectrum towards the shock. As long as the waves do not completely dissipate their energy during this convection, the electromagnetic wave energy accumulates at the location where they arrive at the shock.

The consequence of this accumulation is that the waves affect the shock, cause instability of the shock surface, and reorganise the shock front. This makes quasi-parallel shocks principally non-stationary (or as sometimes said, turbulent) and subject to some kind of irregular reformation. One may thus expect that a supercritical quasi-parallel shock does not represent a solid shock surface over a large area. It consists of a more or less dense ‘patchwork’ of areas which together form a shock but which also appear and disappear in an irregular manner, come and go, and organise the shock in a certain volume where the entropy increases but where a multitude of very difficult to handle processes take place that can be investigated only experimentally or with the help of properly designed numerical simulations. These are, however, more difficult to design than in the quasi-perpendicular

case, because of the greater variability of the conditions at a quasi-parallel shock. Since waves and particles are tied to each other, simulations cannot be discussed separately for waves and for particles. We therefore delay the discussion of the simulation results to the section on quasi-parallel shock reformation. There it will become obvious what the upstream waves are good for and what their role is in the quasi-parallel shock process.

An important question that has not been addressed anywhere in the investigation of the ion-foreshock dynamics concerns the role of the foreshock electron component which, in the large ion-foreshock domain, consists of two populations, the ~ 100 eV upstream electron population which belongs to the quasi-neutral upstream flow, and the hot $\sim$ several keV-electron population which is the product of the shock reflected electrons we are going to discuss in the next paragraphs. This component is isotropic but irregular and might contribute to ion-foreshock instabilities thereby affecting the foreshock turbulence and shock formation. Since these questions are difficult to treat theoretically one has to wait until numerical simulations will become capable of including them into full particle codes in a similar way as has been done [cf. *Matsukiyo & Scholer*, 2003, 2006] for the electron-ion populations in the feet of quasi-perpendicular shocks.

6.2.3 Electron Foreshock

Even though the ion-foreshock occupies the larger part of the foreshock and, which is more important, plays the decisive role in the dynamics of a quasi-parallel supercritical shock, it would be an ignorant attitude not to mention the part that aside of the main flow is populated solely by the reflected higher energy electron component, i.e. the electron-foreshock region. Schematically its extent and relation to the ion-foreshock is shown in Figures 6.3 and 6.16. The electron foreshock has been identified already by *Scarf et al* [1971] in the OGO 5 satellite observations. *Scarf et al* [1971] noted the occurrence of enhanced electron fluxes near the bow shock, which were related to observations of high-frequency electric-field spikes at 30 kHz. They concluded that these spikes resulted from Langmuir waves that had been excited by electron beams arriving from the bow shock along shock-connected magnetic field lines. (A detailed overview of the early observations can be read in *Klimas* [1985].) It could be confirmed that these electron fluxes were magnetic-field aligned and of higher energy than the bulk electrons in the main flow. This was interpreted as electron beams emitted into upstream direction from the shock, even though no mechanism was known that could provide the required shock-electron acceleration/reflection – in fact, a de Hoffman-Teller-frame shock-reflection mechanism had already been proposed by *Sonnerup* [1969] and was reinvented and elaborated on for electrons much later by *Wu* [1984].

Electron Beams

In analogy to the ion-foreshock boundary the electron-foreshock boundary (which is the ultimate upstream boundary of the foreshock) carries narrow bursts of electron beams, which escape into upstream direction from the shock along the magnetic field and which are slightly displaced by the convective flow into downstream direction.

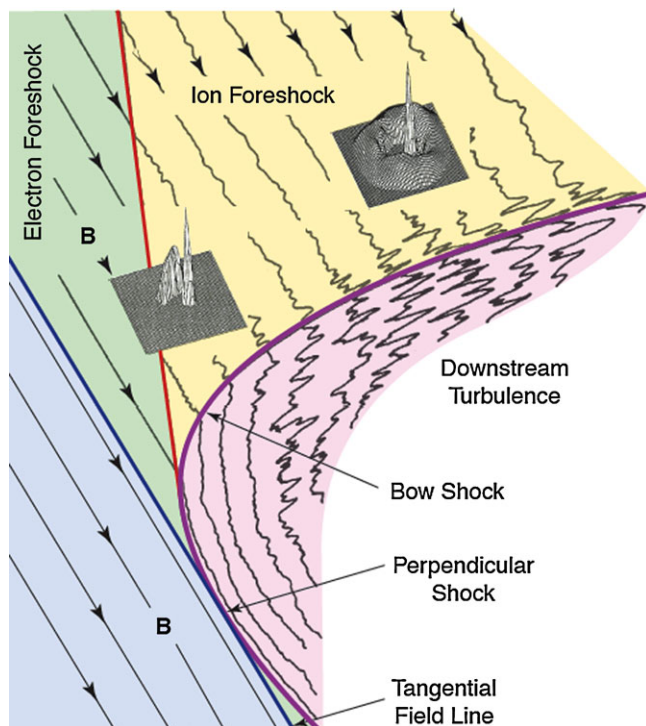


Figure 6.16: Synoptical view of the two foreshocks: the electron and ion foreshocks, respectively.

An example of a narrow energy spectrum of such a beam measured by the IMP-8 satellite is shown on the left in Figure 6.18. The spacecraft was in the electron foreshock boundary for a relatively short time only. Being near apogee, the spacecraft was about standing. Thus the short contact time with the beam most probably indicates that the electron-foreshock boundary is fairly narrow or that the electron beams are bunched into electron bursts. The electron beam also occupies only a small volume in velocity space corresponding to a narrow bump on the distribution [Anderson *et al*, 1981]. Only one energetic electron-gyroradius deeper in the electron foreshock the electron beam becomes depleted and just forms an energetic tail of hot halo electrons on the electron distribution function of the flow.

Figure 6.19 shows electron phase-space observations from the electron instruments on ISEE 1 & 2 during crossings of the electron-foreshock boundary. The left part of the figure [taken from Fitzenreiter *et al*, 1984] is a much higher time and velocity (respectively energy) resolution plot than that in the former figure. It is nicely seen how the electron distribution evolved from a field-aligned (nearly) Maxwellian distribution prior to contact with the foreshock boundary, through a field-aligned-beam-like distribution at contact, into a distribution with an energetic tail along the magnetic field, where the electron beam has been completely washed out. This transition takes place within a time interval of 15 s,

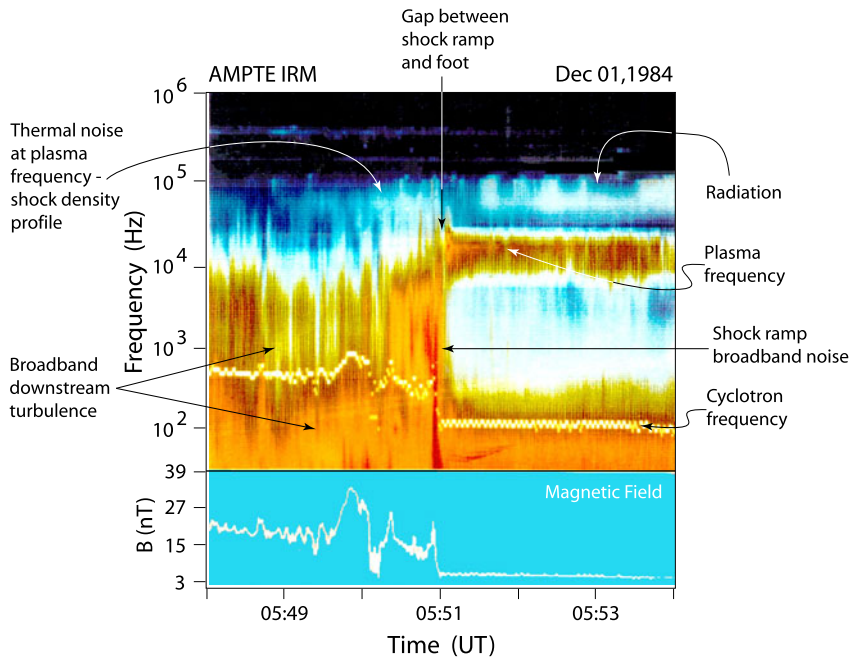


Figure 6.17: AMPTE IRM plasma wave frequency spectrogram during an electron foreshock bow shock crossing [from R. A. Treumann & J. LaBelle, 1986, unpublished]. Note the intense emission in the plasma frequency upstream of the shock transition indicating that the spacecraft is travelling along the electron foreshock boundary. Radiation at the harmonic of the plasma frequency is also emitted. The shock transition and ramp appear as broad regions of intense broadband emission reaching up to the plasma frequency. Across the shock weak noise at the plasma frequency marks the increasing plasma density shock profile.

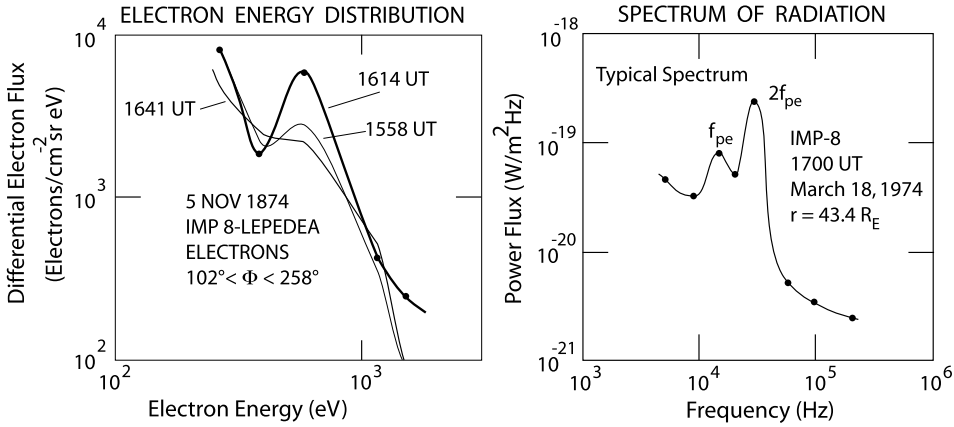


Figure 6.18: *Left:* The electron energy distribution as seen by IMP-8 at three times when crossing the foreshock boundary. *Right:* The high-frequency wave spectrum of the electromagnetic radiation as measured some distance away from the shock near the foreshock boundary. Two types of emissions are visible, the so-called ‘harmonic’ emission $f = 2f_{pe}$ at twice the electron plasma frequency $f_{pe} = \omega_{pe}/2\pi$, and the weaker so-called ‘fundamental’ emission at f_{pe} [after Gurnett & Frank, 1975, courtesy American Geophysical Union].

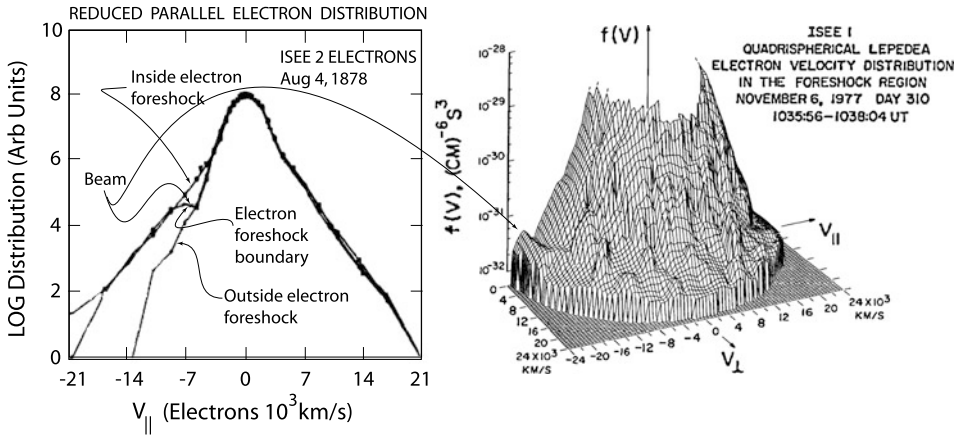


Figure 6.19: *Left:* Three successive reconstructions of the reduced parallel electron distribution function during crossing of the electron foreshock boundary. The distributions are only seconds apart, at 1650:37 UT before touching the foreshock boundary, at 1650:43 UT just crossing it, and at 1650:52 UT being behind it in the electron foreshock. The beam is visible only during the short crossing time. Afterwards the electron distribution shown a heated non-symmetric tail, indicating that the spacecraft sees the depleted beam in the hot electron-foreshock plasma [data from *Fitzenreiter et al., 1984*, courtesy American Geophysical Union]. *Right:* The full two-dimensional electron distribution function during another crossing of the electron-foreshock boundary at the moment when the beam is visible as a narrow enhancement of the distribution in the negative parallel velocities (upstream along the magnetic field). An instant later, when the spacecraft entered into the electron foreshock the beam had disappeared [after *Anderson et al., 1981*, courtesy American Geophysical Union].

just allowing to obtain the three measurements. The pseudo-three-dimensional plot of the electron distribution function on the right [taken from *Anderson et al., 1981*] was obtained on another day when the spacecraft passed the electron-foreshock boundary and a series of distributions were recorded from upstream, across the foreshock boundary, and into the electron foreshock. Here only the foreshock-boundary distribution is shown when the narrow electron beam has evolved flowing into upstream anti-shockward direction along the magnetic field. Note the narrow angular extension of the beam, its weakness, and its comparably large velocity which allows to identify it just outside the bulk of the hot ~ 100 eV background electrons which are much hotter than the beam (indicated by the small beam width in $v_{||}$).

The conclusion that can be drawn from these observations, which in the follow-up time had also been confirmed by measurements of other spacecraft like AMPTE IRM and AMPTE UKS, is that, similar to the ion foreshock, the electron foreshock consists of a very narrow region along about the shock-tangential upstream-magnetic field line (or flux tube) which is populated by a cold, weak, and fast electron beam, and a broad electron region adjacent downstream of this flux tube where the electron distribution exhibits tails and in addition contains a hot, fairly isotropic electron component. The electron beam (like the ion beam at the ion-foreshock boundary) has its origin near the shock-tangential magnetic field line which is connected to the quasi-perpendicular part of the shock. The mechanism of its generation has not yet been unambiguously clarified. Presumably it consists of a combi-

nation of the de Hoffman-Teller frame mechanism of *Sonnerup* [1969] and the bending of the shock surface plus some kind of stochastic acceleration in the electric fields which evolve in the shock transition, ramp and overshoot. Such fields have indeed been reported recently on scales below the electron skin depth $\lambda_e = c/\omega_{pe}$ to be very large, of the order of $\lesssim 100 \text{ mV m}^{-1}$ parallel to the magnetic field and $\lesssim 100 \text{ mV m}^{-1}$ perpendicular to the magnetic field [Bale & Mozer, 2007]. The presence of the hot electron component in the foreshock, on the other hand, is not due to this reflection mechanism but requires another mechanism that heats and accelerates electrons at a quasi-parallel shock over a large area of its surface.

Figure 6.17 shows a crossing of the AMPTE IRM of the electron foreshock boundary in the light of plasma waves. The spacecraft in this crossing is travelling along the electron foreshock boundary all the time observing strong emissions at the Langmuir frequency and at the same time radiation at the harmonic of the plasma frequency. It is interesting that a gap exists between the electron plasma waves and the shock ramp. This gap is probably due to the flux tube along that the electron beam travels at the electron foreshock boundary does not emanate from the shock position where the spacecraft passes the shock ramp. Note also the enormous broadband emission in the ramp which reaches up to the plasma frequency. The extension of the plasma frequency noise to downstream maps the density profile across the shock.

Langmuir Waves

So-called ‘gentle electron beams’, which are just those beams one observes at the electron-foreshock boundary – fast, but not too fast (V_b few times V_1), parallel ($V_b \equiv V_{b\parallel}, V_{b\perp} = 0$), weak ($N_b \ll N_e$), cool ($v_{eb} < v_e \ll V_b$), are known to be the drivers of Langmuir waves with dispersion relation

$$\omega^2(\mathbf{k}) = \omega_{pe}^2 + 3v_e^2 k^2, \quad v_{\text{res}} > \sqrt{3} v_e \quad (6.1)$$

via the resonant kinetic gentle-electron beam instability. The Langmuir resonance condition is $\omega - \mathbf{k} \cdot \mathbf{v}_{\text{res}} = 0$, where $\mathbf{v}_{\text{res}} \sim \mathbf{V}_b$ is the resonant electron velocity. Combining it with the dispersion relation, it is easy to see that in order for the wave number k to be real the resonant velocity must satisfy the condition on the right in the above equation. Below we will discuss the properties of these waves in more detail in connection with observations.

Stability of the Electron Beam. However, there are a number of severe limitations to this instability which theoretically would inhibit its application to the electron-foreshock boundary beam. The most stringent of these limitations is that it is believed that quasi-linear saturation of the gentle-beam instability should quickly, i.e. within a time as short as about $\Delta t \omega_{pe} \sim 10$, deplete the beam and transforming it into a plateau on the distribution function. For an upstream plasma density $N_e \simeq 5 \times 10^6 \text{ m}^{-3}$ the plasma frequency is $\omega_{pe}/2\pi \simeq 20 \text{ kHz}$. Hence, the beam should become depleted within $\Delta t \simeq 0.08 \text{ s}$. A beam of velocity $V_b \simeq 10^4 \text{ km s}^{-1}$ (which satisfies the above condition on the resonant velocity in an inflow of $T_e \sim 100 \text{ eV}$) would thus propagate just 800 km upstream of the shock along the

magnetic field before being depleted, with the implication being that the beam could not escape from its shock-source region. This distance is much less than the distances from the shock at which both the electron-foreshock boundary beams and the beam-excited plasma wave spectra have been observed. Either quasilinear saturation of the Langmuir instability does not work, or the observed plasma waves are no Langmuir waves. The latter is probably not true. So the question arises of how the beam can avoid becoming depleted due to quasilinear effects.

There are several possibilities of inhibiting beam depletion. It has for long time been believed that collapse of Langmuir waves resulting from modulation instability (also known as oscillating two-stream instability) could shift the beam-excited Langmuir waves out of resonance and remove them from reaction onto the beam. This is a very elegant mechanism which is based on the ponderomotive force the Langmuir waves exert on the plasma. The ponderomotive force is the gradient of the wave pressure which is strongest right in the place where the Langmuir waves are most intense, i.e. in the beam region. For the inert ions of the plasma background the intensity gradient $\nabla|\mathbf{e}|^2$ of the fast large-amplitude oscillating Langmuir electric field presents a pressure force which pushes them out of place thus decreasing the plasma density locally. Since the electrons react immediately and follow the expelled ions in order to maintain quasi-neutrality, the wave-ponderomotive-pressure force drills holes into the plasma which are filled with Langmuir waves. The density variation produced

$$\delta N \approx -\frac{\epsilon_0}{4m_i c_{ia}^2} |\mathbf{e}|^2 \quad (6.2)$$

is nothing else but an ion-acoustic wave. The Langmuir waves, when becoming intense enough, can drive ion-acoustic waves of speed c_{ia} , which exist only due to the presence of the Langmuir wave pressure gradient. The broad Langmuir waves in this way organise into a large number of density cavities, and the waves are removed from resonance with the beam electrons. Enhancement of the outside pressure then forces the holes to shrink. This shrinkage of the hole size shortens the Langmuir wavelength, increases the wave number, reduces the phase velocity and, in this way, shifts the waves into the main electron distribution where they can undergo Landau damping and dissipation. Unfortunately, this kind of modulation instability and collapse of Langmuir waves could never be experimentally approved at shocks. It is therefore quite unlikely (though not impossible) that it occurs and stabilises the beam.

Another possibility of stabilising the beam by pushing the Langmuir waves out of resonance is scattering them on thermal ions. This process reads $\ell + i \rightarrow \ell' + i'$. For purely elastic scattering only the moments $\hbar \mathbf{k}_\ell$ of the Langmuir wave ℓ and $\mathbf{p}_i$ of the ion i are changed according to the momentum conservation equation

$$\hbar \mathbf{k}_\ell + \mathbf{p}_i \rightarrow \hbar \mathbf{k}'_\ell + \mathbf{p}'_i \quad (6.3)$$

where the primed quantities are after collision of wave and ion. Loss of momentum decreases the wave number k_ℓ and shifts the wave phase velocity to higher values out of resonance with the beam; gain of momentum increases k_ℓ and brings the wave phase velocity

down into the bulk of the electron distribution where it is again dissipated. The efficiency of this process is, however, low and presumably insufficient for complete stabilisation of the beam. In addition it strongly depends on the available particle number; it is thus most probable to work in the immediate vicinity of the shock only. It does, however, produce a new Langmuir wave ℓ' that plays a role in the generation of radio-radiation from a shock. Because of this capacity it nevertheless remains to be of interest in the physics of collisionless shocks.

Muschiatti & Dum [1991] investigated this nonlinear wave scattering off ions including the polarisation cloud which accompanies each ion. They found that of the ion distribution the ions with velocity $v \simeq 2v_i$ are the most effective scatterers of Langmuir waves. The inclusion of the polarisation cloud increases the scattering efficiency by a factor of three which does not change the above conclusion. However, the wave number of the scattered wave, its direction of propagation and other wave properties change substantially. Nevertheless, the effects are not strong enough to alone explain the long survival of the electron beam over distances very far from the shock.

A review of all the available theories has been given by *Muschiatti [1990]* who performed simulations of a spatially bounded beam. It seems that many effects act together in order to allow the beam to survive. Weak scattering off ions removes some waves from the beam. The remaining waves are only partially reabsorbed because the fast beam front runs them out such that they stay in the trail of the beam only. In this way only the trail part of the beam is depleted and retarded, which affects only the low velocity component of the beam distribution, assuming that the beam has a small but sufficiently large natural beam spread $\Delta v_{eb} \ll v_e$, as is suggested by the observations shown, for instance, in Figure 6.19. If injection is continuous, this implies that the beam starts pulsating, which may be a reason for the observed strong Langmuir wave variations along the electron foreshock-boundary field line.

Finally, in the fast streaming solar wind there opens up another possibility. The Langmuir waves, riding on the beam, are generally much slower than the beam; since $v_{\text{res}}^2 > 3v_e^2$, their group velocity, which carries the energy that can be absorbed, is at most $v_{\ell g} < v_e \sqrt{3}$. So contrary to the beam electrons they are very vulnerable to the downstream convection. The convective $\mathbf{V}_\perp$ -motion of the upstream flow thus sweeps them readily out of the beam region at the electron foreshock-boundary field line removing them from quasilinear reaction, which is probably the main and simple reason for the beam to survive over long distance. Survival of the beam is thus probably simply due to the fast high Mach number upstream flow. All these arguments taken together explain why the beam can survive and, in addition, why the region across the electron foreshock boundary where the beam is detected is as narrow as found in the observations.

Electron Foreshock-Boundary waves. Figure 6.20, taken from *Décrou et al [2001]*, shows a beautiful double passage of the CLUSTER-spacecraft quartet through the Earth's bow shock on December 22, 2000 in the light of the CLUSTER plasma wave spectrum recorded by the WHISPER plasma wave instrument aboard the four CLUSTER spacecraft. Figure 6.21 shows an example of these high-resolution measurements.

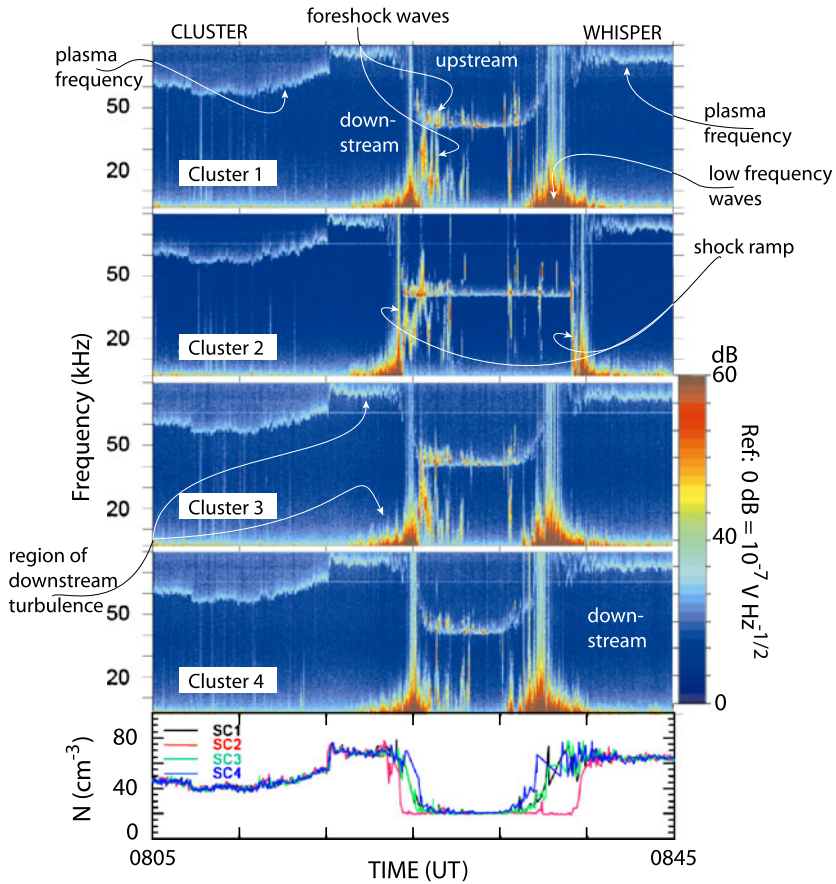


Figure 6.20: A CLUSTER spacecraft passage across the Earth's bow shock region on December 22, 2000. The spacecraft cross the downstream magnetosheath region (*left*) and the shock, enter the upstream solar wind region and back downstream. The figure shows broadband (~ 0 kHz–80 kHz) plasma wave electric field spectra from all four CLUSTER spacecraft (Cluster 1–Cluster 4). The uppermost light blue emissions in each of the panels are at the electron plasma frequency $f_{pe} = \omega_{pe}/2\pi$ which maps the local plasma density N estimates of which are shown in the *lowest panel*. Passages of the shock are signalled by intense broadband waves starting at low frequencies and being correlated with steep drops in N . In the low density upstream solar wind the spots of strong intensification of f_{pe} indicate contact with the electron foreshock boundary beam and excitation of Langmuir waves. These are highly structured, sometimes stretching above, sometimes hanging down below f_{pe} . Also seen are intermediate frequency emissions near the shock [data taken from *Décrou et al.*, 2001].

CLUSTER came from the magnetosphere, entered the downstream region of the shock, passed the shock to upstream, re-entered the shock, and escaped downstream again. The relevant signatures of each region are seen in all four spacecraft. The uppermost light blue trace is the electron plasma frequency $f_{pe} = \omega_{pe}/2\pi$ which maps the local plasma density. The shock appears as the intense broadband emission with maximum intensity in the low-frequency waves. In the density it is mapped as a steep drop in f_{pe} to the low upstream

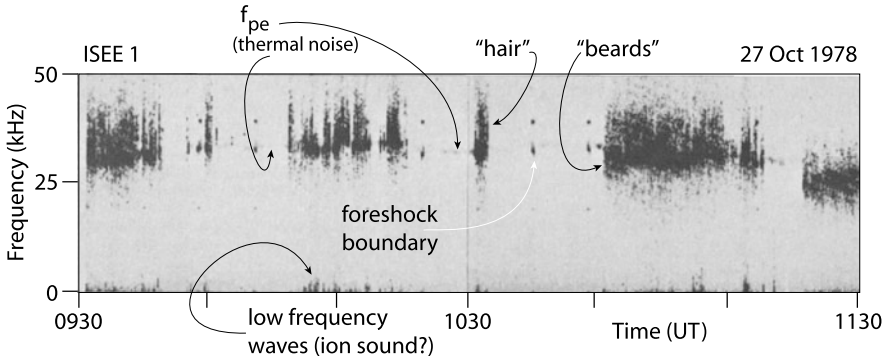


Figure 6.21: An example of ISEE 1 high frequency wave observations during contact and passage of the electron foreshock boundary. The thin faint spotty line around 30–40 kHz is the local plasma frequency $f_{pe} = \omega_{pe}/2\pi \sim 30$ kHz, seen here only as thermal noise with lesser instrumental sensitivity than in the CLUSTER observations. The occasional intense dark spots as the one marked by the white arrow are brief contacts with the foreshock boundary field line when the beam occurs. deeper in the foreshock the spectrum broadens non-symmetrically, evolving ‘hair’ and ‘beards’. Some moving structures can be identified in the hair. The emissions deeper inside the foreshock are correlated with the occurrence of low frequency emissions below a few kHz which are probably ion-sound waves [after *Etcheto & Faucheux, 1984*, courtesy American Geophysical Union].

density values. Close to the shock a number of intensifications in f_{pe} can be recognised. These occur at the time of contact with the electron foreshock boundary when intense plasma waves are excited in the Langmuir mode. It is interesting to note that the intensifications occur in spots and are not necessarily narrow-band. They may exceed f_{pe} , or they may also drop below it: the plasma frequency may have ‘hair’ or ‘beards’. Moreover, sometimes intense lower-frequency emissions are detected half way between the plasma frequency and the low frequency emissions in the foreshock. The most intense of these emissions belong to the region closer to the shock where the density gradient has not yet settled to the upstream values, but weaker emissions of the same kind occur farther away in the upstream low density domain.

A detailed analysis of all these waves detected by CLUSTER has not yet been undertaken and thus is not yet available. It is, however, highly suggestive that the observed intense spots in the plasma frequency, with the ‘hair’ parts exceeding and the ‘beard’ parts hanging down from the local f_{pe} , are related to magnetic field-aligned electron beams emanating from the shock and propagating upstream along the electron foreshock-boundary magnetic field as these are the same signatures as those observed with the ISEE and AMPTE spacecraft, though with much better frequency resolution and higher sensitivity here. For instance, the latter two spacecraft were unable to resolve the plasma frequency which in the CLUSTER data is nicely distinguishable even though it represents just thermal noise. It is most interesting that these contacts with the electron foreshock boundary occur quite irregularly throughout the hole period when CLUSTER was upstream of the shock and that they do not seem to exhibit a one-to-one correlation between the four spacecraft. Taking into account that the upstream plasma frequency away from the shock is fairly constant this implies that the shock-normal angle Θ_{Bn} changes on an irregular and fast time scale,

i.e. that the direction of the upstream magnetic field is highly variable even under conditions of apparent quiescence in the upstream medium. The conclusion drawn from this would not be that the upstream magnetic field is subject to violent variations as in the electron foreshock there is no reason for the magnetic field to be strongly affected. Rather this variability points on the temporal variation of the shock area near the tangential field line. Even a small change in the shock normal there will cause a large variation in the location of the tangential magnetic flux tube at a distance upstream of the shock. Reasons for such a variability have been given in the section on quasi-perpendicular supercritical shocks, but to repeat here, the main reason is that the super-critical shock is neither in thermal nor in thermodynamic equilibrium and is thus by its very nature subject to changes in all its physical parameters.

In addition to the plasma frequency emissions, which hint on contact of the spacecraft with the shock-tangential field line that makes up the electron foreshock boundary, there are a number of most interesting features seen in the wave spectra of Figure 6.20. One of them we merely note at this place because we will return to it farther down. It is the broadband signals that are occasionally seen in many places, in the plasma frequency as signals covering a frequency band of up to ~ 20 kHz, in the lower frequencies near the shock, and in particular in the shock transition itself where the frequency range of the emission $\Delta f > 80$ kHz in all cases exceeds the whole range of the instrument. There are no known plasma waves of comparable bandwidth reaching from ~ 0 kHz deep into the range of free space radiation frequencies. Such spectra are, however, known to be produced from narrow spatially localised electric field structures [Dubouloz *et al*, 1993]. The broader the spectrum the narrower will the structures be. We therefore conclude that in all the places where such broadband emissions are encountered, and in particular in the shock transition, we are dealing with very narrow localised and intense electric fields which have been generated by nonlinear processes. Such structures are presumably solitons and/or phase space holes either of electronic or ionic nature.

The so far most elaborate investigation of the width of the electron-foreshock boundary beam and its relation to the excitation of plasma waves has been performed by *Etcheto & Faucheux* [1984] and subsequently refined and supported by theory [Lacombe *et al*, 1985]. An example of their ISEE 1 observations is shown in Figure 6.21. Since the instrument was much less sensitive than WHISPER the thermal noise at the plasma frequency around $f_{pe} \sim 30$ kHz appears only as a very faint spotty line. Occasional strong narrow-band intensification of this line has been identified as contacts with the electron foreshock-boundary beam by inspecting the simultaneous electron measurements > 300 eV. These emissions were indeed found to be very narrow band, just 1–2 kHz wide at f_{pe} . Deeper inside the foreshock the bandwidth of the emissions broadens substantially in a non-symmetrical way. First, it becomes very noisy, consisting of short emissions clumped together in groups. Second, long ‘hair’ emissions evolve of roughly up to 10 kHz bandwidth, while ‘beard’ emissions are also found being generally weaker, but sometimes they extend to low frequencies. Low frequency waves below a few kHz are also detected in connection with these deep-foreshock emission and were found to smoothly merge from below into the beards.

Etcheto & Faucheux [1984] and subsequently *Lacombe et al* [1985] could show that the spectra changed from very narrowband in the foreshock boundary flux tube to broadband in the foreshock, just a few electron gyroradii away from the foreshock boundary. They also showed that the electron beam distribution flattened over this distance. The waveform of the waves at the plasma frequency in the broadband region away from the foreshock-boundary exhibited modulations, which group the waves into groups of length of a few 10 ms, a frequency roughly corresponding to the frequency of the low frequency waves seen in the dynamic spectrum of Figure 6.21.

The Nature of Electron-Foreshock Waves. The narrowband electron foreshock-boundary waves detected at the foreshock-boundary flux tube are clearly electron-beam driven Langmuir waves of the surprisingly large amplitude of a few mV m^{-1} at the edge of the foreshock and a few % bandwidth in frequency. It is considerably more difficult to infer about the nature of the waves deeper in inside the electron foreshock. These waves reach large amplitudes of a few 10 mV m^{-1} to a few 100 mV m^{-1} and bandwidths of $\gtrsim 30\%$. (Note that the stationary convection electric field in a $\sim 500 \text{ km s}^{-1}$ flow in a $\sim 5 \text{ nT}$ magnetic field is just $E \sim 2.5 \text{ mV m}^{-1}$!) The wave electric field is practically parallel to the upstream magnetic field, and the wavelength is long below and short above f_{pe} . Moreover, these waves are modulated and seem spectrally to connect to the intense low frequency waves. For plasma frequencies of $\sim 30 \text{ kHz}$ the ion plasma frequency is 0.7 kHz . Thus, accounting for the Doppler-broadening of the low-frequency wave spectrum in the fast flow, the low frequency waves can be tentatively identified with ion-acoustic waves, which accompany the high frequency waves at the plasma frequency. Obviously these waves modulate the latter, it is however not clear whether or not they are created by the high-frequency waves via the modulation instability or whether they are generated in a different interaction between the depleted beam and the upstream ion flow via an ion-acoustic instability. We have already argued that the modulation instability is unlikely under the conditions of the electron-foreshock boundary beam. Moreover, the weak modulation of the wave form of the high-frequency wave noted above does not argue in favour of the modulation instability and caviton formation as the waves are not really bundled in localised groups and the wavelength is not changed appreciably. On the other hand, the spiky broadband nature provides a weak argument for some localisation which can, however, be objected owing to the pronounced asymmetry of the spectrum with respect to f_{pe} .

Lacombe et al [1985], ignoring the correlation with the low-frequency waves, interpret the high-frequency waves around f_{pe} on the basis of the kinetic Langmuir wave instability including the beam plasma, i.e. referring to a non-gentle beam of finite temperature and appreciable density. Their simplified dispersion relation then becomes

$$1 = \frac{\omega_{pe}^2}{\omega^2} (1 + 3k^2 \lambda_{De}^2) - \frac{N_b}{N} \left[1 + i \sqrt{\frac{\pi}{2}} \frac{\omega - kV_b}{kv_{be}} \right], \quad \left| \frac{\omega}{kv_e} \right| \gg 1, \quad \left| \frac{\omega}{kV_b} - 1 \right| < \frac{\sqrt{2}v_{be}}{V_b} \quad (6.4)$$

with N_b , V_b , v_{be} the respective beam parameters, density, velocity, and thermal speed, and $\lambda_{De} = v_e/\omega_{pe}$ is the Debye length of the upstream plasma. The first two terms in this

expression are the ordinary Langmuir wave. The second term on the right is the beam contribution. for small gentle-beam densities the contribution to the real frequency can be neglected. However, for larger beam densities we see that N_b/N adds to the unity on the left. Hence the frequency of the wave should decrease for small k . Thus long wavelength waves will have frequencies $\omega < \omega_{pe}$. This agrees nicely with the observation. Short wavelength waves, where the term $3k^2\lambda_{De}^2$ comes into play, will have frequencies $\omega > \omega_{pe}$. This is also in accord with observation. Thus, the inclusion of the dense beam into the dispersion relation does reproduce the basic observed spectral properties of the beam excited waves at least qualitatively in a simple linear way without reference to nonlinear effects like the modulation instability and collapse. The occurrence of the low-frequency ion-acoustic waves is then probably due to the electron-ion velocity difference of the hot electron component and the cold upstream ion flow via the electron-ion acoustic instability.

In a series of one-dimensional particle simulation papers, *Dum* [1990] has investigated the dynamics of beams at the foreshock boundary in an attempt to quantify the above qualitative theoretical conclusions. He considered a pure electron beam on an otherwise neutralising ion background. As expected, a gentle weak beam evolves basically according to quasilinear theory into a plateau which for long times remains unchanged. Even in extremely long simulation times there is practically no evolution in wave energy, and the distribution function remains stable once the quasilinear plateau is formed. This happens approximately after a few hundred plasma periods. For a plasma frequency $f_{pe} \simeq 30$ kHz this amounts to a relaxation time of the order of just ~ 1 ms. Hence, gentle beams at the foreshock should be stabilised within this time, and any nonlinear effects will evolve only at much later times if at all. Excluding the ion dynamics, the simulations have been run until $\sim 3000 \omega_{pe}^{-1}$ without any susceptible change in the distribution and wave level, in physical times until ~ 10 ms. Moreover, the saturation level of the waves is remarkably low, much lower indeed than the theoretical estimates suggested. One does not expect that weak or strong turbulence effects will evolve which could change or prevent plateau formation. For this to happen, of course, ion dynamics should be included into the runs and should produce small k waves, while in the plateau formation the wave number increases because the plateau widens and the beam front proceeds to smaller and smaller velocities until the plateau is completed and Landau damping stops the further evolution. Therefore, based on these simulations, the dynamics of a cold beam cannot explain the richness of the wave observations in the foreshock boundary (but note our earlier remarks on beam stability!).

For a very cold beam the evolution is a bit more complicated as the cold beam initially is not subject to the Langmuir instability but rather to the reactive beam instability which generates a broad spectrum below f_{pe} as inferred above [*Lacombe et al*, 1985]. The maximum growth is close to but below f_{pe} , and the spectrum is sharply cut off just above f_{pe} . Afterwards the beam evolves readily into the kinetic Langmuir stage described above, forming a plateau and stabilising. This transition proceeds due to intermediate electron trapping in the wave, which heats the beam to a temperature when the kinetic instability can take over. One might thus conclude, that the waves observed below f_{pe} indicated the passage of the cold narrow beam front. This, however, is in contradiction to the observation that these waves are not observed in contact with the electron-foreshock boundary but

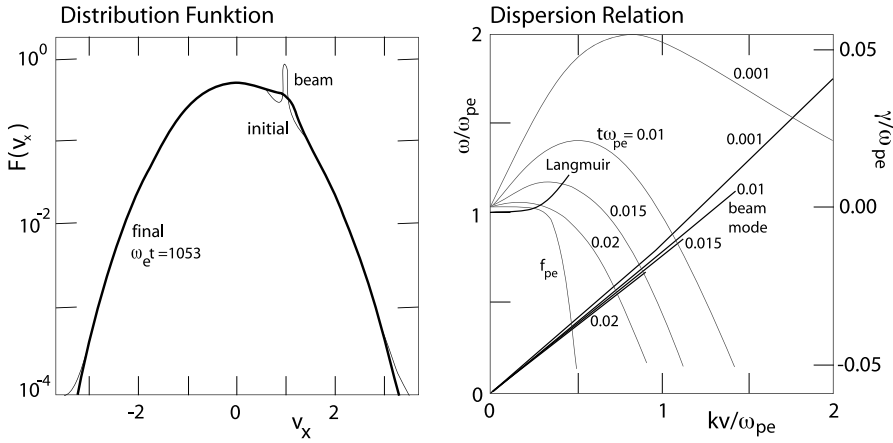


Figure 6.22: Results of a one-dimensional electron-electron beam particle simulation [after Dum, 1990] intended to explain the high-frequency plasma waves in the electron foreshock. *Left:* The electron distribution used. Initially (*thin line*) it consists of a broad (warm) Maxwellian background distribution plus a narrow (cold) electron beam sitting on the main distribution. All quantities are normalised, velocity is normalised to nominal beam velocity. At end time $t\omega_{pe} = 1053$ (*solid line*) the beam has become depleted and heated but has not disappeared. *Right:* Real dispersion relations $\omega(k)$ (*solid lines*) and growth rates $\gamma(k)$ (*thin lines*) at different simulation times $t\omega_{pe}$ as indicated by the numbers. Only the unstable part of the real dispersion relations (corresponding to positive $\gamma > 0$) is plotted. The Langmuir wave is shown separately with its growth rate labelled f_{pe} which is weakly positive only at very long wavelengths. The main unstable waves are shown to propagate in the beam mode (about straight lines starting at origin of ω and k . Maximum unstable frequencies are well below the plasma frequency ω_{pe} . With progressing time the wavelength of the unstable modes increases (decreasing k). The unstable domain in frequency and wave number shrinks. The unstable frequency decreases as a consequence of the decreasing beam velocity which leads to a decrease in the slope of the beam mode.

deeper in the foreshock where no cold beams are present. Hence, another interpretation for the waves is needed.

One solution is to take into account the bulk streaming distribution [Dum, 1990]. This is shown in Figure 6.22. Then the interaction becomes an electron-electron beam mode interaction with the possibility to destabilise the electron acoustic wave mode. This mode has frequency sufficiently far below the plasma frequency, $\omega < \omega_{pe}$. A condition is that the beam velocity spread is small. Initially the beam also excites frequencies substantially above f_{pe} with large growth rate, but these are quickly stabilised. The unstable modes are actually beam modes with resonance condition $\omega \lesssim kv_b$ which are destabilised by Landau damping from bulk electrons, i.e. the beam speed must enter the range of bulk velocities for Landau damping and should thus not be displaced far from the bulk distribution like in the gentle beam case. Plateau formation takes very long time in this case such that the wave can reach quite large intensity of the order of $W/NT_e \lesssim 10^{-3}$. For an upstream density of $N \sim 5 \times 10^6 \text{ m}^{-3}$ and electron temperature of $T_e \sim 100 \text{ eV}$ this yields a rms electric field amplitude $|e|_{rms} \lesssim 1 \text{ V m}^{-1}$ in very good agreement with observation. For larger beam temperatures the general trend as in Figure 6.22 remains valid. It is interesting that the unstable frequency readily drops from higher than ω_{pe} to $\omega < \omega_{pe}$ while the wave number

shrinks, i.e. the wavelength of the maximum unstable waves increases. One should thus observe falling tones in the frequency of the emission. It is also interesting that during the evolution of the instability the slope of the beam mode decreases which simply reflects the retardation of the beam during depletion. The final state is at low frequency and rather long wavelength.

It seems that this calculation explains the observation of electron foreshock wave emissions. However, the problem about this model is that it should work only right at the electron-foreshock boundary where only emission at the plasma frequency is seen like in the classical case of a gentle warm low density beam. The broadband waves below f_{pe} are observed inside the electron foreshock, where no electron beams have ever been detected. Moreover, it does neither explain the high intensities and the broadband nature of the waves above f_{pe} , nor does it explain the connection of the waves below f_{pe} to the low frequency ion-acoustic waves. Therefore we conclude, that probably a two-temperature counter streaming electron-component plasma with both components warm will be more appropriate to the interior of the foreshock, and probably the cold ion component of the bulk upstream flow must also be included. Moreover, since this model ignores the high speed convection of the upstream flow it implicitly assumes that the electron population is extended over a very large area which is the case only in the deep foreshock. Stabilisation of the foreshock boundary beams is due to convective depletion of Langmuir waves.

On the other hand, taking these simulations for serious, the deep electron foreshock region should be filled with many cold and not overwhelmingly fast electron beams propagating along the magnetic field upstream. Possibly the resolution of the current instrumentation is still unable to resolve them. The observations on the right in Figure 6.19 might indeed indicate the presence of such small beamlets that are distributed over the gyro-angle. So far, however, they can only be regarded as measurement fluctuations.

The parallel configuration of the distribution in the foreshock is shown in Figure 6.23 for a warm beam that is the field-aligned part of the diffuse electron component in the electron foreshock. Because of the vanishing-current condition the bulk electron component will be slightly retarded at the small amount $|\Delta V_e| = V_1 N_b / N$. This configuration consists of a hot dense bulk electron component and the warm dilute beam component. In addition to the beam instability it should be unstable with respect to electron-acoustic waves at frequency $\omega \simeq kv_e (N_b / N)^{1/2}$. Electron-acoustic waves are long-wavelength waves below the electron plasma frequency. Thus the low frequency waves find several explanations, none of them completely satisfying though. For the high frequency waves which exceed the plasma frequency the only available interpretation is that they result from localised electric fields. In frequency space these localised waves have a broadband signature. The distributions in the perpendicular direction are non-symmetric halo distributions with the backstreaming electrons populating the halo tails. They still await an in-depth treatment for inferring their contribution to the wave spectra, particle scattering and plasma heating.

We close this section by presenting in Figure 6.24 an average synoptic spectral view of the waves detected in the shock transition as was provided by *Rodriguez & Gurnett [1975]* and *Gurnett [1985]* from consideration of a large number of shock spectra. There

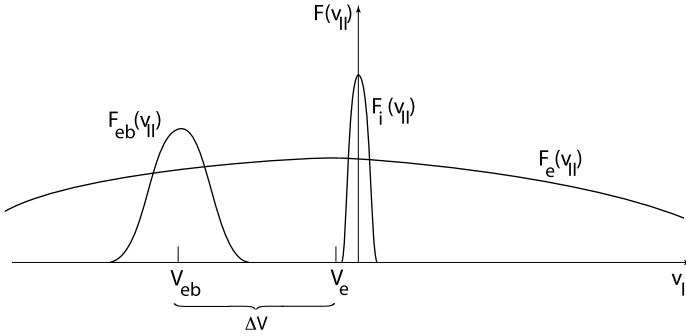


Figure 6.23: Parallel distribution functions in the deep electron foreshock consist of the cold bulk ion distribution and the hot bulk electron distribution on top of which the parallel part of the hot diffuse, i.e. warm upstream electron distribution. This distribution in the upstream frame moves away from the shock. The upstream electron distribution adjusts to the current-free condition by being retarded by the small amount $|V_e - V_i| = V_b N_b / N$. Its maximum is shifted in the direction of the upstream diffuse beam distribution. Such a configuration should be unstable with respect to electron-acoustic waves (dense hot electrons, less dense cool beam electrons).

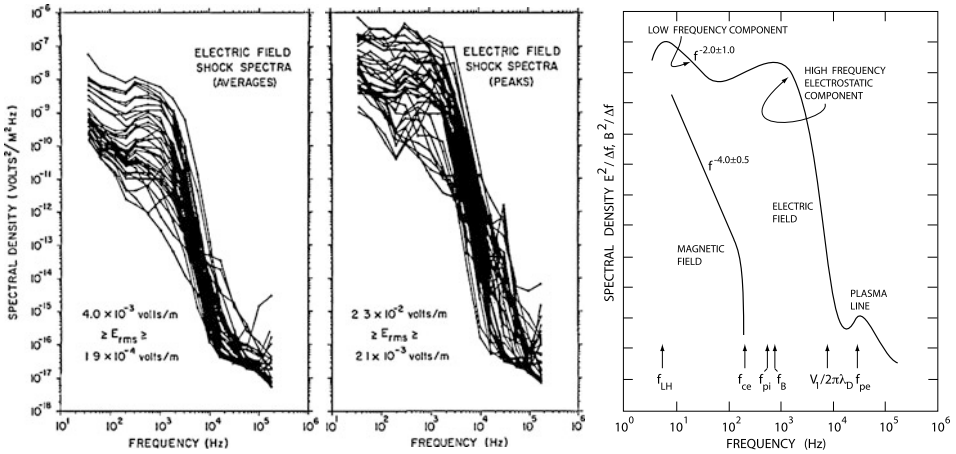


Figure 6.24: A representative distribution of electric shock spectra. *Left:* Average (1 s-averages) spectral values, showing the broad peak in the electric field spectra and the exponential decay of the spectrum toward high frequencies. For intermediate spectral intensities the peak is well developed, while when the spectral intensity is very high it smooths out. *Centre:* The same spectra but in the peak values (30 ms resolution) and not in the averages. The spectra are similar but much more variable and up to two orders of magnitude more intense. This points to the high variability of the electric wave emissions at shocks. In addition indications of the plasma frequency are seen at the high frequency end, suggesting that the emissions in f_{pe} are highly time variable [after *Rodriguez & Gurnett, 1975*]. *Right:* A schematic summary of shock spectra showing the magnetic spectra being cut off at the electron cyclotron frequency f_{ce} , and the electric spectra containing several maxima at the lower hybrid frequency f_{LH} at its low frequency end, around the ion plasma f_{pi} and Buneman two-stream f_B frequencies, with an absorption dip at the electron cyclotron frequency, a steep exponential decay caused by Doppler-shifted Landau damping of the waves in the Buneman and ion-acoustic modes at frequency $f_{pe} V_1/v_e = V_1/2\pi\lambda_D$, and the little bump at the electron plasma frequency [data taken from *Gurnett, 1985*, courtesy American Geophysical Union].

is no new information in this figure except that it summarises at a glance the main features in the higher frequency electric and magnetic field wave spectra. The cut-off of the magnetic spectra at the electron-cyclotron frequency is no surprise. In the electric spectra there is a large temporal variability as can be seen from a comparison of the average and peak values (measured within 1 s measuring time). The variations cover up to two orders in magnitude. Of interest is also that the electric spectra exhibit a maximum at the lower-hybrid frequency, an absorption at the electron cyclotron frequency, and show a broad maximum at the ion plasma/Buneman two-stream frequency before being steeply cut off towards higher frequencies. This cut-off occurs due to Landau damping at the Doppler-shifted frequency of ion-acoustic waves when the waves are shifted into Landau resonance with bulk electrons. In this frequency range all kinds of waves overlap, reaching from ion-acoustic waves and electron-cyclotron harmonics to Buneman waves, localised structures like Bernstein-Green-Kruskal modes and solitons, while the modified two-stream instability provides the large maximum at the lower hybrid frequency. The high variability in the peak values results from the presence of these localised structures which recently have been observed [Bale & Mozer, 2007] in situ but could have been concluded also just from the high variability of these spectra in the frequency range of the Buneman-modes respectively Doppler-shifted ion-acoustic waves. These are nothing but the signatures of many microscopic phase-space holes that obviously accumulate in the shock transition region. Support to such an interpretation has been given long ago by the ISEE measurements of the electron distribution function during shock transition [Feldman *et al*, 1983] which Gurnett [1985] made responsible already for the high variability of the peak spectral values. In these measurements the electron phase-space distribution function transforms from the upstream streaming Maxwellian to the shock-ramp and downstream flat-topped heated electron distribution, which just in the short time interval of crossing the shock -ramp magnetic-overshoot exhibited a clear signature of an electron beam that was sitting on the upstream edge of the flat top of the distribution (as depicted in Figure 5.19). The electron beam in that figure is caused by acceleration in the shock potential and has an upstream directed velocity of a few 1000 km s^{-1} . Compared to the bulk electron temperature this beam is cool. Thus it may excite electron-beam waves and electron-acoustic waves. However, in combination with the bulk ion flow it is also capable of exciting the Buneman two-stream instability which will readily evolve into electron holes and plays a major role in heating the plasma and making the main electron distribution flat-top. One should also remember that the scale of the overshoot is narrow, of the order of a skin depth, and is most probably due to an intense electron current flowing in the ramp or ramp-transition region which, presumably, is also related to the observed beam.

At high frequencies the spectra also exhibit the small peak caused by the beam excited Langmuir electron plasma waves, while the lower frequency foreshock waves are buried in the fat bump of the Doppler-shifted ion-acoustic waves. Of course, more cannot be concluded from a picture like this. The more detailed discussion of the spectra requires much higher spectral resolution. To the extent as it was available at the time of writing, we have given it here. But much more work is to be done until the various modes which can be excited in the foreshock and shock transition will be understood better.

Radiation

Shocks are frequently referred to as sources of radiation. Famous examples are supernova shocks, which are visible in almost all wavelengths, from radio through visible light up to x-rays [e.g., *Dickel & Wang, 2004*], and solar type II shocks with their main radiation signatures seen in the radio waves.

Supernova shocks are relativistic shocks which are not treated here. Their Mach numbers range from weakly relativistic to highly relativistic, but the energy per particle in them remains to be less than the rest energy of an electron, $\lesssim m_e c^2 = 0.511$ MeV, which allows to treat them classically. This does not hold anymore for the central supernova engine which drives the flow and which in some cases results in the generation of ultra-relativistic jets. There the shocks become non-classical, and not only radiation losses but also particle generation must be taken into account in their description.

Observations. The radiation that is occasionally emitted from nonrelativistic shocks is restricted to the radio wave range. They do not generate x-rays or visible emissions because, first, of their comparably low energy per particle, which is less than the rest energy of an electron $\ll m_e c^2$, second, because of their low energy transmission rate (particles are not retarded from flow speed to rest) and, third, because of the low ‘emission measure’ $EM = \int_{los} ds \Delta N_e^2(s)$, i.e. the square number of radiating particles of density ΔN_e integrated along the line of sight that contributes to the emission of x-rays is low. Moreover, since magnetic fields are weak and the ratio $\omega_{pe}/\omega_{ec} > 1$ of electron plasma to electron cyclotron frequency is larger than one, gyro-synchrotron emission is unimportant. Thus, the only means of how free-space radiation of frequency $\omega \gtrsim \omega_{pe}$ can be produced is via plasma wave emission.

Radio emission from collisionless shocks in the heliosphere is a widely studied field including type II solar radio burst, interplanetary type II burst, CME-driven radio bursts and radio emissions from planetary bow shocks. While the solar bursts, because of their high frequency can be observed from Earth, most of the other emissions have been discovered only from aboard spacecraft. Of the enormous wealth of such observations made by the ISEE, AMPTE, Polar, WIND, GEOTAIL and other spacecraft, many of them never published, we pick here just a more recent observation from CLUSTER [*Trotignon et al, 2001*]. Figure 6.25 shows two cases of such observations on December 22, 2000 and three months later on March 13, 2001, when CLUSTER was crossing the bow shock and moved into the electron foreshock. On December 22, 2000 the upstream density was relatively high with plasma frequency near 40 kHz. The intense spots in f_{pe} between 1200–1300 UT indicate touching of the electron foreshock boundary field line as has been discussed above. Broad-band wave modes of around $\sim \frac{1}{2} f_{pe}$ frequency indicate beam modes though no frequency drift is detectable. However, the density during this time is low enough for a faint emission to occur at about ~ 80 kHz, just two times f_{pe} . At this frequency emission can be only in the electromagnetic free space mode. We are thus witnessing local generation of radio emission from the foreshock boundary. At later times when the spacecraft moved deeper into the foreshock – as visible from the widening of the plasma wave spectrum to both sides, up and down from f_{pe} – the radio emission becomes more diffuse and more broad-

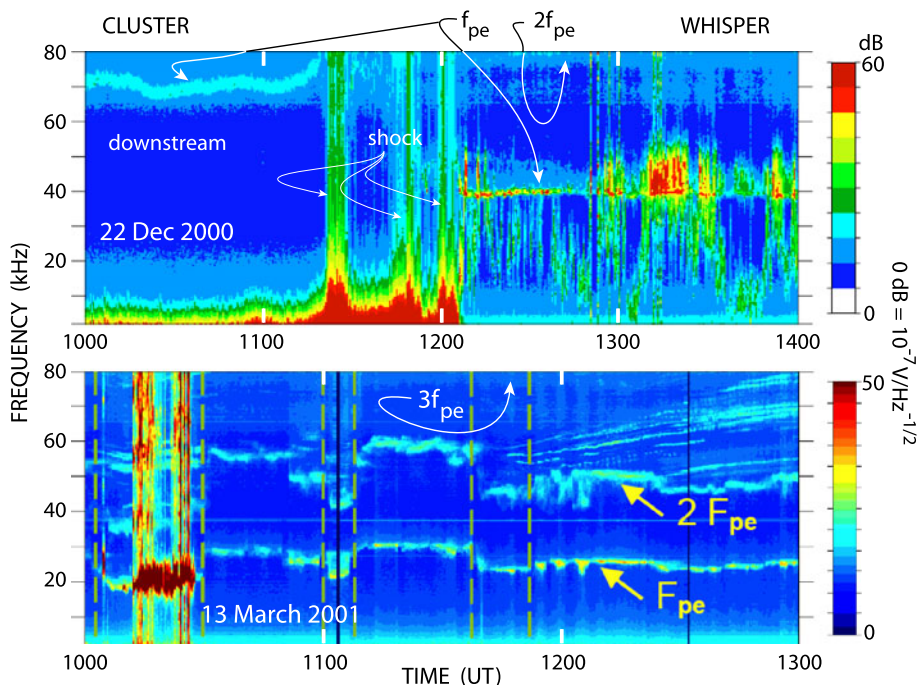


Figure 6.25: Observations of electromagnetic radiation from the electron foreshock, following the multiple CLUSTER shock crossings between 1120–1210 UT marked by intense low frequency noise and its broadband extension through the entire frequency range from 0–80 kHz. *Top:* Typical f_{pe} and lower frequency emissions from 1210–1250 UT at the electron foreshock boundary and $2f_{pe}$ radiation near 80 kHz. Weak emission just above f_{pe} is seen with upper cut-off decreasing in frequency. At later times deeper in the shock broadband electrostatic emissions at f_{pe} occur. These correlate with broadening of the harmonic radiation to lower frequencies. *Bottom:* Three hours of low upstream density when the spacecraft was close to the foreshock boundary. f_{pe} is accompanied by harmonic emission not much less intense than f_{pe} . During broadening of f_{pe} the harmonic emission broadens downward. Near 80 kHz a faint third harmonic band $\sim 3f_{pe}$ can be identified. The band splitting in the harmonic radiation between 1000–1010 UT is of interest. Many narrow-band upward drifting equally spaced in frequency radiation bands after 1140 UT can be identified. They have no counterpart in f_{pe} and may come from a remote source related to sharp changes in density (drop in f_{pe}) at 1140 UT. The dashed vertical lines mark changes in the magnetic field (not shown here), when its direction abruptly turned [after Trotignon *et al.*, 2001, and Trotignon, private communication].

band shifting to lower frequencies than the second harmonic of the plasma frequency. This suggests that possibly here the lower frequency modes close to f_{pe} do actively participate in the generation of emission. On the other hand, radiation from the shock ramp might also contribute to these emissions.

The lower panel in Figure 6.25, recorded on March 13, 2001 during a day of much lower upstream density, shows the more typical case of Langmuir waves at f_{pe} and radio emission at almost precisely $2f_{pe}$. Note that near 80 kHz a faint indication of the presence of an emission at the third harmonic $\sim 3f_{pe}$ can be made out. The dashed vertical lines in this panel indicate the times when the magnetic field direction changed abruptly with the

changes in magnetic field magnitude and density remaining comparably small. The emissions in the plasma frequency and harmonic are well correlated during the entire event. Obviously CLUSTER was constantly close to the foreshock boundary as only intensification in f_{pe} is seen but no ‘hair’ nor ‘beards’ neither evolve except during a period shortly after 1200 UT in the three yellow spots in f_{pe} when the harmonic emission extends to lower frequencies together with the development of a little beard.

Two other interesting features can be read from this panel. At early time in the panel, just before the short active phase of the topside sounder at 1020 UT, splitting of the harmonic emission into two narrow bands is seen, which reminds at the band-splitting in type II bursts. The other interesting feature is the large number of drifting emissions with increasing frequencies, which suddenly evolve right after 1200 UT (following the drop in plasma density at 1140 UT). Both features, the splitting of $2f_{pe}$ and these drifting bursts are not understood yet. The latter might be related to the abrupt changes in density and plasma frequency seen in this panel. These density changes are accompanied also by changes in the magnetic field, which are not shown here. It is worth noting that these banded drifting emissions cannot come from remote simply because the low frequencies arrive first. Radiation from any remote source should be visible first at high frequency. The drifting emissions must be related to a nearby source, most probably the shock or foreshock. Understanding its generation mechanism should provide valuable information about the radiation source.

Interpretation. Theory of shock-emitted radiation is based on plasma processes which under the prevailing collisionless conditions in the shock and foreshock plasmas refer to wave-wave coupling as the main generation mechanism. Direct emission from particles is unimportant, because the energy losses a particle experiences when becoming retarded or reflected at a shock, are not transformed into radiation. In all non-relativistic cases radiation losses can completely be neglected compared with all other energy losses. Nevertheless, the observed radiation is of interest because in many cases, where no measurements are possible to be performed *in situ*, radiation is the only direct and presumably identifiable signature a shock leaves, when seen from remote. The other signature is the generation of energetic particles, which will be treated in the next chapter, but energetic particles are a more diffuse indicator of a shock, because their propagation is vulnerable to scattering from other particles, obstacles and, in the first place, scattering by magnetic fields. They, thus, do not provide an image of the shock as clear as radiation does. Just because of this reason, investigation and understanding of the mechanisms of emission of radiation from shocks enjoys – and deserves – the large amount of attention it receives.

Expecting that – presumably – direct particle involvement into radiation is improbable (maybe with two exemptions, which we will note later) we are left with a small number of possible mechanisms, which all belong to the class of wave-wave interaction in weak plasma turbulence. The most probable of these are three-wave processes. These can be understood as ‘collisions’ between three ‘quasi-particles’. Since only these three are involved, the interaction conserves both, energy and momentum, and can symbolically be written as

$$L + L' \rightarrow T, \quad L \equiv \{\omega_L(\mathbf{k}_L), \mathbf{k}_L\}, L' \equiv \{\omega'_L(\mathbf{k}'_L), \mathbf{k}'_L\}, T \equiv \{\omega(\mathbf{k}), \mathbf{k}\} \quad (6.5)$$

Here L stands for longitudinal, and T for transverse – as the emitted radiation of frequency ω and wave number $\mathbf{k}$ propagates in the free space mode and thus is a transverse electromagnetic wave while its two mother waves are assumed to be longitudinal (i.e. electrostatic) waves of sufficiently high frequency. Moreover, these frequencies $\omega_L(\mathbf{k}_L)$ etc. depend on wave number through the real parts of the electrostatic dispersion relation $D_L(\omega_L, \mathbf{k}_L) = 0$ and may become quite complicated expressions.

In the presence of a plasma there are two free space modes, the ordinary and the extraordinary mode. Naturally, in order to leave the plasma and escape in the form of radiation their frequencies must exceed some lower cut-off frequency. For the ordinary mode

$$\omega^2 = \omega_{pe}^2 + k^2 c^2 \quad (6.6)$$

this is the plasma frequency reached at very long wave lengths $k = 0$. However, since the speed of light $c \gg v_e$ is so large, the radiated wave length is much longer than any of the wavelengths of the longitudinal waves involved. This is immediately recognised when comparing the above ordinary wave dispersion relation with the Langmuir wave relation, with $L \equiv \ell$,

$$\omega_\ell^2 = \omega_{pe}^2 + 3k_\ell^2 v_e^2 \quad (6.7)$$

which is of exactly the same structure. Hence, as long as $v_e \ll c/\sqrt{3}$ we will have $k_\ell \gg k$, and the wave number of the radiated mode is practically zero. Momentum conservation of the three interacting quasi-particles becomes simply $\mathbf{k}_L \approx \mathbf{k}'_L$, implying that the interaction selects counter-streaming electrostatic waves. As for an example, any process that is capable of generating Langmuir waves of comparable wavelengths, propagating in both directions along and opposite to the magnetic field, can in principle contribute to generating escaping radiation. From energy conservation $\hbar\omega_\ell + \hbar\omega'_\ell = \hbar\omega$ of the three ‘quasi-particles’ involved one immediately finds that $\omega \approx 2\omega_{pe}$. This is the origin of and the simplest mechanism for the generation of the $2f_{pe}$ -second plasma harmonic radiation and has been proposed more than half a century ago by *Ginzburg & Zheleznyakov [1958]* as an explanation for the observation of solar type II and type III radio bursts.

In this simple reasoning we have completely neglected not only the contribution of the Langmuir wave number (which turns out not to be important in magnitude, it just shifts the emitted frequency a tiny amount up in frequency) but also the fact that the electric field of Langmuir waves is polarised along $\mathbf{k}_\ell = \pm k_\ell \mathbf{B}/B$ and thus along the ambient magnetic field $\mathbf{B}$, while the electric field of the transverse emitted electromagnetic radiation must necessarily be polarised perpendicular to $\mathbf{k}$ (because of the absence of space charges at frequencies sufficiently higher than f_{pe}). Since the electric field will, after collision and annihilation of the two Langmuir waves involved, remain to oscillate along the magnetic field, the emission is preferably directed perpendicular to the magnetic field. It turns out, then, that it is easier to radiate in the extraordinary than in the ordinary free space mode. The extraordinary mode has a slightly more complicated dispersion relation;

in a dense plasma with $\omega_{pe} \gg \omega_{ce}$ (as is encountered in near Earth space where supercritical collisionless shocks evolve) it has a slightly higher cut-off frequency. But the argument about the smallness of $k \ll k_\ell$ holds also in this case. Radiation at the second harmonic $\omega \approx 2\omega_{pe}$ should therefore be polarised perpendicular to the magnetic field in the extraordinary mode.

Unfortunately, foreshock emission has not been found to show any preference in polarisation [Reiner *et al*, 1996]. Moreover, emission is not only in the second harmonic but has also been detected close to $\gtrsim \omega_{pe}$ and at the third harmonic $\sim 3\omega_{pe}$. These emissions require different waves to be involved for which a number of mechanisms have been proposed [cf., e.g., Cairns, 1988]. So far none of them could be ultimately verified or even agreed upon, each having its merits and pitfalls. Since radiation is energetically negligible, as we have mentioned above, the whole problem could be put aside. However, since the assumptions made in every radiation mechanism contain important information about the source region, the problem of radiation production in collisionless shocks remains to be tantalisingly urgent and awaits resolution.

Since there is no agreement yet about the radiation mechanism, we merely note here some of the different proposals. The first is the above mentioned merging of two counter streaming Langmuir waves. The problems about this simple though suggestive mechanism are numerous. First, Langmuir waves are assumed to be generated by the gentle-beam instability. Ignoring the problem of beam survival during its propagation along the shock-tangential field line, which we have discussed already in detail, gentle beams excite only forward Langmuir waves, which requires some mechanism that back-scatters a substantial percentage of waves and inverts the direction of their wave numbers. There are three known elegant processes that are capable of doing this: modulation instability respectively collapse, scattering of Langmuir waves off thermal ions, and scattering off ion-sound waves, all three proposed long ago [for an early review of the latter two mechanisms cf., e.g., Tsytovich, 1970].

Modulation instability generates ion-sound waves via the ponderomotive pressure force $\mathbf{F}_{pmf} = -(e^2/m_e\omega_{pe}^2)\nabla|\mathbf{e}_\ell|^2$ of the high-frequency Langmuir wave, $\mathbf{e}_\ell(\mathbf{r}, t)$. These waves, when becoming locally large amplitude, structure the plasma into a chain of cavities in which the Langmuir waves become trapped. This process generates long wave lengths. It is described by the Zakharov equations for the combined evolution of the Langmuir wave field and the density variation δN , respectively,

$$\frac{\partial^2 \delta N}{\partial t^2} - c_{ia}^2 \nabla^2 \delta N = \frac{\epsilon_0}{m_i N} \nabla^2 |\mathbf{e}|^2, \quad \frac{\partial \mathbf{e}}{\partial t} + \frac{3\omega_{pe}}{2} \lambda_e^2 \nabla^2 \mathbf{e} = \frac{\omega_{pe}}{2N} \mathbf{e} \delta N \quad (6.8)$$

The first of these equations is a driven wave equation for the density variation which for slow time variations, when the derivative with respect to time is neglected, just gives pressure balance between the ponderomotive pressure on the right and plasma pressure on the left, i.e. proportionality $\delta N \sim -|\mathbf{e}|^2$. In other words, the density variation anti-correlates with the field pressure, which corresponds to caviton formation. The second equation is a nonlinear Schrödinger equation for the evolution of the wave amplitude.

The Langmuir waves trapped in the cavitons must bounce back and forth, which naturally creates counter-streaming waves of equal intensity with opposite wave numbers.

During collapse the cavities shrink in size, the wave numbers and moments of the waves increase, and the wave energy density increases as well because of the shrinking volume. This yields both, the counter streaming Langmuir waves being localised in the same region and, in addition, a large radio emissivity. Unfortunately, we have already noted it, this process – as beautiful as it might be – has not been confirmed experimentally, neither in the observations nor in the full particle simulations. Observed wave intensities are too low in the electron-foreshock boundary and electron foreshock, and the density variations did not indicate the presence of the expected cavities. Simulations, on the other hand support quasilinear evolution and wave scattering off thermal ions. We note, however, that the most recent detection of the very strong electric fields in the shock ramp [Bale & Mozer, 2007] might indicate that it is not the electron foreshock where one should expect caviton and collapse to work and cause the most intense radiation, rather it might be the very shock transition where shock radiation is generated by such processes. It is, in this respect, most interesting to remind of the strange radiative behaviour reproduced in Figure 6.25 that was detected by CLUSTER. We also note that similar observations had been made much earlier with the wave experiment on AMPTE IRM in the Earth foreshock (R. A. Treumann & J. LaBelle, unpublished 1986, see also the trace in the radiation emitted from an Hot Flow Anomaly shown in Figure 6.38). The observed band splitting and high intensities might have been caused by Langmuir caviton collapse [Treumann & LaBelle, 1992].

Other possibilities to produce counter streaming Langmuir waves are scattering of Langmuir waves off thermal ions [Muschiatti & Dum, 1991, investigated this process in full detail numerically including the ion polarisation cloud], a mechanism known since the early sixties. The process reads symbolically $L + i \rightarrow L' + i^*$, where the primed quantities are after the collision, and the star on the ion indicates excitation of the ion as it is too heavy for changing momentum during the collision with the Langmuir wave. It is merely excited while the scattered Langmuir wave has changed direction and lost some of its momentum, i.e. attains a longer wave number and lower frequency. The same process does also work with ion-sound waves as $L + IS \rightarrow L'$. The scattered Langmuir waves then also change direction by absorbing the ion sound. Both processes have been used for radiation generation [Yoon *et al*, 1994].

Radiation at higher frequency, e.g. radiation at the third plasma harmonic can be generated by a four-wave process. This is also favoured by caviton formation and collapse since the waves in this case are all confined to one and the same volume. However, other mechanisms have also been proposed. All these processes are of the kind of wave-wave interactions and thus their efficiencies are proportional to the product of the involved relative wave intensities. Since the latter are usually low, the efficiencies are very small as a rule. An attempt to increase the growth of Langmuir waves has in the recent past been the idea to consider a statistical theory of growth called ‘stochastic growth’ [Robinson, 1995; Krasnoselskikh *et al*, 2007]. This attempt takes advantage of the statistically distributed density fluctuations in the foreshock region like in a random medium. Since the Langmuir-wave growth rate is proportional to $\delta N/N$, an average growth rate over the volume of occupation by the Langmuir wave can be calculated. This might be more realistic than using the linear growth rate. Regions of decreased density contribute strongly to the average Langmuir-wave amplitude. In this way an in the average larger emission efficiency

is obtained. Moreover, the averaging procedure introduces a statistical element which supports the incoherence of the relation between the detected Langmuir waves and radiation, which is in partial agreement with the observation.

Generation of radiation at the fundamental $\omega \gtrsim \omega_{pe}$ in a three wave process requires the presence of a low frequency wave. Ion-acoustic waves are one possibility, other possibilities are lower-hybrid waves, Buneman waves, the modified-two stream instability, various kinds of drift waves, and also electron acoustic waves or electron beam waves. In particular the latter are present in the foreshock region and thus can combine with Langmuir waves to generate fundamental radiation slightly above the plasma frequency.

The two exemptions when particles become involved are the above mentioned scattering of Langmuir waves off ions, and the so-called electron-cyclotron maser instability [for a contemporary review see, e.g., *Treumann, 2006*]. Its advantage is that it operates directly on the free space mode avoiding any intermediate step like three-wave processes or particle scattering. However, it requires a particular form of the electron distribution with a velocity space gradient in the perpendicular direction $\partial F_e(v_{\parallel}, v_{\perp})/\partial v_{\perp} > 0$, a hot electron distribution and low cold electron density. It is barely known whether such distributions are realised at the shock. However, if they are in some place, then the cyclotron maser instability will outrun all other mechanisms and directly feed the free-space electromagnetic radiation modes. Radiation will then be at a harmonic of the electron cyclotron frequency which is a severe restriction if the magnetic field is low and the density high. Therefore, regions of low density and stronger converging magnetic fields are the best candidates for this radiation source.

6.3 Quasi-parallel Shock Reformation

In quasi-parallel supercritical shocks there is not such a stringent distinction between the region upstream of the shock and the shock itself like in quasi-perpendicular shocks. The foreshock, which we have discussed in some detail in the previous section, and the shock itself cannot be considered separately. This is due to the presence of the reflected and diffuse particle components in the foreshock. These, as we have seen are the source of a large number of waves.

The interaction of these waves with the shock is one of the main issues in quasi-parallel shock physics. In the present section this will become clear when we will be dealing with the formation, behaviour and structure of quasi-parallel shocks as it has been inferred less from observation than from numerical simulations. The reason is that the real observations in space do not allow to separate the particles and waves from the shock. They all occur simultaneously and are interrelated and can never be observed in their initial state. The observations to which we will nevertheless occasionally refer will leave the impression of large-amplitude noisy fluctuations. In simulations, on the other hand, it is at least to some degree possible to prepare the system in such a way that a single effect can be studied. For instance, in one-dimensional simulations the direction of wave propagation can be prescribed which allows studying just waves in one direction and their effect on the shock and particles. Moreover, treating the electrons as a neutralising Boltzmannian

fluid suppresses their effect on the ion motion and wave generation. Treating them as an active fluid allows taking account of electron-ion fluid instabilities. Finally, full particle PIC simulations can be performed with low or realistic mass ratios in order to investigate different time and spatial scale dependence and the excitation or coupling to higher frequency waves. Most simulations that have been performed in the past have taken advantage of these possibilities.

6.3.1 Low-Mach Number Quasi-parallel Shocks

It is usually assumed that low-Mach number shocks are stable, i.e. show no substantial time variation or reformation. It is not completely transparent why this should generally be so. Firstly, the critical Mach number has been shown by *Kennel et al* [1985] to become small at narrow shock normal angles $\Theta_{Bn} \rightarrow 0$, in which case even low-Mach number quasi-parallel shocks should become supercritical and reflect ions. Secondly, any fast ions of parallel velocity $v_{\parallel} > V_1$ that have been heated in the shock can in principle escape from the quasi-parallel shock upstream along the magnetic field and should appear in the foreshock where they contribute to wave generation. Therefore, it makes sense to investigate the state of quasi-parallel shocks in view of their stability and wave generation even for low Mach numbers. In addition, any waves that are generated in the shock ramp or transition with upstream directed $\mathbf{k}$ and fast enough parallel phase or group velocities could also escape from the shock in upstream direction. This could, in particular, be possible just for low Mach number quasi-parallel shocks.

To check this possibility *Omidi et al* [1990] have performed one-dimensional hybrid simulations finding that initially the quasi-parallel shock consisted of phase-standing dispersive (magnetosonic) whistler waves with the last whistler wave cycle constituting the shock ramp. As expected, the wave vectors of these phase-locked magnetosonic whistlers are aligned with the shock normal. At later times, backstreaming ions along the upstream magnetic field excite a long-wavelength whistler wave packet upstream of the shock. In the one-dimensional simulation the wave vector is restricted to the shock normal while it is known from theory that the growth rate is largest along the magnetic field. These oblique whistler waves should thus show up in two-dimensional simulation and may be visible at larger distance for sufficiently large upstream phase velocities.

Such two-dimensional (hybrid) simulations with a non-inertial electron fluid have been performed by *Scholer & Fujimoto* [1993] for a Mach number of $\mathcal{M}_A = 2.2$ and angles $\Theta_{Bn} = 20^\circ, 30^\circ, 45^\circ$ and by *Dubouloz & Scholer* [1995] for an angle $\Theta_{Bn} = 30^\circ$ and high Mach number $\mathcal{M}_A = 5$ (Figure 6.26). The lower Mach number simulations show the presence of a substantial number of backstreaming ions which cause an ion-ion instability in the upstream region. However, the excitation and properties of the waves depend strongly on the shock normal angle. Initially, as in the case of *Omidi et al* [1990] phase-locked short-wavelength whistlers appear which are replaced at later times by upstream long-wavelength whistlers with phase velocity directed and amplitude growing towards the shock ramp but upstream directed group velocity, i.e. the shock radiates energy away towards upstream, as one would naively expect, because the shock being supercritical must reject the excess inflow of energy which it can do by both, reflecting particles and emitting

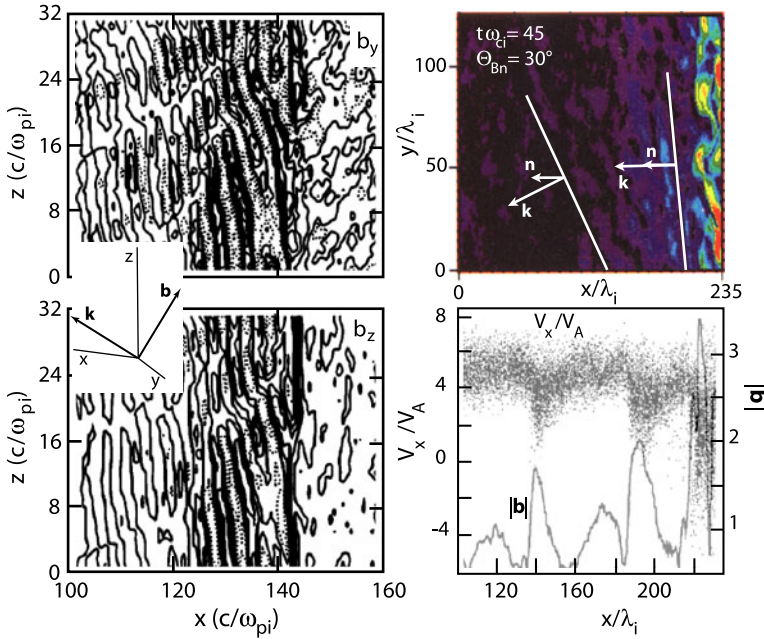


Figure 6.26: Two-dimensional hybrid simulations of the evolution of upstream waves. *Left column:* Upstream wave in a low-Mach number $\mathcal{M}_A = 2.2$ quasi-parallel shock [after Scholer & Fujimoto, 1993, courtesy American Geophysical Union]. The nominal shock is in the (y, z) -plane. Ions escaping to upstream generate the oblique upstream magnetosonic waves. The *contour plot* of the two normalised to the upstream magnetic field components of the magnetic fluctuations shown is taken at relatively early times $\omega_{ci}t = 68$. It shows the nearly plane magnetic wave fronts inclined against the shock in direction x and having wavelengths of $\sim 10c/\omega_{pi}$ in z while being much shorter in x . In the vicinity of the shock the wave fronts turn more parallel to the shock and produce a non-coplanar magnetic component $|b_y|$ which is of same order as the $|b_z|$. Moreover, even though the shock has relatively low Mach number, it is not completely stable but shows structure in z direction which is produced by the presence of the reflected upstream particles and the upstream waves. At places it is impossible to identify one single shock ramp. *Right column:* Two-dimensional hybrid simulations of the evolution of giant magnetic pulsations (SLAMS) in front of a quasi-parallel supercritical shock [after Dubouloz & Scholer, 1995, courtesy American Geophysical Union]. *Top:* The simulation plane showing the structure of the (normalised) magnetic fluctuation field $|b|$ at time $t\omega_{ci} = 45$ and shock normal angle $\Theta_{Bn} = 30^\circ$. The accumulation of the growing wave fronts at the shock transition, their increasing amplitudes, and their turning towards becoming parallel to the shock is clearly visible from the rotation of the two wave fronts and their $\mathbf{k}$ vectors shown in white. Away from the shock transition the angle between $\mathbf{k}$ and the shock normal $\mathbf{n}$ is large. Close to the shock the two vectors are about parallel. The magnetic field is in the wave front, so Θ_{Bn} is close to 90° here. *Bottom:* Pulsation amplitude and ion phase space. The fluctuations evolve into large amplitude pulsations when approaching (and making up) the shock. The strong retardation of the upstream flow by the pulsations is visible in the shock-normal velocity component (Mach number $\mathcal{M}_A$). In the hybrid simulations this slowing down is accompanied by some ion heating.

waves into upstream direction. These waves are excited by the backstreaming ion component in a strongly nonlinear interaction process because of the evolving steep ion-density gradient, which is of the same scale as the whistler wavelength. The $\mathbf{k}$ -vector turns away from the magnetic field having comparable components parallel to $\mathbf{B}$ and parallel to the

shock normal $\mathbf{n}$. For small Θ_{Bn} a remnant of the initial phase-locked whistlers survives but disappears at $\Theta_{Bn} = 30^\circ$. Close to the shock, where the backstreaming ion density is high, the waves have short wavelengths, and $\mathbf{k}$ is almost parallel to $\mathbf{n}$. In the high Mach number simulations no shock is produced but instead reflected ions were artificially injected with same Mach number as the incoming flow but with much higher temperature $v_i = 14.1 V_A$, forming a spatially uniform ion beam. The intention was to investigate the effect of the hot reflected ions. This is shown on the right in the above figure. The result resembles the former one where a shock was generated by reflection at a wall, but the effect in the injected beam case is stronger because of the higher Mach number. Hence it is the hot reflected ion component that is responsible for the wave dynamics and the shock dynamics.

All this can be seen from the two-dimensional intensity contours of these waves in the foreshock, which are plotted at a relatively early time in the shock evolution in the simulations in Figure 6.26. On the right in this figure the geometry is given, with the magnetic field fluctuation vector $\mathbf{b}$ in the (y, z) -plane. The bottom panel on the left shows contours of the b_z fluctuations in the (x, z) -plane. The nominal shock ramp is at $x \approx 145 \lambda_i$ ion inertial lengths $\lambda_i = c/\omega_{pi}$ at this time. The upper panel shows the non-coplanar component b_y -contours in the same representation. Behind the shock the fluctuations are irregular and disorganised. However, in front of the shock a clear wave structure is visible with strongly inclined wave fronts and of roughly $\sim 10 \lambda_i$ wavelengths in z parallel to the nominal shock surface. The wavelength in x is about three times as short.

These waves are seen in both components, b_z, b_y , are low amplitude at large distance from the shock but reach very large amplitudes simultaneously in both components during shock approach while, at the same time, bending and assuming structure in z -direction that is different from the regular elongated shape at large distance. This deformation of wave front may be due to the residual whistlers near to the shock, but it implies that the shock has structure on the surface in both directions x and z and is not anymore as planar as was initially assumed. The shock becomes locally curved on the scale of the shock-tangential wavelength. The waves deform the shock and, in addition, being themselves of same amplitude as the shock ramp, become increasingly indistinguishable from the shock itself. *The shock is, so to say, the last of the large-amplitude magnetic wave pulsations in downstream direction, and the shock-magnetic field is not anymore coplanar*, because the waves have contributed a substantial component b_y that points in the direction out of the coplanarity plane.

We have emphasised this phrase, because it expresses the importance of the low-frequency upstream magnetic waves in quasi-parallel shock physics. Contrary to quasi-perpendicular shocks where the reflected gyrating ions in combination with the reflected-ion excited modified-two stream instability were responsible for the shock dynamics and different kinds and phases of shock reformation, quasi-parallel shock reformation and much of its physics is predominantly due to the presence of large-amplitude and spatially distinct upstream waves. These are the generators of the shock and, due to their presence, the shock changes its character. It is highly variable in time and position along the shock surface and is – close to the shock transition on a smaller scale – ‘less quasi-parallel’ (or more perpendicular, i.e. the shock-normal angle Θ_{Bn} has increased on the scale of the upstream waves). The latter is due to the out-of coplanarity-plane component

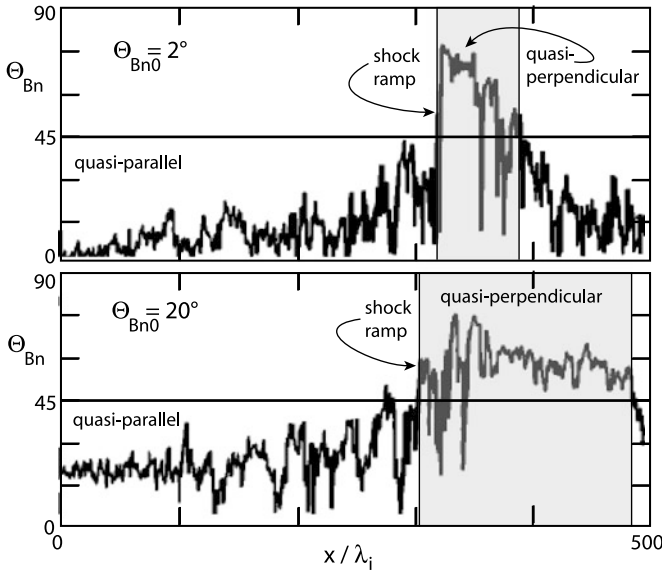


Figure 6.27: The evolution of the shock normal angle Θ_{Bn} on distance from the shock in two-dimensional hybrid simulations for two initial quasi-parallel shock-normal angles $\Theta_{Bn0} = 2^\circ$ and $\Theta_{Bn0} = 20^\circ$, respectively [simulation results taken from *Scholer & Fujimoto, 1993*, courtesy American Geophysical Union]. The horizontal line at 45° is the division between quasi-perpendicular and quasi-parallel shock normal angles. In both cases Θ_{Bn} evolves from quasi-parallel direction into quasi-perpendicular direction. The shaded areas identify the quasi-perpendicular domains.

of the upstream waves. In spite of concluding this from a hybrid simulation, this conclusion remains basically valid also in full particle simulations. It had been suggested already earlier on the basis of ISEE 1 & 2 observation of the evolution of the upstream ultra-low frequency wave component [*Greenstadt et al, 1993*].

The gradual evolution of the shock normal angle Θ_{Bn} has been demonstrated in other hybrid simulations by *Scholer & Fujimoto [1993]* and *Dubouloz & Scholer [1995]* who investigated the evolution of the shock normal angle in dependence on distance from the shock. This is shown in Figure 6.27 for two-dimensional hybrid simulations with initial shock-normal angles $\Theta_{Bn0} = 2^\circ$ and $\Theta_{Bn0} = 20^\circ$, respectively, which we anticipate here. In both cases Θ_{Bn} evolves from quasi-parallel to quasi-perpendicular angles. Qualitatively there is little difference between the two cases. At the shock ramp Θ_{Bn} is deep in the domain of quasi-perpendicular shocks. The only difference is that for the nearly parallel case the angle jumps to quasi-perpendicular quite suddenly, just before approaching the shock ramp, while the evolution is more gradual for the larger initial Θ_{Bn} . In both cases the evolution is not smooth, however, which is due to the presence of large-amplitude foreshock waves. Transition to quasi-perpendicular occurs for the initially nearly parallel case at the nominal shock ramp while for the initially quasi-parallel case it occurs at an upstream distance of about $\lesssim 100\lambda_i$ from the shock. This is sufficiently far away for the upstream flow ions to feel the change in the shock normal.

One notices that this transition is on the ion scale, implying that in the region close to the shock the ions experience the shock occasionally (because of the large fluctuations in Θ_{Bn}) – and when ultimately arriving at the shock – as quasi-perpendicular. It is thus not clear, whether the electrons do also see a quasi-perpendicular shock, here (which cannot be concluded from these hybrid simulations, however). The ISEE measurements of the electron distribution function by *Feldman et al* [1983] at the shock do not show a difference between quasi-perpendicular and quasi-parallel shocks. This fact suggests, in addition that, close to the shock transition, quasi-parallel shocks behave like quasi-perpendicular shocks as well on the electron scale, which is just what we have claimed.

We finally note that the behaviour of the shock normal angle gives a rather clear identification of the location of the shock transition in the quasi-parallel case, as indicated in Figure 6.27 by shading. Three distinctions can be noticed:

- first, at larger initial shock-normal angles the transition to quasi-perpendicular angles occurs earlier, i.e. farther upstream than for nearly parallel shocks. This is due to the stronger effect of the large amplitude upstream waves in this case;
- second, at larger initial shock normal angles the quasi-perpendicular shock transition is considerably broader than for nearly parallel shocks, i.e. it extends farther downstream before the main quasi-parallel direction of the magnetic field in the downstream region takes over again and dominates the direction of the magnetic field;
- third, at an initial shock-normal angle of 20° , this region is roughly $\sim 150\lambda_i$ wide, implying that the magnetic field direction behind a quasi-parallel shock remains to be quasi-perpendicular over quite a long downstream distance measured from the shock ramp. For the nearly parallel shock this volume is only about $\sim 50\lambda_i$ wide.

This observation must have interesting implications for the physics downstream of quasi-parallel shocks. For instance, applied to the Earth's bow shock, where $\lambda_i \sim 10^3$ km, both distances correspond to regions wider than the order of $>5R_E$ which is larger than the nominal width of the magnetosheath! Thus, behind the bow shock a substantial part of the magnetosheath plasma should behave as if the bow shock would have been a completely quasi-perpendicular shock.

6.3.2 Turbulent Reformation

When speaking about turbulent reformation we have in mind that a quasi-parallel supercritical shock is basically a transition from one lower entropy plasma state to another higher entropy plasma state that is mediated by a substantially broad wave spectrum. Such a transition has been proposed by *Schwartz & Burgess* [1991] based on the detection of the large amplitude magnetic pulsations (SLAMS) in the foreshock.

Figure 6.28 on the left shows their model assuming that somewhere upstream in the foreshock magnetic pulsations have been excited which become convected downstream toward the shock by the convective flow, grow in amplitude and number and accumulate at the shock transition to give rise to a spatially and temporarily highly variable transition

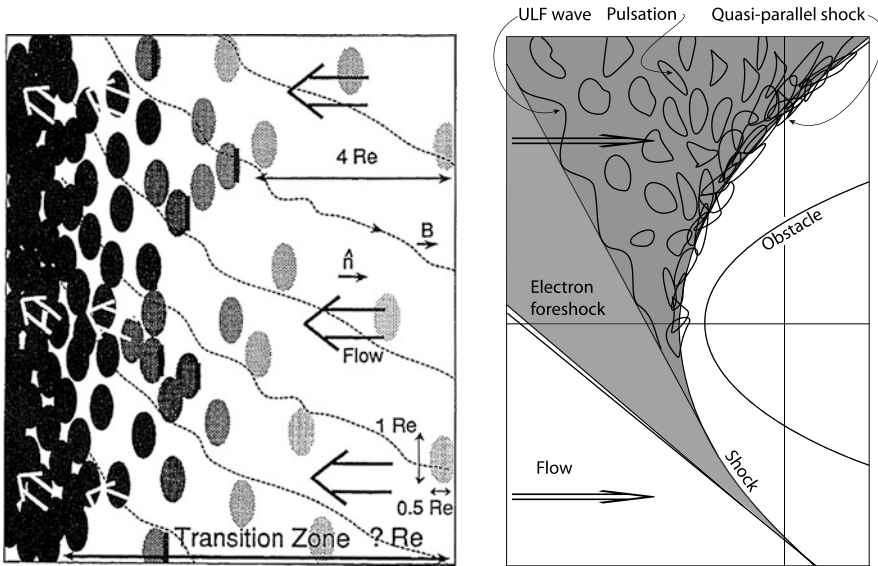


Figure 6.28: The patchwork model of *Schwartz & Burgess [1991]* of a quasi-parallel supercritical shock mentioned earlier. *Left:* Magnetic pulsations (SLAMS) grow in the ion foreshock and are convected toward the shock where they accumulate, thereby causing formation of an irregular shock structure. Note also the slight turning of the magnetic field into a direction to the shock normal that is more perpendicular, i.e. the magnetic field is more parallel to the shock surface with the shock surface itself becoming very irregular [after *Schwartz & Burgess, 1991*, courtesy American Geophysical Union]. *Right:* The same model with the pulsations being generated in the relatively broad ULF-wave-unstable region in greater proximity to the ion-foreshock boundary. When the ULF waves evolve to large amplitude and form localised structures these are convected toward the shock, grow, steepen, overlap, accumulate and lead to the build up of the irregular quasi-parallel shock structure which overlaps into the downstream direction.

from upstream of the shock to downstream of the shock. An important clue in this argument was the observation that, first, the pulsations grow in amplitude when approaching the shock and that, second, they slow down. This slowing down is effectively an increase in their upstream directed velocity on the plasma frame with growing amplitude such that their speed nearly compensates for the downstream convection of the flow.

On the right of the figure, which is suggested by the observations of *Kis et al [2007]*, a larger volume is seen. Here the pulsations are the result of growing ultra-low-frequency waves which are generated in a volume inside the foreshock but relatively close to the ion-foreshock boundary. These waves grow to large amplitudes until evolving into pulsations which the flow carries toward the shock. Growth, slowing down, and accumulation then lead to the pile up of the pulsations at the shock location and formation of the turbulent shock structure. Of the magnetic field (in this figure on the right) we plotted only the shock-tangential upstream field line. In the left part, several field lines are schematically shown exhibiting the fluctuations imposed by the background level of ultra-low-frequency fluctuations. Moreover, a certain bending of the field lines is included here in approaching the shock transition with the field lines turning more perpendicular the

closer they come to the shock. This bending is what we claim to be a parallel shock turning quasi-perpendicular at a scale very close to the shock. In this schematic drawing, however, there would be no reason for the magnetic field to turn this way. What closer observations and simulations show is that the turning of the field is the result of the presence of the large amplitude magnetic pulsations. This will become clearer below.

Observations

Of course, the model shown in Figure 6.28 is a schematic model only which, however, has some merits in explaining the observations. The signature of a quasi-parallel shock in the magnetic field is indeed quite different from that of a quasi-perpendicular shock. We have already seen in the electric recordings reproduced in Figure 6.25 that the quasi-parallel shock appears in the electric wave spectrum as a broadband emission of highest spectral intensities at the low frequency end. The magnetic signature of a quasi-parallel shock is quite similar in that it lacks a clear location of the shock front. Rather one detects a broad region of very large amplitude compressive oscillations in magnetic magnitude and in the direction of the magnetic field that subsequently is recognised as a passage across the quasi-parallel shock.

An example is shown in Figure 6.29 as measured by the CLUSTER spacecraft. This figure shows eight hours of observation by CLUSTER in the immediate vicinity of the quasi-parallel shock. It is difficult to say where in the figure the shock transition is located as the large fluctuations in the magnetic field magnitude and directions mask the various back and forth passages across the shock that are contained in the data. Clearly, at the beginning near 1400 UT the spacecraft was in the downstream region. The fluctuations show that during almost the entire sequence the magnetic field exhibits compressive fluctuations. These belong to the shock transition. At the same time large fluctuations in the direction of the magnetic field are also observed. In the compressions of the magnetic field buried are also upstream pulsations (SLAMS), and many of the changes in direction belong to the ultra-low-frequency waves present at and near the shock. The changes in direction indicate that the shock does not behave like a stationary flat surface. Instead, it shows structure with highly fluctuating local shock normal directions.

Lucek *et al* [2002] have checked this expectation by determining the local shock-normal angle Θ_{Bn} and comparing it with the prediction for Θ_{Bn} estimated from magnetic field measurements by the ACE spacecraft which was located farther out in the upstream flow. The interesting result is that during the checked time-interval of passage of the quasi-parallel shock the prediction for the shock normal was around $20\text{--}30^\circ$, as expected for quasi-parallel shocks. However, this value just set a lower bound on the actually measured shock normal angle. The measured Θ_{Bn} was highly fluctuating around much larger values and, in addition, showed a tendency to be close to 90° . This is a very important observation. It strengthens the claim that quasi-parallel shocks are locally, on the small scale, very close to perpendicular shocks, a property that they borrow from the large magnetic waves by which they are surrounded. In fact, we may even claim that locally, on the small scale, quasi-parallel shocks are quasi-perpendicular since the majority of the local shock normal

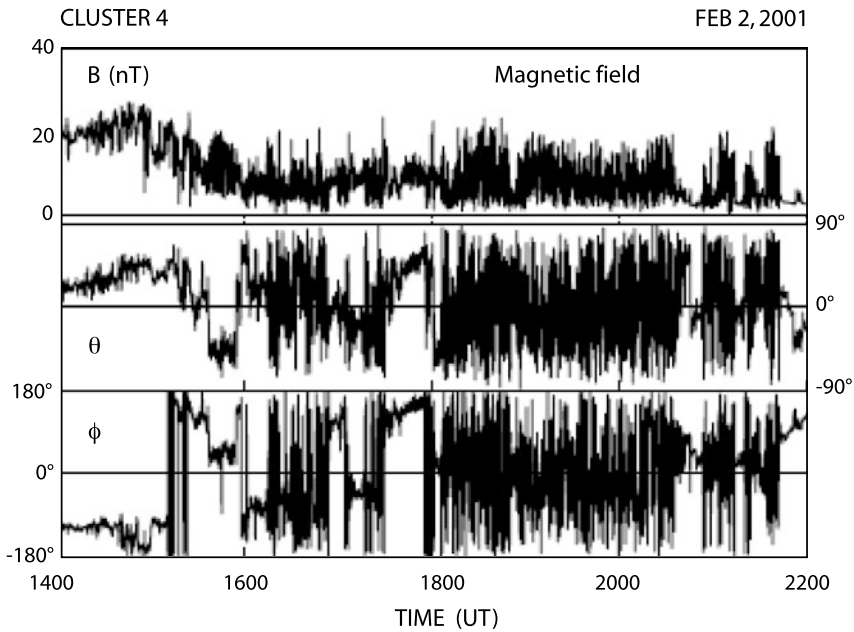


Figure 6.29: Eight hours of CLUSTER magnetic field data during a long passage near and across the quasi-parallel supercritical (Alfvénic Mach number $\mathcal{M}_A \sim 12\text{--}13$, ion inertial length $\lambda_i \sim 140$ km) bow shock. The time resolution was 4 s. The *top panel* shows the variation in the magnitude of the magnetic field. The *two lower panels* are the respective elevation and azimuthal angles θ, ϕ in a GSE coordinate frame [data taken from *Lucek et al., 2002*]. Large variation in the magnetic compression and direction can be seen to be associated with this quasi-parallel shock crossing. Buried in these large variations on this highly time-compressed scale are many magnetic pulsations (SLAMS). The compressive large amplitude fluctuations in the *upper panel* are typical for a quasi-parallel shock transition.

angles was $>45^\circ$. By small-scale a length scale comparable to a few times the ion inertial length or less is meant here.

The data suggest that, indeed, the quasi-parallel shock is the result of a build-up from upstream waves which continuously reorganise and reform the shock. Figure 6.30 shows three representative examples of such upstream waves which are far from being continuous wave trains. The upper panel is taken from the large density fluctuation region in the shock transition. This region turns out to consist of many embedded magnetic pulsations (SLAMS) of very large amplitudes. In the present case amplitudes reach $|\mathbf{b}| \sim 25$ nT. These pulsations have steep flanks and quite irregular shape, exhibit higher frequency oscillations probably propagating in the whistler mode while sitting on the feet or shoulders of the pulsations. It is most interesting that the different CLUSTER spacecraft – at spacecraft separation <1000 km – do not observe one coherent picture of a particular pulsation. This implies that the pulsations in the shock are of shorter scale than spacecraft separation: the different spacecraft observe different structures respectively different pulsations. In addition, the magnetic field directions (not shown in the figure) are very different from

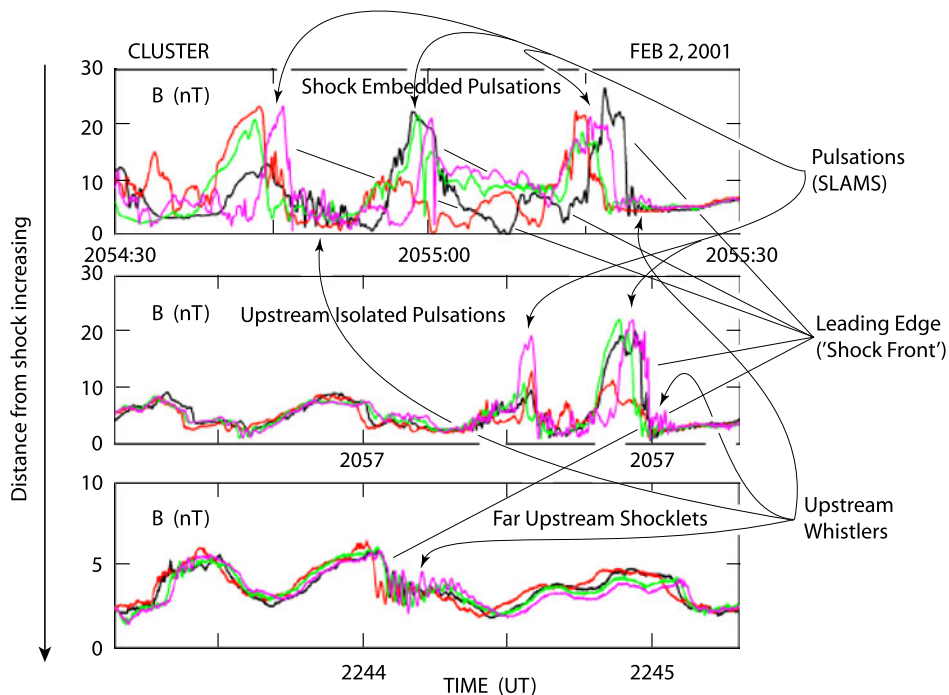


Figure 6.30: CLUSTER magnetic field measurements of magnetic pulsations (SLAMS) near and remote from a quasi-parallel shock on February 2, 2001 [data from *Lucek et al., 2002*]. Only the magnetic field magnitude is shown for all four CLUSTER (colour coded) spacecraft. The spacecraft separation was between a few 100 km and 1000 km. *Top:* Clustered pulsations in the shock transition. Three events of large amplitudes are shown. These structures are very irregular with steep fronts. Note that in spite of the small spacecraft separation the shapes of the structures differ strongly from spacecraft to spacecraft. Moreover, determination of the pulsation fronts and normals (not shown) indicates high variability over the spacecraft separation distance. Thus the structures are of relatively small scale and large amplitude. *Middle:* Isolated pulsation at greater distance from shock outside compression region. The structures are seen almost simultaneously at the spacecraft and thus must be of larger size. Compared to the embedded pulsations the amplitudes are lower, and the structures are more regular. *Bottom:* A shocklet observed outside the pulsation region in the domain of ultra-low-frequency waves. The steep shock-like front is well expressed with the attached whistler waves it carries with it. Note the much lower amplitude than the pulsations.

spacecraft to spacecraft and thus from pulsation to pulsation, and even for one pulsation at its front edge and trailing edge different magnetic field directions are observed. The directions of the magnetic normal across a pulsation change on very short spatial scales. The quasi-parallel shock front has thus a rather irregular shape, which will be bent locally with changing direction of its normal.

The second panel shows an isolated pulsation farther away from the shock transition. This pulsation exhibits a much more coherent way on the four different CLUSTER spacecraft. It seems as if the pulsation is still in evolution as three of the spacecraft see a nearly coherent structure while the fourth which is farther away sees it in a different state. Con-

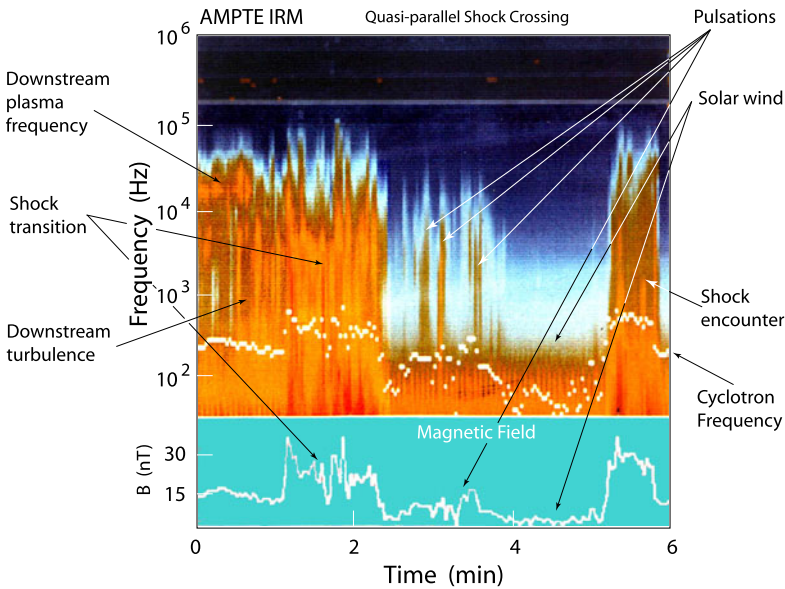


Figure 6.31: An AMPTE IRM quasi-parallel shock crossing in the plasma wave picture. The shock plasma wave spectrum is broadband. The magnetic field variation in the lower part exhibits the presence of large amplitude pulsations SLAMS. In plasma waves these appear as broadband signals on the background of moderate foreshock wave activity. Pulsations are thus very active in exciting plasma waves, resembling shock elements [from R. A. Treumann & J. LaBelle 1986, unpublished].

cluding from this event, isolated pulsations seem to have larger dimensions and lower amplitudes, which would be consistent with the assumed solitary properties of pulsations.

The third panel shows a shocklet, i.e. a structure which presumably has little in common with pulsations. It is embedded into long wavelength ultra-low-frequency wave trains, evolves into steep front and drives whistlers attached to this front across the flow. These waves were already observed by *Russell et al* [1971]. Their properties indeed resemble those of sub-critical little shocks which propagate against the flow, though with slower speed such that they effectively are slowly convected towards the shock. Observations of plasma wave spectra by AMPTE IRM are shown in Figure 6.31. There the large amplitude SLAMS pulsations appear as broadband wave signatures resembling narrow shock structures.

Behlke et al [2003] report another interesting property of magnetic pulsations in the shock transition region where they overlap to form the quasi-parallel shock. Measurement of the electric cross-SLAMS potential identify a substantial unipolar drop in the electric potential of several 100 V, corresponding to a potential ramp, when passing from upstream to downstream across the pulsation. Such a drop signifies the presence of an electric field in one direction across the pulsation. Taking the mean size of a pulsation to be roughly 1000 km, the mean electric field is $\langle E \rangle \simeq 400 \text{ mV m}^{-1}$. However, this field drops mainly at the leading edge of the pulsation. Such a field presumably corresponds to a steep pressure

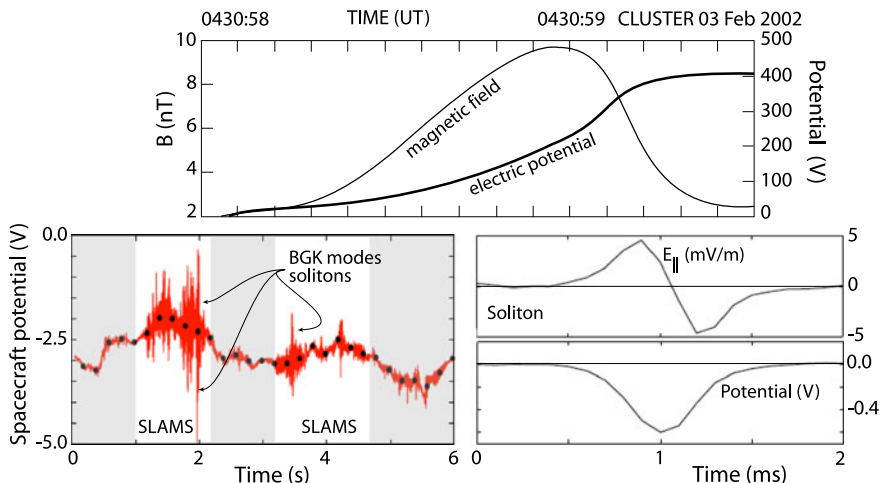


Figure 6.32: Observation of electric field structures in large magnetic pulsations (SLAMS) in the quasi-parallel shock transition region [data taken from *Behlke et al*, 2003, 2004, courtesy American Geophysical Union]. Structures on three different time scales are shown, corresponding also to three different spatial scales. *Top*: CLUSTER passage across on (moderately large amplitude) magnetic pulsation in the shock transition. The (smoothed) magnetic field structure is a slightly steepened magnetic bump. The stationary parallel electric potential field across this structure shows a potential ramp with steep gradient at the leading edge of the pulsation. The potential drop of ~ 400 V corresponds to an electric field of ~ 0.47 V m $^{-1}$. Note that the time scale in this panel is 90 s. *Bottom left*: Six seconds of a CLUSTER passage through the shock transition. The black dots show the spacecraft potential variation which maps the local density variation. Overlaid is the high frequency WHISPER trace of the plasma frequency. In the magnetic pulsation regions (white) the plasma frequency exhibits huge excursions to both sides similar to those seen in the overview Figure 6.25 on wave observations. These excursions trace the BGK (non-symmetric) modes and (symmetric) solitons. *Bottom right*: One example of one of the solitons on a 2 ms time scale. It is nicely seen how symmetric the parallel potential trough and the corresponding bipolar parallel electric field shape look like in the solitary wave structure.

gradient in the pulsation. It could (but is unlikely) also be generated by an anomalous collision frequency. This remains to be tested by further observations and data analysis.

The measurements of *Behlke et al* [2003] anticipated the more recent report of strong electric fields in the shock by *Bale & Møzer* [2007]. In addition to this observation it was found [*Behlke et al*, 2004] that the single pulsations were subject to a fairly large number of high frequency/Debye scale structures in the electric field seen in the WHISPER recordings (bottom panel), which belong to electron holes or solitons which form in the pulsation gradient regions as shown in Figure 6.32. The bipolar electric field and unipolar potential across one – indeed very symmetric – soliton is seen in this figure. These observations suggest that the pulsations are indeed the main constituents of a quasi-parallel shock with the dynamics on the micro-scale of a quasi-parallel shock going on mainly in the single pulsations of which the shock transition is built.

The occurrence of these intense nonlinear electrostatic electron plasma waves at the quasi-parallel shock transition is intriguing. It forces one to draw another very important main conclusion from these observations (and other related observations like those of *Bale*

& Mozer [2007]): that quasi-parallel shocks are sources of electron acceleration into electron beams, which are capable to move upstream along the magnetic field over a certain distance and excite electron plasma waves at intensity high enough to enter into the non-linear regime, forming solitons and electron holes (BGK modes). This is possible only if quasi-parallel shocks are quasi-perpendicular as well on the electron scale.

Indeed, electron reflection into beams which generate Langmuir waves is presumably possible only when in the supercritical quasi-parallel shock transition region the magnetic field changes from quasi-parallel to quasi-perpendicular on the electron scale $\sim \lambda_e$. Indications of such a change on the ion scale λ_i have been noted above at a number of occasions, but the detection of solitary structures in the electron plasma waves in relation to quasi-parallel shock transitions provides a very strong argument for this to be true on a scale which is well below the ion scale. Only if this is the case, there will be ample reason for electrons to become reflected and accelerated into beams from the transition region in a quasi-parallel shock. As we already noted, we may, therefore, expect that quasi-parallel supercritical shocks on the electron scale are not anymore quasi-parallel but change to become locally quasi-perpendicular, while on the larger ion scale they still maintain properties of quasi-parallelity. Because of this reason high Mach number quasi-parallel shocks will be embedded into a diffuse high energy electron component that emits radiation.

If this conclusion will turn out to be true and will sustain future more sophisticated experimental tests, it will have important consequences for collisionless shock physics. Supercritical collisionless – nonrelativistic – shocks will, in fact on the electron scale, always behave quasi-perpendicularly – and it may be suspected that this conjecture will also hold for relativistic shocks though probably for other reasons (like the generation of transverse magnetic fields by the Weibel instability, which becomes dominant in relativistic shocks [see, e.g., Jaroschek *et al.*, 2004, 2005]). This implies also that the true quasi-parallel shock physics cannot be properly elucidated when ignoring electron effects as is, for instance, done in hybrid simulations.

Simulations of Quasi-parallel Shock Reformation

Nevertheless, before turning to the very few more realistic full particle PIC simulations with different mass ratios from small to about realistic, we are first going to discuss here hybrid simulations in one and two dimensions of the formation and behaviour of quasi-parallel shocks. Initial particle [Quest *et al.*, 1983; Quest, 1985, 1988] and hybrid simulations [Burgess, 1989; Thomas *et al.*, 1990; Winske *et al.*, 1990; Scholer & Terasawa, 1990; Scholer & Burgess, 1992] did already illuminate some of the particular interrelations between the dynamics of quasi-parallel shocks: energy dissipation by short wavelength whistlers at the shock transition, the presence of a diffuse hot ion component upstream of the shock, and the importance of upstream waves. Most of these have been the subject of reviews [the interested reader might be directed to the papers by Burgess, 1997; Burgess *et al.*, 2005; Lembège *et al.*, 2004, which more or less systematically discuss selected aspects of these interactions]. The selection of results we will provide here is guided by the progress that has been achieved in the understanding of the quasi-parallel shock physics and in its relation to the observations.

The first arising question is, whether nearly parallel shocks are at all capable of reflecting ions. This question is neither nonsensical nor academic. It makes sense, because the investigation of the critical Mach number by *Kennel et al* [1985] becomes unreliable at small angles, say $\Theta_{Bn} \lesssim 30^\circ$, because it is based on linear whistler dispersion and does not take into account the completely modified plasma conditions near the shock ramp. Thus the simple conclusion that the critical Mach number approaches $\mathcal{M}_{\text{crit}} \rightarrow 1$ for $\Theta_{Bn} \rightarrow 0$ might be an unjustified extrapolation. Theoretically, the reflection of particles at small Θ_{Bn} must become entirely due to the electrostatic shock potential drop with the magnetic part of the Lorentz force being obsolete. The potential drop, however, should decrease with $\mathcal{M}$, and thus the reflection of particles should cease, while being of vital importance for the generation of a collisionless quasi-parallel shock. Hence, the question is also not purely academic.

Of course, we do already know from the observations that quasi-parallel shocks exist at small Θ_{Bn} with their foreshocks being populated by a diffuse ion component that excites upstream waves and mediates the beam-generated upstream ion-foreshock boundary waves. The impossibility for this diffuse component of being entirely due to shock reflection in the quasi-perpendicular part of the shock, immediately proves that the quasi-parallel (or even the nearly parallel shock) must be able to reflect particles upstream. Hence, either a quasi-parallel shock is capable of generating a large cross-shock potential, or it is capable of stochastically – or nearly stochastically – scattering ions in the shock transition region in pitch angle and energy in such a way that part of the incoming ion distribution can escape upstream, or – on a scale that affects the ion motion – a quasi-parallel shock close to the shock transition becomes sufficiently quasi-perpendicular that ions are reflected in the same way as if they encountered a quasi-perpendicular shock.

Observations suggest that the latter is the case, while observations also suggest that large potential drops occur in the large-amplitude magnetic pulsations (SLAMS) where they accumulate in the shock ramp [*Behlke et al*, 2004]. Hence, reflection of ions will be due to the combination of both effects, the electric potential drop and the magnetic deflection. In fact, this can be a quite complicated process for an ion passing across a number of magnetic pulsations, in each of which it is being retarded and at the same time deflected by a small angle until its normal velocity component is decreased sufficiently that a further deflection in pitch angle suffices to let it return into the upstream region.

Hybrid Simulations in 1D. In agreement with what is known today, first one-dimensional hybrid simulations in an extended simulation box [*Burgess*, 1989] suggested that the reformation of quasi-parallel shocks is about cyclic and is caused by the impact of large-amplitude upstream waves. *Scholer & Terasawa* [1990], using $\Theta_{Bn} = 20^\circ$ and $\mathcal{M}_A = 3.5$ with an upstream ion thermal velocity $v_i = v_A$ (corresponding to $\beta_i = 1$) in one-dimensional hybrid simulations (with small numerical resistivity) showed that the reflected ions are not coming from the core of the incident upstream ion distribution but originate in the shell of this distribution, having initial velocities $v \gtrsim 1.7 v_i$. These ions escape from the shock quite far upstream and excite ultra-low frequency waves with upstream directed velocity of $\sim 1.3 V_A$ at distances up to $> 300 \lambda_i$, which are convected downstream to reach the shock.

In this one-dimensional hybrid simulation the only mode in which they can propagate is the compressive fast magnetosonic mode.

These waves are in fact what in observations has been identified as pulsations (SLAMS) but is not yet recognised as such, here. During downstream convection the waves grow and slow down in the interaction with the foreshock ion component. When approaching the shock they generate a large amount of new reflected ions. These slow the incident ion population down and steepen the wave crest, which becomes the new shock front. In the time between the arrival of the compressive waves the shock is about stationary and develops phase-locked upstream whistlers which the arriving next wave crest destroys. From these simulations it could not be concluded what process produced the reflected ions, however, as one-dimensional simulations among suffering from other deficiencies select only one particular direction of wave numbers and are thus not general enough for drawing final conclusions.

The nature, generation and effects of the large-amplitude upstream waves have been further investigated in more detail in one-dimensional [Krauss-Varban & Omidi, 1991; Scholer & Burgess, 1992; Scholer, 1993, among others], and in two-dimensional hybrid simulations [Krauss-Varban & Omidi, 1993; Scholer, 1993; Dubouloz & Scholer, 1995]. Since shocks are three-dimensional, it is clear that two-dimensional numerical simulations at same resolution come closer to reality. However, they suffer from restrictions in size of the simulation box and simulation time. Since reality does not confront us with an initial state, large boxes and long times are needed. However, for the investigation of particular questions, simulations have the great advantage of providing the possibility to prepare them for answering just those questions. Concerning the propagation of upstream waves in one direction with respect to the shock normal, one-dimensional simulations are just good enough.

In order to identify a particular wave mode, the dispersion of the wave must be investigated. This dispersion relation depends on the frame in which it is taken, because the energy/frequency of a wave is not invariant with respect to coordinate transformations; in a medium moving with velocity $\mathbf{V}$ it is Doppler shifted according to $\omega' = \omega(\mathbf{k}) - \mathbf{k} \cdot \mathbf{V}$, where $\omega(\mathbf{k})$ is the dispersion relation in the rest frame of the flow. While the Doppler shift at high frequency is negligible, it completely changes the dispersion of ultra-low frequency waves at large Mach numbers. Figure 6.33 on its left shows the deformation of the fast mode dispersion relation in the shock frame at large Mach number $\mathcal{M} > 1$ and for waves propagating upstream in the plasma frame. The deformation causes negative frequencies of the waves which imply downstream directed phase velocities, which is nothing else but the intuitive downstream convection of the waves by the flow. However, the minimum in the dispersion relation implies that waves of a particular frequency and wave number have zero group velocities. In the shock frame the energy of these waves is stationary. Smaller wave numbers have energy moving downstream, larger wave number have energy moving upstream away from the shock. The right part of the figure shows simulations of upstream waves according to one-dimensional hybrid simulations by Krauss-Varban & Omidi [1991] for $\Theta_{Bn} = 30^\circ$ and a Mach number $\mathcal{M}_A = 2.8$. The entire dispersion of the simulated waves is negative. The waves are all convected toward the shock as their Mach number is less than the streaming Mach number. Near the shock most of the wave energy

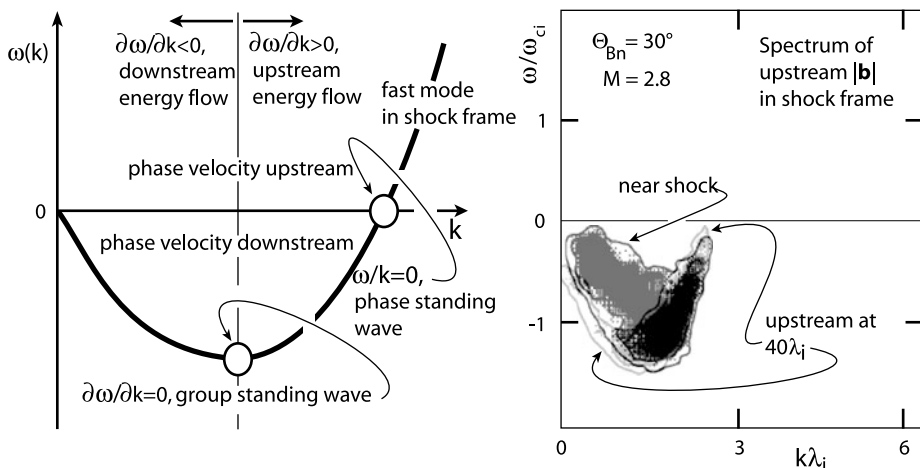


Figure 6.33: Fast mode dispersion relation in simulations in the shock frame. *Left:* The Doppler shifted fast mode dispersion relation in a supercritical flow in the shock frame. The Dispersion relation assumes negative frequencies corresponding to the downstream convection by the flow of Mach number $\mathcal{M}_A = 2.8$ and at $\Theta_{Bn} = 30^\circ$. Waves at zero group velocity have energy at rest in the shock frame. Negative group velocities imply downstream transport of energy, positive group velocities imply upstream transport. *Right:* Simulated upstream wave dispersion spectra near the shock and upstream of the shock. Near the shock wave energy accumulated around standing and downstream transport. Away from the shock the wave energy still moves upstream [data from Krauss-Varban & Omidi, 1991, courtesy American Geophysical Union].

moves downstream and will cross the shock. Still some shorter wavelength waves (large k) move in energy upstream in the shock frame. Farther away from the shock most of the wave energy encountered is seen to move upstream.

Scholer [1993] investigated these waves further in one-dimensional hybrid simulations performing several numerical experiments on them, taking away the shock and instead injecting a diffuse ion component from downstream. The main finding is that the large amplitude upstream magnetic pulsations (SLAMS) evolve out of the ultra-low frequency wave spectrum in the interaction with the diffuse ion component. In accord with observation the pulsations move upstream in the plasma frame. Thereby their upstream leading edge steepens and is right-hand circularly polarised like required for whistlers. However, dispersion is unimportant; the main cause of the evolution of large pulsations is nonlinearity when the wave interacts with the diffuse ion distribution. This distribution has a steeper shock directed density gradient than the pulsation wavelength. Moreover, the flow becomes decelerated at the leading edge of the pulsation (as is seen in Figure 6.26), and here the velocity difference between the flow and the diffuse ion component drastically decreases, which shifts the k vector of the resonant wave to larger values, and the wavelength decreases during convection of the pulsation toward the shock. (Note that no resonant ion beam-whistler interaction exists as the beam is hot and diffuse.) The standing whistlers in the leading edge are simply generated by the current flowing in the edge. It is thus concluded that it is the gradient in the hot diffuse ion component over a length of the

same order as the wavelength which produces the pulsations. Ultimately these pulsations cause a quasi-periodic reformation of the shock, as we have described earlier, which thus has been proven by one-dimensional hybrid simulations. The same result is obtained when the simulation starts right away without a shock but with an injected hot beam (which is no surprise as the generation of the shock, before it was removed in the former simulations in order to keep with the wave field, was due to the plasma flow-reflected ion beam interaction).

Hybrid Simulations in 2D. The two-dimensional evolution of the pulsation (SLAMS) was studied later by *Scholer & Fujimoto* [1993] and *Dubouloz & Scholer* [1995] (see Figure 6.26). It basically confirmed the conclusions drawn from one-dimensional simulations with the following three important modifications: (1) the wave fronts of the pulsations (SLAMS) rotate into a direction that is more parallel to the shock thereby increasing the shock-normal angle locally to become quasi-perpendicular; (2) the pulsations have short wavelength in shock normal direction, but are of substantially longer but finite lengths in the direction tangential to the shock, which provides structure to the shock in tangential direction; and (3) shock reformation is a result due to the steeping and accumulation of the pulsations and is a quasi-periodic process, but the downstream structure of the shock over some distance is caused by the downstream convection of the old shock front, i.e. the bulk of the pulsations that had accumulated at the location of the former shock transition. It is interesting to note that from one reformation cycle to the next rather large fluctuations in the magnetic field and density exist in the transition from upstream to downstream which are capable of trapping particles. The two-dimensional simulations do also confirm the conclusion that the diffuse upstream ion component is responsible for the growth of the pulsations (SLAMS). The two-dimensional simulations by *Dubouloz & Scholer* [1995], in particular, followed the same scheme as the one-dimensional simulations, injecting hot diffuse ions into upstream in order to control the interaction between the diffuse ion component and large amplitude pulsations (SLAMS).

Full Particle PIC Simulations. So far we dealt just with hybrid simulations where the ions are macro-particles while the electrons represent a charge-neutralising background of zero mass. Clearly, such simulations are unrealistic if whistlers become involved. This is, however, the case, as we have discussed above, whenever large amplitude pulsations evolve at the leading edge to which (partially phase-locked) whistlers are attached. The question what role the electrons play in the evolution of the pulsations can only be answered by full particle simulations. These require large simulation boxes and at the same time high temporal and spatial resolutions. So far they could therefore only be performed in one dimension [*Pantellini et al*, 1992; *Scholer et al*, 2003; *Tsubouchi & Lembège*, 2004]. *Pantellini et al* [1992] used a mass ratio of $m_i/m_e = 100$. They had too small a box (just $\sim 30\lambda_i$ in the upstream direction) for following the evolution of upstream waves but stressed the importance of whistlers in shock reformation. One decade later it became possible to substantially enlarge the box and at the same time to switch to a larger mass ratio while staying with one dimension only. *Scholer et al* [2003], using the same mass ratio at $\Theta_{Bn} = 30^\circ$ and

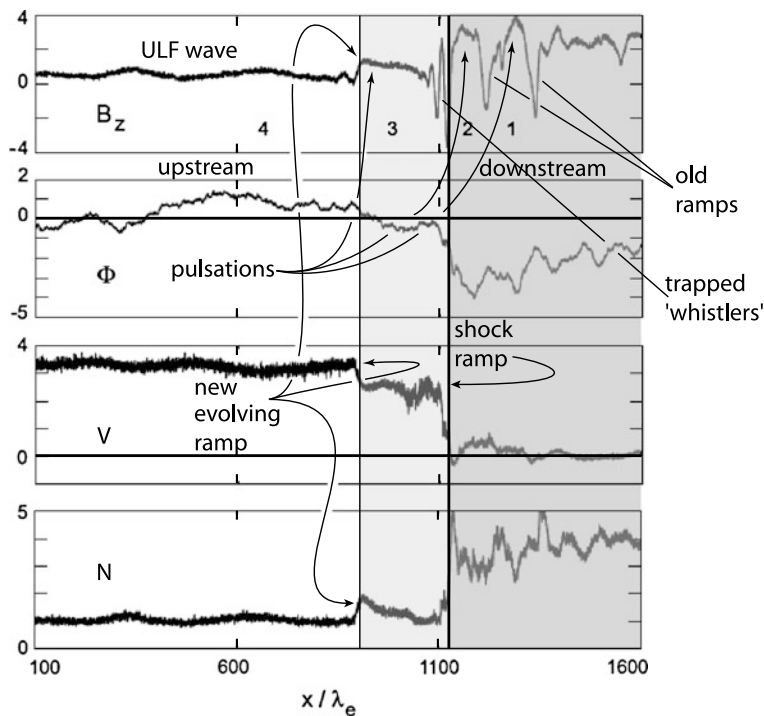


Figure 6.34: Full particle PIC simulations of the evolution of a quasi-parallel shock in one dimension only [simulation data from *Scholer et al., 2003*, courtesy American Geophysical Union]. From top: main magnetic field component B_z , electric potential Φ , bulk plasma flow velocity V , density N , all in simulation units. The *numbers* indicate three pulsations (SLAMS). Pulsation 1 was the old shock. Pulsation 2 is the actual shock coinciding with the drop in V to zero and the steep increase in density and potential. Pulsation 3 is just evolving. It already has a steep leading upstream edge and decelerates the plasma flow. It will become the next shock ramp. Number 4 indicates a bump in the upstream waves that will become a pulsation. The actual shock ramp has some phase-locked whistlers attached to it which are not well resolved on the scale shown (see the next figure).

$\mathcal{M}_A \simeq 4.7$, had an upstream extension of $\sim 200 \lambda_i$ and could follow the shock evolution for a time $t\omega_{ci} \sim 100$. Similar simulations with mass ratio $m_i/m_e = 50$ have been performed by *Tsubouchi & Lembège [2004]*. Here we discuss in greater detail the most informative simulations of *Scholer et al [2003]*. An overview of their results is shown in Figure 6.34 in the fixed lab-frame for the main magnetic field component B_z , electric shock potential Φ , bulk stream velocity along the shock normal V , and density N , all as functions of distance x (measured in electron inertial lengths λ_e).

Far upstream from the quasi-parallel shock the magnetic field exhibits long wavelength ultra-low frequency waves (4) which, when approaching the shock, start steeping at their leading edges (note that in the plasma frame these waves are moving upstream, as also does the shock, i.e. to the left in the figure). The amplitude of the wave increases (3), and the wave becomes a pulsation (SLAMS). Its amplitude is large enough to already substantially brake the upstream bulk flow, which causes a drop in V and in Φ , and an increase

in density. The pulsation has slowed down the flow to a velocity $\sim 2.8 V_A$ already here. In fact the leading edge of the pulsation behaves like a ‘baby shock’, which later will become the real ‘adult shock’. The shock itself is formed further downstream at the local position of the leading steep edge of the previous pulsation (2). In front of this edge (i.e. the genuine shock at this instant and location) between it and the trailing edge of pulsation (3) a standing (possibly phase-locked) large amplitude whistler has evolved. This whistler is spatially damped by the approaching pulsation (3). An indication of such a whistler has already been seen in front of the leading edge of pulsation (3) as well. Pulsation (2) (the instantaneous shock) has a substantial downstream extension. Further downstream of it the ‘old shock’ is seen, which was formed at an earlier time by pulsation (1); and even farther downstream a remainder of earlier shock ramps (pulsations) is recognised in the trace of B_z . The instantaneous shock ramp (pulsation 2) is high enough to locally completely brake the upstream flow, the velocity of which drops to zero while the density steeply increases and forms a dense wall.

Figure 6.35 gives an impression of the shock evolution in higher temporal and spatial resolution. It shows instantaneous magnetic profiles in a box of lengths $400 \lambda_e$ including the shock ramp with time axis running upward. In this simulation frame the shock moves upstream to the left. The heavy step-like line indicates the approximate location of the nominal shock ramp. Also shown is an upstream ultra-low frequency wave that at time $t \omega_{ci} = 82.5$ has begun to evolve into a pulsation while convectively approaching the shock, growing in amplitude and developing a steep leading (upstream) ramp in front of which magnetosonic whistlers start growing. When approaching the shock at $t \omega_{ci} = 92.5$, the pulsation kills the phase-locked whistlers that were waiting in front of the ramp by damping them out. At this time the pulsation takes over the role of the shock, and the nominal shock position jumps ahead to upstream by roughly a distance of the width of the pulsation $\sim 100 \lambda_e$. This process repeats itself at $t \omega_{ci} = 102.5$ showing that the quasi-parallel (on this time scale) shock is not stationary but undergoes nearly periodic reformation which is mediated by the arrival of large amplitude pulsations.

In order to illuminate what happens in the shock ramp, the ion phase space data have been overlaid at time $t \omega_{ci} = 100$ just before a new reformation. The shock ramp coincides with the location where the upstream flow is completely braked and the hot downstream ion distribution begins. This location corresponds to the leading edge/ramp of the first broad downstream pulsation. On the other hand, the upstream whistlers that are attached to its ramp are seen to trap ions, which form large phase space vortices (holes) with ions bouncing back and forth in the associated potential wells. These are not shown in the figure but resemble the magnetic oscillations, being minimum in the centres of the holes. However, what is really interesting to note is that there is little indication of any back streaming reflected ion beam in this simulation. Instead, a very dilute but broad completely diffuse ion component is seen all over the simulation box, which is denser downstream than upstream but of similar velocity dispersion. Somehow this component has been generated by the shock, and it is which is responsible for the growth of the pulsations.

The evolution of the proton and electron distribution functions is shown in Figure 6.36. The two proton (ion) distributions are taken in the phases when fully developed whistlers with trapped ions and phase-space vortex formation exist, and when these whistlers have

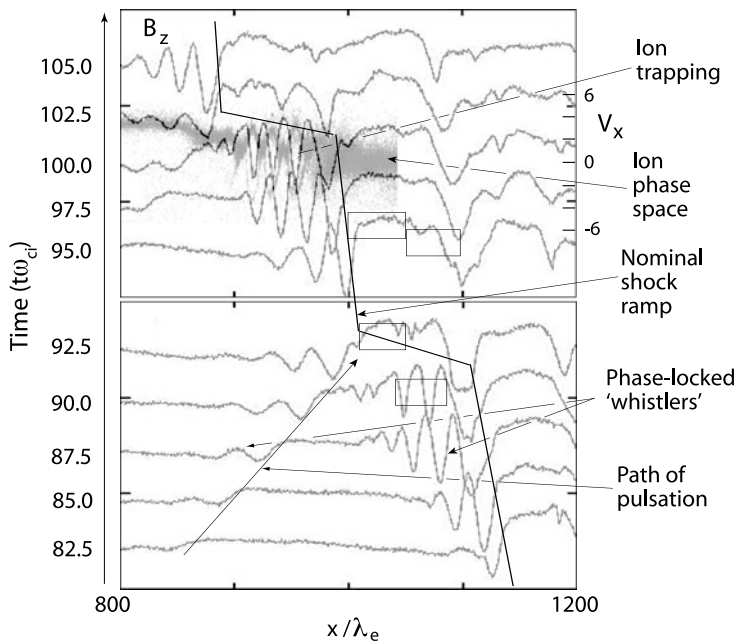


Figure 6.35: Full particle PIC simulations of the evolution of a quasi-parallel shock in one dimension [simulation data taken from *Scholer et al., 2003*, courtesy American Geophysical Union]. The profile of the main magnetic component B_z is shown for subsequent simulation times shifted by $\Delta t\omega_{ci} = 2.5$ upward. The representation is in the simulation frame, i.e. the shock moves to the left into the upstream direction. The simulation shows the reformation of the quasi-parallel shock resulting from the exchange of the shock with an incoming upstream wave which has steepened to become a pulsation (SLAMS). The magnetic field trace at time $t\omega_{ci} = 100$ has been overlaid on the ion phase space at this time. The heavy steps show the location of the nominal shock ramp (where the flow is stopped). It moves slowly upstream until a new pulsation arrives and when it suddenly jumps forward by roughly $100\lambda_e$. Also shown is the fast approach of an upstream pulsation starting at $t\omega_{ci} = 82.5$ and arriving at the shock at $t\omega_{ci} = 92.5$ to take over the role of the shock. Note that in the minima of the ‘whistler’ field fluctuations (at $t\omega_{ci} = 100$) ions are trapped, oscillating back and forth and forming hole vortices in phase space centred around local minima of the electric potential Φ (not shown). The little boxes indicate where particle (ion or electron) phase space distributions have been determined.

been completely damped by the trapped component. In the former case at time $t\omega_{ci} = 90$ the (perpendicular) ion distribution shows substantial heating with respect to the upstream ion distribution, has been slowed down by a factor of roughly ~ 8 in velocity and, in addition, exhibits a non-thermal top-flat part in the upstream direction which is due to the presence of the trapped bouncing ions in the phase space vortices. At time $t\omega_{ci} = 95$ when the whistlers are completely damped and the pulsation has taken over to become the shock, the trapped vortices have disappeared, as has been shown above, and the distribution evolved into a broad and hot distribution at about zero flow velocity. Thus the vortices are entirely due to the interaction between the trailing edge of the shock-arriving pulsation and the whistlers with the latter being compressed, amplified in amplitude until they trap ions, while the pulsation itself is free of those trapped ions. This was already clear from Figure 6.35 but its effect on the distribution is nicely seen here.

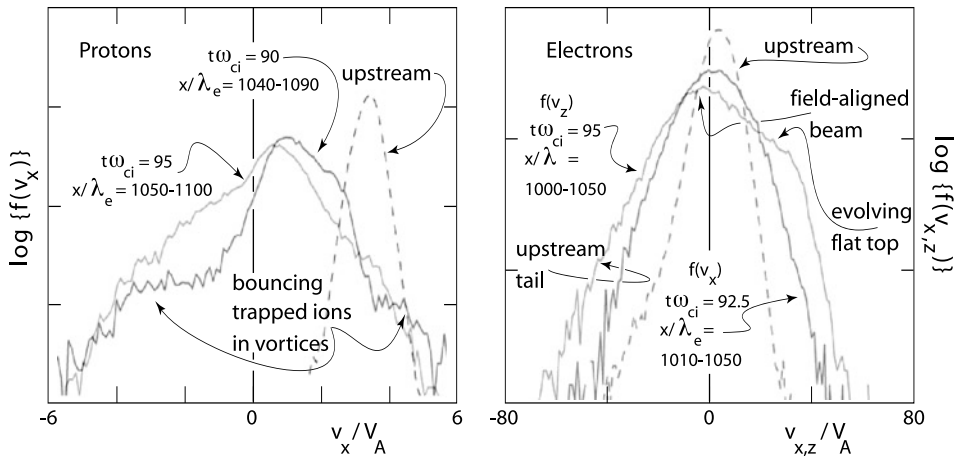


Figure 6.36: Velocity distribution functions $f_i(v_x)$ for the full particle PIC simulations of Figure 6.35 during different phases of shock evolution as indicated by the boxes in that figure [simulation data taken from *Scholer et al., 2003*, courtesy American Geophysical Union]. *Left:* Proton distribution functions at times $t\omega_{ci} = 90$ and $t\omega_{ci} = 95$. The first period is in the arriving pulsation before it takes over to become the shock. Shown is the part of the pulsation where it encounters the whistlers. The upstream distribution has changed here into a heated distribution exhibiting a substantial upstream directed tail which is due to the trapped ion component. The whole distribution assumed a large velocity spread and is strongly heated. At time $t\omega_{ci} = 95$ the trapped ions have damped the whistlers. The distribution has now become very hot but has a two-temperature structure with high-energy tail in the leading part of the pulsation. At this stage the incident plasma is completely slowed down. *Right:* Two electron velocity distributions $f_e(v_x)$ at time $t\omega_{ci} = 92.5$ and $f_e(v_z)$ at time $t\omega_{ci} = 95$, respectively. The former is right in the new shock built up of a fresh pulsation showing a hot but otherwise structure-less perpendicular electron distribution, the latter is in the well-developed shock but along the main magnetic field. This parallel distribution is strongly non-symmetric, heated, exhibits an upstream tail, indication of the formation of a top flat distribution and a remaining upstream-beam-like part similar to those distributions as had been measured by *Feldman et al [1983]*.

The right part of Figure 6.36 shows two cases of electron distributions. The first in the perpendicular direction, the second in the parallel direction taken in the arriving pulsation when it becomes the shock, i.e. after the reformation. The boxes are given in Figure 6.35. The perpendicular distribution $f_e(v_x)$ simply shows the heated, about symmetric electron distribution in the shock. However, the parallel distribution $f_e(v_z)$ exhibits a number of interesting features. It is also hot, in fact substantially hotter than the perpendicular distribution, exhibiting non-thermal tails in both directions along the magnetic field with a stronger upstream tail. In addition, on the downstream side it shows the evolution of a broad flat top on the distribution with a remainder of an non-flattened upstream directed beam on the upstream side. This is interesting as it is almost exactly the type of distributions that had been reported by *Feldman et al [1983]* from crossings of the super-critical quasi-parallel bow shock [see also *Feldman, 1985; Gurnett, 1985*].

Conclusions. Full particle simulations provide much deeper insight into the shock physics than did the hybrid simulations. Still, they have been only one-dimensional and thus

are restrict to waves propagating solely into shock normal direction. Moreover, the mass ratio is still not realistic as it competes with the length of the box, and it is obvious that the length of the box is crucial in quasi-linear shock physics for properly investigating the evolution of the waves which build the shock, and for their interaction with the particle component. It is, in particular, of considerable interest that, even though the magnetic field should allow the streaming of particle upstream in a quasi-parallel shock, no up-streaming beams have been seen in the simulations. Electrons exhibit a signature of an up-streaming beam, but ions do not. The diffuse component which is homogeneously distributed over the entire region depending only on the distance from the shock is obviously participating in a diffusive process and is not injected into the upstream region. This sheds light on the acceleration and confinement process acting in quasi-parallel shocks. These simulations have demonstrated that quasi-parallel supercritical shocks are non-stationary on the time scale of the simulations. They are subject to reformation. However, this reformation is quite different from the reformation of a quasi-perpendicular shock as it is not due to the presence of gyro-bunched reflected ions in the foot of the shock. Rather it is caused by the accumulation of large amplitude magnetic pulsations which are not separate entities but grow out of the upstream ultra-low frequency waves in the range where the wavelength of the wave is of the same order as the density-gradient scale of the diffuse ion population in the foreshock. *Scholer & Fujimoto [1993]*, *Dubouloz & Scholer [1995]* and *Scholer et al [2003]* have shown that extracting the diffuse ion component results in a lack of upstream pulsations, while injecting additional diffuse ions speeds the evolution of large amplitude pulsations up. Thus, pulsations are the result of the fine-tuning of the resonance of ultra-low frequency upstream waves with the diffuse ion component. These pulsations play an important role in quasi-parallel shock. Still, their dynamics has not been completely clarified and, in addition, is subject to some controversy. Closer investigation of pulsation dynamics is therefore of vital interest.

To a certain extent this has been the subject of the one-dimensional full particle PIC simulations performed by *Tsubouchi & Lembège [2004]* who use a modest mass ratio $m_i/m_e = 50$ [see also *Lembège et al, 2004*] which, however, allows them to extend the simulation time scale to times as long as $t\omega_{ci} \sim 38$ and to enlarge the simulation box to $\sim 700\lambda_i$. Their simulation is based on the magnetic piston method instead of the reflecting wall that has been used by *Scholer et al [2003]*. In general these simulations confirm the previous results, adding some facts about the structure, size and life times of isolated pulsations whose sizes are determined to lie in the interval between $(10-20)\lambda_i$ with the tendency of a gradual shrinkage during the reformation process. This shrinkage can simply be attributed to the dynamical compression of the pulsations at the shock transition where many embedded pulsations compete for the available shock volume.

Also confirmed is the rotation of the magnetic field from quasi-parallel to quasi-perpendicular in the reformation process. *Large amplitude pulsations (SLAMS) behave like localised supercritical quasi-perpendicular shocks of which the quasi-parallel shock is constructed.* We stress the importance of this conclusion here again. Moreover, their evolution is accompanied by the generation of phase-locked whistler precursors at the leading edge of a pulsation, as has been shown above to be the case [*Scholer et al, 2003*]. *Tsubouchi & Lembège [2004]* report in addition the observation of large-amplitude spiky

electric fields in the leading edge of a pulsation. In view of the observations by *Behlke et al* [2004], their investigation, however, will require even much better time resolution on the electron plasma scale. Structures of this kind require the inclusion of the full electron dynamics and imply accounting for realistic mass ratios in the simulations.

Inclusion of particle dynamics is one of the main problems in quasi-parallel shock physics. We have seen that the evolution of the building blocks of a quasi-parallel shock, i.e. the large amplitude pulsations (SLAMS), is completely determined by the interaction of the ultra-low upstream wave spectrum with the diffuse ion component. On the other hand, the pulsations build up the quasi-parallel shock transition region which is responsible for the generation of the diffuse ion component in a way which has not yet been satisfactorily clarified. The current knowledge will be discussed in the next paper in this series.

One idea is that the quasi-perpendicular subshocks of which the quasi-parallel shock is constructed – and which are nothing else but the various pulsations of which the shock consists on the small scale of a few λ_i – reflect sufficiently many ions back upstream. However, because of the complicated magnetic structure of the shock transition region, these ions cannot simply escape along the magnetic field to upstream but remain trapped for long enough time in the shock transition, being scattered in pitch angle and energy until ultimately picking up enough energy in order to leave the shock either upstream or downstream in a diffusion process that works in the shock and partially also in the foreshock regions. The observation of a nearly isotropic diffuse ion distribution both upstream and downstream of the shock in the simulations and the lack of observation of ion beams, both in the measurements and in the simulations, provides a strong argument for a diffusive mechanism to act in the shock transition. The observations by *Kis et al* [2007] also support this conclusion as they show that the diffuse ion component exhibits an exponential density gradient towards the shock with the shock being their source. Such a gradient is typical for a diffusive process.

Whether electrons undergo a similar process, remains an unresolved question that will be attacked with the increasing capacities of computers when the full electron dynamics can be included into the simulations. There are strong hints on the importance of electron dynamics in the presence of the observed spiky large amplitude electric fields in the pulsation and shock transitions. These electric field structures point on the generation of electron beams. Since no such beams have been observed in the deep upstream foreshock, the structure of the shock on the small scale presumably inhibits the escape of electrons to far upstream in the form of beams. This can be due to either a chaotic magnetic field configuration on the electron scale which makes the escape of electrons along the magnetic field impossible. It can also be due to the trapping of electrons in the phase space holes which are related to the smallest-scale electric field structures. These, on the other hand, require strong electron currents to flow in the shock transition and, if present, will stochastically heat the electrons. Moreover, they might be responsible for the electromagnetic radiation generated in the shock ramp, a process that has been observed in solar and interplanetary type II radio bursts but never found a convincing explanation by any known mechanism. The finding in the PIC simulations of the formation of top-flat electron distributions and a rudimentary upstream directed field aligned beam in the shock transition [*Scholer et al*, 2003] might indicate that the electrons are accelerated in the embedded

pulsations by electric fields and can form a field aligned upstream beam which apparently cannot leave far upstream but possibly has sufficient free energy in the electrons for driving either a Buneman or modified two-stream instability which causes electron plasma waves, electron holes, and radiation – and at the same time scatters, heats, and confines the electrons.

6.4 Hot Flow Anomalies

Would the collisionless upstream flow be stationary at constant Mach number $\mathcal{M} > M_{\text{crit}}$, we could safely terminate the discussion of the supercritical shock transition regions at this place and focus our attention on the downstream part of the shock transition. Unfortunately, this is not the case. The upstream flow is usually in a highly disturbed fluctuating state with the smallest disturbances in it readily evolving into structures, which can appear in various form, either as turbulence on almost all scales, from the electron scales up to scales of a substantial fraction of the macroscopic dimension of the flow, as localised large amplitude waves, as shocklets, or as various kinds of boundaries in the flow, which appear as discontinuities and current layers.

Then the question arises, what happens when any such forms impact on the shock and interact with it. As long as we are dealing with the interaction of a single wave with the shock one needs to investigate the conditions for wave reflection, refraction and transmission across the shock. This is most easily done for wavelength much longer than the shock transition scale (including the foreshock extension!) and leads to modifications of Snell's law when taking into account the Rankine-Hugoniot jump conditions at the shock, for instance, when the MHD approximation holds. For shorter wavelengths the continuous change in the wave propagation conditions as given in the dispersion relation can be parametrically treated. The shock may then act as a spectral filter depending on which waves it allows to pass.

6.4.1 Observations

Of quite a different kind of problems is the interaction of a current sheet or discontinuity boundary with the shock. Current sheets and discontinuities are large amplitude distortions of the upstream flow and cannot be treated like waves interacting with a shock. Observations near the Earth's bow shock have shown that they lead to severe distortions of the shock, which have been termed 'Hot Flow Anomalies' (HFAs, the term now commonly used even though it does not appropriately refer to the physics involved), 'Active Current Sheets', 'Diamagnetic Cavities' or '3D-Plasma Structures with Anomalous Flows'. They have been discovered in spacecraft transitions of the Earth's bow shock from analysing AMPTE and ISEE measurements [Schwartz *et al*, 1985, 1988; Thomsen *et al*, 1986; Woolliscroft *et al*, 1986, 1987; Fuselier *et al*, 1987; Paschmann *et al*, 1988].

The first observation of the HFA phenomenon was made by Schwartz *et al* [1985] in the plasma and field observations of AMPTE UKS who detected an unusual very hot upstream plasma event resembling a broad layer of plasma in which the upstream magnetic field

performed a large rotation as it is known from current sheets. The most surprising observation was, however, that the upstream flow inside this current sheet was sharply deflected by a substantial angle close to $\sim 90^\circ$ from its original upstream direction, a deflection far in excess of the deflection the bow shock would provide and in addition not expected to occur in front of the shock in its upstream region. The upstream flow was about stopped from its original flow direction. This deflection was also different from the one predicted by the magnetic stresses across the boundary of the hot flow anomaly, and it was clear that it required a violent momentum exchange between the upstream flow, the shock, and some (unknown) agent. It was also found that the structure was not in its final state but was evolving. Hence, most probably it did not have come all the way along from the source of the upstream flow, the Sun, and it was speculated that it was either caused by an upstream-flow current sheet interacting with the supercritical bow shock or also could have had its origin in a violent change of the configuration of the obstacle, i.e. the magnetosphere in this case, an idea that was advocated later on [Paschmann *et al.*, 1988]. It took about one decade of measurement and theoretical investigation until a consensus was reached that the HFA phenomenon was a shock phenomenon and was not caused by the obstacle. Clearly, violent changes in the obstacle configuration do also affect the shock [Paschmann *et al.*, 1988; Sibeck *et al.*, 1999] but most probably to a lesser extent than it happens to be observed in hot flow anomalies.

The signatures of a Hot Flow Anomaly in the upstream foreshock plasma are: (1) a strong deflection of the bulk flow velocity, (2) reduced magnetic field strength, (3) reduced density, and (4) considerable heating of the plasma. The region where all these changes in the plasma parameters are observed is usually flanked by walls of enhanced plasma density and magnetic field strength when both, the upstream plasma and upstream magnetic fields pile up.

All these properties are nicely seen in the most recent observation of a Hot Flow Anomaly by THEMIS plotted in Figure 6.37. Similar observations by the CLUSTER spacecraft have been reported by Lucek *et al.* [2004], but while all four CLUSTER spacecraft were upstream of the shock, the THEMIS spacecraft were both upstream and downstream such that they could follow the evolution of the HFA from upstream to downstream.

The most spectacular effect in the observation of a Hot Flow Anomaly is seen in the fifth panel on the left of this figure. The velocity component V_x turns positive in the HFA, which corresponds to a complete stopping or even reversal of the flow in the downstream direction. The other two flow velocity components are enhanced simultaneously, indicating a deflection of the flow in the (y, z) -direction.

Inspecting the magnetic field, one finds that its magnitude drops to a small value. However the traces of its transverse components become dipolar, which is the sign of a current that is flowing inside the HFA. To both sides of the HFA current both the magnetic field and density pile up. It is interesting that this pile up is not symmetric but is stronger on the upstream side. This should be so since the HFA must stop the upstream flow, which causes a stronger compression on the upstream than on the downstream sides. In the figure the two pile up regions are shown in shading. They flank the current region where the plasma is dilute and hot, as is seen from the strong increase in temperature and from the ion energy flux in the uppermost panel which exhibits a broadening of the distribution that

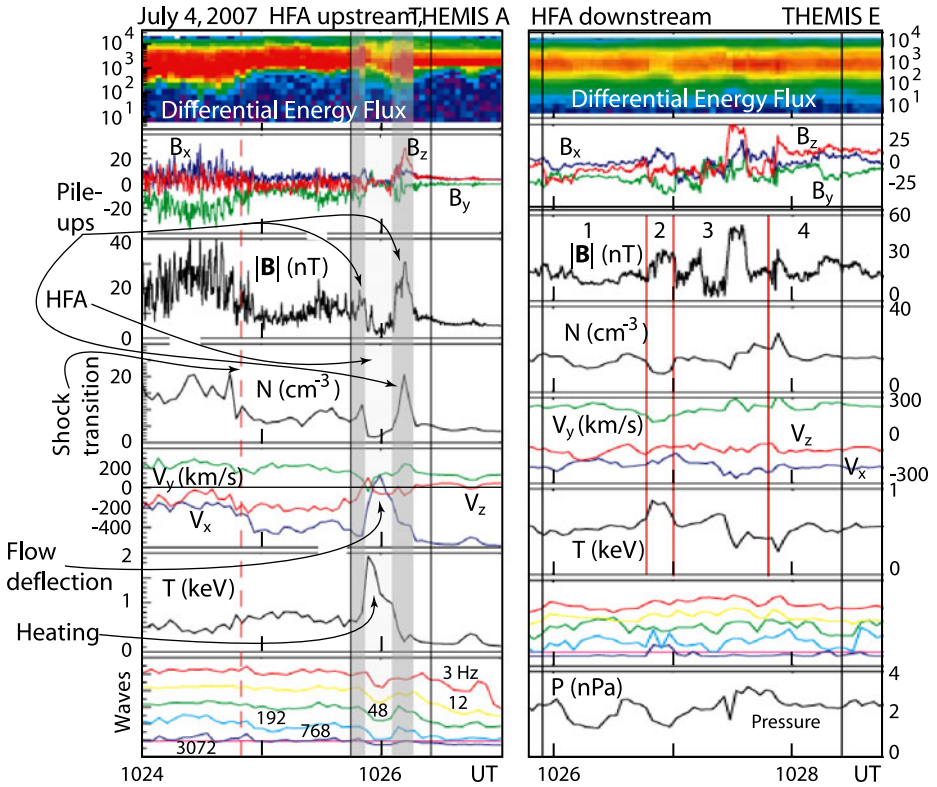


Figure 6.37: THEMIS spacecraft measurements of a Hot Flow Anomaly upstream and downstream of the quasi-parallel Earth's bow shock wave [data taken from *Eastwood et al, 2008*, courtesy American Geophysical Union]. *Left:* The THEMIS A observations in the foreshock close to the shock transition. Three regions of the HFA have been shaded: darker shading refers to the pile-up regions at the HFA edges, The magnetic depression region in the centre of the HFA is shaded light. The quasi-parallel shock transition is about 1 min before the HFA encounter. The deflection (retardation in V_x) of the upstream flow is clearly seen. It coincides with a drop in density and magnetic field and strong plasma heating. Bipolar shape of B_z , B_y indicates the presence of a current sheet. Note also the asymmetry of the HFA pile-up edges. The upstream edge which is exposed to the flow is much stronger piled up. These regions are themselves shocks. The *lowest panel* is the power spectral density of magnetic waves at different frequencies. *Right:* The simultaneous THEMIS E observations downstream of the shock behind the shock transition. The signature of the HFA is very irregular here. The later onset confirms that the HFA is an upstream phenomenon.

is comparable to the shock transition at the beginning of the panel. However, though the temperature is high, the mean energy of the ions is low, reflecting the strong deceleration of the flow.

The right part of the figure shows simultaneous measurements downstream of the shock. These indicate that the HFA downstream is very complex, causing a strong disturbance in the downstream flow. It, moreover, sets on later than in the upstream region which is a confirmation for the HFA being an upstream phenomenon and is not caused by

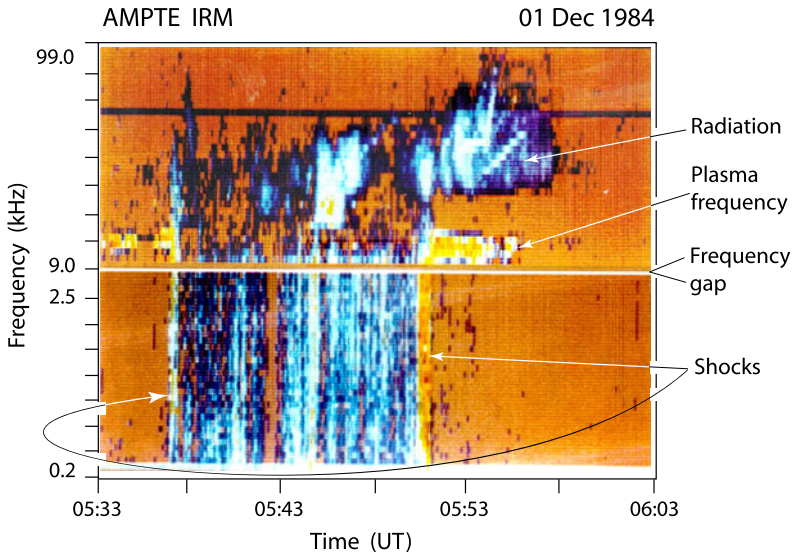


Figure 6.38: AMPTE IRM plasma wave measurements during a rare passage of an HFA in the frequency range >200 Hz. (Note the instrumental frequency gap between 2.5 kHz and 9 kHz.) The transition from the surrounding to the interior of the HFA is quite sharp in the plasma wave spectrogram and is reminiscent on a shock transition. In particular the broadband emission that flanks the interior is similar to shocks, and the occurrence of the plasma frequency at ~ 20 kHz is typical for electron reflection and acceleration. Note the broadband high activity of plasma waves in the interior. In addition intense radio waves are generated by the structure exhibiting strange traces which drift in frequency [from R. A. Treumann & J. LaBelle 1986, unpublished].

the obstacle or downstream of the shock. This rules out a number of models that had been proposed in the past.

We will not go into the details of this downstream observation other than point on the much more diffuse shape of the HFA after its “passage” across the shock to downstream. It has become dissolved into a number of different regions which still belong together while each differs from the others. The larger of these regions are numbered from 1 to 4 in the figure. Some of them showing very clear signatures (correlated magnetic plus density compressions, regions 1 and 4) of fast modes, others of slow expansion modes (anti-correlated magnetic and plasma variations plus plasma heating, region 2). In any case, it becomes clear that the shock has split the HFA into different filaments. The HFA has lost the compactness, consistence, and coherence it possessed in the foreshock and during the interaction with the shock. It would contradict the imagination (and the second law) if such a diffuse downstream structure would combine to make up a coherent HFA of the kind we see in front of the shock. In particular, its later downstream than upstream onset, which causes a time delay between the upstream and downstream disturbances, provides a very strong causal argument against a downstream origin of the HFA, at least in this particular case and for the class of HFAs that exhibit similar properties.

For completion we also show plasma wave measurements during an AMPTE IRM passage of an HFA on December 01, 1984 in Figure 6.38. The wave activity outside the HFA

is surprisingly low. It is concentrated to the interior of the HFA where broadband noise dominates the spectrum below $\lesssim 9$ kHz. Note also the very sharp transition in the spectrum to both sides of the HFA from the outside to the interior which is strongly reminiscent on a shock transition. Moreover, just in front of the structure intense electron plasma waves at the plasma frequency ω_{pe} are seen (20–30) kHz signalling electron reflection and acceleration. These emissions are accompanied by intense and surprisingly broadband radiation of radio waves above ~ 40 kHz exhibiting substructure that drifts upward in frequency. Some of this radiation is also seen to originate in the interior of the HFA [from Treumann & LaBelle 1986, unpublished].

We note that *Thomsen et al* [1993] investigated a large number of Hot Flow Anomalies and determined the directions of their normal vectors, applying the minimum variance method. They found that the normal vectors are in good agreement with the assumption that upstream tangential discontinuities are involved into HFAs. This assumption since dominates the interpretation and theory of Hot Flow Anomalies.

6.4.2 Models and Simulations

The mechanism of hot flow anomalies is believed to be related to the interaction of the shock with an upstream current sheet [*Burgess & Schwartz, 1988; Thomas et al, 1991*], most probably a tangential discontinuity (TD), in the near-shock quasi-parallel foreshock. (It has also been proposed that rotational discontinuities in the upstream flow (RD) would cause a violent change in the shock structure.) It is this model that is most strongly supported by the recent THEMIS observations shown in Figure 6.37 and also by the CLUSTER observations reported by *Lucek et al* [2004]. The current sheet (tangential discontinuity) deflects the upstream reflected ion component, channels the ions back and focusses them along the current sheet. This produces a highly localised hot ion population which starts expanding, blowing off the upstream flow to both sides, deflecting the flow from its original direction, and piling the upstream plasma up to the sides of the current sheet where it causes new shock waves.

Tangential discontinuities have the property that they are convected with the upstream flow. The model is shown schematically in Figure 6.39. There is no flow across the discontinuity. Thus, in the upstream frame of reference the velocity component V_n perpendicular to the discontinuity is strictly zero; in the shock or laboratory frames it is the same on both sides of the discontinuity. The normal component of the upstream magnetic field $B_n \equiv 0$ vanishes in all frames. The flow velocity $\mathbf{V}_T$ and the upstream magnetic field $\mathbf{B}_T$ are both tangential to the discontinuity to both of its sides but can be rotated by an arbitrary angle and can have different magnitudes on both sides, while the total pressure $P = NT + B^2/2\mu_0$ is constant across the discontinuity in the co-moving upstream frame. Thus, the density, tangential magnetic field, and tangential velocity can change arbitrarily across the discontinuity. The temperature change is then fixed by the continuity of the pressure (we do not consider here an anisotropic pressure, however).

Any change in the magnetic field, which is a rotation of the field in the discontinuity plane plus a stretching of the tangential component, corresponds to an electric current flowing in the discontinuity. (Note that different magnitudes of the tangential magnetic

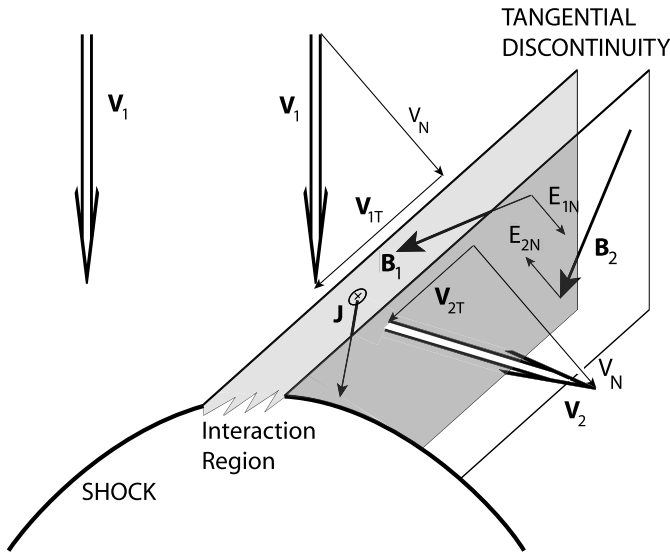


Figure 6.39: Schematic of the interaction of an upstream tangential discontinuity with the shock. The tangential discontinuity is convected along with the normal component V_n of the upstream flow which, in the shock frame, is thus the same on both sides. However, the tangential velocities and magnetic fields are different on both sides of the discontinuity. The latter is due to the inclined current flow $\mathbf{J}$ in the discontinuity. Here the rotation of the magnetic field is shown in such a way that the normal components of the electric field E_n point inward on both sides of the discontinuity thus confining any shock-reflected ions to the interior of the discontinuity which here is shown as an extended region. Due to the presence of the discontinuity and interaction with the shock the upstream flow is deflected from its original direction, because the tangential velocity has changed across the discontinuity. For other current direction the tangential magnetic field would rotate differently. E_n can then have other directions, pointing away from the discontinuity on both or also only on one side.

fields to both sides of the discontinuity simply mean that the current flowing in the discontinuity is not distributed homogeneously over the width of the discontinuity, while a rotation of the field across the discontinuity implies an inclined current.)

In addition, tangential discontinuities have one property that turns out to be of vital importance for the formation of a Hot Flow Anomaly. Tangential discontinuities possess a non-vanishing normal electric field component $\mathbf{E}_n = -\mathbf{V}_T \times \mathbf{B}_T$. This field vanishes only when $\mathbf{V}_T$ and $\mathbf{B}_T$ are either parallel or anti-parallel. However, its direction can be along the external normal to the tangential discontinuity (pointing away from the discontinuity) or anti-parallel to the external normal (pointing toward the discontinuity). In the first case, ions will be removed from the discontinuity by this field, while in the second case ions will be returned to the discontinuity. One expects that only in the second case a Hot Flow Anomaly is created in the interaction with the shock, while the first case will cause only a minor distortion of the shock. One should, however, keep in mind that near to the interaction region the character of the tangential discontinuity is lost as the observation of the HFA being bounded by shocks suggests. This behaviour can be interpreted as being due to a deformation of the shock and thus is a structure of the shock

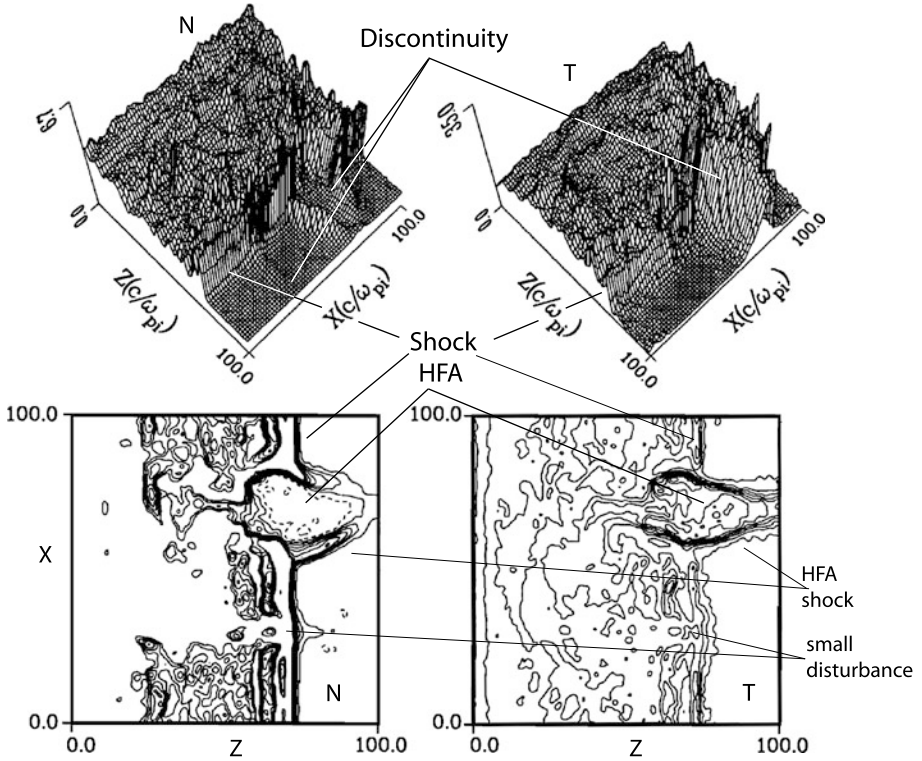


Figure 6.40: Two-dimensional hybrid simulation of the formation of a Hot Flow Anomaly (HFA) in the interaction of a tangential discontinuity with a quasi-perpendicular shock [after Thomas *et al.*, 1991, courtesy American Geophysical Union]. The idea is that the tangential discontinuity first touches the shock at the quasi-perpendicular site because the upstream magnetic field is contained in the discontinuity plane, initially interacting with the shock-reflected ions. Two discontinuities are used. One with normal electric field $E_n = E_n^+$ pointing toward the discontinuity, the other with $E_n = E_n^-$ pointing away from the discontinuity. Only the former generates an HFA by capturing the shock reflected ions. The other causes only a minor disturbance of the shock. The figure shows the density N and plasma temperature T in the simulation plane (x, z) . *Top:* Stacked profiles of density and temperature. The E_n^+ -discontinuity produces a major disturbance in N and T at the shock surface, consisting of two dense and hot walls surrounding a diluted region. *Bottom:* Iso-contours of N and T showing the extension of the HFA disturbance on the shock profile.

and not of the tangential discontinuity. Hybrid simulations show that this is indeed the case.

Two-Dimensional Hybrid Simulation. An example is given in Figure 6.40 which shows the results of a two-dimensional hybrid simulation of the interaction of two discontinuities with a shock [Thomas *et al.*, 1991]. The only difference between the two discontinuities is that the normal electric fields $\mathbf{E}_n = -\mathbf{V}_T \times \mathbf{B}_T$ to both sides of the discontinuity at $x = 25$ points away from the discontinuity while to both sides on the discontinuity at $x = 75$ it points into the discontinuity. The effect on the interaction between the disconti-

nuities and the shock is dramatic. The $x = 25$ discontinuity causes only a minor distortion on the shock while the discontinuity at $x = 75$ digs a big hole into the shock and at the same time produces a violent effect in the upstream region which shows all the signs of an HFA. Note that because of technical reasons the simulation was performed assuming that the discontinuity impinges on a perpendicular shock, which is a reasonable assumption for the initial time of the interaction, because the discontinuity, being a tangential discontinuity and therefore containing in its plane the full upstream magnetic field, first touches the shock in the area of the perpendicular shock where the upstream field is tangential to the shock. Only later it is swept over the shock surface by the convective flow sitting for the longest time in the larger quasi-parallel region where HFAs are usually observed.

The explanation of the simulation result is quite simple even though the complete physics of the interaction is still poorly understood. The quasi-perpendicular supercritical shock reflects ions. When $E_n = E_n^-$ points away from the discontinuity all reflected ions in the interaction region between shock and discontinuity are channelled away from the shock. This has the effect, that shock reformation is inhibited by the discontinuity, and the shock becomes nearly stationary with the exception of minor distortion around the intersection point, as is seen in the lower part of the figure around $x = 25$. There is a slight increase in density N at the interaction point in front and a dilution behind the shock surface along the direction of the discontinuity where all the fluctuations are damped out by the discontinuity, and the shock front becomes narrow, weekend, and stable. A similar effect is seen in the temperature T .

The more interesting violent distortion in the case $E_n = E_n^+$ around $x = 75$ is caused by the channelling of the shock-reflected ions into the discontinuity. Inside of the tangential discontinuity the ions become approximately non-magnetised due to the presence of the current (which generates the discontinuity) and the correspondingly low internal magnetic field. The isotropic ions accumulate and cause a hot dilute ion cloud that compresses the shock in the downstream direction, excavate the upstream flow from the region, deflect it into tangential direction, and cause the density and magnetic field to pile up at the boundaries of the discontinuity. The shock is hereby violently deformed both in the upstream and downstream directions. In addition, upstream shocks are formed at the discontinuity boundaries (in the non-symmetric case only one shock boundary would form); downstream a HFA bubble structure is caused with sharp, dense and hot walls, as seen from the figure. These simulations are hybrid and thus not completely reliable as the effect of the electrons is not taken into account. Here the electron fluid is simply heated. Moreover, they use a perpendicular shock, i.e. the simulation is relevant only for the initial state of HFA formation.

‘Global’ Hybrid Simulations. The question arises whether the HFA survives or is modulated when moving into the quasi-parallel shock region. This question has been investigated with the help of a *global* two-dimensional hybrid simulation [Omidi & Sibeck, 2007] which takes account of the entire obstacle, in this particular case the magnetosphere in its interaction with the supersonic solar wind. In such a simulation, because of the finite transverse size of the obstacle (the blunt magnetosphere), the bow shock assumes its natural curvature around and distance from the blunt obstacle. A tangential discontinuity that

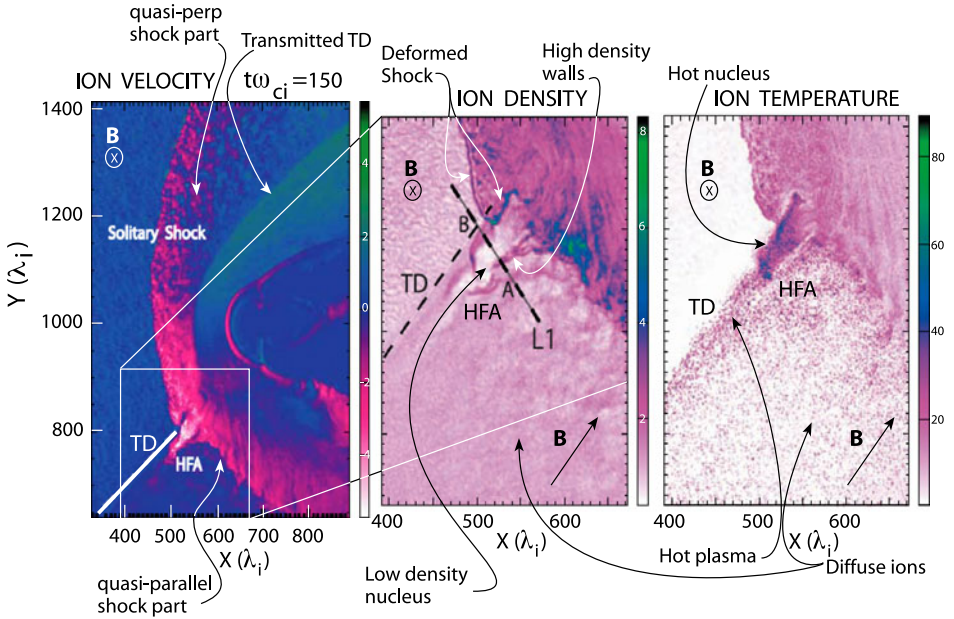


Figure 6.41: “Global” hybrid simulations of the interaction of a tangential discontinuity in the upstream medium (solar wind) with a supercritical shock (bow shock) [data taken from *Omidi & Sibeck, 2007*, courtesy American Geophysical Union]. The obstacle is the magnetosphere. *Left:* The bulk velocity in colour coding. The simulation plane is the noon meridian plane (North-South). The magnetosphere is on the right, upstream is on the left. Originally the magnetic field had a shock normal angle $\Theta_{Bn} = 45^\circ$, i.e. the upper half of the shock is quasi-perpendicular, the lower half quasi-parallel. The tangential discontinuity (TD) is parallel to the field (*white fat line in the lower left corner*). Approaching the shock from left it moves down in the box from upper left corner to its position shown here at time $t\omega_{ci} = 150$. Behind the TD the magnetic field points into the simulation plane. Hence E_n points away from the TD behind and into the TD in front of the discontinuity. The case is non-symmetric. *Centre:* Enlargement of the white box in the panel on the left. Shown is the ion density. *Right:* Ion temperature in the enlarged box.

is convected toward the obstacle thus starts interacting with the self-consistent shock first in the quasi-perpendicular region before reaching the quasi-parallel domain. On the other hand, global simulations suffer, in addition to being hybrid only, from the largeness of the scale that can be resolved. In these simulations the cell size is $\sim 1\lambda_i \times 1\lambda_i$. Moreover, global simulations are not completely collisionless, which introduces some non-realistic element through the resistive scale $\ell_\eta = \lambda_i \sqrt{v_c/\omega}$, where v_c is the numerical collision frequency, and $\tau = \omega^{-1}$ is the time scale of the variation of the field quantities.

Figure 6.41 summarises the results of the global simulation. On the left of this figure a large part of the simulation box is shown at time $t\omega_{ci} = 150$ when the tangential discontinuity (TD) has arrived in the quasi-parallel shock region. Shown is the bulk flow velocity in colour coding. Upstream directed velocities are from white to red, downstream directed from green to blue. The two panels in the centre and on the right are enlargements of the white box on the left as indicated. They show, in higher spatial resolution, the

ion densities and ion temperatures, respectively. The original upstream magnetic field (in front, i.e. downstream of the tangential discontinuity TD) was directed at a shock normal angle $\Theta_{Bn} = 45^\circ$ (as shown in the lower right corners of the central and right panels). Thus the original shock was quasi-perpendicular in the upper half of the left panel, while in the lower half it was quasi-parallel. Behind the TD, i.e. upstream of it the magnetic field points into the simulation plane, thus being rotated by 90° in the discontinuity plane. The TD being inclined like the original magnetic field at an angle of 45° arrived at the shock from above in the left panel touching it first at its quasi-perpendicular part, and then moved down along the shock into the quasi-parallel domain. The first and very interesting observation in this simulation was that during the entire pass along the quasi-perpendicular bow shock no HFA was created. This observation is somewhat disturbing in view of the simulations by *Thomas et al* [1991] who used a perpendicular shock. Possibly this is because of the asymmetry of the normal electric field in the “global” simulation which inhibits accumulation of reflected ions on the backside of the TD.

When the TD arrives at the quasi-parallel shock a Hot Flow Anomaly (HFA in the figure) is immediately formed. This is shown for the time $t\omega_{ci} = 150$ in the figure. The left panel plots the local ion velocity. White-to-red colouring indicates upstream directed velocities, conversely green-to-blue colouring indicates downstream directed velocities. The magnitude of the velocity has been colour coded (as given by the bar on the right of the left panel in relative simulation units). The Hot Flow Anomaly appears only on the front side of the TD and is seen as an extended region of fast upstream directed ions. These ions are concentrated and flowing along the TD, which has caused a distortion of the quasi-parallel shock.

The central and right panels are enlargements of the white box region in the left panel, but showing now the density and temperature variations, respectively. In both cases white is low, green and blue high. Again, in the density and temperature the HFA appears along the front side of the TD as a very dilute hot plasma region that is bounded by dense plasma walls which are approximately parallel to the TD.

The bow shock is heavily distorted and deformed by the presence of the TD in a way similar to what had been found in the local hybrid simulations of *Thomas et al* [1991], bending into the downstream direction, away from its original shield shape. On the downstream side of the TD there is no sign of any distortion of the medium other than that the bow shock has become a broad transition region of upstream directed velocities with a sharp, dense upstream edge, which *Omidi & Sibeck* [2007] identify as a “solitary shock”. This sharp edge is also seen in the density and temperature plots.

Most interesting is the strong effect of the TD on the ion population. Upstream of the TD, i.e. behind it, the plasma is cold and dilute and of high downstream directed velocity. Downstream, in front of the TD, the plasma is hot and dense, being composed of the hot diffuse ion component that populates the upstream foreshock of the quasi-parallel shock. The TD sweeps this diffuse ion component along when moving with the flow, forming a very sharp boundary between the two regions. It thus seems that the diffuse upstream ion component at the quasi-parallel shock is heavily involved into the formation of the HFA. These ions are absent at the quasi-perpendicular shock. However, the mechanism is not clear in this case, and while it is very interesting to see these differences between the local

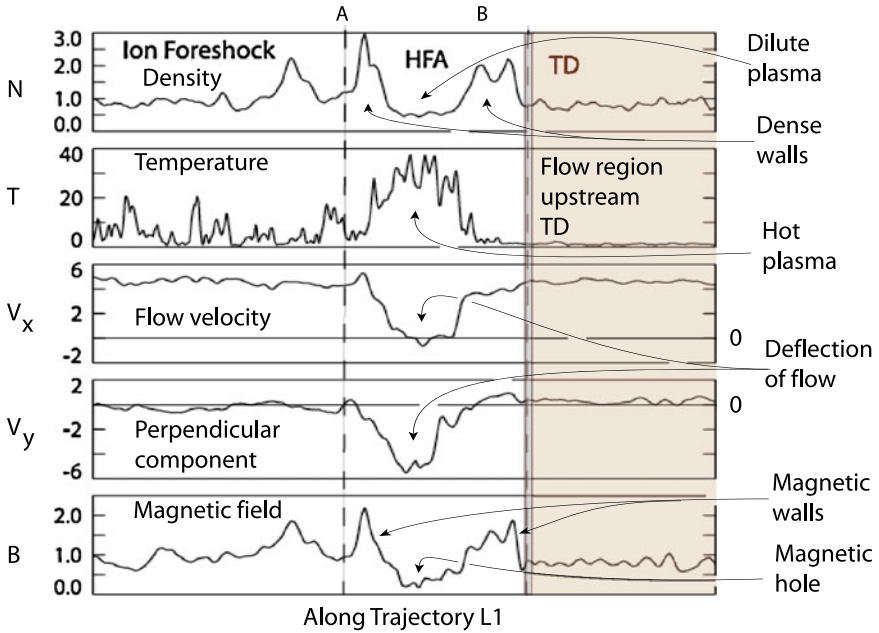


Figure 6.42: Variation of the plasma parameters and fields along the trajectory L1 of a fictive spacecraft passing the HFA in Figure 6.41 [data taken from *Omidi & Sibeck, 2007*, courtesy American Geophysical Union]. The simulated HFA shows all the signs of an observed HFA.

and global simulations, the global simulations differ so much from the local simulations that it is very difficult to compare them on the same safe grounds.

In order to compare the simulation results with observation *Omidi & Sibeck [2007]* let a fictive spacecraft fly along the line L1 in their simulations. Figure 6.42 shows the changes in the plasma and field parameters along L1. On the left in this figure is the ion foreshock region (up to point A, the TD being at point B) with the usual fluctuations in the plasma and field quantities that are typical for the foreshock. The fictive spacecraft encounters the HFA about at point A, detecting a dense wall and strong magnetic field that separate the hot HFA interior from the environment. This interior has the property that the flow velocity V_x has dropped to zero while the transverse flow velocity V_y has taken over the entire moment of the flow. Note that outside the HFA this component was zero. The magnetic field inside the HFA is weak, either corresponding to a magnetic hole or a current sheet. Leaving the HFA close to the tangential discontinuity TD, density and field are again enhanced, and the flow returns to its original direction. All this happens only on the front side of the TD. After passage across the TD, the fictive spacecraft finds itself located in the upstream flow domain well outside the foreshock without any signature of an HFA. The plasma is dilute, cold and quiet and the magnetic field is weak, here.

Most of these signatures agree quite well with the observations of Hot Flow Anomalies. As we have noted, the discrepancy is mainly the boundary property of the TD between the

foreshock and the quiet upstream flow. Moreover, the HFA appears only on the front side of the HFA lying entirely inside the ion foreshock on the quasi-parallel shock. Whether this is an effect of the asymmetry of the normal electric field at the tangential discontinuity is an open question that cannot be decided at the present time but needs more and specialised simulations. Concerning the tangential discontinuity we note that the ‘global’ simulations seem to indicate that on the perpendicular side of the shock the tangential discontinuity simply passes the shock and survives in a modified form behind the shock in the downstream flow. This seems to be the case in the upper part of the left panel in Figure 6.41 which advocates such a conclusion.

6.4.3 “Solitary Shock”

Omidi & Sibeck [2007], in their ‘global’ hybrid simulations, discovered a new type of a thin shock, which they call a ‘solitary shock’ (because of its narrowness). The signature of this solitary shock is seen in Figure 6.41 as the sharp boundary between downstream and upstream directed flows behind the tangential discontinuity.

Omidi & Sibeck [2007] investigated this new type of shock more closely, finding that it is completely unrelated to the presence (or absence) of the tangential discontinuity being a proper phenomenon of formation of a bow shock in front of a blunt obstacle in a magnetised flow (note that the ‘global’ simulations are not completely collisionless, even though in investigating the ‘solitary shock’ *Omidi & Sibeck [2007]* have tried to determine the dependence of its formation on the resistive scale!). Simulations without a tangential discontinuity also show the formation of a ‘solitary shock’.

The condition under that it is produced seems to be related to the direction of the motional upstream electric shock field only. If this has a component that is directed upstream (in the shock frame *away* from the shock), then ‘solitary shocks’ seem to form spontaneously. This motional electric field direction determines the drift and acceleration of the shock-reflected ions. Nonetheless, this observation does not yet lead to a complete understanding of the very mechanism that leads to the formation of this new kind of a shock that under the condition of the upstream directed motional electric field component might evolve and is attached to a bow shock.

Figure 6.43 gives an example of the ‘solitary shock’ taken from the former simulations (still including the TD) plus a variation of the plasma and field quantities along the cross section L in the figure. The ‘solitary shock’ appears as a narrow line in the density which bounds the new sheath region sharply, separating it from the upstream flow. The plasma and field data on the right show the sharpness of the ‘solitary shock’ (grey shading), the spike in the electric field related to it, and the fluctuating new sheath formed behind it all being separated from the old sheath.

Simulations performed for a perpendicular shock without the presence of a tangential discontinuity but otherwise under exactly the same conditions with just different directions of the motional upstream electric field are compared in Figure 6.44. The ‘solitary shock’ on the left behaves very similar to the perpendicular shock on the right. The two exceptions are the initial assumptions that the direction of the electric field in the ‘solitary shock’ is opposite to the direction of the electric field in the perpendicular shock. Related to this is

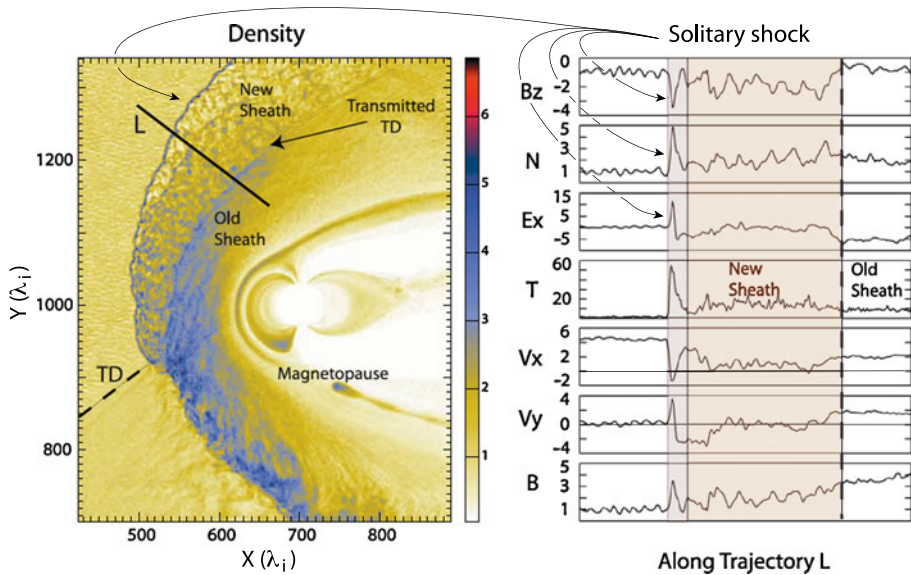


Figure 6.43: ‘Global hybrid’ simulation of the formation of a ‘solitary shock’ at the bow shock [data taken from *Omidi & Sibeck, 2007*, courtesy American Geophysical Union]. *Left:* The solitary shock as a replacement of the bow shock in the quasi-perpendicular bow shock region appears as a narrow sharp boundary with a new downstream region behind it (labelled ‘New sheath’). The ‘Old sheath’ is seen as a dense region. *Right:* Plasma data taken along the fictive spacecraft trajectory indicated in the *left part* of the figure by L. The solitary shock is shaded, while the new sheath region is light coloured.

the opposite direction of the flow deflection in V_y . In addition, the solitary shock transition is substantially narrower than the transition region of the perpendicular shock.

We do not intend to discuss these observations/simulations in further detail as neither additional simulations or investigations in depth are currently available, nor is a comparison at hand to observations in real shocks like Earth’s bow shock. Moreover, local simulations have demonstrated the importance of electron dynamics for shock formation and reformation. This dynamics is completely ignored in the global hybrid simulations. It is therefore not certain whether this finding will survive the experimental test in space or full particle PIC simulations with large mass ratios. Such ‘solitary shocks’ should, if they exist, be found at the bow shock, because the bow shock is a three-dimensional structure. Hence, at some location on its surface the motional electric field should have the correct direction to generate a ‘solitary shock’. So far no indication of such a phenomenon has been seen, a lacking result which might be due to having escaped recognition.

6.5 The Downstream Region

All shocks possess a downstream region which is located between the obstacle and the shock or the shock and the slow stream which has been overturned by a fast stream. This

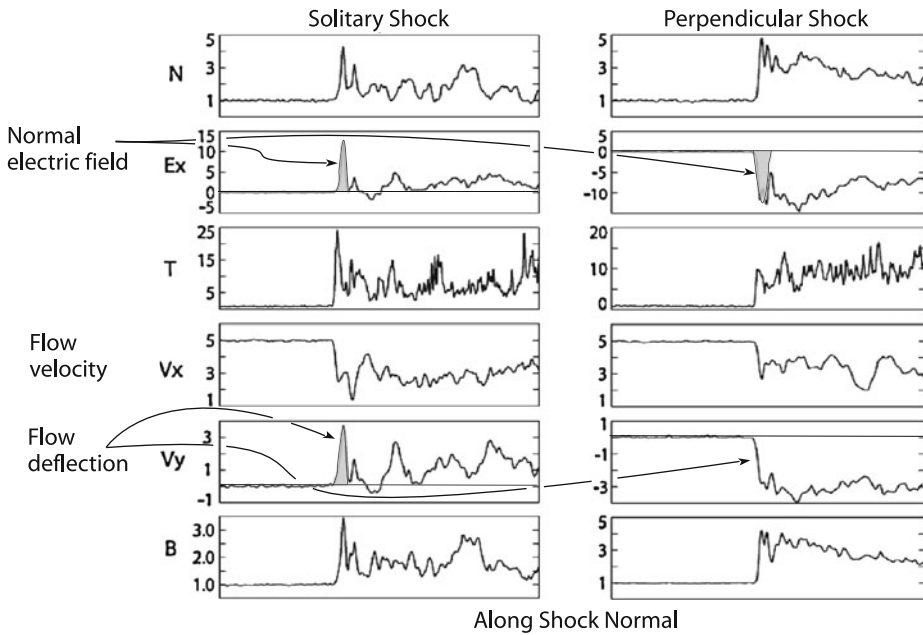


Figure 6.44: Variation of the plasma parameters and fields along a direction normal to the ‘solitary shock’ and a model ‘perpendicular shock’ [from *Omidi & Sibeck, 2007*, courtesy American Geophysical Union]. Indicated are the main differences: the motional electric field and the deflection of the flow which are both in opposite directions.

downstream region belongs to the shock transition with its properties being determined by and large by the shock. The downstream region has two boundaries, the shock and the (blunt) obstacle in the wide sense of the meaning, being a solid body, an unmagnetised atmospheric gas, a magnetosphere, a magnetic piston, the driver of a blast wave, the volume of a Coronal Mass Ejection (CME), or simply the slow stream flow. In each case the properties of the downstream region will be different because of the differences in the downstream boundary conditions at the obstacle. Therefore, the properties of the downstream region cannot be considered without including the obstacle. Nevertheless, a few general conclusions can be noted as far as they concern the effect of the shock boundary on the downstream region.

It is clear that in the average downstream of the shock, whether quasi-perpendicular or quasi-parallel, the flow is decelerated to Mach number $\mathcal{M} \lesssim 1$, deflected to flow around the obstacle (independent on whether it can penetrate the obstacle in some cases up to a certain percentage like in the presence of a magnetosphere and diffusive processes or reconnection), the magnetic field direction changes, the plasma and magnetic fields are compressed, and the temperature and pressures are increased. Moreover, the pressure and temperature anisotropies which were very moderate upstream of the shock change across the shock because of some adiabatic heating in the shock and other effects. The precise

processes and the degree to what extent this heating, change in anisotropy and changes in the parameters are predicted for the different kinds of shocks, shock normal angles, and Mach numbers are not precisely known, but the general trend is well described by the above list of changes. Thus, when asking for the effects of the presence of the shock on the downstream region we are less interested in these average quantities than in the dynamic behaviour of the downstream region as a function of the nature of the shock.

The dynamic behaviour meant is the turbulent behaviour of the plasma downstream of the shock. What are the turbulent properties? Can turbulence develop downstream of a shock? Is the turbulence, if it exists, a consequence of the presence of the shock, i.e. is it generated by the shock, or does it evolve locally in the downstream region, and the shock provides just the background properties for it such as, for instance, plasma pressure anisotropies? Are there differences in the properties of the turbulence between quasi-parallel and quasi-perpendicular shocks? Of what modes is this turbulence made of? What is the shape of the possible turbulent spectra?

These are some of the main questions that arise. Few of them can currently be answered less, because there would be insufficient observational material but, because turbulence theory is still incomplete and has not been developed for regions of such a limited transverse extent like the transition region between a shock and the obstacle in a high Mach number supercritical collisionless plasma flow. As usual, the best investigated example is the Earth's magnetosheath plasma. But even here no consensus has been reached so far about the state of the turbulence and whether the state of the magnetosheath plasma can at all be called turbulent or not. The evolution of turbulence requires time. In a streaming plasma like the magnetosheath, which is only a few Earth radii in diameter, there is barely sufficient time available for the turbulence to develop up to a stationary state. In particular, waves of large wavelength where the energy input is expected to occur at the shock, might not have enough time to cascade down to form a turbulent spectrum. Hence any spectrum will be cut off at wave numbers corresponding to the scale of the inverse diameters of the magnetosheath. Another reason is that the downstream flow is highly inhomogeneous on all scales larger than a few λ_i . This implies that the theory of homogeneous turbulence does not apply, and it is not known very well what to make out of it, because at the large scales the effect of the presence of the downstream obstacle boundary cannot be neglected anymore, which inhibits to make general conclusions about the turbulent state of the downstream plasma.

6.5.1 Sources of Downstream Fluctuations

Let us – for the moment – assume that the downstream plasma is indeed in a turbulent state, an assumption we may correct later. What are the sources of this turbulence?

Upstream Waves

The first source that may contribute to the downstream fluctuations and the possible evolution of downstream turbulence are the upstream waves. If the upstream waves vanquish the shock barrier, they can enter the downstream region, interact with the downstream plasma

population until either being damped or amplified; they can decay into other waves which the downstream plasma allows to propagate and, if of sufficiently large amplitude, they can cascade down into a broad spectrum of turbulence. As we have seen, there are several types of upstream waves: periodic ultra-low frequency waves, shocklets, huge pulsations, several types of whistlers from upstream propagating Alfvén ion-cyclotron waves to phase-locked whistlers standing in front of the shock and its elements, the pulsations; there are waves which are generated in the ramp, high-frequency electron waves, BGK modes, and the waves in the feet of the quasi-perpendicular shock and, since the pulsations have turned out to behave very similar to quasi-perpendicular shocks, the waves in the feet of pulsations, viz. Buneman and modified-two stream modes, depending on the conditions and plasma composition. Of course, in the quasi-parallel shock upstream region the plasma conditions are modified by the presence of the diffuse ion component, which will change the conditions in the foot region of the pulsations. These questions have not yet been attacked properly, at least not to the extent as they have been investigated for quasi-perpendicular shocks [as was done in the papers by *Matsukiyo & Scholer, 2003, 2006*].

ULF Waves. Quasi-parallel shocks possess a large spectrum of upstream waves that propagates mostly in the upstream direction away from the shock, when seen from the plasma frame but is convected into the shock by the flow. These ultra-low frequency waves propagate in the fast, Alfvén-ion cyclotron, and whistler modes.

In the high-Mach number upstream flow they have no chance of propagating far ahead of the shock. However, arriving at the shock or passing it, the downstream convection becomes sub-magnetosonic and the wave speed has a chance to compensate for the convection. In this case the passing waves will accumulate in the shock transition just behind the ramp. However, the ramp is nonstationary. It reforms quasi-periodically and jumps ahead upstream. In a real system like that of the bow shock where the shock is found at about the same position this reformation implies that the shock ramp oscillates back and forth around its nominal position. It consists of an accumulation of pulsations, and when a new pulsation takes over to become the shock, the old pulsation is expelled downstream with all the upstream waves that have accumulated in it and are attached to it.

Reformation of a quasi-parallel shock thus implies that quasi-periodically a pulsation and a bunch of upstream waves is added from the shock to the downstream medium where they start their own shock-independent life. They then contribute to the downstream wave population that is in some kind of turbulent state and to wave dynamics, may participate in the resonant or nonresonant interaction with the particle populations or with other waves, or they contribute to the turbulence via the turbulent cascade. This is seen in the simulations of Figure 6.34 where the region downstream of the quasi-parallel shock ramp contains a few old pulsations which fake deep magnetic holes respectively large magnetic peaks which can easily be mistaken for mirror modes but are nothing else but the old shock ramps that have been ejected to downstream after arrival of new pulsations. It can also be inferred from the AMPTE IRM wave observations downstream of the quasi-parallel shock (on the left in Figure 6.31) where the broadband emissions indicate the presence of spatially well-separated shock-ejected downstream pulsations.

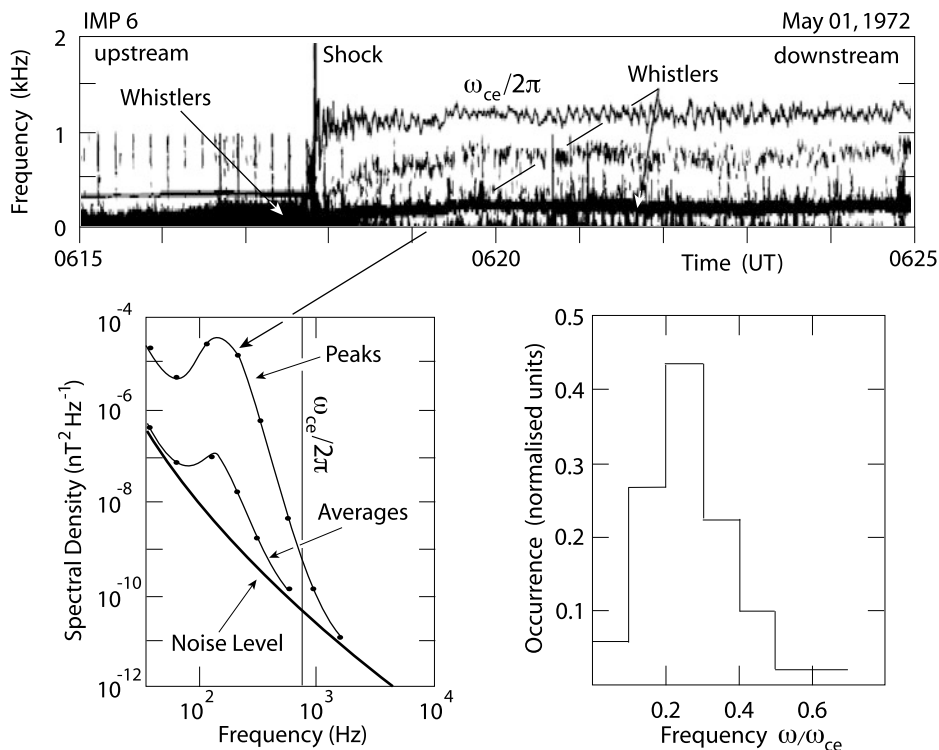


Figure 6.45: Connection between upstream (foot) whistlers and downstream (transmitted/excited) long duration whistlers behind the (quasi-perpendicular) shock [data taken from *Rodriguez, 1985*, courtesy American Geophysical Union]. The whistlers form a well distinct band in the downstream region. *Top panel:* An early dynamic spectrum observed by the IMP 6 spacecraft in 1972. The shock is quasi-perpendicular in this case. The *vertical lines* are time marks. The *horizontal line* upstream at 300 Hz is the electron gyro-frequency, downstream of the shock it fluctuates strongly around 1 kHz. Waves at the local whistler frequency appear in the shock foot. Some of these waves are whistlers. Behind the shock only the higher frequency part of this whistler band survives. *Bottom left:* The downstream magnetic spectrum showing the spectral power density above instrumental noise level. The *two curves* correspond to few second averages, the *upper curve* shows the 30 ms peak values within the average measuring time. Both curves show the presence of whistler waves at a fraction of the electron cyclotron frequency. Peak values are up to three orders of magnitude higher indicating the high variability of the downstream waves. *Bottom right:* The occurrence frequency of whistler versus frequency. The peak is at about 25% of the electron cyclotron frequency.

Whistlers. The higher frequency branch of the upstream waves propagates in the whistler mode. An example of an observation of (what is believed to be) whistler waves across a (quasi-perpendicular) shock is shown in Figure 6.45. The upstream whistlers in the foot of the shock form a broad intense fluctuation band well below the electron cyclotron frequency (the about stationary horizontal line upstream of the shock ramp). There is a continuous connection between this whistler band across the shock to downstream, now with the low frequency part of it cut off. Only the higher frequencies, which are now found far below the electron cyclotron frequency, survive in the downstream region. It is, however,

not known whether these waves have propagated from upstream across the shock (here the bow shock) to reach the downstream region or whether they are excited about locally in the downstream region (here the magnetosheath). The continuity of the whistler band from upstream to downstream suggests the former: probably most of the downstream spectrum leaks in from upstream. If this is the case, then the lower cut-off of the spectrum is due to the propagation direction of the whistlers.

The upstream whistler waves are attached to the leading edge of the shock, respectively the pulsation – in the case of a quasi-parallel shock. We already saw that the arrival of a new pulsation in the quasi-parallel case damps these whistlers out, in the simulations apparently killing them completely. The ejected ‘old shock front’ that has been the ‘old pulsation’ shifts to downstream, but the whistlers which had been attached to it, have gone in the simulation. Hence, any whistlers which occur in the downstream region should have been produced in a different way, either by the downstream particle population, by wave-wave interaction, or in some way by the cascading of the large amplitude pulsation into smaller amplitude waves which build up a turbulent spectrum. Neither of these mechanisms has so far been explored. On the other hand, if the damping of the whistlers is not complete – which is highly probable in reality –, the pulsation which is ejected to downstream will carry the surviving whistlers along with it and add them to the downstream wave population. In any case the continuity of the whistler band across the shock in Figure 6.45 suggests that the downstream whistlers are not independent of the upstream whistlers, at least at the quasi-perpendicular shock. If this is the case, it should not be much different in the quasi-parallel case in particular that everything in the simulations points on the quasi-parallel shock behaving quasi-perpendicular the closer to the shock ramp.

Local generation of whistlers downstream of the shock requires large electron temperature anisotropies $T_{e\perp}/T_{e\parallel} > 1$ and electron resonant energies $\mathcal{E}_{e,\text{res}} > \frac{1}{2}m_e V_{Ae}^2$ from linear theory, where $V_{Ae} = B/\sqrt{\mu_0 N m_e}$ is the electron Alfvén velocity, here taken in the downstream region. Figure 6.47 below indeed indicates an anisotropy in the electron temperature downstream of the shock. It also indicates the presence of counter streaming electron beams in the downstream region. Both, the anisotropy and the beams, might contribute to the excitation of whistlers locally. Whether this is the case has not yet been checked numerically.

High Frequency Waves. Plasma waves like Langmuir waves, Buneman modes and ion-acoustic waves have a rather different behaviour at the shock. The sharp cut-off of Langmuir waves at the local plasma frequency prevents them from passing across the steep density increase at the shock ramp. Only such Langmuir waves can be convected downstream which are generated at the top of the density overshoot in the shock. Langmuir waves from the electron foreshock when being ultimately convected into the shock ramp, assuming that they have not been damped by the diffuse electron population in the foreshock, will stop long before arriving at the shock and will being cut off.

Ion acoustic waves behave differently. As long as they approximately satisfy the linear dispersion relation the increase of the density when approaching the shock and in the shock ramp means that the upstream ion acoustic waves move down on the ion acoustic dispersion curve, i.e. their wave number decreases and they become long wavelength

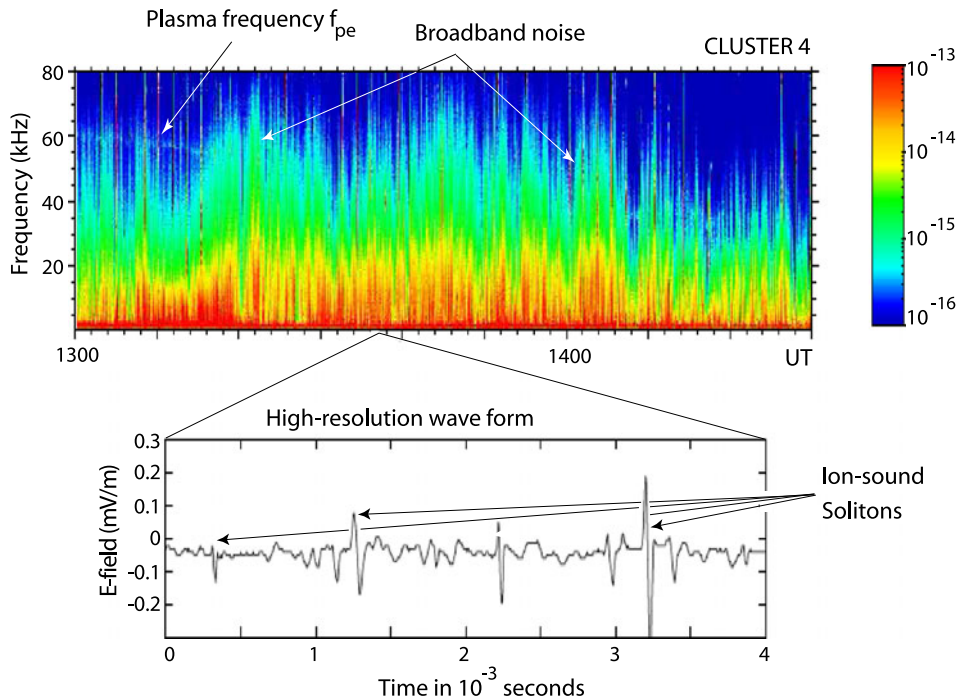


Figure 6.46: Downstream magnetosheath high frequency waves observed by the CLUSTER spacecraft [data taken from *Pickett et al., 2004*]. The dynamic spectrum shows indication of the plasma frequency and broadband noise of very spiky nature, similar to that observed earlier by *Rodriguez & Gurnett [1975]* with the ISEE spacecraft. The high-time resolution wave electric field given in the *bottom panel* shows that the latter noise is (as expected) produced from the occurrence of many bipolar electric field structures which indicate the presence of ion acoustic solitons in the downstream region.

waves which can pass into the downstream region where ion Landau damping in the warm downstream ion population will absorb their energy and contribute to heating the downstream ions stronger than electrons. It is, however, not clear what happens to the spiky ion acoustic waves and ion BGK-modes or holes which have been observed in the foreshock. Whether these waves survive the passage across the shock cannot be answered yet.

Figure 6.46 shows an example of high frequency electric wave spectra downstream of the bow shock as have been measured by the CLUSTER 4 spacecraft. The dynamic spectrum shows a weak indication of the plasma frequency in Langmuir waves which must be excited locally either as thermal noise or unstably by weak electron beams. The main signals are, however, the broadband short time emissions which do not show the slightest indication of a dispersive drift and in frequency go even beyond the plasma frequency in several cases. At a plasma frequency of ~ 60 kHz taken from the dynamic spectrum, the ion plasma frequency, where the ion acoustic wave branch flattens out, is between 1 and 2 kHz which corresponds to the intense red line at the bottom of the dynamic spectrum.

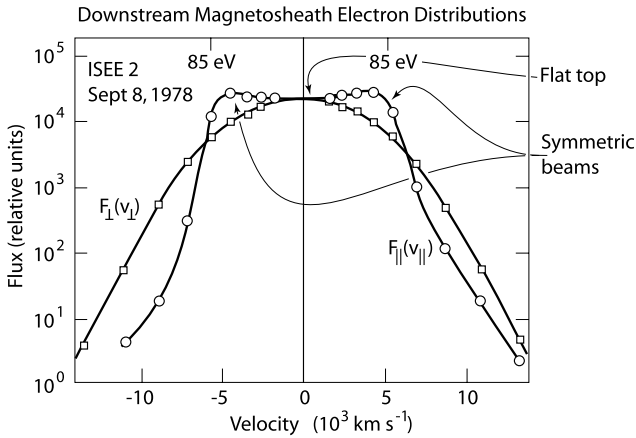


Figure 6.47: Downstream magnetosheath electron distributions as measured by ISEE 2 [after *Feldman et al*, 1983, courtesy American Geophysical Union]. The perpendicular distribution $F_{\perp}(v_{\perp})$ is a hot symmetric Maxwellian, while the parallel distribution $F_{\parallel}(v_{\parallel})$ is cooler, flat top and shows the indication of antiparallel electron beams.

Such broadband signals can only be produced by highly localised wave packets in the streaming plasma. At a streaming velocity of $\sim 100 \text{ km s}^{-1}$ this implies wavelengths of the structures of the order of a few meters. In fact the extremely high time resolution of the wave electric field given in the lower panel confirms the presence of these local wave fields. They form (sometimes non-symmetric) bipolar signals of amplitudes of some $\sim 10^{-4} \text{ V m}^{-1}$. There are ~ 1 or 2 such structures per ms. At a streaming velocity of $\sim 100 \text{ km s}^{-1}$ this implies that the structures are of transverse size of $< 50 \text{ m}$, i.e. they are Debye scale structures of the kind of solitons or BGK modes propagating on the ion-acoustic wave branch. These waves and structures may have passed the shock and entered the downstream magnetosheath region. They possess similar or even slightly larger amplitudes than the similar waves in the foreshock, so they should have been amplified in the interaction with the shock or have originated in the shock ramp. Downstream amplification implies downstream free energy in the form of electron currents. Thus their origin is rather in the shock transition than downstream. It is, however, clear that they contribute to the high frequency turbulence in the downstream region.

Shock Generated Waves

The shock barrier can permit a limited number of upstream waves to pass downstream, it can refract the waves and it can transform them into other wave modes. In the latter case the shock barrier will itself act as a wave generator becoming a source of downstream turbulence. One type of those waves that is known from the observation and simulation of quasi-parallel shocks are the large amplitude pulsations (SLAMS). Even though they are not generated in the shock, seen from the downstream region the shock is their source because they reach their largest amplitudes in the shock from where they spread downstream.

Other proposals are that the shock itself generates low frequency waves due to interface instabilities. Streaming instabilities have been proposed which may arise due to the sheared velocity tangential to the shock. Different proposals concern the generation of interface waves due to the interaction between the downstream hot ion component and the upstream ion flow overlap in the shock ramp region. Typically those waves have long wavelength parallel to the shock surface and short wavelength of the order of the shock width normal to the shock. Along the shock they are included into ripple formation. These interface modes could be swept downstream by the flow. Whether this can happen depends on their cross-shock wavelength. This is quite short such that their upstream directed phase speed probably compensates for the downstream convection speed confining them to the shock region. Only during ejection of pulsations they will together with the pulsation pass into downstream. Quasi-parallel shock reformation by pulsations is thus one kind of those oscillations downstream of the shock surface.

The shock transition is of sufficient width for short wavelength electron waves to be excited in the shock by the current flowing in the shock, and by the free energy contained in the shock heated electron distribution. Such waves will be radiated downstream from the shock. However, as we have noted above, Langmuir waves can enter only if their frequency exceeds the high downstream plasma frequency, while ion acoustic waves excited in the shock can be convected downstream.

Passing Particles

The free energy stored in the downstream particle distribution is another source for the excitation of downstream turbulence. In the first place it is the energetic ions which carry most of the free energy. They are shock-reflected either in the quasi-perpendicular shock or in parts of the quasi-parallel shock. After having been accelerated they can pass the shock ramp and move downstream where they occur as the energetic gyrating ion component which causes a high perpendicular excess in the ion pressure and thus a pressure anisotropy $P_{i\perp}/P_{i\parallel} > 1$. These ions excite propagating waves on the transverse electromagnetic ion cyclotron (EMCI) branch with frequencies below the ion cyclotron frequency $\omega < \omega_{ci}$ and phase velocities $\omega/k < V_A$. This kind of anisotropy is largest close to the shock where the gyrating ion component has not yet merged into the plasma background. Further down, the plasma flow between shock and obstacle further increases the anisotropy adiabatically. Under certain conditions zero-frequency mirror modes can then be excited which require that $P_{i\perp}/P_{i\parallel} - 1 > \beta_i^{-1}$. Both types of waves have been observed, indeed.

The competition between these two modes poses an interesting problem. Both require that the anisotropy $A = P_{i\perp}/P_{i\parallel} - 1 > 0$. This has frequently given rise to controversy about the existence or non-existence of mirror modes in regions like the magnetosheath [e.g. Gary, 1992]. Excitation of the mirror mode was attributed to the presence of heavy ions which may consume the free energy in the anisotropy in favour of the mirror mode [see Gary *et al.*, 1995]. Indeed, the heavy ions resonantly absorb the energy of the EMCI. However, the particles contributing to both instabilities are separate. The mirror mode resonates with ions of zero parallel velocity $v_{\parallel} \sim 0$ while ion cyclotron waves resonate predominantly with the high parallel energy component of the ion distribution. Thus the con-

ditions for resonance partially exclude each other. Restricting for the purpose of demonstration to the case $\beta \lesssim 1$, the ion cyclotron instability condition is $A \gtrsim \omega/(\omega_{ci} - \omega)$. The resonance condition $kv_{\parallel} = \omega - \omega_{ci}$ combined with the ion cyclotron dispersion relation $k^2 V_A^2 / \omega_{ci}^2 \simeq \omega^2 / \omega_{ci}(\omega_{ci} - \omega)$ is easy to show to yield the (minimum) parallel resonance energy $\mathcal{E}_{i\parallel} \gtrsim \mathcal{E}_B / A^2 (A + 1)$, with $\mathcal{E}_B = m_i V_A^2 / 2$.

In order to obtain a viable threshold for this mode one needs to express the resonant energy on the left of this expression through the anisotropy and plasma $\beta_{\perp}$. This can be done calculating the average resonant particle energy above $\mathcal{E}_{i\parallel}$ by multiplying both sides of the above inequality by the parallel Maxwell distribution $(N/\sqrt{\pi}) \exp(-v_{\parallel}^2/v_{\parallel\text{th}}^2)$ and integrating from $\mathcal{E}_{i\parallel}$ to ∞ . The left-hand side then yields the partial resonant (parallel) pressure, while the right-hand side just becomes proportional to the fractional resonant particle density. Dividing by $\mathcal{E}_B$ and noting that $T_{\parallel}/\mathcal{E}_B = \beta_{\parallel} = \beta_{\perp}/(A + 1)$ then yields the wanted threshold condition for the ion cyclotron instability as

$$A^2 \beta_{\perp} > 1/F(A, \beta_{\perp}) \quad (6.9)$$

where the function on the right-hand side results from integrating the Maxwellian and is given by

$$F(A, \beta_{\perp}) = 1 + \frac{2}{\sqrt{\pi}} \frac{\sqrt{\xi} \exp(-\xi)}{1 - \operatorname{erf} \sqrt{\xi}}, \quad \xi = \frac{A + 1}{\beta_{\perp}} \quad (6.10)$$

This threshold condition holds in a low β plasma. The difference between mirror and EIC waves is indeed small. For large A the exponential factor makes $F \sim 1$. Exactly along the same lines one can (numerically) obtain a threshold condition for the ion-cyclotron instability in high β plasma.

The condition for the mirror instability is simply $A\beta_{\perp} > 1$, subject to the limitations mentioned above. Under the conditions that $F \sim 1$ there is a range where both instabilities exclude each other. Still there is a large domain where both instabilities coexist and where the one with the larger growth rate will win. On the other hand, numerical simulations in three dimensions have demonstrated that the mirror mode, not depending on the presence of high energy resonant ions, gains from the increase of dimensionality as it provides a higher number of degrees of freedom and allows for coalescence of mirror bubbles to form larger bubbles. This is, however, a nonlinear effect that is not contained in the linear consideration given here. Anisotropies of the order of $A \sim 10$ are not unusual in the magnetosheath behind Earth's bow shock (cf. Chapter 10). Kinetic [Pokhotelov *et al*, 2000] effects slightly modify the marginal stability boundary and growth rates of the mirror instability. Nonlinear effects show in addition the saturation of the mirror instability by trapping of low energy ions. These ions have the effect of flattening the shape of the magnetic field inside a mirror bubble [Pokhotelov *et al*, 2008]. In addition it has been claimed that the mirror mode may as well generate peaks instead of bubbles. The cited nonlinear theory does not confirm this claim, however, but there is still controversy about this fact as one-dimensional simulations seem to confirm the evolution of peaks, and other nonlinear treatments [Kuznetsov *et al*, 2007] suggest that plasma collapse would be responsible for peaks. The situation is quite unclear as observations in the magnetosheath indeed show the

presence of peaks behind the bow shock which one might be tempted to attribute to mirror modes.

On the other hand, the observed peaks are probably not caused by an instability. Rather they are the remains of old quasi-parallel shock ramps which have been ejected downstream from the quasi-parallel shock after reformation when a new upstream pulsation arrived at the shock. The proximity of peaks in the outer magnetosheath provides a strong argument for this to be the case.

For an estimate of the mirror growth we consider magnetosheath conditions. The angular ion cyclotron frequency in the magnetosheath is of the order of $\omega_{ci} \sim 10 \text{ s}^{-1}$. With a realistic assumption on the mirror growth rate $\gamma \sim 0.01 \omega_{ci}$, flow velocity of $V \sim 100 \text{ km/s}$, and the width of the magnetosheath being of the order of $\sim 2R_E \gtrsim 10^4 \text{ km}$, a transverse convective crossing of the magnetosheath just corresponds to one e-folding time of 100 s for a magnetic mirror hole, if the hole would form right behind the bow shock. However, the anisotropy must first evolve from the bow shock along the convective path to reach into the mirror unstable domain. Behind the shock one has $\beta_{\perp} \sim 1-2$. Hence, an anisotropy of $A \sim 5-10$ is needed for the mirror mode. Anisotropies of this magnitude evolve roughly at about half the distance between the shock and the magnetopause. Observations showed that at about this distance the state of low frequency wave excitation in the magnetosheath changes about abruptly [Hill et al, 1995], indeed, from peaks to mirror dips. This apparently reduces the accessible e-folding time.

However, the flow is not crossing the magnetosheath straight from the shock to hit the magnetopause. At the contrary, it is deflected and is turned around the magnetopause. In addition, observationally, the strongest mirror activity is found close to the magnetopause at either the flanks or at higher latitudes. Thus the flow moves roughly a distance of $5-8R_E \gtrsim 6 \times 10^4 \text{ km}$ until the mirror modes have evolved to the observed amplitudes. This yields a growth time of $\tau \sim 600 \text{ s}$ or six e-folding times corresponding to a factor $e^6 \sim 400$ in the magnetic mirror amplitude. The observations show that mirror modes in their evolved state, when they are observed in the magnetosheath, have hole amplitudes $\delta B/B \simeq -(50-80)\%$ (less than) the surrounding ambient field value. The latter is of the order of $B \sim 30 \text{ nT}$, which suggests that the initial infinitesimal mirror wave amplitude at about $1R_E$ behind the bow shock was of the order of $|\delta B| \sim (0.1-0.3) \text{ nT}$, which at these low mirror frequencies $\omega \sim 0$ is in the measurement noise. This estimate is consistent with our assumptions on the growth of mirrors.

Electrons, on the other hand, do not appear as a transmitted gyrating component in the downstream plasma. They are heated and accelerated in the shock and are seen downstream as a single hot electron component of lesser temperature than the ions. Reasons for this difference in temperature have been given above. Accordingly, electrons do in the average provide a warm background component of little anisotropy. This might, of course, locally be incorrect, when current sheets are formed as the consequence of shock behaviour or by internal processes in the downstream plasma like acceleration in approaching current sheets, current dissipation processes, the internal processes acting in mirror modes, and circumstantially sometimes even reconnection in narrow current sheets. Figure 6.47 shows the typical electron distribution in the downstream region as had been measured by the ISEE spacecraft [Feldman et al, 1983]. Two cuts through this distribution are given,

the perpendicular cut is a symmetric Maxwellian $F_{e\perp}(v_{\perp})$ of slightly larger temperature than the parallel component. The parallel distribution $F_{e\parallel}(v_{\parallel})$ has a flat top and shows an indication of the presence of two nearly symmetric but counter streaming electron beams. These measurements are taken close to the shock and are of interest for what concerns the electron wave dynamics of the near shock downstream region. Such weak beams might be involved into the generation of the plasma frequency close behind the shock which is seen in Figure 6.46 as well as for the formation of the BGK modes. They are one of the sources of the high-frequency turbulence in the downstream region.

The formation of the particular form of the electron distribution shown in the figure has not yet been satisfactorily understood. Electrons are heated in the shock mostly parallel to the field while also being energised adiabatically in the perpendicular direction in the compressed magnetic fields when crossing the shock. Electron distributions of this kind can well be responsible for the local excitation of whistlers, as we have mentioned above. However, farther downstream of the shock transition the electron distribution looks different, much more isotropic with flat tops in the parallel and perpendicular directions and temperature anisotropies near zero [as seen in the CLUSTER measurements analysed by Masood *et al*, 2006]. In the following we discuss observations of downstream fluctuations causing turbulence and compare them with simulations.

6.5.2 Downstream Fluctuations and Turbulence

Let us now check whether in the downstream region turbulence can evolve and which are the dominant modes that can be observed. This can be done in two ways, either by referring to observations or by referring to simulations.

Observations

Figure 6.48 shows an example of recent CLUSTER observations of magnetic field fluctuations downstream of the Earth's bow shock wave in the deep magnetosheath [Sahraoui *et al*, 2004]. The top frame is a plot of the normalised fluctuation amplitude for a time period of 160 s and for all 4 CLUSTER spacecraft which were at separation of a few 100 km. The fluctuations measured by the different spacecraft are colour coded, and the main field has been subtracted (for our purposes it is not necessary to know which colour represents which spacecraft).

The important observation is, however, that on a time scale longer than ~ 10 s there is practically no difference in the signals between the different spacecraft. Thus, on this time scale and the corresponding separation distance which maps to the wave number, the downstream plasma is not in a turbulent state. Any ultra-low frequency waves having entered the magnetosheath or have evolved in the magnetosheath remain as such and have not cascaded down into turbulence while having evolved to nonlinearity, as is signalled by the large relative amplitudes. On the other hand, on the much shorter time scales the signals are incoherent among the spacecraft, and the high frequency fluctuations might have entered some turbulent state.

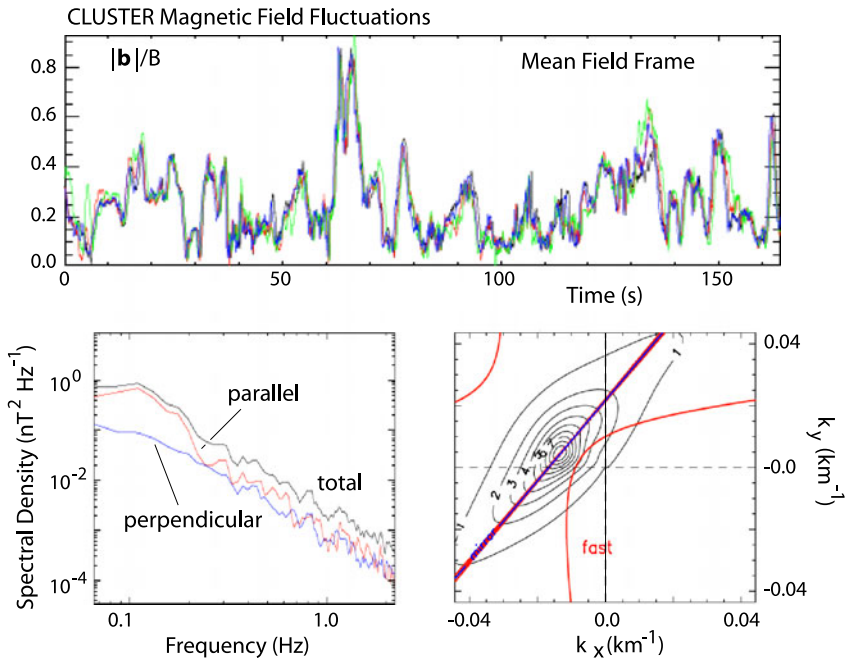


Figure 6.48: CLUSTER data of magnetic fluctuations downstream of the bow shock [data taken from *Sahraoui et al., 2004*]. The *top panel* shows the normalised fluctuation modulus in the mean magnetic field frame for the four spacecraft which are colour coded. Note that on the separation of the spacecraft there is no difference in the fluctuations for any times longer than ~ 10 s. Hence, on these time scales the fluctuations are practically coherent, and there is no turbulence neither on this time scale nor on the corresponding spatial scale. The *lower panels* show the power spectrum of fluctuations (*bottom left*) exhibiting some incoherent power law shape that might indicate the presence of turbulence. Based on the spacecraft separation the spectral density in the wave number plane (k_x, k_y) has been determined (*bottom right*). The power peaks at long wavelengths corresponding to ‘mirror modes’ as indicated by the superimposed projection of the mirror and fast dispersion curves.

This is reflected by the power spectrum in the lower left panel which evolves into a power law tail at frequencies higher than ~ 0.2 Hz. Whether or not this tail signals real turbulence remains to be answered. The time domain power read from this figure is $\sim -2.3 \pm 0.2$. If real, it might indicate fractal-dimensional intermittent turbulence, however. The panel on the right bottom in the figure is a reconstruction of the k -space for ultra-low frequency waves, the only ones for which it can be done on the basis of the spacecraft separations. The spectral power peaks at very small values of k_x, k_y typical for long wavelength, but the structure does not lie on the fast or Alfvén modes, they fit better to the mirror mode. Note also that the shape of the wave region in k -space is elongated along the mirror dispersion curve.

The method of determination of these curves depends strongly on the spacecraft separation and the related resolution and thus remains unaffected by short wavelength turbulence. *Walker et al [2004]* tried to identify the contribution of modes to higher frequency turbulence. Choosing a frequency of 0.61 Hz they found that the spectrum at this frequency

was composed of the superposition of an Alfvén wave, a weak contribution of a slow mode and some remainder of the mirror mode. It thus seems that the zero frequency mirror mode slowly decays into other wave modes.

Narita & Glassmeier [2005] analysed the plasma rest frame dispersion relation of downstream ultra-low frequency waves in dependence on the downstream distance from the shock (using CLUSTER observations across the bow shock when the spacecraft were on short separation of few 100 km). They confirmed the finding that close to the shock the waves seemed to be a mixture of Alfvén ion cyclotron and mirror modes. Farther away from the shock closer to the obstacle (the magnetosphere) the mirror mode dominated over the Alfvén ion cyclotron mode. *Narita et al [2006]* extended this analysis to many observed events for the regions downstream of the quasi-parallel as well as quasi-perpendicular bow shock. They found a downshift in the frequency ω/ω_{ci} normalised to the local cyclotron frequency and a decrease in $k\lambda_i$ from upstream to downstream in both cases (see Figure 6.50). Generally, the downstream region of a quasi-parallel shock is more disturbed than the downstream region of a quasi-perpendicular shock, i.e. the number of fluctuation events is larger. Moreover, from upstream to downstream the direction of wave propagation changes from along the background (main) magnetic field to perpendicular to the magnetic field, upstream being in the fast-Alfvén ion cyclotron mode, downstream showing properties of the slow or mirror modes.

From this it might be concluded that the upstream waves (mainly in the case of a quasi-parallel shock) when transmitted downstream [*Narita et al, 2006*] change direction which is due to a change in wave number. The wave vector direction downstream seems to support this conclusion and contradicts the two-dimensional hybrid simulations of quasi-parallel shocks [*Krauss-Varban & Omidi, 1993*] which are shown in Figure 6.49. These simulations suggest that waves can be transmitted and that some of the downstream waves propagate like the upstream waves in the Alfvén-ion cyclotron mode in upstream direction and some in the fast mode in downstream direction.

Simulations

Krauss-Varban & Omidi [1993] in their two-dimensional hybrid simulations consider a quasi-parallel shock at different Mach numbers from $\mathcal{M}_A < 2$ to $\mathcal{M}_A > 3$. As expected at the small Mach numbers upstream standing whistlers are initially attached to the shock. Later the backstreaming ion component adds long wavelength waves to them which are excited by the right-hand resonant ion-ion instability and propagate along the magnetic field. For the low Mach numbers they can escape upstream, but when the Mach number is large they are convected back to the shock and can be transmitted downstream. During transmission they convert into upstream propagating (directed along the shock normal) Alfvén-ion cyclotron modes. Because of the refraction of the magnetic field their wave vector assumes a substantial perpendicular component (Figure 6.49). These results have essentially been confirmed by two-dimensional simulations at quasi-perpendicular shocks also shown in Figure 6.49. To study the transmission, *Scholer et al [1997]* performed one-dimensional high-Mach number hybrid simulations at a nearly parallel $\Theta_{Bn} = 5^\circ$ shock. These authors wanted to infer about the role of the shock transition region in the transmis-

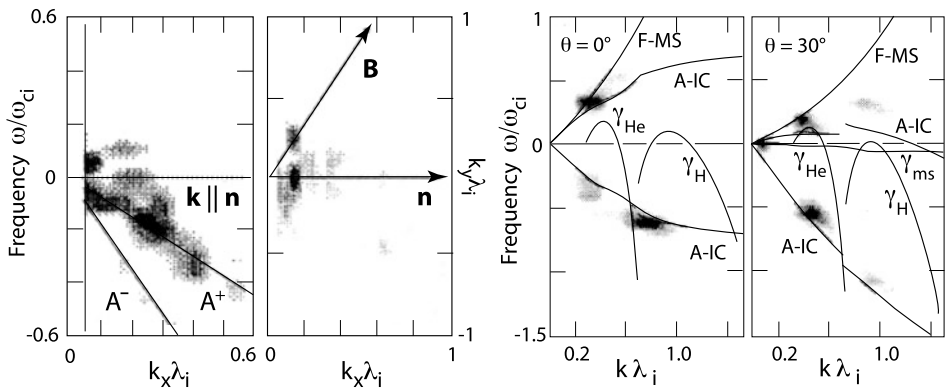


Figure 6.49: Left two panels: Dispersion relation of downstream low frequency waves in two-dimensional hybrid simulations at Mach number $\mathcal{M}_A = 3$ [data taken from Krauss-Varban & Omid, 1993]. Shown is the normalised power in the B_z fluctuations plotted versus wave number and frequency. On the left is the dispersion relation in the shock frame indicating the two polarisation branches of the waves, on the right is the wave power in the (k_x, k_y) -plane. The waves are mainly in the A^+ mode (directed along the shock normal upstream on the plasma frame) splitting into propagation parallel to the magnetic field $\mathbf{B}$ and the shock normal $\mathbf{n}$ as seen in the right panel. The waves are thus Alfvén-ion cyclotron waves into which the upstream waves had been converted when passing to downstream. Right two panels: A two-dimensional confirmation of these results for two wave propagation angles $\theta = 0^\circ, 30^\circ$ and two ion populations, protons (H^+) and Helium nuclei (He^{++}) [data taken from McKean et al, 1996, courtesy American Geophysical Union].

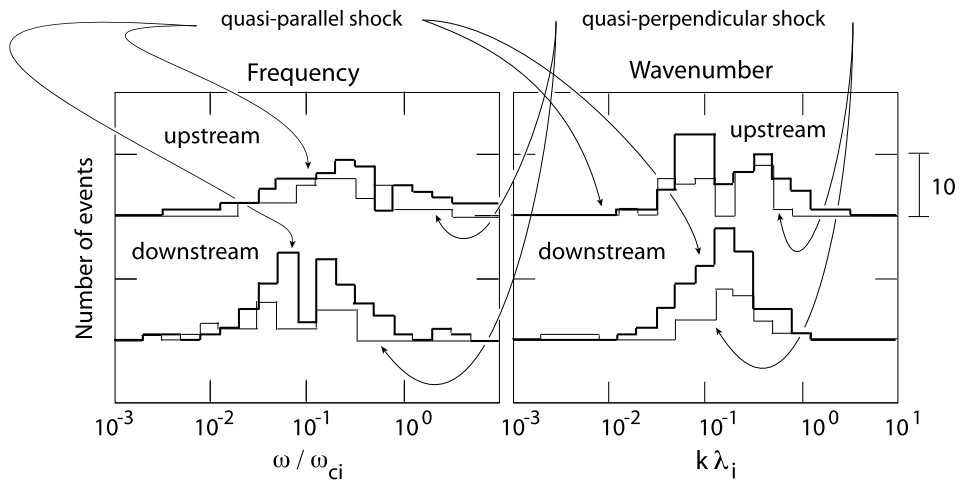


Figure 6.50: Statistical behaviour of normalised plasma rest frame frequency $\omega/\omega_{ci} = (\omega_{sc} - \mathbf{k} \cdot \mathbf{V})/\omega_{ci}$ (left panel) and wave number $k \lambda_i$ (right panel) from upstream to downstream in quasi-parallel (solid lines) and quasi-perpendicular (light lines) shocks [data taken from Narita et al, 2006, courtesy American Geophysical Union]. Note the shift in frequency to lower frequencies which is most prominent in quasi-parallel shocks suggesting that a substantial part of upstream waves is transmitted with about constant frequency. Since ω_{ci} increases, the ratio will decrease. The wave numbers increase instead.

sion process. They found that the shock transition plays a role at medium super-critical Mach numbers $\mathcal{M}_A < 8$ similar to what had been described by *Krauss-Varban & Omidi* [1993] (who did not see the formation of pulsations because of the too small number of diffuse ions in their simulations). For larger Mach numbers right hand polarised waves of both positive and negative helicity were found behind the shock ramp which could have been excited by the resonant ion-ion and the non-resonant right hand ion-ion instabilities in the very shock transition region.

Linear theory predicts such a mechanism in the overlap region of the incoming fast ion beam and the dense downstream slow ion component in the very ramp. Here the relative velocities are large between the two ion components, and the ambient downstream ion density is high, as it should be for the instability to work. Hence the shock transition behaves actively in generating waves in addition to those which are already present.

At high Mach numbers these waves have larger wavelengths and, because their phase velocities are large, they are weakly damped and can survive when propagating downstream. Downstream they may undergo an inverse cascade to generate longer wavelength waves in the course of a parametric decay instability. In the high Mach number regime $\mathcal{M}_A > 10$ no damping of these waves has been observed. This implies that for such high Mach numbers the waves will survive undamped over the entire downstream transition region between shock (for instance the bow shock) and obstacle (for instance the magnetosphere). One would expect to observe monochromatic waves all over the transition region in this case.

Figure 6.51 shows the negative helicity waves seen in this simulation at high Mach number. Clearly the waves are generated in the shock ramp as is visible from the upper left panel. The waves assume quite large amplitude and are moving upstream in the downstream frame, however much slower than the shock. They are thus flowing downstream toward the obstacle. With downstream propagation they broaden and attain longer wavelengths. The spectrum on the right shows that the wave power and wavelength increases with Mach number. At large Mach number the spectrum exhibits an exponentially decaying tail which suggests that some shorter wavelengths are generated but that the waves do not participate substantially in a cascade.

The effects observed in the latter simulations were not present in those of *Krauss-Varban & Omidi* [1993] because of their much larger shock normal angles and further restrictions on the simulations. It thus seems that quasi-parallel shocks can themselves produce waves in their transition regions. It is, however, difficult to identify these waves in real observations if waves from upstream are mixed in into the transmitted waves. This will necessarily be the case for more oblique shocks that generate large amplitude pulsations upstream which themselves become the shock ramps and during reformation expel the old pulsation-old ramp downstream where it becomes downstream fluctuations, respectively it becomes large amplitude coherent downstream waves with broad low frequency spectrum and may under conditions of a sufficiently extended downstream region evolve into turbulence.

In addition, whistlers might be attached to these pulsations and escape from them simply because the downstream Mach number is small enough for letting them go. Hence, the downstream turbulence is rather in a mixed state of some large amplitude waves than

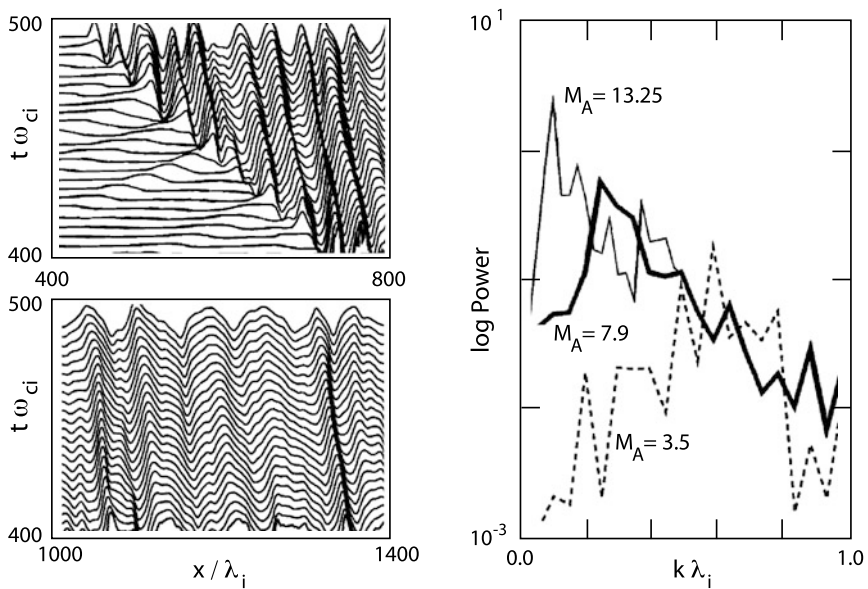


Figure 6.51: Shock generated downstream waves as inferred from one-dimensional hybrid simulations at high Mach numbers [simulation data taken from Scholer *et al*, 1997, courtesy American Geophysical Union]. *Left:* Wave profiles of the negative helicity component of the magnetic field B_y for high Mach number $\mathcal{M}_A = 13.25$ and two different regions $400 < x/\lambda_i < 800$ and $1000 < x/\lambda_i < 1400$ in the downstream frame as function of time (time is upward on the ordinate). The first domain contains the shock location. It is seen that the waves are generated in the shock, are moving in upstream direction but much slower than the shock such that they effectively are moving downstream toward the obstacle. They are large amplitude and practically undamped. The *lower panel* is taken in the downstream region. The waves have flattened and attain a longer wavelength. Some shorter waves have been produced such that the wave profile is less sinusoidal but still exhibits coherence. *Right:* Shock interface generated wave power spectra as function of wave number for three Mach numbers. The wave power increases with Mach number. At the same time the interface waves assume longer wavelengths with the spectral peak shifting to small k . Note also that for large $\mathcal{M}_A \gtrsim 5$ the spectrum has an exponentially decaying tail towards larger k . Hence there is no remarkable cascading.

in a turbulent state, with the mixed state consisting of pulsations, detached whistlers, and interface waves from the shock transition, as well as waves which are excited by the instability of the diffuse downstream ion component in the downstream plasma background. These downstream diffuse ions should as well interact with the downstream wave component, either damping them, or exciting them, or stimulating their decay in a nonlinear wave-particle interaction process. Turbulence, on the other hand, will hardly have time to evolve in the narrow transition region between the shock and the obstacle for high Mach numbers, because the transition time for the waves will in general be shorter than the decay time of the waves or the typical time to become involved into a turbulent cascade. Therefore any conclusions about real turbulence must be taken with care.

Another interesting conclusion that can be drawn from these considerations and simulations concerns the possibility of mirror modes near the shock. The simulations have all shown that the release of pulsations from the shock to downstream is accompanied by

a sequence of large amplitude magnetic fluctuations and magnetic holes with the latter containing groups of hot nearly isotropic ions which compensate for the pressure balance. These holes are similar to mirrors but are in fact generated not by the mirror instability but by the shock process of accumulation of pulsations and trapping of ions between the pulsations. They are in fact enforced structures that have not evolved through an instability but in the violent impact process of an upstream pulsation onto the shock ramp. Their observation downstream of the shock implies that the pulsations can survive for quite a long distance downstream.

We are not entering into any discussion of any processes in the downstream region farther away from the shock close to the obstacle. Such processes are strongly affected by the nature of the obstacle and have relatively little in common with the shock which for them provides an input source and a boundary condition. The latter consists for any waves generated near the obstacle in the inhibition of propagation across the shock into the region upstream of the shock. As we will see in the section on particle acceleration such a propagation is possible only for fast particles and, of course, for high frequency radiation. Fast particles are produced in almost all shocks to some degree. However, in non-relativistic shocks no efficient mechanism is known that would be capable of producing electromagnetic radiation downstream of the shock.

6.6 Summary and Conclusions

Supercritical quasi-parallel shocks behave completely different from quasi-perpendicular shocks. They are subject to a very high fluctuation level with the shock ramp often very hard to distinguish from the environment in observations. They possess extended foreshocks which are divided into electron and ion foreshocks. Electron foreshocks contain a shock-reflected and accelerated upstream electron population that is flowing in upstream direction. These electrons are believed to have their origin at the quasi-perpendicular part of a bent shock. So far it remains unclear how these electrons are generated and reflected.

The shock potential is attracting and not repulsing for electrons. Therefore, the mechanism of generating reflected electrons is more likely kind of a mirror mechanism in converging magnetic fields. Reflected electrons return into the upstream flow along the tangential field line as narrow fast electron beams. These can generate a spectrum of plasma waves via several instabilities but can survive due to the combined action of convection and escape from the self-excited wave population. Probably only part of the electron distribution is subject to quasi-linear depletion.

The ion foreshock is less extended upstream than the electron foreshock. Similar to electrons an ion beam is flowing along the ion-tangential field line and is responsible for the generation of ultra-low frequency waves. These waves fill the ion foreshock. In addition, the ion foreshock is filled with a diffuse energetic ion population which interacts with the wave spectrum. The upstream waves are instrumental in shaping the quasi-parallel shock. During their interaction with the upstream ion population they evolve into ‘shocklets’ and large amplitude ‘pulsations’, so-called ‘SLAMS’, which have a particular structure consisting of upstream leading fronts and downstream trails. In the plasma frame

they move upstream at slow speed. In the shock frame they are convected toward the shock front, interact with the diffuse ion background gradient in a resonant way, steepen, become very large amplitude and, when arriving at the shock ramp, clump together with other pulsations, accumulate, retard to zero speed and form the new shock ramp. They are responsible for quasi-linear shock reformation, which is completely different from quasi-perpendicular shock reformation.

Quasi-perpendicular shocks reform due to bunched gyrating foot ions. Quasi-parallel shocks, on the other hand, reform due to accumulation of pulsations that have formed upstream in the nonlinear resonant interaction of low frequency upstream waves with the diffuse upstream ion component and become convected downstream towards the shock. During reformation the old pulsation/ramp is expelled downstream in this accumulation process. Whistlers that had been standing in front of it are as well convected downstream and become partially damped. These and the old pulsations form the initial sources of the downstream “turbulence” behind a quasi-parallel supercritical shock.

Thus this “turbulence” consists to a large part of the overlap of the old large amplitude shock pulsations present in the downstream convective flow which, not being supported anymore as strongly by nonlinear interaction as upstream of the shock, start decaying and cascading to shorter wavelength waves. In addition, other contributions to turbulence downstream are shock-interface generated waves, leaking whistlers, and locally excited plasma waves due to anisotropy in the particle distribution, due to the plasma inhomogeneity involved into the presence of the pulsations, and due to small scale currents flowing in the rudimentary shock ramps of the residual downstream pulsations.

Two important facts concerning quasi-parallel shocks should be noted here. The first is that the generation of pulsations in the foreshock implies that the upstream magnetic field is turned about into tangential direction to the nominal shock surface. This happens on the scale of a few ion inertial lengths or less. On the large scale the quasi-parallel shock remains to behave like a quasi-parallel shock. However, on the short scale the quasi-parallel shock becomes locally quasi-perpendicular, a fact that should be very important in considering particle dynamics, in particular the reflection and acceleration of electrons. It should be noted in this context that no ion beams have been seen in front of the quasi-parallel shock neither in observations nor in the available simulations. This suggests that quasi-parallel shocks to some extent remain to behave quasi-parallel for ions, which implies that the scale on that the shock becomes quasi-perpendicular leaves not enough space for ion reflection – or is sufficiently strongly fluctuating such that ions become scattered but not really reflected to form upstream beams. Instead they become diffusively accelerated. No consensus has yet been reached on this fact.

The second fact is that the foreshock is populated not by beams but by energetic diffuse ions the density of which increases exponentially towards the shock. There is a close relation between the upstream wave growth and this density scale. It is most interesting that this diffuse component has maximum density at the shock ramp and is otherwise about homogeneously distributed both towards upstream and downstream of the shock. This implies that it is generated in the shock transition.

All these observations leave a large number of question unanswered of which the production of diffuse energetic ions is only one. How are these ions generated? How are they

injected? Why does the quasi-parallel shock (apparently) not generate ion beams, why not on the short scale? How does a quasi-parallel shock look like in two or three dimensions? It cannot behave stationary because it reforms continuously. It jumps back and forth around its nominal position that is governed by the arrival of new pulsations at the shock. How is the ejection of the old pulsation/ramp to downstream going on? We have seen that in between ion holes are formed and the magnetic field shows minima. What kind of structures are these? Are they related to mirror structures, as is naively believed, or aren't they rather violently produced ion holes instead of growing from an instability? What kind are the different particle distributions just upstream and just downstream of the shock ramp? Which waves can they drive unstable? What is the mechanism that generates the electron holes and spiky electric field structures that have been observed in quasi-parallel shocks? Is the generation of radiation at the shock ramp possible at all?

The list of these question could be extended. It is clear that their answering requires considerably more observational work, mostly on the small scales. In addition, it requires performing full particle PIC simulations with realistic mass ratio in two or even three dimensions in order to include the variability of the shock along its surface and investigating the reaction of the various particle components to the shock ramp variations.

6.7 Update – 2012

A small number of papers of relevance to this chapter have been published during the period after its preparation in 2008. Some of them, selected according to their interest for quasi-parallel shock physics, are briefly brought to the reader's attention below.

Lucek et al [2008] have used CLUSTER data at small separation in the quasi-parallel bow shock in order to investigate the role of magnetic pulsations in the shock process. Pulsations are shown to grow within only a few seconds. They also found that the shock thickness was rather narrow, in the bow shock the order of 10^3 km only while the pulsations had comparable transverse extension while being coherent along the shock over lengths the order of 10^3 to 1.5×10^3 km. Observations of this kind are quite consistent with our claim that the large pulsations are the generators of the quasi-parallel shock over some distance along the shock and that the quasi-parallel shock forms by accumulating and rejecting pulsations downstream.

CLUSTER magnetic field and energetic ion data in the foreshock have been used by *Kis et al* [2008] to determine the power in the low-frequency fluctuations in the ion-cyclotron wave spectrum caused by the energetic particle component in view of an extension of their former investigation of the dependence of the e-folding distance on the wave power spectrum. They confirm the approximate constancy of the power in the transverse fluctuations as function of distance along the magnetic field from spacecraft location to field-line foot point on the shock, whereas the power of the compressive fluctuations increases exponentially with decreasing distance. This can be interpreted as confirmation of growth of the shock-relevant compressive fluctuations (pulsations) which finally will make up the shock when arriving at the shock, as argued in this chapter when discussing quasi-parallel shock reformation. These authors also find that the fluctuations resonant with the field-aligned

shock reflected ion beam is about constant as function of distance from the shock. Thus the power of the transverse fluctuations in the foreshock having the same property is nothing else but the convected power of the beam excited transverse wave spectrum. The compressive spectrum, on the other hand, is – contrary to the view of these authors – probably not caused by the same mechanism but results from the interaction of the energetic diffuse ion component in the foreshock which is the agent of shock reformation via the growth of pulsations.

In continuation of these investigations, *Kis et al* [2010] report simultaneous CLUSTER observations of diffuse ion events upstream of the quasi-parallel Earth's bow shock in view of ion scattering in the foreshock region as function of interplanetary magnetic field and plasma conditions, finding that during quiet interplanetary times the intensity of the field-aligned shock-reflected ion beam is maximum yielding strong excitation of low-frequency foreshock waves up to large distances from the shock. These waves are convected toward the shock and cause scattering of low-energy ions.

Hietala et al [2009] investigated the downstream region of a quasi-parallel shock based on CLUSTER data finding that the flow behind the shock is highly structured containing bulk flows with high kinetic energy jets. This observation was preceded by two similar earlier observations [*Savin et al*, 2008; *Amata et al*, 2009] of such high energy jets throughout the entire downstream magnetosheath shock region. The mechanism of generation of these jets remains unclear. *Savin et al* [2008] excluded reconnection as a viable source while *Hietala et al* [2009] suggested an undefined shock interaction and formation of secondary shocks in the magnetosheath downstream as sources. Presumably there are several mechanism which can produce such intermittency in the flow and plasma properties, local and non-local. For instance reflection from the magnetopause of ion beams may cause sufficient inhomogeneity, instability and heating for modulating the downstream flow. Otherwise internal intermittency in the turbulent downstream flow similar to the generation of large amplitude waves on the ocean may cause such stream or jets. Closer to the shock presumably the most probable mechanism is that fast upstream ions of sufficient energy generate local heating and acceleration of the flow. Perforation of the shock during reformation may be another mechanism because on the small scale and for short times the flow may be considered here as highly time variable and the shock consists of many parts of highly variable different rigidity for particles and flows.

Kajdič et al [2011] using CLUSTER observations of two extended cavitons in the Earth's bow shock foreshock where they appear deep inside the foreshock being immersed into a sea of ULF waves and supra-thermal ions. In the upstream reference frame they move away from the shock. There is a pronounced difference between these types of cavitons and the well-known normal cavitons in the upstream medium, while no mechanism is known of how they have been produced and whether the shock is involved or not.

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Particle Acceleration

Abstract. Shocks are believed to be the primary accelerator of charged particles in space. This chapter reviews the basic theory and observations in near-Earth space. It discusses the Fermi and diffuse ion acceleration mechanisms including the self-consistent quasilinear theory. Some remarks on super-diffusion are included, stressing that super-diffusion takes place only at times shorter than binary collision time. Substantial space is provided for the discussion of particle injection based on observation and numerical simulation. Electron acceleration at shocks fills a long subsection. The nonlinear generation of high-energy tails on the electron distribution is discussed referring to shock particle simulations.

7.1 Introduction

Interaction of fast particles with shocks has been proposed by *Enrico Fermi* [1949] as a simple natural mechanism of pushing *charged*-particle energies up to very high values, such high that the respective particles in the Universe can appear as Cosmic Rays.

Fermi's mechanism is based on the insight that the particle energy while being a constant of the particle motion is not an invariant with respect to coordinate transformations. For instance, when sitting on a fast (classical) particle one does not feel its velocity. What is felt is just its *Zitterbewegung*, i.e. its jitter motion caused by its internal oscillations, and the acceleration and deceleration it experiences when interacting with fields that are external to it, large scale external electric or magnetic fields, and the proper fields of the surrounding particles. The latter interaction is what causes its thermal velocity $v_{e,i}$, or in terms of energy, its temperature $T_{e,i}$, taken in energy units – J or eV –, where the indices e, i refer to either electrons or ions.

However, seen from the rest frame, one adds to these energies the translational energy of the particle. Thus, a particle that passes from downstream across the shock to upstream, is reflected there, and returns to downstream, gains the velocity difference between the upstream and downstream media and thus has gained an energy that (non-relativistically) is of the order of $\Delta\mathcal{E} = 2m_{e,i}V|v_{\parallel}| \sim 2m_{e,i}V_{A1}\mathcal{M}|v_{\parallel}|$, and $\Delta|v_{\parallel}| = 2V$, where V is the velocity of the scatterer ('mirror') [Jokipii, 1966]. The second part of this expression holds only in an approximate sense as we have not included the angular dependence of the energy conserving scattering process and the obliqueness of the particle velocity. A more precise argument [originally given by Bell, 1978] can be found in Drury [1983]. A downstream particle of initial relativistic velocity $\beta \cos \theta_{vn} > V_2/c$ and pitch angle θ in the downstream region will cross the shock upstream into the region moving with speed V_1 , become scattered in the flow (thereby conserving energy), and returns downstream with pitch angle θ'_{vn} .

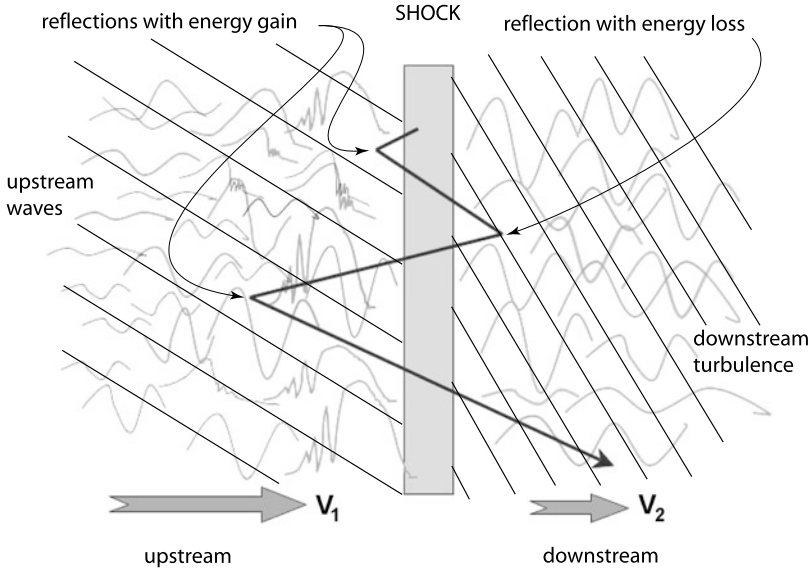


Figure 7.1: Schematic of the acceleration mechanism of a charged particle in reflection at a quasi-parallel ($\Theta_{Bn} < 45^\circ$) supercritical shock. The upstream plasma flow (left, $V_1 \gg V_2$) contains the various upstream plasma modes: upstream waves, shocklets, whistlers, pulsations. The downstream (right) is turbulent. The energetic particle that is injected at the shock to upstream is reflected in an energy gaining collision with upstream waves, moves downstream where it is reflected in an energy losing collision back upstream. It loses energy because it overtakes the slow waves, but the energy loss is small. Returning to upstream it is scattered a second time again gaining energy. Its initially high energy is successively increased until it escapes from the shock and ends up in free space as an energetic Cosmic Ray particle. The energy gain is on the expense of the upstream flow which is gradually retarded in this interaction. However, the number of energetic particles is small and the energy gain per collision is also small. So the retardation of the upstream flow is much less than the retardation it experiences in the interaction with the shock-reflected low energy particles and the excitation of upstream turbulence.

In the downstream frame its total (relativistic) energy gain $\Delta\mathcal{E} = \mathcal{E}' - \mathcal{E}$, with $\mathcal{E} = pc/\beta$ the relativistic energy, is

$$\frac{\Delta\mathcal{E}}{\mathcal{E}} = \frac{\Delta V}{c} \frac{\beta \cos \theta_{vn} - \beta' \cos \theta'_{vn}}{1 + \beta' \Delta V \cos \theta'_{vn}/c}, \quad \Delta V = V_1 - V_2 \quad (7.1)$$

which is positive because for returning particles $\pi/2 < \theta_{vn}$ and $0 < \theta'_{vn} < \pi/2$. Here ΔV is the velocity difference between the upstream and downstream scatterers (note that these, strictly speaking, usually move at velocities that are different from the respective flows!). Thus the particle energy increases monotonically for multiple shock crossings, as had been suggested by *Fermi* [1949].

Clearly, in one single reflection cycle back upstream and back downstream the particle does not pick up a large amount of energy and momentum. Thus, in order for the acceleration to be efficient, a large number of shock crossings and reflections back and forth is required. This is schematically shown in Figure 7.1. Shock particle acceleration therefore depends on the scattering process which is clearly a stochastic process, depending on the

presence of scattering centres upstream and downstream and on the probabilistic nature of the changes in the scattering angle. It also assumes that the scattering conserves energy, i.e. the scattering is assumed elastic such that the scattered particles are practically not involved into any form of dissipation of their energy of motion. In particular they should not become involved into excitation of instabilities which consume part of their motional energy. The only actual dissipation that is allowed in this process is dissipation of bulk motional energy from where the few accelerated particles extract their energy gain. On the other hand ‘dissipation’ is also attributed to direct particle loss by either convective transport or the limited size of the acceleration region. This mechanism therefore works until the gyro-radius of the accelerated particle becomes so large that it exceeds the size of the system or, otherwise, in a very extended plasma – for instance in the environment of an interstellar travelling shock –, until the energy of the particle becomes so large that the back-scattering of the particle becomes ineffective.

The dependence on the gyro-radius imposes a severe limitation on the acceleration mechanism. In order to experience a first scattering, i.e. in order to being admitted to the acceleration process, the particle must initially already possess a gyro-radius much larger than the entire width of the shock transition region. Only when this condition is given, the shock will behave like an infinitesimally thin discontinuity separating two regions of vastly different velocities such that the particle when crossing back and forth over the shock can become aware of the bulk difference in speed and take an energetic advantage of it. This restriction is indeed very rigid as it rules out any particles in the core of the upstream inflow distribution from participation in the acceleration process. In fact, in order to enter the Fermi shock-acceleration mechanism a particle must be pre-accelerated or pre-heated until its gyro-radius becomes sufficiently large. This condition poses a problem that has not yet been resolved and to which we will return several times in the present chapter.

Being aware of this difficulty, *Leverett Davis* [1956] extended and modified Fermi’s mechanism to include a medium with many diffusely distributed scattering centres in relative motion. It is clear that in this process the resulting accelerated particle distribution will have been scattered into all angles; it will be about isotropic or “diffuse”. Thus this mechanism is now known as the “diffusive” or “first order” Fermi-acceleration mechanism (leaving the original shock-acceleration “second-order” Fermi acceleration), which we will discuss in application to the shock environment. However, this mechanism suffers from the same disease as Fermi’s second-order mechanism as acceleration of thermal particles is inefficient as well. Reviews of the shock Fermi acceleration mechanisms working under different space and astrophysical conditions are found among many others in the papers and books of *Jokipii* [1971], *Drury* [1983], *Forman & Webb* [1985], *Blandford & Eichler* [1987], *Jones & Ellison* [1991], and *Schlickeiser* [2002].

The vivid interest of astrophysics in shock acceleration has in the past six decades led to an enormous activity in this field and the publication of an uncountable number of papers, which liberates us from the obligation of producing an exhaustive review of the various approaches and achievements in shock acceleration theory and application. In addition, Cosmic Rays are of much higher energy than can be generated in the heliosphere referring to the non-relativistic shocks it contains. Figure 7.2 shows the Cosmic Ray spec-

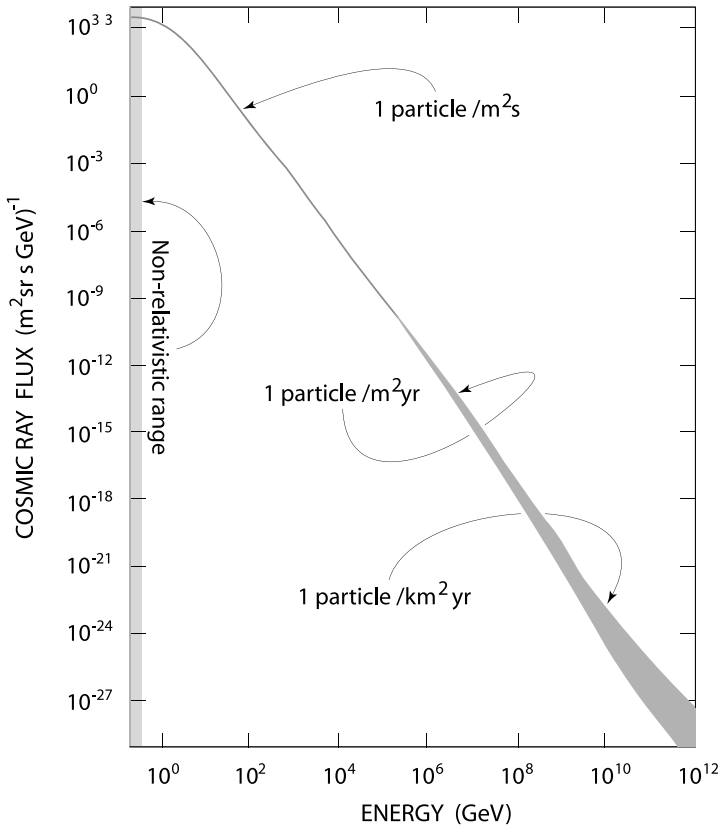


Figure 7.2: The Cosmic Ray spectrum. The nonrelativistic range is indicated in shading at the *far left* of the figure [after *Cronin et al.*, 1997]. The spectrum is approximately power law with breaks in the power at several energies. At the highest energies the count rate is very low and the shape of the spectrum is therefore quite uncertain. However, the impression is that the spectrum follows an approximate power law over ten decades.

trum as is known today from various observations. The belief in Cosmic Ray acceleration by shocks is largely fuelled by the spatial isotropy of Cosmic Rays as well as from its approximate power law shape over wide ranges of the spectrum even though the spectrum exhibits several breaks in this shape (see the figure) and becomes quite uncertain at extremely high energies. However, Cosmic Rays require highly relativistic or even ultra-relativistic shocks for being the product of shock acceleration [cf., e.g., *Waxman*, 2006]. Thus the contribution of heliospheric shock acceleration is quite naturally restricted to the range of weakly relativistic particles and to the investigation of particle acceleration by measuring energetic particle spectra *in situ* the shock environment. These measurements can then be compared with theory and in the first place numerical simulations in order to select the relevant acceleration models for medium energy particles ($\sim$ few GeV ions and $\sim$ few MeV electrons).

In addition, because of the availability – or at least the occasional availability – of collisionless shocks in space, like planetary bow shocks, travelling interplanetary shocks, corotating interaction regions, coronal shocks and the Heliospheric Termination Shock, one of the most interesting questions in shock acceleration theory can be treated. This is the above mentioned complex of questions that are related to the so-called shock particle *injection problem*: Which of the various mechanisms is capable of accelerating ions and electrons out of the main streaming thermal plasma distributions to energies high enough that they can become injected into the cycle of the shock-Fermi acceleration machine? Theory has so far been unable to ultimately answer this question. However, a number of sub-processes acting in the shock have in the past been proposed of which it is believed that some of them are indeed capable of contributing to answering this question. This problem does not directly stimulate astrophysical interest as it is believed that in the huge astrophysical objects with the available high energies sufficiently many particles will always have sufficiently high energy for initiating the Fermi process. Here another problem awakens attention even when the shocks are non-relativistic: this is the question what happens to a shock, if it is exposed to a substantial density of energetic particles, particles that have undergone Fermi acceleration and fill all the space upstream and downstream of the shock. These particles are believed to modulate the shock, transforming it into an energetic particle (or Cosmic Ray) *mediated* shock wave. We are not going to treat this problem here as in the heliosphere there is presumably only one single shock that may be subject to weak modulation by the Anomalous Cosmic Ray component that is present in the heliosphere, the Heliospheric Termination Shock, which we will briefly treat in passing in the second part of this volume.

7.2 Accelerating Ions when They Are Already Fast

When dealing with the acceleration of particles by shocks, the physics of the shock stands back and is not of large interest. The shock appears as a boundary between two independent regions of different bulk flow parameters which are filled with scattering centres for the particles as sketched in Figure 7.1 (see also the cartoon in Figure 7.3). These scattering centres are assumed to move at speeds not very different from the speeds of the bulk flows to both sides while being capable of scattering even high energy particles in angle. The scattering process is in most cases modelled by some dependence of the mean free path of the particles on momentum, or it is modelled according to some self-consistent quasilinear velocity space diffusion coefficient that takes account of the particle distribution function. Either method is not completely self-consistent since the formation of the shock is not taken into account and since the injection process remains unsolved. We will therefore only briefly go through the theory and later refer to the simulational approach to the shock acceleration problem.

In referring to multiple scattering of particles (mainly ions) one has in mind that the medium upstream and downstream of the shock is turbulent. There is still large uncertainty about this turbulent state and whether it is indeed capable of sufficient particle scattering, in particular when the particles are of high energy. Turbulence in the foreshock is stronger

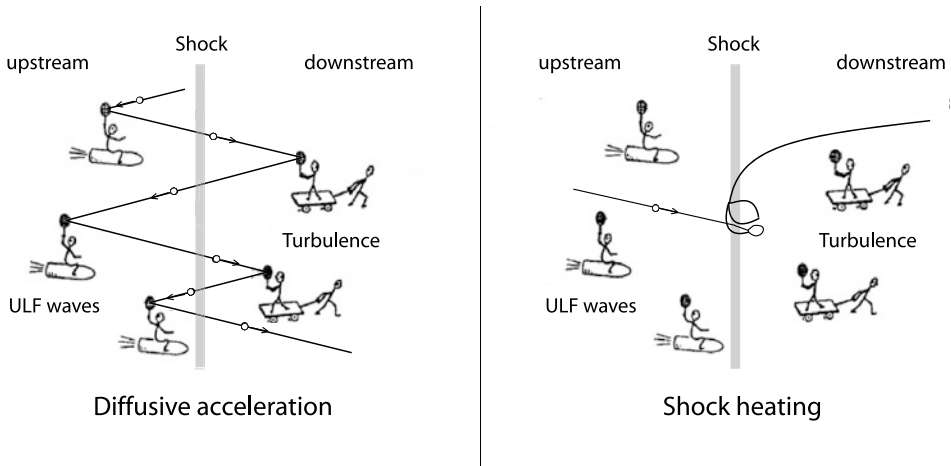


Figure 7.3: Cartoon of the diffusive shock acceleration (*left*) and shock heating mechanisms [after Hoshino, 2001, after an original sketch by M. Scholer]. In diffusive shock acceleration the particle is scattered around the shock being much faster than the shock. The requirement is the presence of upstream waves and downstream turbulence or waves. In shock heating the particle is a member of the main particle distribution, is trapped for a while at the shock and thereby thermalised and accelerated until leaving the shock.

than in the nearly quiescent upstream flow. However the foreshock is limited in extension, and the turbulence is confined to the domain close to the shock.

One may, therefore, ask whether charged particles of any energy might not couple in a *non-scattering* way to the flow. This is indeed the case in an upstream *magnetised* plasma streaming with velocity $\mathbf{V}$. After returning to upstream, all energetic particles perform a (large) gyro-circle upstream of the shock when having been back-scattered in the downstream medium. At the apex of their gyration the particles are at rest with the flow where they experience the upstream convection electric field $\mathbf{E} = -\mathbf{V} \times \mathbf{B}$ and become accelerated along $\mathbf{E}$. More precisely, the particles experience the convection field all along their upstream half-gyrocircles. This is well-known from the foot of the shock as well as from pick-up ions. In this process they get some amount of energy which *increases* their energy and gyro-radius. If they are capable to return many times from downstream to upstream they will each time accumulate some gain in energy. This process causes only gains and no losses in energy. It depends on the number of returns (downstream back-scatterings) the particle experiences.

Assuming (for generality) relativistic flows¹ with gamma-factor $\gamma^{-1} = \sqrt{1 - V^2/c^2}$, the magnitude of the electric field is $E = -cB\sqrt{\gamma^2 - 1}/\gamma$. Multiplying with the ion charge Ze , velocity $V_{\perp} = V \sin \omega_{ci}t$ and time $t = \pi/\omega_{ci}$ yields an ion energy gain $\Delta\mathcal{E}_i$ in one single upstream gyration

¹We do not treat relativistic shocks in the present treatise. Inclusion of relativistic effects changes the structure of shocks substantially as also the acceleration of particles. Relativistic flows are already high speed with velocity close to c , thus acceleration implies gain in momentum and energy, and readily leads to particle creation and other effects which do not allow to separate anymore between fields and particles.

$$\Delta\mathcal{E}_i \simeq -\frac{\pi ZeBc^2\gamma}{\omega_{ci}} \left(\frac{\gamma^2 - 1}{\gamma^2} \right) \int_0^\pi d\phi \phi \sin \phi = \pi^2 m_i c^2 \left(\gamma - \frac{1}{\gamma} \right) \rightarrow \begin{cases} \pi^2 m_i c^2 \gamma, & \gamma \gg 1 \\ \frac{1}{2} \pi^2 m_i V^2, & \gamma \rightarrow 1 \end{cases} \quad (7.2)$$

The relativistic cyclotron frequency $\omega_{ci} \rightarrow \omega_{ci}/\gamma$ has been used, gyro-radius $r_{ci} = V/\omega_{ci}$, and the shock has been assumed as infinitely extended and plane, in which case the particle performs just half a gyro-circle upstream of the shock, and the integration is with respect to time $t\omega_{ci} \equiv \phi$ in the interval $[0, \pi]$. For a bent shock of curvature radius R_s one would have to add a small correction factor $\propto r_{ci}/R_s$ (not given here). In the limit of highly relativistic flow the energy gain increases as $\Delta\mathcal{E}_i \propto \gamma$, while for $\gamma \rightarrow 1$ one recovers the non-relativistic result $\Delta\mathcal{E}_i \approx \frac{1}{2} \pi^2 m_i V^2$ that the particle picks up about four times the flow energy. It is easy to see by induction that a highly relativistic particle that starts with initial energy $m_i c^2 \gamma$ picks up a factor $(1 + \pi^2)$ in energy in each new round such that its energy after ξ returns to upstream will have risen to the quite substantial final value

$$\mathcal{E}_i^{\text{fin}} = (1 + \pi^2)^\xi m_i c^2 \gamma \approx 10^\xi m_i c^2 \gamma, \quad \xi \in \mathbb{N} \quad (7.3)$$

where it is assumed that the particle does not loose (or gain) energy in the downstream (elastic) back scattering process and during the shock passage, and ξ is a natural number, $\xi = 1, 2, \dots$. This simple assumption is subject to improving via taking into account the diffusive nature of the scattering process. This will be done below.

Thus any particle which returns from downstream to upstream is accelerated in the upstream flow, even in the absence of any upstream turbulence and upstream scatterings. The entire acceleration in this case depends on the possible number ξ of elastic downstream scatterings, i.e. the availability of scattering centres *downstream* of the shock and of a sufficiently large *extension* of the downstream region such that the particles, after having achieved a large upstream gyro-radius, still find themselves inside the downstream region and do not immediately escape into space after having passed the shock transition. The presence of upstream scattering centres increases the probability of acceleration but, for the basic process of picking up upstream energy, they are not absolutely necessary provided that the upstream medium is *magnetised* and is sufficiently extended to host the upstream gyration orbit. Magnetisation, even very weak, is *most important*, while in non-magnetised media (e.g. in the models of non-magnetic relativistic shocks used in astrophysics) upstream turbulence and upstream scattering are vitally necessary for any shock acceleration mechanism. This distinction is frequently missed in high Mach number shock acceleration models where the effect of magnetic fields is neglected thus suppressing the main cause of acceleration.

7.2.1 Diffusive Acceleration

The multiple transitions of a fast particle back and forth across the shock under the assumption that the scattering is a stochastic process suggests a diffusive approach to the acceleration mechanism. Before briefly going into detail of the diffusive treatment, we follow *Bell* [1978] in giving a simple heuristic argument for the resulting particle distribution being power law in momentum or velocity.

The Heuristic Argument

Assume that the particle that is to be scattered has already sufficiently high initial velocity $\Delta V \ll v_{\text{in}}$ or initial momentum p_{in} . In this case the relative change of the particle momentum p_j in one scattering j is $\Delta p_j/p_j = (\cos \theta_j - \cos \theta'_j)\Delta V/v_j$, which by ensemble averaging over the scattering angle in the interval $0 \leq \theta \leq \frac{\pi}{2}$ simply produces a factor $\frac{2}{\pi}$ per each cosine, yielding $\langle \Delta p_j/p_j \rangle = \frac{4}{\pi}\Delta V/v_j$ for the average change in momentum during the scattering process. (Note that averaging over the angular distribution implies complete isotropy, which is a crucial assumption in diffusive acceleration theory leaving us just with the momentum dependence!) Now assume that the particle of initial momentum p_{in} has indeed experienced a large number J of scatterings before being injected away from the shock into space. Then the final particle momentum becomes

$$\frac{p_J}{p_{\text{in}}} = \prod_{j=1}^{j=J} \left(1 + \frac{4}{\pi} \frac{\Delta V}{v_j} \right) \simeq \exp \left(\frac{4}{\pi} \sum_{j=1}^{j=J} \frac{\Delta V}{v_j} \right) \quad (7.4)$$

where the expression in brackets is understood as the expansion of the exponential for small argument. When the loss of particles from the downstream region 2 is mainly convective, then the probability of loss is the ratio of downstream flux to entering flux from upstream, $NV_2/\frac{1}{4}Nv$, and the probability for a scattering to happen before the particle is convectively transported away is $P_j = 1 - NV_2/\frac{1}{4}Nv_j$. Thus the J s probability, i.e. the total probability at the last scattering in the Fermi cycle is

$$P_{\geq J} = \prod_{j=1}^{j=J} \left(1 - \frac{4V_2}{v_j} \right) \simeq \exp \left(-4 \sum_{j=1}^{j=J} \frac{V_2}{v_j} \right) \quad (7.5)$$

Writing this in terms of the momentum by eliminating the sum in the exponential with the help of the above expression for p (dropping the large arbitrary index J), the probability of a particle to reach its final momentum is found to be

$$P_{\geq p} = \exp \left\{ -\frac{\pi V_2}{\Delta V} \ln \frac{p}{p_{\text{in}}} \right\} = \left(\frac{p}{p_{\text{in}}} \right)^{-\frac{\pi V_2}{\Delta V}} \propto N_{\geq p} \quad (7.6)$$

This expression is a power law for the probability, and the power is a constant which is given by the ratio of the convective loss probability to the change in $\ln p$ per scattering and depends just on the shock compression ratio. This result is also the number density $N_{\geq p}$ of particles with moments larger p . Since this can be expressed as the integral $N_{\geq p} = 4\pi \int_p^\infty p'^2 F(p') dp'$ over the isotropic distribution function $F(p)$, one obtains for the latter

$$F(p) = -\frac{1}{4\pi p^2} \frac{dN_{\geq p}}{dp} \simeq \frac{\pi V_2}{\Delta V} \frac{N_{\geq p}}{4\pi p_{\text{in}}^3} \left(\frac{p}{p_{\text{in}}} \right)^{-\frac{\pi V_2}{\Delta V} [1 - (1 - \frac{3}{\pi}) \frac{V_2}{V_1}]} \quad (7.7)$$

The momentum distribution is a power law $F(p) \simeq (p/p_{\text{in}})^{-\alpha}$ with $\alpha \approx \pi V_1/\Delta V$ [Bell, 1978; Drury, 1983; Forman & Webb, 1985]. (Fermi's second-order acceleration yields

essentially the same formula, however with $\alpha = 1 + \tau_{\text{acc}}/\tau_{\text{esc}}$ being the ratio of acceleration time to escape time. The relative energy gain in one collision is then $\sim |V/c|$ for the scatterer moving with speed V , while energy-increasing collisions are favoured and introduce another factor V/c . Thus the acceleration time is of second order $\tau_{\text{acc}} \sim c^2/V^2$ – from where the term “second-order” Fermi originates.)

The above power law solution of the distribution function has two deficiencies: it is valid only for large momenta $p > p_{\text{in}}$ and, thus, does not describe the low energies nor the way how particles can cross the p_{in} -boundary in momentum space in order to become injected; in addition, however, it does not describe particle losses through a distant boundary. The former problem is of principal nature and cannot be solved by a simple acceleration theory; its solution must take care of the entire shock formation process and particle dynamics, which is not included in a simple acceleration theory like the one described here. The loss problem can be accounted for by simply assuming a probability of loss. This can be done in two ways, either assuming that the particles, when crossing a distant boundary located at $x = \Delta_d$ are considered lost and do not return. This probability is simply $P_L = r_{ci}/\Delta_d$, i.e. the ratio of gyroradius to width. Or, one also accounts for an additional convective return to the shock. Then it becomes the product of the loss and convective return probabilities $P_L = (r_{ci}/\Delta_d)(r_{ci}v/V_1\Delta_d)$, which is a reduction of the escape probability [Jones & Ellison, 1991]. The effect of the loss is an exponential cut-off of the distribution which is most severe if the convective return is neglected:

$$F(p) \rightarrow F(p) \exp\left(-\frac{3p^2/2m_i}{ZeB\Delta V\Delta_d}\right) \quad (7.8)$$

which shows that the cut-off depends on the energy to charge ratio $p^2/2m_iZe$.

Fokker-Planck Diffusion Equation

Whatsoever the nature of the scatter is, the assumption of a stochastic process implies that the basic equation that governs the process is a phase space diffusion equation of the kind of a Fokker-Planck equation

$$\frac{\partial F(\mathbf{p}, \mathbf{x}, t)}{\partial t} + \mathbf{v} \cdot \nabla F(\mathbf{p}, \mathbf{x}, t) = \frac{\partial}{\partial \mathbf{p}} \cdot \mathbf{D}_{\mathbf{pp}} \cdot \frac{\partial F(\mathbf{p}, \mathbf{x}, t)}{\partial \mathbf{p}}, \quad \mathbf{D}_{\mathbf{pp}} = \frac{1}{2} \left\langle \frac{\Delta \mathbf{p} \Delta \mathbf{p}}{\Delta t} \right\rangle \quad (7.9)$$

where $\Delta \mathbf{p}$ is the variation of the particle momentum in the scattering process which happens in the time interval Δt , and the angular brackets indicate ensemble averaging. $\mathbf{D}_{\mathbf{pp}}$ is the momentum space diffusion tensor. It is customary to define $\mu = \cos \alpha$ as the cosine of the particle pitch angle α and to understand among $F(p, \mu)$ the gyro-phase averaged distribution function, which depends only on $p = |\mathbf{p}|$ and μ .

The scattering centres are of course waves of phase velocities $v_{\text{ph}} \equiv \omega/k \ll v_{\parallel} = v\mu$. In this case the diffusion tensor has only the three components $D_{\mu\mu}, D_{\mu p} = D_{p\mu}, D_{pp}$. These can all be expressed through the pitch-angle diffusion coefficient $D_{\mu\mu}$ and are given by the expressions

$$D_{\mu\mu} = \frac{\pi \omega_{ci}}{2} \frac{|k| W_k}{B^2} \Big|_{\text{res}} (1 - \mu^2), \quad \frac{D_{p\mu}}{D_{\mu\mu}} = \left(\frac{pv_{\text{ph}}}{v} \right), \quad \frac{D_{pp}}{D_{\mu\mu}} = \left(\frac{pv_{\text{ph}}}{v} \right)^2 \quad (7.10)$$

Here the spectral power density W_k of the waves that are involved into the scattering has to be taken at the resonant wave number $k_{\text{res}} = \mu v / \omega_{ci}$, and the ion cyclotron frequency is given by its relativistic expression $\omega_{ci} = ZeB/\gamma$.

One immediately realises that pitch angle scattering is the main process at small phase velocities as long as p becomes not too large, because the energy diffusion D_{pp} in scattering is quadratically small in the ratio $p v_{\text{ph}}/v$ becoming more important only at large momenta p . One may use the pitch-angle diffusion coefficient to define a pitch-angle scattering frequency as $\nu_\mu = 2D_{\mu\mu}/(1 - \mu^2) \ll \omega_{ci}$ which, of course, is much less than the cyclotron frequency. The mechanism of acceleration is, hence, a (fast) multi-pitchangle scattering process followed by slow energy diffusion, just as we have suspected from simple logic that multiple crossings of the shock (or otherwise multiple reflection from moving waves) would be required in order to push the particles to higher energy.

Wave Spectra Involved

The dependence on the wave spectral power density W_k is important. It implies that the entire diffusive acceleration process depends crucially on the assumption of which waves are involved serving as scatterers in the upstream and downstream regions of the shock. By the nature of the diffusive approximation of the Fokker-Planck equation these waves must be distributed isotropically in space, a condition that is barely satisfied in the vicinity of curved shock waves like planetary bow shocks in a high Mach number stream like the solar wind (or stellar winds). If the waves are related to the shock then it is clear that the turbulence will be inhomogeneous. The turbulent power should maximise close to the shock and decay with distance from the shock. This implies that scattering probabilities decrease with increasing distance from the shock.

A digression into our previous investigation of supercritical shocks reveals that upstream of quasi-perpendicular shocks susceptible wave power is found only in the shock foot region. The transverse size of shock feet is of the order of one (or at most a few) thermal upstream ion gyroradii. Energetic ions entering the upstream region from downstream having been scattered there into the upstream direction presumably completely ignore the foot since it belongs to the shock transition which in this kind of theory is assumed infinitesimally narrow. Both, observations and simulations have shown that farther away from the quasi-perpendicular shock the upstream flow only contains the weak upstream turbulence and possibly a few whistler precursors which have managed to escape some short distance into the upstream flow before being damped out. If the upstream turbulence is incapable of providing a high enough turbulence level, quasi-perpendicular shocks will not become effective diffusive accelerators. In the case of the solar wind, the upstream solar wind turbulence is about Kolmogorov with spectrum $W_k \propto k^{-5/3}$ but is weak [Tu & Marsch, 1995; Goldstein et al, 1995; Biskamp, 2003], in particular when the solar wind is fast, coming from the coronal holes. Higher turbulence levels are observed in slow solar wind streams when it contains a component of travelling discontinuities.

On the other hand, upstream of the supercritical quasi-parallel shock there is a substantial level of wave activity, low frequency upstream waves, Alfvén-ion cyclotron, whistler waves, shocklets and pulsations (SLAMS). These waves produce a comparably intense

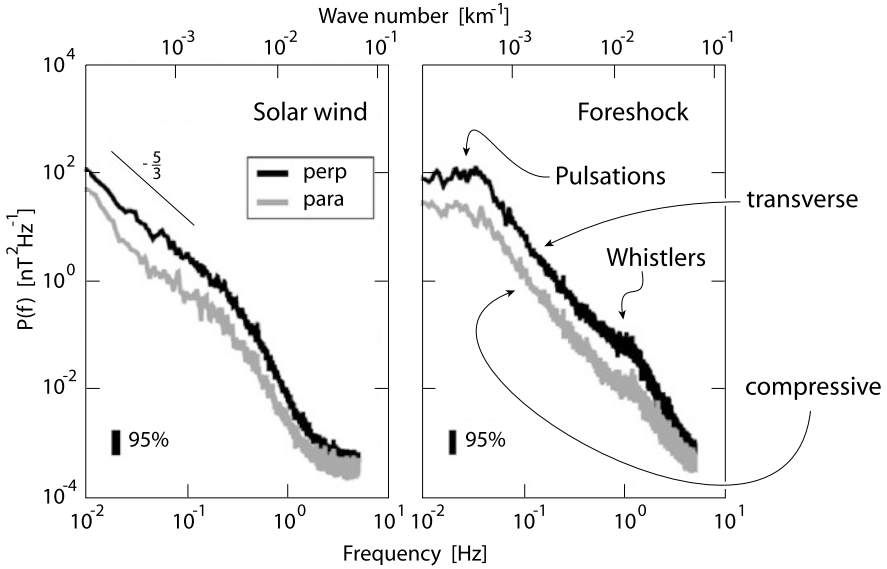


Figure 7.4: Power spectral densities for the magnetic fluctuations measured by CLUSTER in the solar wind and foreshock versus time t and wave number k . The spectra distinguish between perpendicular (transverse) and parallel (compressive) fluctuations indicating the strong anisotropy of the fluctuations both in the solar wind (*left*) and the foreshock (*right*). The perpendicular fluctuations are more intense than the parallel fluctuations. In the solar wind the spectra decay approximately like Kolmogorov for low frequencies. In the foreshock the spectra are flat at low frequencies according to the presence of shocklets and pulsations (SLAMS), exhibit a maximum at roughly ~ 0.05 Hz which is the frequency range of pulsations corresponding to transverse wavelengths of $\lesssim 10^3$ km. At higher frequencies and wave numbers the spectra are again similar to Kolmogorov though the range of power law is short and the uncertainty might hide a steeper decay [after Narita *et al*, 2008].

though anisotropic turbulence level [Narita *et al*, 2006, 2008] which, in addition, decays with increasing distance from the shock. The assumption of an isotropic and homogeneous wave distribution upstream of a quasi-parallel shock is therefore barely realistic. Theoretical spectra are not very reliable as they cannot take into account the entire complexity of the interaction between the diffuse particle component in the wave generation process as well as the wave cascading or nonlinear interactions that will take place if the power is injected in a certain wave number band and reaches high enough values. The best is to refer to *in situ* observations which again can be performed only in the vicinity of Earth and otherwise must be inferred from remote using indirect methods. Figure 7.4 shows spectra measured by CLUSTER in the solar wind and the foreshock of the Earth's bow shock wave. In both cases the anisotropy is quite pronounced with the solar wind behaving like Kolmogorov at low frequencies/large wavelengths. The foreshock spectrum shows the large contribution of pulsations as well as the contribution of whistlers. The range of power law is shorter than in the solar wind and is possibly also steeper.

The turbulence in the foreshock is sufficiently intense (see also Figure 6.15) for scattering ions into a nearly diffuse component. However, the theoretical assumption of a sin-

gle wave component (for instance of purely Alfvénic turbulence) is certainly unjustified. Neither shocklets nor pulsations are of this kind of waves; they are localised large amplitude irregular structures the contribution of which to the turbulence is largely unknown. Moreover, their spatial evolution is not accounted for in any of the acceleration theories. On the other hand, because of their large amplitudes and their effect on the magnetic field magnitude and direction they must be rather efficient in the scattering of particles.

The expectation is thus that the scattering is strong at an upstream distance close to the shock and levels out into the ordinary upstream flow scattering efficiency. But the most severe restriction is the narrowness of the downstream region which downstream of the Earth's bow shock is only one to two Earth radii thick. Once the particle reaches a gyro-radius larger than the width of the downstream region the probability for being scattered back upstream must drastically decrease. This spatial restriction must naturally impose severe limits on the backscattering of ions from downstream to upstream thereby cutting off the Fermi cycle. The gyroradius is $r_{ci} = p/ZeB$. For a proton gyroradius of $2R_E$ in a field of 30 nT the momentum is $p \simeq 10^{-19}$ Js/m yielding an energy $\epsilon_i/m_i c^2 \simeq 1.002$, which corresponds to a maximum energy of $\epsilon_{i\max} \sim (1-2)$ MeV. On the other hand the high level of wave activity downstream of the shock compensates partially for its narrowness. Spectral wave energy densities downstream of the quasi-parallel shock are at least two orders of magnitude above those in the foreshock [cf., e.g., the recent CLUSTER observations by *Narita et al*, 2008]. Since these enter the diffusion coefficients linearly the scattering probability behind the shock is much higher than from the foreshock.

Hence the question arises, what the shock is good for in acceleration if the downstream region limits the acceleration by being too narrow and the upstream region either containing weak turbulence as in the case of the quasi-perpendicular shock or containing inhomogeneous and anisotropic turbulence like in the case of quasi-parallel shocks. The answer is that shocks should serve as pre-accelerators for energetic particles. Probably they are the sources of the energetic seed particle component for further acceleration. For this they must trap ions and electrons for long enough time in the shock transition. Before discussing the related injection problem we will return to the diffusive shock acceleration, however, in order to briefly complete the picture.

7.2.2 Convective-Diffusion Equation

Since convective losses enter the above quasi-linear diffusion equation, the equation can be rewritten into a diffusion equation in real space by taking into account that convection produces a streaming density term $\mathbf{s}(p) = \int \mathbf{v} F(p, \mu) (1 - \mu^2) d\mu$ that is added to the diffusion equation. For Alfvén waves propagating in both directions along the magnetic field with Alfvén speed V_A this term can be written as $\mathbf{s} = [\mathbf{K} \cdot \nabla - \frac{1}{3} \mathbf{w} p \nabla_p] F(p, \mu)$, where $\mathbf{w} = \mathbf{V} \pm V_A \mathbf{B}/B$ is the velocity relative to the Alfvén waves which are the presumable scatterers. The new quantity $\mathbf{K}$ that appears in this contribution is the spatial diffusion tensor that describes the spatial diffusion of the energetic particles as consequence of the combined pitch angle scattering and convection. In the simple model where the waves

propagate just parallel and antiparallel to the magnetic field this diffusion tensor is given by

$$\mathbf{K} = \begin{pmatrix} \kappa_{\parallel} & 0 & 0 \\ 0 & \kappa_{\perp,1} & -\kappa_A \\ 0 & \kappa_A & \kappa_{\perp,2} \end{pmatrix}, \quad \kappa_{\parallel} \approx \frac{1}{3}\lambda v, \quad |\kappa_A| \approx \frac{1}{3}vr_{ci} \quad (7.11)$$

These expressions hold under the assumption that $v_{\mu}/\omega_{ci} \ll 1$, and then $|\kappa_A|/\kappa_{\parallel} \ll 1$ as well. Usually the perpendicular diffusion coefficient $\kappa_{\perp}$ is also neglected.

It is clear that this kind of theory is only approximative. The assumption that the scattering waves are Alfvén waves is very problematic when referring to our present knowledge on the kind of waves that are present in the vicinity of the shock both upstream and downstream. It is also problematic in view of the solar wind turbulence which is far from being purely Alfvénic. In addition, the assumption of purely parallel propagation is barely correct. Pulsations do not propagate parallel to the magnetic field. The waves in the quasi-perpendicular shock foot, which are versions of the electromagnetic modified two stream instability, propagate essentially perpendicular (i.e. highly obliquely) to the field. Jones [1990] has extended the theory for including oblique propagation while still restricting to Alfvénic type waves. The analytical expressions obtained are very involved in this case inhibiting an analytical treatment anyway. Finally the neglect of the transverse diffusion of energetic particles is questionable as well. Transverse diffusion is first order in $p v_{ph}/v$ as we have seen and can hardly be neglected, therefore. However, since all these considerations affect only the acceleration of relatively low energy particles, which are affected by the presence of the shock, the deficiencies of the theory are relatively unimportant for high energy particles in which astrophysics is interested. But when dealing with shock mediation and the injection problem they should become crucial.

In the above approximation and with the given expressions for the streaming density function $\mathbf{s}$ and the spatial diffusion tensor $\mathbf{K}$ the isotropic diffusive particle acceleration equation becomes

$$\frac{\partial F}{\partial t} + \nabla \cdot \mathbf{s} + \frac{1}{p^2} \nabla_p [p^2 J(p)] = S - L, \quad J(p) = \left[\frac{1}{3} \mathbf{V} \cdot \nabla - D_{pp} \nabla_p + \dot{p} \right] F(p) \quad (7.12)$$

This equation is still a quasi-linear equation with the quasilinear phase space streaming current density function $J(p)$. It includes the convective streaming which comes in simply because of the difference in the velocities to both sides of the shock which inhibits transformation to a comoving coordinate system. Such an equation had first been given by Parker [1965] and Jokipii [1971]. The functions on its right include possible sources S and losses L of energetic particles other than convection.

For more elaborate approximative convective diffusion equations and their numerical or analytical solutions with inclusion of various other wave modes one may consult the extended literature. However, the problem lies not in the mathematical solution of those more complicated equations. It is buried in the fact that the acceleration of Cosmic Rays to high energies can be understood only on the basis of relativistic or even ultra-relativistic shock theory. Non-relativistic shock theory on the other hand should serve to illuminate

the initial injection of energetic particles into the acceleration process. Whether this can be treated with the help of approximative diffusion equations is a question that might be doubted. One probably needs to refer to numerical simulation work in two or three spatial dimensions of self-consistent shock generation including the acceleration of particles.

One-Dimensional Case

The above convection-diffusion equation including transverse diffusion [following Jones, 1990] taken in one-dimensional real space and under stationarity (i.e. for asymptotic times $t \rightarrow \infty$) becomes

$$\frac{\partial}{\partial x} \left[V(x)F(x, p) - \kappa \frac{\partial F(x, p)}{\partial x} \right] = \frac{1}{3} \left[\frac{\partial V(x)}{\partial x} \right] \frac{\partial}{\partial p} [pF(x, p)] \quad (7.13)$$

where $\kappa = \kappa_{\parallel} \cos^2 \Theta_{Bn} + \kappa_{\perp} \sin^2 \Theta_{Bn}$ is the projection of the diffusion coefficient onto the shock normal direction. When integrating this equation with respect to the spatial coordinate from upstream (index u) infinity $x \rightarrow -\infty$ to downstream (index d) infinity $x \rightarrow \infty$ one may apply the boundary conditions $F(x \rightarrow -\infty, p) = F_u(p)$, $F(x \geq 0, p) = F_d(p)$, and $\partial_x V(x) = -\Delta V \delta(x)$ where as usual $\Delta V = V_u - V_d$. Then

$$V_d F_d(p) - V_u F_u(p) = -\frac{1}{3} \Delta V \nabla_p [pF_d(p)] \quad (7.14)$$

It is convenient to define the shock compression ratio $\rho_{sh} = V_u/V_d > 1$ with the help of which this equation becomes an inhomogeneous ordinary differential equation for the distribution $F_d(p)$ far downstream $p(\partial F_d/\partial p) + \alpha F_d = (\alpha + 2)F_u$, where $\alpha = (\rho_{sh} + 2)/(\rho_{sh} - 1)$. The solution of this equation is simply

$$F_d(p) = (\alpha + 2)p^{-\alpha} \int_{p_{in}}^p p'^{(\alpha-1)} F_u(p') dp' + C p^{-\alpha} \quad (7.15)$$

C is an arbitrary integration constant of the homogeneous equation which is put to zero in order to prevent an infrared catastrophe at $p \rightarrow 0$. The lower bound p_{in} on the integral is required already large in order to satisfy the condition $v \gg V$ of validity of the entire diffusive acceleration theory. p_{in} is the momentum of the particles which are injected into the acceleration cycle. Assuming that the injection function $F_u(p) \sim F_u(\varepsilon) \sim \delta(\varepsilon - \varepsilon_{in})$ is a δ -function, one obtains from here the differential energy flux of particles as

$$\frac{dJ(\varepsilon)}{d\varepsilon} = \frac{1}{4\pi} \frac{(\alpha + 2)N_u V_u}{[\varepsilon_{in}(\varepsilon_{in} + 2)]^{\frac{1}{2}}} \left[\frac{\varepsilon(\varepsilon + 2)}{\varepsilon_{in}(\varepsilon_{in} + 2)} \right]^{-\frac{1}{2}\alpha} \frac{\text{particles}}{\text{m}^2 \text{s ster}} \quad (7.16)$$

In this expression $\varepsilon \rightarrow \varepsilon/m_i c^2$ is the particle energy normalised to the particle rest energy $m_i c^2$, and N_u is the density of particles with $p \geq p_{in}$ far upstream. It is remarkable that in this special solution with these special boundary conditions the asymptotic distribution of the particles is independent of κ . The shock plays a role only in providing the shock compression ratio ρ_{sh} , viz. offering the particles the possibility of changing border from

downstream to upstream and vice versa many times until a finite state is reached. Accounting for the diffusion is, however, possible [cf., e.g., *Jones & Ellison, 1991*]. It changes the power in the spectrum, which now becomes

$$\alpha' = 1 + \frac{3V_u}{8\kappa} \frac{\alpha + 1}{\alpha} \Delta \quad (7.17)$$

Δ is the width of the shock transition layer, which had been assumed to be infinitesimally thin before. In the case of accounting for the diffusion it cannot be neglected anymore. However, following our previous discussion, this is still not correct simply because of the limiting extent of the downstream region which would cut off the integration in x at some $x_{\max} = \Delta_d$, where $\Delta_d \gg \Delta$ is the width of the downstream region. This can be included by adding the exponential cut-off factor derived in the former section which accounts for losses through the distant boundary.

The neglect of the perpendicular diffusion coefficient implies that the distribution has a spatial dependence only along the magnetic field. The upstream spatial dependence is roughly an exponential decay $\propto \exp(-x/\ell_{\parallel})$ away from the shock with e-folding distance $\ell_{\parallel} \approx \kappa_{\parallel u}/V_u$. This distance depends on energy, remembering that the parallel diffusion coefficient $\kappa_{\parallel} = v\lambda_{\text{mfp}}$ also depends on energy through both the velocity v and the mean free path λ_{mfp} in collisionless plasma, the latter being a complicated functional that is determined through the collisionless interaction between particles and wave spectrum. Inclusion of perpendicular diffusion, which presumably is important for acceleration at quasi-perpendicular shocks but, as numerical simulations show, plays a role also in quasi-parallel shocks is not yet contained in this theory. It violates the assumption of one-dimensionality and thus brings with it a severe complication of shock acceleration theory.

Miller's Non-selfconsistent Approach

Figure 7.5 shows a non-selfconsistent calculation of the particle spectra that are expected to be produced by a power law wave spectrum. Here protons are injected continuously (as assumed by the shock) at energy 1 MeV into the spectrum. Solution of the diffusion equation produces an exponential spectrum which evolves in time until a stationary state is reached which depends on the acceleration efficiency.

This not self-consistent theory has been developed by *Miller et al [1990]* in view of application of stochastic (second-order Fermi) acceleration to solar flare proton acceleration and electron heating. This approach uses constant escape times and a given wave spectrum and assumes that the wave spectrum decays like a power law $W_k \propto k^{-s}$. The resulting ion spectrum is an exponential, similar to those spectra obtained by *Lee [1982]* in the quasilinear self-consistent diffusive (first-order Fermi) calculations (see below). However, in this way it becomes possible to check which waves contribute most to the acceleration. The result suggests that it are indeed Alfvén waves that play the dominant part in the acceleration through the cyclotron resonance, as was assumed in all the above theories. Magnetosonic waves, which participate in acceleration through Landau resonance, turn out to be much less important.

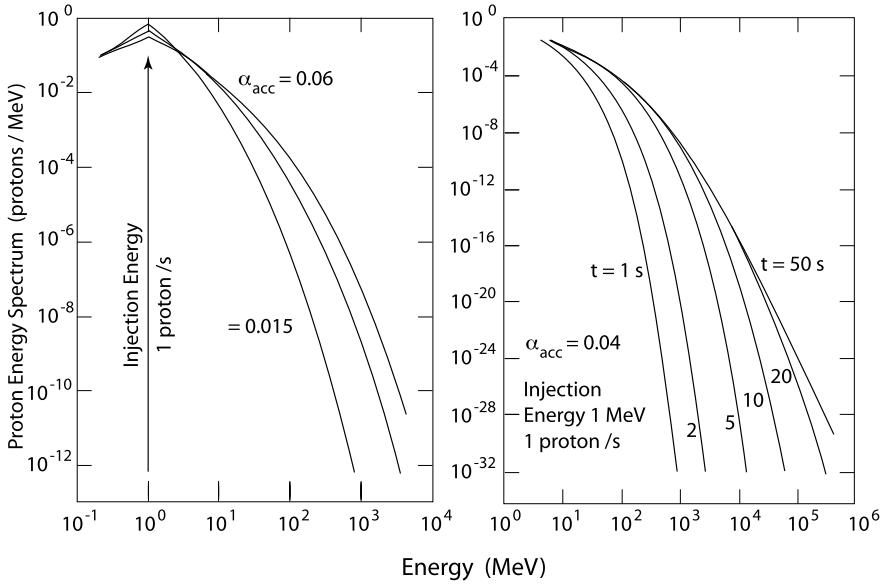


Figure 7.5: Proton energy spectra in first order (stochastic) Fermi acceleration in Alfvénic turbulence. The wave spectrum has been assumed power law $W_k \sim k^{-s}$. *Left:* Dependence on the acceleration efficiency $\alpha = (V^2/c\lambda_{\text{mfp}})\tau_{\text{esc}}$ in final state. The escape time has been assumed 1 s. Protons are continuously injected at energy 1 MeV. The spectrum is no power law but an exponential function. *Right:* Time evolution of the spectrum after continuous injection of protons for one given efficiency. The spectrum evolves into harder and harder exponential shape [after Miller *et al.*, 1990].

In spite of the unrealistic assumption of the power law wave spectrum and the non-selfconsistent treatment the Miller [1991] approach goes one step ahead of the below self-consistent Lee [1982] quasilinear calculation as it in addition to quasi-linear wave particle interaction accounts for nonlinear wave-particle interactions – which are also familiar as ‘nonlinear Landau damping’ – according to the elementary resonance condition $\omega_1 - \omega_2 = (k_1 - k_2)v_{\parallel}$, between two Alfvén waves (k_1, ω_1) and (k_2, ω_2) and low energy protons of parallel speed $v_{\parallel}$. This leads to the claim that nonlinear wave-particle interaction does indeed efficiently heat the low energy protons and thus might be capable of lifting thermal ions above the initial energy threshold level for stochastic acceleration. If this is the case, then nonlinear Landau damping could help solving the injection problem explaining how thermal particles can become hot enough in shock interaction to enter the acceleration cycle.

The reservation one might have with this theory is that, in contrast to the approach of Lee [1982], which will be discussed below, this approach is not self-consistent, assumes continuous injection of protons (or ions) at a given relatively high energy (in Figure 7.5 at the substantial value of kinetic energy of 1 MeV) as well as a fixed power law spectrum of waves. Observation of waves in the foreshock does not necessarily support this latter assumption even though a power law range has been found, as we have discussed above.

However, the suggestion of inclusion of nonlinear wave particle interaction into Lee's theory might provide an improvement if the interaction time will be shorter than the time for escape from the narrow downstream region.

7.2.3 Lee's Self-consistent Quasilinear Shock Acceleration Model

Prior to proceeding to simulations and observational evidence for shock particle acceleration at real shocks in interplanetary space we are going to briefly discuss an analytical selfconsistent model calculation of shock acceleration that has been developed by *Lee* [1982].

The Quasilinear Equations of Particle and Wave Dynamics

Lee [1982] assumed scattering of ions at circularly polarised Alfvén waves, in which case the quasilinear pitch-angle diffusion coefficient becomes

$$D_{\mu\mu} = \frac{\pi\omega_{ci}^2}{2B^2} \frac{1-\mu^2}{v|\mu|} W_k(k_{\parallel}, z) \Big|_{k_{\parallel}=\omega_{ci}/v\mu}, \quad W_k = W_k^+ + W_k^- \quad (7.18)$$

where, as usual, the power in the magnetic field fluctuations is given by $\langle |\mathbf{b}|^2 \rangle = \int W_k dk$, and the signs $\pm$ on $W_k^{\pm}$ refer to Alfvén waves propagating in the direction $\pm z$ parallel or antiparallel along the magnetic field $\mathbf{B} = B\mathbf{e}_z$, depending on magnetic field direction and wave polarisation being circularly left or right handed. This analytical expression allows to express the parallel diffusion coefficient in terms of wave power

$$\kappa_{\parallel}(v, z) = \frac{v^2}{4} \int_{-1}^1 \frac{(1-\mu^2)^2}{D_{\mu\mu}} d\mu = \frac{m_i^2 v^3}{2\pi e^2} \left\{ \sum_{\pm} \left[W_k \left(\frac{\pm\omega_{ci}}{v}, z \right) \right] \right\}^{-1} \quad (7.19)$$

The perpendicular diffusion coefficient is modelled according to the classical diffusion model $\kappa_{\perp} = r_{ci}^2/\tau_{\text{diff}}$, and the diffusion time is $\tau_{\text{diff}} = \lambda_{\text{mfp}}/v$. With $r_{ci}\omega_{ci} = v$ and $\kappa_{\parallel} = \frac{1}{3}v\lambda_{\text{mfp}}$ this yields $\kappa_{\perp} \sim v^4\omega_{ci}^{-2}/3\kappa_{\parallel}$. This essentially implies that the perpendicular diffusion increases in proportion to the wave power. The diffusive convection equation then reads

$$\nabla_z [\kappa_{\parallel}(v, z) \nabla_z F(v, z)] - (\xi_1/a)^2 \kappa_{\perp}(v, z) F(v, z) + VF(v, z) = 0 \quad (7.20)$$

In this form the diffusion equation takes account of a radius a of a magnetic flux tube from where the particles can freely escape by convection perpendicular to the magnetic field. This had been introduced by *Eichler* [1981a, b] to account for the free escape boundary. In the case of the bow shock this radius would amount to the radius of the bow shock assuming that the entire bow shock is embedded into this flux tube of radius $a \sim 15 R_E$. In such a case $r/a \ll 1$ and the original distribution function has been expanded in cylindrical symmetry with respect to Bessel functions of which only the zero order function $J_0(\xi_1 r/a)$ has been taken into account, and $\xi_1 \approx 2.4$ is the first zero of J_0 . [Note that such an approach

does not account for the finite width of the magnetosheath although one could interpret it in such a way assuming that the entire flux tube that is embracing the bow shock is draped around the magnetosphere behind the bow shock. In this case the cross section of the flux tube would become $\sim a\Delta_d$, which allowed to approximate $a \rightarrow \sqrt{a\Delta_d} \sim 5R_E$.]

For a self-consistent treatment that is still quasi-linear the above convection-diffusion equation must be accomplished by an equation for the evolution of the wave power in dependence on the presence of the fast diffuse ion component if it is assumed that the diffuse ions are responsible for the excitation and damping of the waves. Such an equation, for Alfvén waves under stationary conditions, can be written as the (truncated) wave kinetic equation

$$-(V \mp V_A) \nabla_z W_k^\pm = 2\gamma_\pm^A W_k^\pm \quad (7.21)$$

where only convective losses and linear excitation of waves with growth rate $\gamma_\pm^A$ have been retained. The boundary condition on this equation is simply $\lim_{z \rightarrow \infty} W_k^\pm = W_{0k}^\pm$, which means that is that at infinity the wave spectrum is the interplanetary wave spectrum. With the resonant Alfvén growth rate $\gamma_- = -\gamma_+ \equiv \gamma_A(k, W_k^\pm, z)$ being a known functional of the distribution function $F(v, z)$ one may redefine the spatial variable $z \rightarrow \zeta = \frac{1}{2} \int_0^z dz' / \sum_\pm W_k^\pm$. Then the above wave kinetic equation, on neglecting the small Alfvén speed $V_A \ll V$ against V , is rewritten into

$$\frac{\partial W_k^\pm}{\partial \zeta} = \mp \frac{2\gamma_A}{V} \frac{W_k^\pm}{\sum_\pm W_k^\pm} \quad \text{or} \quad \frac{\partial}{\partial \zeta} (W_k^+ - W_k^-) = -\frac{2\gamma_A}{V}, \quad \frac{W_k^+ W_k^-}{W_{0k}^+ W_{0k}^-} = 1 \quad (7.22)$$

Together with the above kinetic equation and appropriate boundary conditions at the distribution function $F(v, z)$ these equations form a closed quasilinear system for the distribution function $F(v, z)$ and spectral density of the Alfvén waves $W_k^\pm$ that has been solved analytically by Lee [1982] for a given set of plasma and shock parameters in view of application to the bow shock, but the solutions being valid quite generally.

In this approach the acceleration times for particles to reach a certain energy or momentum p can be formally estimated [Drury, 1983] from

$$\tau_{\text{acc}}(p) \simeq \frac{3}{\Delta V} \int_{p_0}^p \frac{dp'}{p'} \left[\frac{\kappa_u(p)}{V_u} + \frac{\kappa_d(p)}{V_d} \right] \quad (7.23)$$

where κ_u, κ_d are the respective upstream and downstream spatial diffusion coefficients along the shock normal direction (when neglecting the perpendicular diffusion coefficient these are the projections $\kappa_\parallel \cos \Theta_{Bn}$). Recently Giacalone & Neugebauer [2008] provided upper and lower bounds for this diffusion time (for infinitely extended regions) in terms of the shock density compression ratio $\rho_{sh} = N_d/N_u > 1$ as

$$\frac{3\rho_{sh}}{(\rho_{sh} - 1)} \frac{\kappa_u(p)}{sV_u^2} < \tau_{\text{acc}}(p) < \frac{3\rho_{sh}(1 + \rho_{sh})}{(\rho_{sh} - 1)} \frac{\kappa_u(p)}{sV_u^2} \quad (7.24)$$

on the assumption that $\kappa_u(p) \propto p^s$ is a power law function of the particle momentum p . The upper bound is obtained for constant κ which corresponds to no additional downstream turbulence. The lower bound follows for a zero downstream diffusion coefficient.

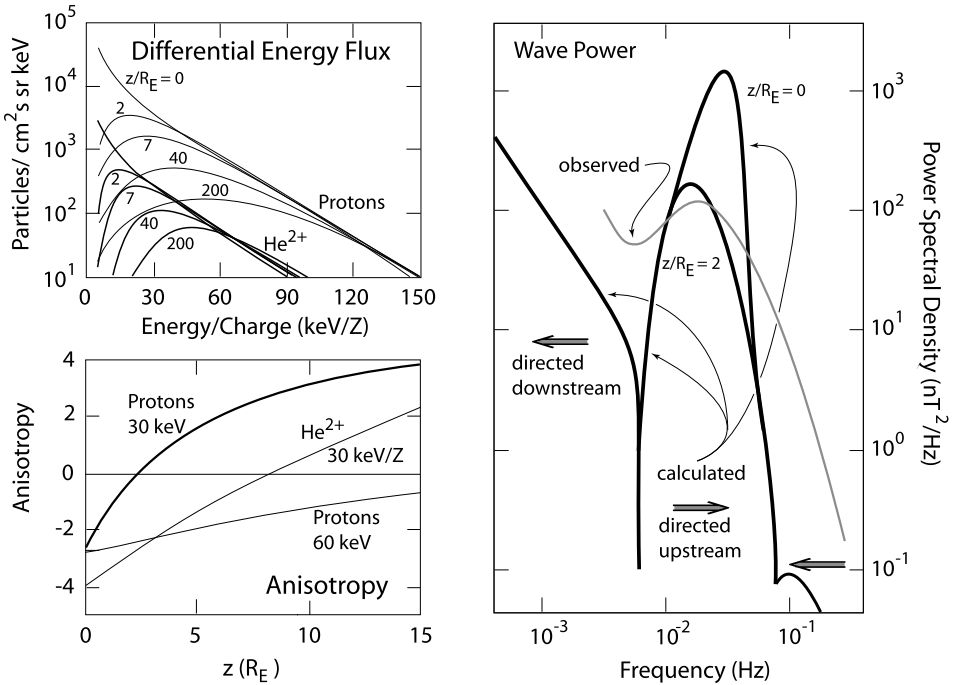


Figure 7.6: Self-consistent particle and wave spectra determined from quasi-linear convection-diffusion theory [after Lee, 1982, courtesy American Geophysical Union]. *Left:* Theoretical dependence of differential energy flux (top panel) on energy per charge number at different upstream distances z from the shock. The sample shock is the Earth bow shock to which all parameters are scaled. Therefore z is measured in Earth radii. It is clear that the particle spectra for both protons and Helium ions (He^{2+}) are exponential functions. The spectra drop towards low energies because the theory becomes invalid there, i.e. the diffusion coefficient is increasingly incorrect here. The lower panel shows the particle anisotropy as function of distance from the shock for protons and Helium ions and for two energy channels. *Right:* Self-consistent power spectral density of parallel and antiparallel Alfvén waves at two distances z from the bow shock exhibiting a resonant peak at ~ 3 mHz. The wave intensity decreases with distance. These waves are directed upstream (in the plasma frame). At higher and lower frequencies the waves flow downstream. Their high frequency part is the continuation of the low frequency spectrum. There is some similarity between the calculated and measured spectra.

Self-consistent Particle and Wave Spectra

Figure 7.6 given an impression of the power of this self-consistent theory. It shows the calculated differential energy flux for protons and Helium ions as function of energy per charge and the self-consistently generated anisotropy in the particle fluxes in dependence on distance from the shock. On the right of the figure the self-consistent wave power spectral density is shown for two distances. It is obvious that the interaction between the diffuse ions and the waves generates a large peak of upstream propagating Alfvén waves.

The remarkable result of this theory is that the solutions predict ion energy spectra close to the shock which are exponential in energy ε per charge Ze above a certain low energy cut-off with the same e-folding coefficient $\sim 20 \text{ keV}/Ze$ for all species, which is

similar to observation. Moreover, near the shock the ion intensities decrease with distance from the shock with energy dependent scale length which is $\sim 5R_E$ for 30-keV protons. The solutions also produce an anisotropy in the particle energies which may become quite high, up to $\lesssim 50\%$.

What concerns the wave spectra so the interplanetary spectrum is strongly modified near the shock by the resonant interaction between waves and diffuse particles. It maximises close to the shock in the resonant frequency range and decreases with distance together with the ion density. It is interesting enough that the waves are found to be polarised essentially linearly suggesting that both polarisations contribute comparable parts to the wave intensity, which is a little an unexpected result that is explained, however, when referring to the assumed complete isotropy of the distribution with respect to the parallel and anti-parallel propagation along the magnetic field. Thus, there are equal numbers of particles flowing in both directions that resonate with both types of waves. It is important to note that observations do not support this symmetry in the wave spectra.

Energetic Particle Diffusion Coefficients

In spite of the above noted reservations the theory of *Lee* [1982] represented a substantial progress in the analytical approach to the self-consistent particle acceleration problem. Apart from the simplifications used in order to make the problem treatable, its weakest point is the poorly known diffusion coefficient. Diffusion coefficients have been calculated under the assumption of quasi-linear theory [first introduced by *Jokipii*, 1966] of a random walk model of magnetic fields. These models have been shown to be crude for describing what is observed both in space and in numerical simulations [*Giacalone & Jokipii*, 1999]. The diffusion coefficient is anisotropic (see Eq. (7.11)) and contains three components, $\kappa_{\perp}$, $\kappa_{\parallel}$, κ_A , where κ_A is the Alfvénic part used before, resulting from scattering in Alfvén waves, and the perpendicular diffusion coefficient is assumed to be related to the parallel diffusion coefficient by

$$\kappa_{\perp} = \kappa_{\parallel} \left[1 + \left(\frac{\lambda_{\parallel \text{mfp}}}{r_{ci}} \right)^2 \right]^{-1} \quad (7.25)$$

and perpendicular transport is usually neglected because the mean free path is much larger than the gyroradius. In this expression the poor knowledge of the mean free path is the big problem as, for instance, there is no reason for having the same mean free path for all particles. *Matthaeus et al* [2003] therefore went deeper into statistical theory and proposed an integral equation form for the perpendicular diffusion coefficient

$$\kappa_{\perp} = \frac{av^2}{3B^2} \int \frac{S(\mathbf{k}) d^3\mathbf{k}}{k_{\perp}^2 \kappa_{\perp} + k_{\parallel}^2 \kappa_{\parallel} + v/\lambda_{b\parallel} + \gamma(\mathbf{k})} \quad (7.26)$$

where a is a constant of proportionality that is determined from simulations, $\lambda_{b\parallel}$ is the parallel correlation length of the fluctuating magnetic fields, and $S(\mathbf{k})$ is defined through

$S(\mathbf{k}, t') = S(\mathbf{k})\Gamma(\mathbf{k}, t')$, with $\Gamma = \exp[-\gamma(\mathbf{k})t']$, being the spatial correlation of the transverse magnetic fluctuations

$$\langle b_x[\mathbf{x}(0), 0] b_x[\mathbf{x}(t'), t'] \rangle = \int R(\mathbf{y}, t') P(\mathbf{y}|t') d\mathbf{y} \quad (7.27)$$

$R(\mathbf{y}, t')$ is the two-point, two-time correlation, and $P(\mathbf{y}|t')$ is the probability density for a particle at time t' being displaced by the amount $\mathbf{y}$ in transverse direction. Moreover, the two-time parallel-velocity auto-correlation is modelled by the isotropic assumption $\langle v_{\parallel}(0)v_{\parallel}(t') \rangle = (v^2/3)\exp(-v_{t'}/\lambda_{b\parallel})$, and it is assumed that $\gamma = 0$. With P being a symmetric Gaussian distribution for the particle trajectory (assuming that the displacement is at all times diffusive), the average over the exponential wave factor $\langle \exp[i\mathbf{k} \cdot \mathbf{x}(t')] \rangle$ which is introduced through the magnetic field fluctuations, yields just $\langle \exp[i\mathbf{k} \cdot \mathbf{x}(t')] \rangle = \exp[(-k_{\perp}^2 \kappa_{\perp} - k_{\parallel}^2 \kappa_{\parallel})t']$.

This Non-Linear Gyro-Center (NLGC) diffusion model is closer to reality as it takes account not only of the random walk of the fluctuating magnetic field lines but also of the distortion of the particle orbits in this random walk. Diffusion of particles is caused by decorrelation of their orbits from the fluctuating magnetic field. A further development in this theory by using not only Alfvénic turbulence but taking into account also the contribution of compressive modes has been presented by Zank *et al* [2004], but the deviations from the NLGC model are very small and can be neglected.

Figure 7.7 shows the used turbulent fluctuation model of the magnetic field and the comparison of the NLGC diffusion coefficient with the simpler models. The agreement between the NLGC diffusion model and the two-dimensional numerical simulation based on the two-dimensional magnetic field turbulence is very good while all other models show severe discrepancies between the simulations of the diffusion and theory. However it is not only this agreement which is of interest. More interesting is that the perpendicular diffusion coefficient is considerably less than quasilinear theory predicts and that the deviations from quasilinear theory become susceptible already at $r_{ci} \lesssim 0.1\lambda$. Still this model is not yet self-consistent as it ignores the feedback of the particles on the magnetic field turbulence which – at least to some degree – is a function of the presence of the energetic particles, as suggested by the self-consistent theory of Lee [1982]. Determination of the diffusion coefficient at this stage is a necessary intermediate step. A complete solution of the problem including parallel and perpendicular diffusion for the relevant waves as well as energy diffusion can be expected to come only from three-dimensional full-particle PIC simulations of shock formation including particle diffusion and acceleration to high energies.

Non-classical Diffusion ('Super-Diffusion')

The above diffusion coefficients are all based on the assumption of classical diffusion. Since shocks are narrow transitions from one plasma state to another one particle interactions in the vicinity of shocks might be subject to statistics of extremes and not to classical statistics. In this case the diffusion becomes time-dependent as no final state is reached in the process of particle scattering. The probability $P(\mathbf{y}|t') \rightarrow P(\mathbf{y}|t'; v)$ is not anymore Gaussian but has a long tail of power $-v$ that extends toward the rare large excursions and

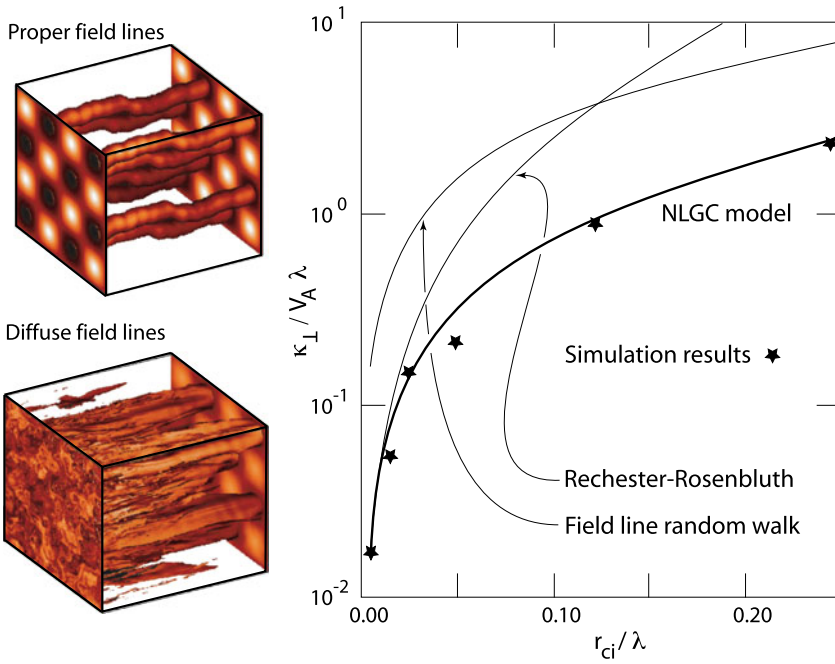


Figure 7.7: Models of the perpendicular spatial diffusion coefficient $\kappa_{\perp}$ as function of the normalised to λ ion gyroradius r_{ci} [after *Matthaeus et al, 2003*]. *Left:* The turbulent magnetic field used in the NLGC diffusion model. The *upper part* shows a couple of magnetic flux tubes that do not undergo random walk. The field lines remain unbroken. The *lower part* has 80% of the magnetic energy in turbulent two-dimensional transverse modes. The magnetic field is highly mixed. This model is used in the numerical NLGC calculations. *Right:* The perpendicular diffusion coefficient for several models, the *Rechester & Rosenbluth* [1978] model, the ordinary field line random walk (quasi-linear) model, and the NLGC model. The *stars* show the values obtained from the numerical simulation using the turbulent magnetic field on the left.

long waiting times between the excursions. The theory of such distributions goes back to *Lévy* [1954]. More contemporary developments based on fractal theory have been given by *Schlesinger et al* [1993] and *Metzler & Klafter* [2000], and a statistical mechanical argument has been developed by *Treumann & Jaroschek* [2008]. An early application to the derivation of the perpendicular diffusion coefficient [Treumann, 1997] yields a time dependent (parallel) mean square particle displacement $\langle z(0)z(t) \rangle = \kappa_{\parallel}(d, t; \nu)t$ which leads to a time dependent parallel (or non-magnetised) particle diffusion coefficient

$$\kappa_{\parallel}(d, t; \nu) = \kappa_{\parallel \text{cl}} [\nu / (\nu - d/2)] (v_{\text{an}} t)^{\alpha}, \quad \text{where} \quad \alpha = (4\nu - 2d - 1)^{-1} \quad (7.28)$$

Here d is the dimension of the system, v_{an} is an anomalous collision frequency that is usually nonzero, and $\kappa_{\parallel \text{cl}}$ is the classical diffusivity based on v_{an} . The condition that P is a positive valued probability distribution is that $\nu - d/2 > 1$. The diffusion coefficient then scales as $\kappa_{\parallel} \propto T_i v_{\text{an}}^{\alpha-1} t^{\alpha}$, which for the scaling with the collision frequency gives a scaling exponent $\frac{2}{3} < |\alpha - 1| < 1$. In order to find the scaling of the perpendicular diffusion coef-

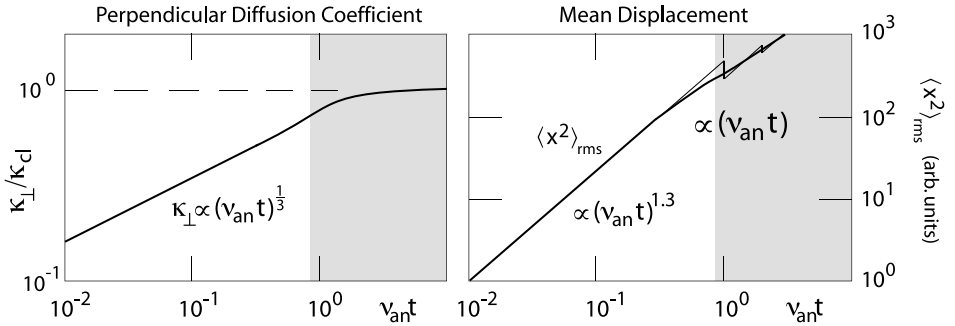


Figure 7.8: Evolution of the anomalous perpendicular diffusion coefficient under conditions of Lévy statistics. *Left:* The diffusion coefficient is a power law increasing with time t with anomalous collision frequency ν_{an} . Its maximum slope is $+\frac{1}{3}$. Approaching collision time ν_{an}^{-1} it merges into the constant classical diffusion coefficient (shaded region). *Right:* Evolution of the rms perpendicular displacement $\langle x^2 \rangle$ as function of time. The maximum slope of the power law is ~ 1.3 . When classical diffusion takes over the average displacement evolves linearly with time while between collisions being anomalous (thin line).

ficient one can use the ordinary classical formula $\kappa_{\perp} = \kappa_{\parallel} / (1 + \omega_{ci}^2 / \nu_{an}^2)$ (a more precise theory should use the integral equation for $\kappa_{\perp}$ derived in the previous section instead of this classical approximation) which gives

$$\kappa_{\perp} = \frac{v}{v - d/2} \left(\frac{\nu_{an}}{\omega_{ci}} \right)^2 \frac{T_i}{m_i \nu_{an}} (\nu_{an} t)^{\alpha} \quad (7.29)$$

Under the condition that $\nu_{an} \propto \omega_{ci}$ this yields the scaling $\kappa_{\perp} \propto \kappa_B (\omega_{ci} t)^{\alpha}$ with $0 \leq \alpha < \frac{1}{3}$. It is very important to realise that the entire physics of deviation from purely stochastic diffusion processes is contained in this extraordinarily narrow range of exponents. Very precise determination of this exponent is therefore crucial for elucidating the actual physics of diffusion and hence also diffusive particle acceleration. Below we will discover in numerical simulations of energetic particle diffusion near shocks that the non-classical limit might indeed be realised there.

This kind of transverse diffusion increases with time (see Figure 7.8) and thus the displacement of the particles grows at a faster rate than any classical transverse displacement and even faster than under Bohm diffusion κ_B . However, if the system would have sufficient time for reaching a final state it would end up in the limit of classical diffusion with linear growing root mean square displacements. This means that under non-classical conditions the system evolves diffusively faster than it would evolve under classical conditions, while given sufficient time it would settle in classical diffusion. A diffusion of this kind is called super-diffusion. In an infinitely extended system it will be realised only temporarily in the initial state for times shorter than the typical classical collision time.

In a spatially limited system like the shock transition the final stationary diffusion state might not be reached, however, since the particles have not sufficient time to undergo classical collisions. In this case the above classical theories require correction for time dependent diffusion, and a stationary state is achieved only when the time dependent dif-

fusion is balanced by losses at the moment when the mean displacement exceeds the size Δ_d of the downstream region. From the definition of $\kappa_\perp = \Delta_d^2/\tau_D$ one may estimate the limiting diffusion time τ_D as

$$v_{\text{an}} \tau_D = \left[\frac{\Delta_d^2}{r_{ci}^2} \left(1 - \frac{d}{2v} \right) \right]^{\frac{1}{\alpha+1}} \Rightarrow \left(\frac{2}{d} \frac{\Delta_d^2}{r_{ci}^2} \right)^{\frac{3}{4}} < v_{\text{an}} \tau_D < \frac{\Delta_d^2}{r_{ci}^2} \quad (7.30)$$

The left limit holds for the extreme case $v - d/2 = 1$, the right for $v \rightarrow \infty$ (the latter corresponding to a Gaussian probability distribution). For the bow shock with the downstream region being the magnetosheath of width $\Delta_D \sim 2R_E$ the relative diffusion times will be roughly $v_{\text{an}} \tau_D \gtrsim 2 \times 10^4$. Measuring the limiting diffusion time τ_D of the energetic particles provides an opportunity to determine the anomalous collision frequency v_{an} that governs the diffusive interaction.

Injection Threshold

So far we have been very vague in talking about the energy or velocity threshold $v > v_{\text{F-inj}}$ a particle has to overcome for being injected into the first order Fermi acceleration process. Our argument was that the width $\Delta_s \ll r_{ci}$ of the entire shock transition must be infinitesimally narrow compared with the ion gyro radius. This condition is still difficult to treat as, for instance in view of the quasi-parallel shock structure the whole extension of the foreshock is to be considered as the shock transition because the upstream flow loses streaming energy in this region long before arriving at the shock ramp. Moreover, the shock ramp is built up of the upstream large amplitude waves which are part of the foreshock long before reaching the shock. Therefore, making the gyro radius larger than the shock transition implies making it larger than the extension of the foreshock. It remains true that this is required if considering the particles which undergo Fermi acceleration without having any effect on the shock using the shock just as an infinitely thin boundary between two media of different speed. However, in determining the energy threshold for finding at what energy a particle can be considered as being subject to Fermi acceleration cannot be considered as strictly. There is an intermediate range when a particle has high enough energy for being accelerated while still belonging to the component that cannot be neglected when the very shock physics is to be considered. We will later from full particle PIC simulations see that this threshold is reached at energies and gyro radii that are substantially smaller than the extension of the foreshock.

A threshold condition for a particle of becoming injected into the first order Fermi cycle can be derived from the diffusive particle acceleration approach assuming that the components of the diffusion tensor are known [cf., e.g., *Giacalone & Jokipii, 1999*] as

$$\frac{v_{\text{F-inj}}}{V_1} \approx 3 \left\{ 1 + \frac{(\kappa_A/\kappa_\parallel)^2 \sin^2 \Theta_{Bn} + (1 - \kappa_\perp/\kappa_\parallel)^2 \sin^2 \Theta_{Bn} \cos^2 \Theta_{Bn}}{[(\kappa_\perp/\kappa_\parallel) \sin^2 \Theta_{Bn} + \cos^2 \Theta_{Bn}]^2} \right\}^{\frac{1}{2}} \quad (7.31)$$

an expression showing that the injection threshold is larger than $> 3V_1$, three times the upstream flow velocity or roughly 10 times the upstream flow energy as seen from the

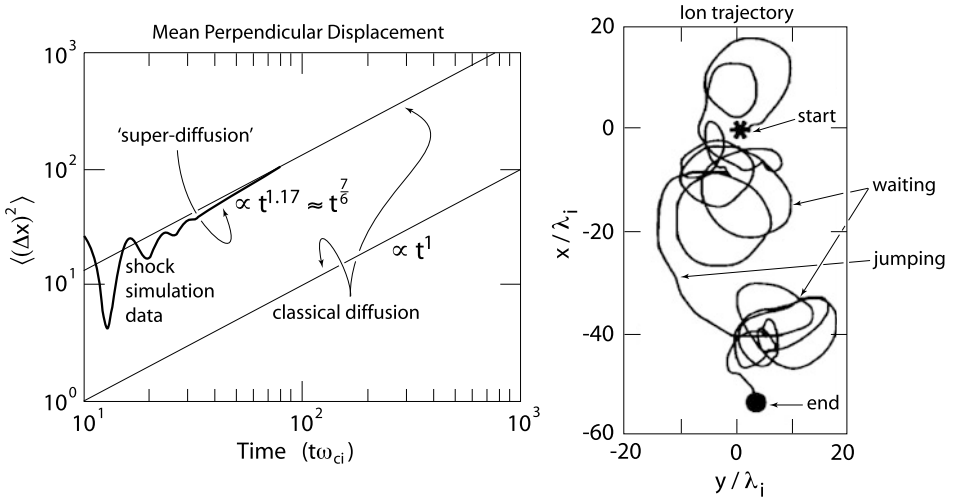


Figure 7.9: Two-dimensional numerical simulation result of the mean downstream perpendicular displacement of quasi-perpendicular ($\Theta_{Bn} = 87^\circ$) supercritical ($\mathcal{M}_A = 4$) shock-accelerated ions as function of simulation time [simulation data from Scholer *et al.*, 2000, courtesy American Geophysical Union]. *Left:* The displacement performs an initial damped oscillation before settling in a continuous diffusive increase at about $t \sim 40\omega_{ci}^{-1}$. The further evolution deviates slightly from classical (linear) increase, following a $\langle \Delta x^2 \rangle \propto (t\omega_{ci})^{1.17}$ power law. This is very close to a power of $\frac{7}{6}$, suggesting that the particle diffusion process is in fact ‘super-diffusive’ [Treumann, 1997]. *Right:* Late time trajectory of an arbitrary ion projected into the plane perpendicular to the mean magnetic field. The ion orbit in the superposition of the ambient and wave magnetic field is not a smooth stochastic trajectory. It consists of waiting (trapped gyrating) parts and parts when the ion suddenly jumps ahead a long distance as is typical for rare extreme event statistics.

shock frame. Hence, the injection threshold velocity varies between the limits $v_{F-inj,para} = 3V_1 < v_{F-inj} < v_{F-inj,perp} = 3\sqrt{2}V_1 \approx 4.25V_1$ for the parallel and perpendicular shocks, respectively.

Shock Simulation of Particle Diffusion

In order to confirm the diffusive nature of the ion acceleration process corresponding to the region downstream of a quasi-perpendicular ($\Theta_{Bn} = 87^\circ$) shock (the Mach number was $\mathcal{M}_A = 4$), Scholer *et al.* [2000] performed three-dimensional numerical hybrid simulations measuring the root-mean-square high energy ion displacement $\langle \Delta x^2 \rangle$ in the direction perpendicular to the magnetic field in the downstream region as a function of simulation time. Due to restrictions of computing power they could follow the process just up to roughly $90\omega_{ci}^{-1}$.

The result of these simulations is shown in Figure 7.9 where $\langle \Delta x^2 \rangle$ has been plotted as a function of time $t\omega_{ci}$. Initially the displacement performs a large amplitude damped oscillation until the diffusive equilibrium is attained at about $t\omega_{ci} \sim 40$. For later times the particle displacement increases continuously. However, the exponent of the increase is found not to be unity as was expected for classical diffusion. It is rather larger, being

1.17, which is very close to $\frac{7}{6}$ identifying the diffusion process as ‘super-diffusion’. (Note that, because of the large number $\sim 6.3 \times 10^6$ of simulation macro-particles used of which 525000 have high energies and contribute to the determination of the mean displacement, and because of the high time resolution, the statistical error of the measurement is less than the line width!) *Scholer et al* [2000] note this discrepancy while, nevertheless, insist on interpreting their simulations in terms of the classical diffusion picture. Above it has been shown that the exponent of the mean displacement in anomalous or ‘super-diffusion’ is $1 + \alpha$ with $\alpha = (4\nu - 2d - 1)^{-1}$, as given in Eq. (7.28) and $0 \leq \alpha \leq \frac{1}{3}$. The lower bound corresponds to classical diffusion (the linear increase of the displacement in Figure 7.9). In the simulated case $\alpha \approx \frac{1}{6}$ lies clearly in the permitted range of exponents. We may use this value together with the known dimensionality $d = 3$ of the simulation in determining the value of the anomalous parameter $\nu = \frac{13}{4} = 3.25$. (If using the exact value $\alpha = 0.17^{-1}$, we obtain $\nu = 3.22$.) This value is sufficiently far above the marginal value of $\nu = \frac{3}{2}$ for the three-dimensional case (or $\nu = 1$ for the two-dimensional case). It is, however, also far enough below the classical diffusive limit $\nu \rightarrow \infty$ thus identifying the diffusion process as anomalous super-diffusive and non-stochastic (non-Markovian). Since these simulations are completely collision-free, particle diffusion is entirely determined by anomalous processes that are mediated by the self-consistently excited wave spectrum downstream of the shock.

In view of this reasoning it is instructive to look at the trajectory an arbitrarily chosen fast ion performs in the magnetic fluctuation field. The right part of Figure 7.9 shows such an orbit in the plane perpendicular to the magnetic field at a late time $100 < t\omega_{ci} < 160$, when no average diffusive displacement was determined anymore. Clearly, the trajectory is far from being a smooth diffusive particle orbit. It rather consists of a sequence of gyrating sections and sections when the particle breaks out of gyration. When performing the former sections, the particle is ‘waiting’ at its guiding centre location. In this regime it is probably not interacting with waves. During the break-out sections it suddenly jumps out in order to occupy another waiting position. The break-out is probably caused by a brief intense wave-particle interaction. The instantaneous location of the chosen particle is progressing only in x -direction. Other particles also break out into direction y . On a short time scale such a process cannot really be described anymore as being stochastic.

The non-stochastic behaviour is reflected in the time dependence of the perpendicular diffusion coefficient. However, as we have noted above, this time dependence will last only as long as the time remains to be shorter than the collision time, when the diffusion coefficient will assume its constant classical value. This reasoning implies that *the actual diffusion in the collisionless regime is very slow*. This has been concluded (with surprise, because referring to classical diffusion) by *Scholer et al* [2000] from the simulations. The diffusion is in fact much less than classical diffusion even though the increase of the displacement is substantially faster than under classical diffusive conditions. Using classical diffusion in diffusive acceleration grossly overestimates the diffusive effect. The term ‘super-diffusion’ does not apply to the strength of diffusion but just to its time-dependence.

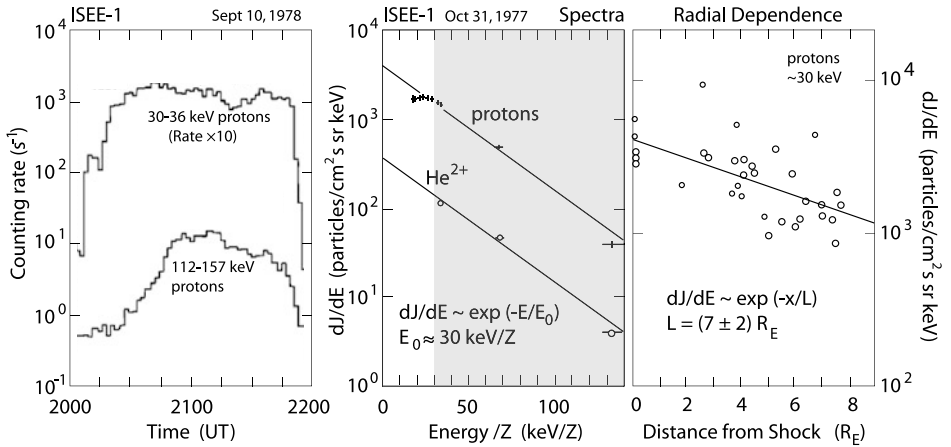


Figure 7.10: Energetic ion spectra upstream of the Earth's bow shock [data taken from *Ipavich et al., 1981a, b*, courtesy American Geophysical Union]. *Left:* The count rate in the upstream region of the bow shock for protons in the channels of 30–36 keV and 112–157 keV. Note the delay of ~ 40 min between the rises of low and high energy fluxes. *Centre:* The energy per charge number spectra decay exponentially (*shaded region*) with same energy-per charge number scale factor for all ions, suggesting energy diffusion. *Right:* Spectral variation with shock distance. The intensity of energetic particles decays exponentially with distance from shock, which is a strong argument for the shock source and for a diffusive process.

7.3 Observations

Observation of energetic ions around shocks has a long history starting in the late sixties of the past century [for a relatively recent review of inner heliospheric observations see *Reames, 1999*]. Since the ISEE spacecraft it became possible to determine both the spectra of energetic particles and their dependence on distance from the shock [in this case again the Earth's bow shock; for a review of energetic ions in the foreshock see *Fuselier, 1995*]. Figure 7.10 shows examples of those observations upstream of the shock. In the left part the count rates for protons in two energy ranges 30–36 keV and 112–157 keV have been plotted against time when the spacecraft was in the upstream foreshock region observing an energetic event that was connected with the shock [*Ipavich et al., 1981a, b*]. For two hours the shock injected high fluxes of energetic protons into the upstream region. It is interesting to note that the higher energies occur later.

Even though it is not known whether this is a purely spatial or a temporal effect, the delay in the energetic particles points on the action of an energy filter which naturally would be given by energy diffusion. The central panel shows the energy spectrum of energetic particles, this time protons and helium nuclei (α particles) plotted per charge number. Note that the plot is half-logarithmic. Hence, in both cases the spectra are exponentials, exhibiting the same scale ~ 30 keV/ Z in unit of energy per charge number Z , again reflecting the energy diffusion process. The panel on the right shows the spatial dependence of the differential energy flux for protons of 30 keV energy. Though the scatter is substantial (the correlation coefficient is 0.62), which is due to the time variability of the fluxes, in

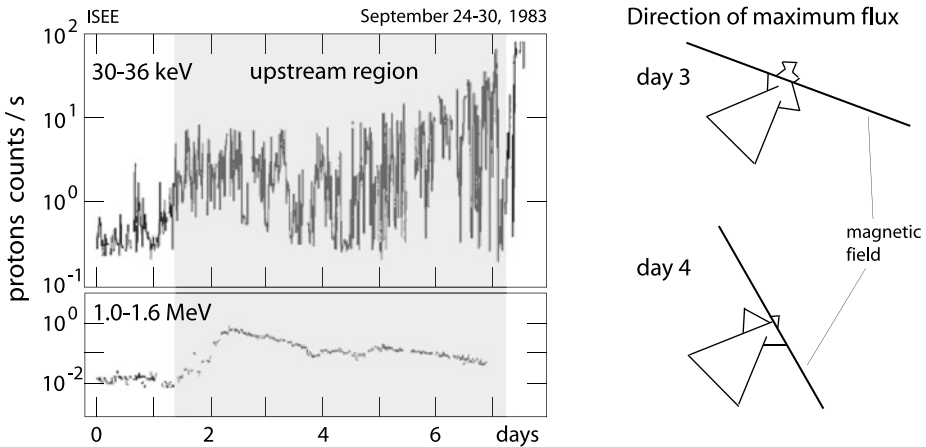


Figure 7.11: ISEE observations of energetic ions upstream of the Earth's bow shock [data taken from *Terasawa et al., 1985*, courtesy American Geophysical Union]. *Left:* The count rate in the upstream region of the bow shock (shaded) for protons in the channels of 30–36 keV and 1–1.6 MeV. The enhanced fluxes in the upstream region are well expressed. Moreover, on this long time scale the fluxes are highly variable in time in the lower energy channel. At high energies the time resolution is considerably less, so the variability is suppressed stronger. *Right:* The direction of arrival of the maximum 30–36 keV proton flux as seen in the spacecraft frame and relative to the instantaneous direction of the magnetic field for two occasions, day 3 and day 4 of ISEE observations. The counts are anisotropic, preferably distributed perpendicular to the magnetic field.

the mean the dependence of flux on distance from the shock is described quite well by an exponential decay with shock distance, implying that the energetic particles undergo a spatial diffusion process away from the shock. Hence, the shock is the source of the energetic particles.

Terasawa et al. [1985] used ISEE-3 measurements in order to determine the dependence of the upstream energetic ion fluxes on the direction of the magnetic field. Figure 7.11 shows these observations over a time period of seven days. There is a high variability of the energetic proton fluxes in the 30–36 keV channel that can be traced back to the variation in the magnetic field direction. On the right of this figure two cases of direction measurements of the fluxes are given. Clearly the maximum flux is directed perpendicular to the interplanetary magnetic field. Thus the fluxes are highly anisotropic which is partially affected by the particles being convected in this direction by the flow. At higher energies >1 MeV there is less variability in the proton fluxes which may be due to larger isotropy of the high energy particles while being affected by the lesser resolution. *Terasawa et al. [1985]* also conclude that the main acceleration of the protons occurs at the shock nose part of the upstream region which is probable as the shock is strongest in this part.

These bow shock observations have in the past been confirmed by other spacecraft [see, e.g., the AMPTEIRM observations in Figure 7.12 reported and successfully modelled by Monte-Carlo test particle simulations by *Ellison et al., 1990*]. One might argue that these exponential spectra would be a peculiarity of the conditions at the bow shock, therefore. However, energetic particle spectra downstream of interplanetary travelling shocks

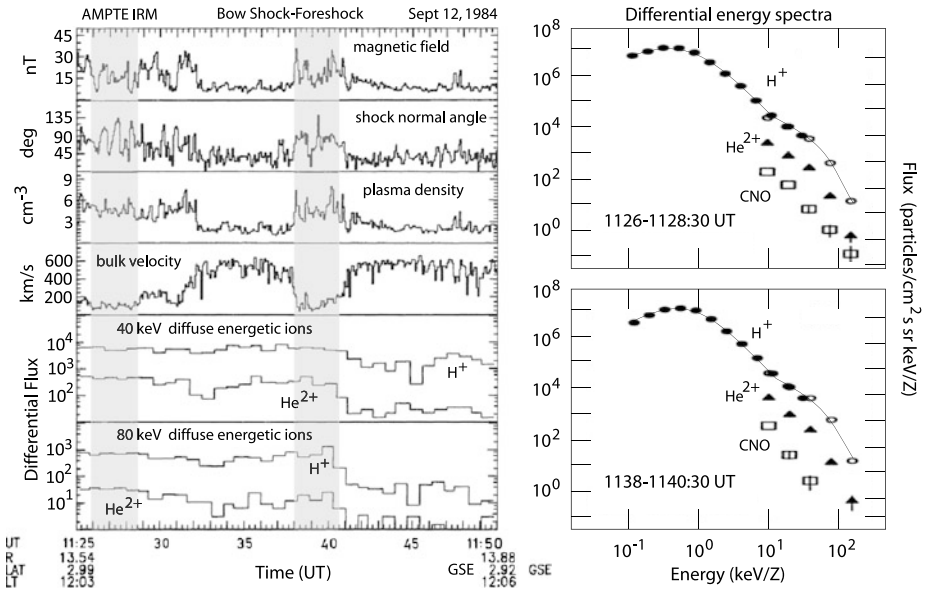


Figure 7.12: Energetic AMPTE IRM ion observations and spectra in the downstream region of the Earth's bow shock [data taken from *Ellison et al., 1990*]. This time interval was chosen because the spacecraft was at the nose of the bow shock and the interplanetary magnetic field was nearly radial providing about stationary quasi-parallel shock conditions. *Left:* Plasma and high energy particle data (the lower two panels showing proton and Helium ion fluxes at 40 keV and 80 keV). *Right:* Differential energy spectra obtained during the shaded time intervals on the left for protons, helium ions (triangles) and CNO ($Z = 6$; squares). Full and open circles are from two different instruments. A high energy tail evolves on the distribution function at >10 keV/Z. Within the measuring error the distributions are identical for all ions.

do also decay exponentially with energy. This is shown in most recent observations by the ACE spacecraft (see Figure 7.13) and has been confirmed for the same events by GEOTAIL near Earth [*Giacone & Neugebauer, 2008*]. Within the certainty of the measurement the slope on both spacecraft was the same. The intensity difference may have had many reasons, being affected by the travel time between the two spacecraft (ACE was more than 1 Million km upstream in sunward direction of GEOTAIL) and the geometry of the shock and environment. Obviously the acceleration physics is similar at bow shocks and at traveling shocks. Both kinds of shocks accelerated ions to higher energies when being exposed to scattering centres to both sides of the shock. Such scattering centres are abundant in the foreshock region and downstream of the shock due to the self-consistently excited wave spectra that are present in the close vicinity of the shock. The intensity of these waves decays at large distances from the shock and the scattering process will then rely solely on the general turbulence present in the stream.

The selfconsistent theory and observations near Earth demonstrate that the accelerated particle spectra are exponential. Hence, the acceleration mechanism generating these particles is of the kind of the first-order Fermi mechanism (or diffusive acceleration). This is in contrast to the power spectra observed at higher energies in cosmic rays. This discrep-

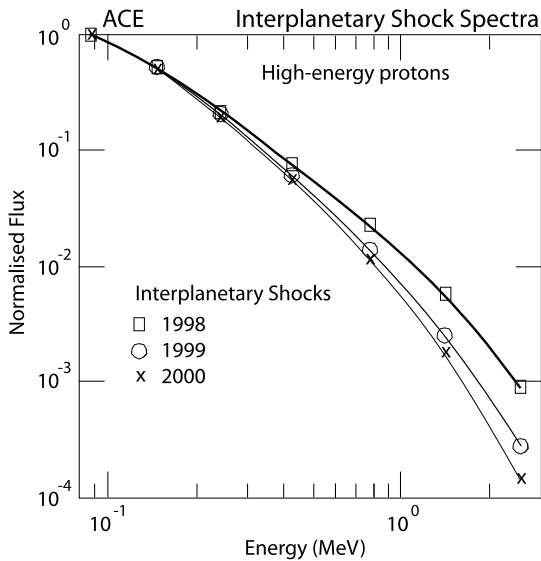


Figure 7.13: Energetic post-shock ion spectra at three travelling interplanetary shocks observed by the ACE spacecraft far upstream (about 1–2 Mill. km) in the solar wind [after *Giacalone & Neugebauer, 2008*]. Nearly identical spectral slopes though of varied intensity were observed for the same shocks during their passage near Earth by GEOTAIL. The spectral shape of the accelerated protons at the interplanetary shocks is as well exponential.

ancy has not yet been resolved while there are indications that the higher energy particles do in fact behave more like test particles and do not become so much involved into the selfconsistent shock formation/acceleration process. Moreover, at higher energies the relativistic nature of shocks comes into play thereby completely changing the picture one naively has of the process of acceleration. Shock acceleration is then closer to the original second-order Fermi process with the result that the spectrum becomes power law over some extended range in energy the extension of which depends on the width of the acceleration region and the loss mechanism. At the end of this range an exponential cut-off will limit the spectrum. The cosmic ray spectrum of Figure 7.2 suggests however that this cut-off for the energetic cosmic rays lies at extraordinarily high energies. Nothing similar applies to the shocks in the heliosphere. It thus seems that these are in the majority located in the regime of first order Fermi acceleration and do not generate a clear power law.

7.4 The Injection Problem

One of the most tantalising problems in shock acceleration physics is the origin of the energetic seed ion population that is required for starting the acceleration cycle. How are the relatively large numbers of ions generated by the shock that have energies above the injection threshold? In anticipating the result of the following discussion of a few models

that have been proposed we note that there is no solution in sight yet at this instant. Since the problem is extraordinarily complex it can be investigated only on the way of numerical simulation studies and, at its best, presumably only with the help of self-consistent full particle PIC simulations adopting real ion to electron mass ratios.

There are two ways for treating the complicated process of shock particle acceleration by numerical simulation: test particle simulations and selfconsistent simulations of shock formation. The classical way has for long time been performing non-selfconsistent test particle simulations. Because the diffusive process of particle acceleration is a slow process, selfconsistent simulations require long times and large simulation boxes. Accounting for both has partly become possible only in most recent times. In the following we mention test particle simulations only in passing before turning to a deeper rooted recollection of recent results on selfconsistent mostly hybrid and low-dimensional simulations, even though test particle simulations may come close to reality in extraordinarily large systems. Such systems are available in astrophysics for the small fractions of accelerated particles in Cosmic Rays. The space physics environment, on the other hand, requires self-consistency because the shock scales are finite, energies of shock-accelerated particles are comparably low, and the numbers of accelerated particles still substantial. Under such conditions the interaction between the accelerated particles and the plasma cannot be neglected.

7.4.1 Ion Shock Surfing

Since quasi-perpendicular shocks with their shock potential reflect low energy ions out of the upstream thermal ion population one may expect that this ion reflection might be capable of starting the Fermi mechanism under the condition that the ions are not reflected by upstream turbulence but experience scattering in the shock foot waves, i.e. they surf along the shock front [Sagdeev, 1966; Sagdeev & Shapiro, 1973; Ohsawa, 1985a, b]. This reflecting shock electric potential field can be estimated from the electron momentum equation, neglecting the motional electric field,

$$e\phi(x) \approx \int_{-\infty}^x dx \frac{1}{N} \frac{\partial}{\partial x} \left(\frac{B^2}{2\mu_0} + P_e \right) \quad (7.32)$$

Making use of the constancy of the total pressure $[B_2^2/2\mu_0 + P_{e2} + m_i V_2^2] = 0$ and neglecting the electron pressure in first approximation (see the above discussion) this can be rewritten, yielding

$$e\phi(x) \approx \frac{B_1}{\mu_0 N_1} \int_{-\infty}^x dx \frac{\partial B}{\partial x} = \frac{1}{\mathcal{M}_A^2} \frac{\Delta B(x)}{B_1} m_i V_1^2 \quad (7.33)$$

where $\Delta B(x) = B_2(x) - B_1$ is the difference between the downstream and upstream magnetic fields taken at position x . For ions this potential implies that ions with energy $\frac{1}{2}m_i v_x^2 < e\phi(x)$ will be reflected at position x because the potential is positive. As discussed in Chapter 5 on quasi-perpendicular shocks these reflected ions are partially responsible for the quasi-periodic reformation of quasi-perpendicular shocks. When gyrating back into the upstream plasma they form the shock foot, retard the inflow, cause a foot current, and

most important experience the motional electric field in which they become accelerated along the shock surface until gaining sufficient energy to pass the reflecting shock potential. Clearly, their acceleration is on the expense of the upstream motional electric field, and consequently the inflow is slightly retarded in this acceleration process.

The maximum energy gain of the accelerated ions in this multiple upstream reflection before passing the shock, which we can call ‘ion surfing’, can be estimated from the balance between the Lorentz and the electric forces on the ion, yielding the simple upper limit for the tangential to the shock ion velocity $v_{y,\max} \sim E_x/B_z$, where E_x is the cross shock reflecting electric field based on the above potential. Combination with the above result then produces the maximum energy

$$\mathcal{E}_{i,\max} \sim \frac{m_i}{2} v_{y,\max}^2 \sim \frac{m_i}{2} \left(\frac{1}{\mathcal{M}_A^2} \frac{m_i V_1^2}{e B_1 \Delta_s} \frac{\Delta B}{B_1} \right)^2 \quad (7.34)$$

an ion can gain in this multiple reflection process, where Δ_s is the thickness of the cross shock ramp electric field region. Hence the maximum energy gain depends crucially on the narrowness of this region. Since the shock thickness is not well known and the effect of the overshoot field is difficult to estimate, one assumes that in the overshoot region the magnetic field can be described approximately by a solitary structure yielding an estimate $\Delta B/B_1 \sim 2(\mathcal{M}_A - 1)$ for the maximum overshoot amplitude. With this estimate the maximum energy gain becomes very large

$$\mathcal{E}_{i,\max} \sim 2m_i V_A^2 \mathcal{M}_A^2 (m_i/m_e) = 2m_i V_1^2 (m_i/m_e) \quad (7.35)$$

In fact, in the solar wind the streaming energy of ions is of the order of 1 keV, implying that the maximum energy an ion could gain in surfing along the bow shock amounts to the order of $\mathcal{E}_{i,\max} \sim 7.5$ MeV quite high. This energy would be sufficient for entering the Fermi acceleration cycle, while the observations suggest a steep cut off of the proton spectrum at the bow shock above a few 100 keV. This could be due to the curvature of the shock that does not allow a proton to surf for sufficiently long time. But the question remains unresolved whether or not the bow shock and thus any shock can by itself accelerate ions by surfing to exceed the energy threshold of the Fermi cycle.

7.4.2 Test Particle Simulations

Test particle simulations assume a shock that is given for granted and is not (or very little) affected by the accelerated particles. This holds as long as the density of the high energy particles remains much less than the ambient plasma density in the flow. The shock can be prescribed by the Rankine-Hugoniot conditions, assuming a given shock compression ratio (thus implicitly assuming a Mach number), and the particles are subject to some stochastic scattering mechanism in a given turbulent wave spectrum after having been injected at some high enough energy such that the shock can be safely considered infinitesimally thin with no substructure. The wave spectrum is usually modelled by a power law $W_k \propto [1 + (|k|\lambda_{corr})^\xi]^{-1}$ of purely isotropic magnetic fluctuations. Here k is the wave number,

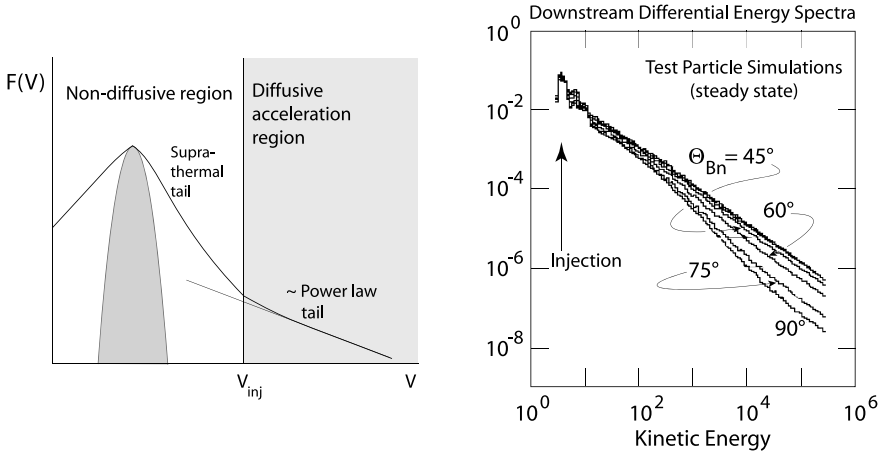


Figure 7.14: Test particle simulations of shock acceleration injecting protons into a pre-existing shock at different shock normal angles Θ_{Bn} [after Giacalone, 2005a]. *Left:* Schematic of shock acceleration. Below injection velocity V_{inj} it is assumed that the selfconsistent shock formation process produces a non-thermal tail. Sufficiently energetic ions will be capable of crossing the injection velocity threshold at the lower bound for the diffusive shock acceleration cycle. Above injection velocity a power law tail evolves due to the acceleration process. Test particle simulations are performed in this regime injecting test particle ions at an energy that is sufficiently far above the injection threshold. *Right:* Test particle spectra obtained for different Θ_{Bn} . Energy is in units of the upstream flow kinetic energy. The hardest spectra with clearest power law tail are obtained for small Θ_{Bn} , but the dependence is weak.

λ_{corr} the wave correlation length, and ζ the wave spectral index. The wave correlation length is modelled as $\lambda_{corr} \propto V_1/\omega_{ci}$ proportional to the gyroradius of the flow. Taking the constant of proportionality sufficiently large, being of the order of $10^3 - 10^4$ ensures that the injected ions perform many gyro-circles before getting out of contact with the waves. In such a spectrum of waves particles ‘diffuse’ across the magnetic field by following the meandering magnetic field in the magnetic fluctuations.

The part of the distribution below the injection threshold is not affected by the presence of the test particles, and to good approximation it is assumed that the Rankine-Hugoniot jump relations remain intact during the acceleration process. Ions are injected at energy sufficiently high above the injection threshold and are assumed to be reflected back and forth across the infinitesimally thin shock in the wave turbulence until moving out to higher energies and creating a power law tail in the high-energy part of the ion distribution. This tail is cut off exponentially at the loss boundary as has been discussed before. We repeat that the particles are accelerated only in the upstream scattering by the upstream motional electric field $\mathbf{E} = -\mathbf{V} \times \mathbf{B}$. In the downstream scattering process the particles are overtaking the waves and thus loose energy.

Figure 7.14 on its right shows an example of the power law tail spectra obtained in such test particle simulations [performed by Giacalone, 2005a]. The spectra are given in units of the upstream flow energy. The injection energy is taken as 1 MeV, and the final steady state spectra are given for different shock normal angles Θ_{Bn} . One notes that the flattest

spectrum is obtained for small Θ_{Bn} in the quasi-parallel shock. The spectra show that irrespective of the shock normal angle particles are accelerated assuming a final steady state spectrum. This spectrum becomes power law above roughly 10^3 -times the upstream flow energy for the quasi-parallel and 10^4 -times for the perpendicular shocks. Above about this energy the slopes of the spectra are comparable for both quasi-perpendicular and quasi-parallel shocks. The difference is only in the flux with the quasi-parallel shock generated flux at given energy about one order of magnitude larger than the quasi-perpendicular shock generated flux. Nevertheless, time dependent test particle simulations of the same kind show that even though the fluxes are low the quasi-perpendicular shock accelerates particles to higher energy at a given time than the quasi-parallel shock. In other words, quasi-perpendicular shocks show a higher acceleration rate than quasi-parallel shocks.

Giaccalone [2005a] also investigated the dependence of the acceleration efficiency on the strength of the magnetic fluctuations. He found that with decreasing strength the fluxes are lower but the slope of the distribution remained about the same. On the other hand, the presence of very long scale waves turns out to be important. The longer the wave and correlation lengths the more efficient is the acceleration; this holds in particular at quasi-parallel shocks while quasi-perpendicular shocks are relatively insensitive to this variation. This effect is probably caused by the variation in the shock normal angle Θ_{Bn} close to the quasi-parallel shock that is related to the presence of long scale waves. Near the shock transition these transverse waves which in the spectrum have the largest amplitudes provide the quasi-parallel shock a less parallel (or more quasi-perpendicular) character thereby enhancing the acceleration efficiency.

7.4.3 Self-consistent Shock Acceleration Simulations

While at the largest energies the dependence of the acceleration is rather insensitive to the character of the shock, self-consistent numerical simulations are the only means of receiving theoretical information about the acceleration efficiency and the very mechanism of acceleration at the shock. The lower energies are very sensitive to the nature and structure of the shock and, vice versa, the shock depends strongly on the properties of the low energy particles as long as their population forms part of the main flow. This is the case for all observed deformed particle distributions in the heliosphere where the energetic tails make up a non-negligible fraction of the main distribution. Their generation requires performing particle PIC simulation, hybrid if just the ions are the primary subject of interest, or full particle PICs if the dynamics and acceleration of electrons and their modification of the ion dynamics are in question. In both cases, for investigating reflection and acceleration, the simulations should be higher dimensional, firstly because the particle motion consists of the superposition of the two-dimensional gyro-motion and the parallel particle dynamics, secondly because of the two-dimensionality of the diffusion process occupies a central place in the acceleration of particles and their escape from the shock into space.

In the present section we review the available more recent numerical simulations of particle acceleration at shocks restricting ourselves to the non-relativistic subset of these

simulations. Our main purposes of doing so is to receive information about the theoretical possibilities of shock-injection of particles into the Fermi acceleration process. We, moreover, treat the acceleration or heating of ions and electrons separately paying attention first to ion acceleration and injection.

Downstream Ion Heating – Leakage from Downstream to Upstream

Quite early *Tanaka et al* [1983] suggested a mechanism for ion heating and injection from downstream by wave scattering. The idea of this mechanism is triggered by the observations of *Sckopke et al* [1983] that shock reflected ions from a quasi-perpendicular shock surface after having been accelerated in the upstream convection electric field (which, we remind, is geometrically parallel to the ion gyration in the quasi-perpendicular shock foot; note also that is these ions that contribute to quasi-perpendicular shock reformation!) gain high enough energy to overcome the shock ramp and enter the downstream region. This may take several cycles of ion reflection in gyration in the quasi-perpendicular shock foot.

When having passed the shock ramp and overshoot and entered the downstream region, these ions possess a large excess thermal energy in the perpendicular motion, $T_{i\perp} \simeq (3-5)T_{i\parallel}$, which is sufficient to excite unstable electromagnetic ion-cyclotron waves in the downstream region just behind the shock. Observations do indeed show that these waves form part of the downstream electromagnetic spectrum [*Narita et al*, 2006, 2008]. Quasi-linear wave particle interaction between the ion-cyclotron waves and the energetic ions scatters the ions in pitch-angle thereby approaching isotropy. Some of the ions then will escape upstream along the magnetic field adding to the upstream energetic ion component.

The simulations of *Tanaka et al* [1983] suggest that the escape is preferentially for $\Theta_{Bn} \sim 45^\circ$, and the ions have roughly parallel energy two times the upstream flow energy. This is still much too small for entering the acceleration cycle. However, these simulations assume a plane shock at fixed Θ_{Bn} . In a curved shock like the bow shock, the ions will escape upstream at the quasi-parallel part of the shock where escape is easier. Arriving upstream they will again become accelerated in the perpendicular direction by the convection electric field up to a perpendicular energy roughly a factor of two higher. Passing the shock another time the cycle will be repeated with self-scattering in pitch angle until some of the particles may reach energy high enough for entering the first-order Fermi acceleration cycle. This might be one possibility for feeding first order Fermi acceleration. However, there is only qualitative information available and no selfconsistent simulation exists so far which could support this claim. Moreover, this process is slow and probably not fast and efficient enough for ion injection. A sufficiently intense downstream wave spectrum is clearly required for pushing ions back across the shock barrier upstream by scattering. But this scattering itself is insufficient of arming the ions with the required jump in energy to overcome the injection threshold even though the ions when re-entering the shock from downstream feel the shock potential which further accelerates them in the upstream direction helping them to pass the shock back upstream against the flow.

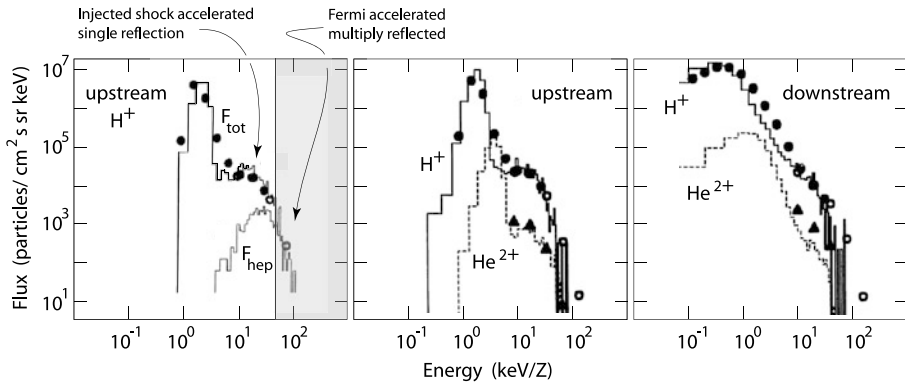


Figure 7.15: One-dimensional hybrid simulation with ion splitting showing the acceleration of ions up to the injection boundary for multiple reflection (first order Fermi) acceleration [data from *Scholer et al., 1992*, courtesy American Geophysical Union]. *Left:* The simulated upstream proton spectrum compared with the observations of *Ellison et al. [1990]*. The total proton flux F_{tot} agrees satisfactorily well with the observations. The distribution exhibits the Maxwellian upstream flow and a broad shock-accelerated exponential bump at energies >10 keV, typical for shock acceleration. The distribution F_{hep} contains the population that has been accelerated by the shock but is scattered back towards the shock thus becoming multiply scattered. Above ~ 50 keV it makes up the total distribution identifying the shaded region as the Fermi acceleration domain. *Centre:* Similar upstream simulation for protons and Helium nuclei. *Right:* The downstream distribution for the two particle components fitted to the observations. Here the data points are the same as in Figure 7.12. The agreement between observation and simulation is reasonable.

Hybrid Simulations

One-dimensional and two-dimensional hybrid simulations of self-consistent shock formation with the focus on particle acceleration have been performed for the past roughly two decades with increasing resolution [*Burgess, 1987; Scholer, 1990; Kucharek & Scholer, 1991; Scholer & Kucharek, 1999; Giacalone, 2004, 2005b*, and others].

Scholer et al. [1992] tried to reconstruct the AMPTE IRM spectra observed by *Ellison et al. [1990]* with the help of self-consistent hybrid simulations of acceleration at a quasi-parallel $\Theta_{Bn} = 30^\circ$ shock in order to check the validity of the test particle approach used by those authors and to infer about the possible self-consistent ion injection mechanism. Since the amount of self-consistently produced shock-accelerated ions in the simulations was very small they applied the trick of splitting each reflected ion into 30 parts thereby generating a factor of 30 more ions than the simulation produced while not changing the total charge and mass of the ion component. This splitting was applied only to the reflected ions once an ion was recognised as being reflected. Otherwise the simulation was one-dimensional and computer-time limited. There were a number of further restrictions that were related to the upstream boundary of particle escape which was assumed fixed in space such that the distance shock-escape boundary continuously shrunk during the simulation, and a finite resistivity had to be assumed making the simulations not fully collisionless.

Figure 7.15 shows the results of these simulations for a pure proton run (left) and a run including protons and helium nuclei (centre and right) compared to the observations

of *Ellison et al* [1990]. The fits of the simulated and observed data are satisfactory for both the upstream and downstream particle energy fluxes. The spectra obtained are exponentials and no power laws. Similar simulations in one dimension in a larger domain but assuming the presence of a given upstream wave turbulence spectrum [*Scholer et al*, 1999] and with better resolution yield also exponentials for both particle species, protons and helium nuclei.

The important physics is, however, contained in the left panel which shows the results for a pure proton hybrid simulation with a large number of particles. The two flux curves in this panel belong to the total particle flux upstream which consists of the Maxwellian flow and a shock accelerated exponential bump on the tail of the distribution which is completely self-consistently produced by the shock. The simulations do not tell in which way these protons have been accelerated. However, since in the simulation the accelerated particles had been given a flag, it was possible to plot the distribution of those particles separately which had been scattered back towards the shock from upstream. These particles have the distribution F_{hep} and obviously will, after reaching the shock, have another chance to be reflected again and further accelerated. In fact, above ~ 50 keV (in these simulations which have been tailored for bow shock conditions) the total high energy flux is completely built up of these multiply reflected particles. Hence this is the domain of first order Fermi acceleration where the particles become accelerated.

We may therefore conclude that the quasi-parallel shock is capable of accelerating ions into an exponential energetic bump. Part of the flux in this energetic bump becomes reflected back to the shock and has the chance of becoming accelerated in the Fermi mechanism. Parallel shocks are thus capable of providing the pre-accelerated particle component that becomes the seed population for shock Fermi acceleration. Concluding from these simulations the quasi-parallel shock itself is capable of generating a seed population for further first order Fermi acceleration. It does not need to be fed by particles that are reflected at the quasi-perpendicular part of the curved shock surface. Further hybrid simulations by *Scholer et al* [1998] support this conclusion by showing that the quasi-parallel shock-generated ion spectrum extends into the energy range of a few ten times the upstream ram flow energy. This is a conclusion of vital importance in application to astrophysics.

However, the very mechanism of how the parallel shock is able to accelerate the particles in order to inject them into the Fermi cycle remains an open problem. It is not answered by these simulations except that they show that the energetic particle component is not a population that is leaking out from the downstream region into the upstream region as proposed in the mechanism of *Tanaka et al* [1983]. The seed particles are part of the thermal upstream distribution that has become accelerated by the quasi-parallel shock in a first step followed by a second step in the acceleration chain.

Similar simulations in one dimension but a much larger spatial simulation domain $x \geq 28000\lambda_i$ and simulation times $t\omega_{ci} = 4000$ have been performed by *Giacalone* [2004] for a parallel shock of Mach number $\mathcal{M}_A = 6.4$. Like in the simulations of *Scholer et al* [1992], the differential energy flux spectrum exhibits a shoulder at high energies. However in this simulation, because of the large simulation domain and long simulation times, these

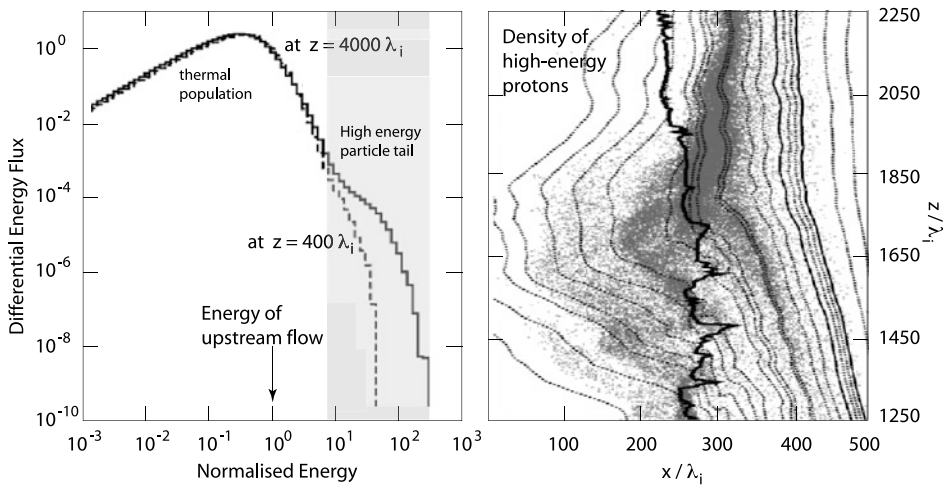


Figure 7.16: Two-dimensional self-consistent hybrid simulations of generation of energetic protons in an extended simulation box around a nearly perpendicular shock [after *Giacalone, 2005b*]. *Left:* The spectra at two different locations. The high energy tail that has evolved at the distribution is more expressed as the distant location. *Right:* The energetic proton distribution in the x, z -plane and the field lines around the shock at time $t\omega_{ci} = 150$. Plotted are only the protons in the tails of the distribution (shaded on the left at two locations). The shock ramp is the *dark line*, the field lines are the *light lines*. Energetic particles concentrate almost exclusively inside the looping magnetic field lines. It seems that here the strongest acceleration takes place.

shoulders extend out to energies up to 200 times the initial upstream plasma ram energy (in application to the solar wind an acceleration up to 200 keV).

Giacalone [2005b] has recently investigated the complementary problem of acceleration at a nearly perpendicular shock in his half-self-consistent numerical simulations. Here the ‘half’ refers to the fact that a turbulent magnetic wave spectrum is superimposed on the magnetic field with spectral index $\zeta = \frac{5}{3}$ simulating Kolmogorov wave turbulence. His results are shown in Figure 7.16. The peculiarities of this simulation are the two-dimensionality and the spatial extension along the mean magnetic field, which is about parallel to the shock surface. Such a simulation geometry allows including long wavelength magnetic perturbations. In the case of field line meandering this should enhance any perpendicular particle diffusion. The assumed Mach number is $\mathcal{M}_A = 4$. As common in hybrid simulations a small anomalous resistivity is included. The imposed turbulence together with the self-consistent shock dynamics affects the shock surface which becomes rippled. The longest wavelength is of the size of the simulation box which was also assumed to be the turbulent correlation length.

The right part of Figure 7.16 shows the instant $t\omega_{ci} = 150$ in part of the simulation box of length $\Delta z = 1000\lambda_i$. The shock surface is located at the dark line undulating between $x = 200\lambda_i$ and $x = 300\lambda_i$ along z . The light lines are the magnetic field lines. The dots are the instantaneous locations of the energetic particles with energy $> 10 \times \frac{1}{2}m_i V_1^2$, i.e. larger than ten times the upstream ram energy. The upstream density of these ions closely follows

the magnetic field, which forms a large loop whose both ends are located at the shock at two different locations in z , and the flux of supra-thermal ions is non-uniformly distributed over the shock surface. The conclusion *Giocalone* [2005b] draws is that the ions are specularly reflected and subsequently accelerated in these loops bouncing back and forth along the magnetic field from the shock. They thus do not need an upstream reflection; since they move along the loop magnetic field they are automatically returned to the shock between the bounces. The particles are acceleration because during their many bounces in the magnetic loop they experience the perpendicular upstream motional electric field that accelerates them in the direction perpendicular to the magnetic field and the flow. It is interesting to remark that this mechanism should in principle be independent of the presence of an upstream turbulence as the reflected ions anyway experience this acceleration when participating in perpendicular shock reformation. If there are many reflected ions, which is the case at high Mach number, their thermal pressure in a two or three dimensional geometry will locally cause the expansion of the magnetic field into the upstream direction and form magnetic loops along the shock surface. This expansion is not seen in one-dimensional simulations; however the local magnetic field maximum at the upstream edge of the shock foot region signals that the magnetic field there is enhanced by first, the retarding effect on the flow causing pile up of the field and, second, because of the foot current field that adds to the upstream field. In two and three dimensions this effect takes place locally along the shock front causing a look like structure of the upstream field. The upstream turbulence and the self-excited waves in the shock foot amplify this effect until in some locations along the shock surface large loops are formed in the interiors of which the particles are temporarily trapped and accelerated to high energies.

The left part in Figure 7.16 shows the particle differential energy flux as function of energy normalised to the upstream flow energy. Two spectra are shown both taken in the interval $335 < x/\lambda_i < 490$ but for different domain sizes $z_{max} = \lambda_{corr}$ in z . These flux spectra show the formation of a high energy population on the particle fluxes up to energies of 100 times the upstream flow energy. The spectra are very steep, however, and it requires very large correlation lengths and extensions of the shock in the second dimension in order to flatten them substantially. Nevertheless, the value of this simulation is that it demonstrates the possibility of a perpendicular shock to accelerate particles to high energies under the condition that the flow is high Mach number and magnetic loops form in the upstream part of the shock either because of prevalent turbulence (as assumed in this simulations) or because of any other reasons as, for instance the noted foot field contribution in two dimension or the working of instabilities like the Weibel instability which also causes the formation of magnetic loops. The importance of this mechanism is that it is capable of accelerating shock-reflected ions which are citizens of the thermal population.

In fact, similar ideas by *Balogh & Erdős* [1991] and *Erdős & Balogh* [1994] have anticipated these simulations long before. They were stimulated by spiky energetic particle injections observed during six ISEE 3 passages of nearly perpendicular ($\Theta_{Bn} = 84^\circ$ for the 30 November 1979 event, for instance) interplanetary shocks [*Balogh & Erdős*, 1983, 1985; *Erdős & Balogh*, 1990]. These authors found that close to the shocks the upstream angular distribution of particles surprisingly was bidirectional with peak proton flux in all energy channels from 35 keV to 1 MeV right at shock crossing. The higher

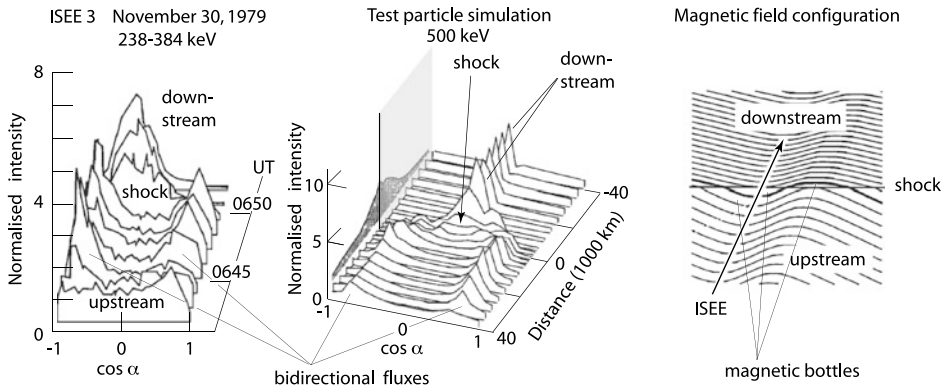


Figure 7.17: Particle acceleration in magnetic loops upstream of an interplanetary shock [after *Balogh & Erdős, 1991*, courtesy American Geophysical Union]. *Left:* Time stacked ISEE 3 pitch angle distributions of energetic upstream protons on 30 November 1979 near a quasi-perpendicular interplanetary shock showing the bidirectional proton fluxes at pitch angles close to parallel and anti-parallel to the upstream solar wind magnetic field. The shock normal angle was $\Theta_{Bn} = 84^\circ$. *Centre:* The 500 keV test particle simulation of fluxes of bottle accelerated protons which reproduce the bidirectional proton fluxes measured by ISEE 3. *Right:* The assumed magnetic field model forming magnetic loops with their foot points on the shock. Note that there are loops upstream and loops downstream of the shock. Acceleration takes place in this case in the upstream loops. This is similar to what is observed in the two-dimensional simulations by *Giacalone [2005b]*.

energies close upstream to the shock exhibited a well defined loss cone in both directions, toward and away from the shock. This bidirectional character of the distributions immediately disappeared when crossing the shock to downstream. This observation led *Balogh & Erdős [1991]* propose that upstream of the interplanetary shock the magnetic field formed loops that had both footpoints on the shock and trapped particles. This was supported by test particle numerical calculations of particle motions in an assumed magnetic loop (or bottle) structure that basically confirmed the assumptions. Figure 7.17 compares the observed and calculated pitch angle distributions as function of distance from the shock for the magnetic field model. However, the interesting point in this idea is that the trapped particles bounce back and forth in the magnetic bottle during shock approach and are scattered in pitch angle. The particle population thereby effectively gains energy and assumes a bidirectional pitch angle distribution. Particles that cross the shock are back-scattered and energised in the bottle field configuration such that their fluxes increase. We will later, in the application to the heliospheric termination shock in Chapter 11 make use of these ideas in order to explain the observations of acceleration of anomalous cosmic rays in the downstream heliosheath region of the termination shock. Acceleration there becomes possible downstream because of the concave geometry of the termination shock such that particles trapped in the spiral field in the downstream heliosheath become accelerated when the local position of the termination shock fluctuates.

It is of interest to compare both, the self-consistent hybrid simulations of *Scholer et al [1992]* and *Scholer et al [1998]* and the half-self-consistent simulations of *Giacalone [2005b]* with the theoretical predictions of the injection threshold given in Eq. (7.31). The

latter predicts that the injection threshold energy lies in the interval $9\mathcal{E}_u < \mathcal{E}_{\text{F-inj}} < 18\mathcal{E}_u$, where $\mathcal{E}_u = \frac{1}{2}m_i V_1^2$ is the upstream flow energy in the shock frame. In the solar wind in the bow shock frame we have $\mathcal{E}_u \approx 1$ keV. Hence the injection threshold should be somewhere between 10 keV and 20 keV, depending on the quasi-parallel or quasi-perpendicular parts of the shock. The simulations of *Scholer et al* [1992] in Figure 7.15 which have been confirmed by the more extended simulations of *Scholer et al* [1998] for the bow shock show that the threshold lies slightly higher near an energy of ~ 50 keV which is not in disagreement with the theoretical prediction. The bow shock is capable of producing ions in this range. The more optimistic two-dimensional half-self-consistent simulations by *Giacomini* [2005b] in Figure 7.16 suggest a somewhat lower threshold. Note, however, that in these simulations no distinction has been made between the particles such that it is not known which of the particles in the tail have actually passed the injection threshold and participate in the Fermi cycle.

Terasawa's Mechanism of Trapped Particle Acceleration

In an attempt of finding a mechanism that explains the extraction of ions from the thermal particle population in the interaction with the shock inferred by *Scholer* [1990], *Kucharek & Scholer* [1991], *Scholer et al* [1992] and *Scholer & Kucharek* [1999] from their hybrid simulations, *Sugiyama & Terasawa* [1999] suggested a simple mechanism which could efficiently accelerate ions trapped in a large amplitude upstream wave when the wave is crossing the shock. The mechanism exploits the fact that the flow velocity of an upstream (monochromatic) wave decreases when the wave is crossing the shock. This is obvious when remembering that the velocity of upstream waves is essentially the Alfvén speed. During the time of passage of the wave across the narrow shock transition, the trapped particle in this case bounces back and forth across the shock transition between a fast and a slow propagating wave and has a chance to pick up the phase velocity differences. The particle moves along the lines of constant energy when jumping from one frame to the other. The process gradually increases the particle energy. For a high bounce frequency of the particle it experiences many reflections and can thus be speeded up to an energy which might become high enough for injection into the Fermi mechanism. The action of the mechanism is schematically shown for a monochromatic wave in Figure 7.18. This mechanism has been elaborated by *Sugiyama et al* [2001]. The particle orbit and energy gain is shown on the right in the figure. The process is essentially diffusion in energy space which for sufficiently large wave amplitudes becomes chaotic. The stepwise increase in particles energy is well illustrated.

A restriction of this process is the assumption of monochromatic waves. Simulations have shown that quasi-parallel shocks are built of large amplitude and irregularly shaped pulsations. These pulsations are fast before becoming the shock, i.e. before reaching the shock. It is also seen in the simulations that ions are not only retarded but also trapped between the pulsations and the old shock front. During the approach of the pulsation they may bounce many times between the old and new shock, a process that is similar to Terasawa's wave acceleration while at the same time being fundamentally different. In the case of pulsations the ions are trapped between two approaching large amplitude magnetic

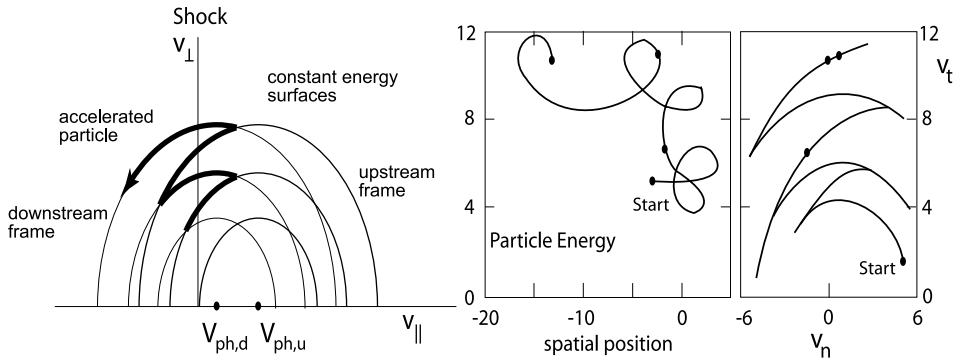


Figure 7.18: Particle injection from a quasi-parallel shock by trapping acceleration across the shock [after Sugiyama *et al*, 2001]. *Left:* Schematic of the mechanism. The upstream and downstream waves move at different phase velocities. The curves show circles of constant energy for a trapped particle. The ion moving back and forth across the shock along the different lines picks up the energy between the upstream and downstream wave frames exactly like in Fermi acceleration, until having sufficient energy to escape and enter the real Fermi process. *Right:* Particle orbits in this process. The energy increases while the particle is moving with the waves. The *outer panel* shows the bouncing in normal and tangential velocity coordinates.

waves like between two mirrors and experience a strong acceleration by bouncing back and forth between the mirrors as seen from the system of inertia. This is a kind of Fermi acceleration the ions experience. However, at the same time the ions experience the convection electric field of the upstream flow which is along the magnetic shock ramp since the magnetic field is locally about quasi-perpendicular. Hence the ions become accelerated due to two effects, the bounce motion between the approaching mirrors of the two pulsations (the new and old shock ramps) and the acceleration in the convection electric field along the shock. Both effects together will lift the energy of the trapped ions and case a diffuse isotropic high energy ion component. This process has so far escaped attention due to the difficulty of simulating it. It is, however, quite reasonable to assume that it might accelerate a fraction of the trapped ions in a quasi-parallel reforming supercritical shock to high energies.

7.5 Accelerating Electrons

Unlike in astrophysics, shock acceleration of electrons is of secondary importance in the Heliosphere. The first and obvious reason – with few exceptions – is the accessibility of most of the shocks in the Heliosphere which allows for the performance of measurements *in situ* the shock environment. The second and less obvious reason is that in non-relativistic collisionless shocks the dynamics of the shock, shock formation and reformation is mainly governed by ions which are the main carriers of the kinetic energy in the primary supercritical Mach number flow. This is also the reason for the great success of hybrid simulations in describing the shock dynamics.

It has only recently been realised that electrons do also affect the shock dynamics at high Mach numbers through the excitation of instabilities in which ions are involved and through small scale currents flowing on the microscopic scale $\sim \lambda_e$. Otherwise electrons serve the remote sensing and diagnostics of shocks from remote distances. Such inaccessible shocks are the shocks in the solar atmosphere, from the chromosphere up into the high corona, long before some of the shocks that have been generated there manage to escape into interplanetary space and becoming visible as CME driven shocks, shocks in corotating interaction regions (CIR shocks) and, closest to the Sun, the solar type II burst shocks.

As long as the shocks remain to be hidden in the hot corona where spacecraft cannot enter, such shocks are realised only in the signature of the radio radiation they emit, as well as in X-rays and sometimes even Gamma-rays. This radiation is commonly used for diagnostic purposes that are based on theories of its excitation. Here the acceleration of electrons becomes of interest as well, even though some of the observed accelerated electrons and their signatures in the radiation are definitely not shock generated. The exempt applies to solar type III radio bursts and possibly also to some classes of prompt solar X-ray and Gamma-ray flares, which are believed to be caused by reconnection rather than shock acceleration, or by the direct action of localised strong electric potential drops along the local magnetic fields. In the following we, however, restrict ourselves to a concise description of what is currently known about the acceleration of electrons by non-relativistic collisionless shocks.

7.5.1 Introductory Remarks

Electron acceleration in non-relativistic shocks meets at least two severe problems. The first problem is that a super-critical shock must reflect ions in order to get rid of the excess energy, simply because ions at a given Mach number carry the bulk of the kinetic energy. In order to do this the shock potential is positive. A simple explanation for the generation of a positive potential in a high Mach number flow at the quasi-perpendicular shock transition can be based on the large discrepancy between the convective electron and ion gyroradii: at same perpendicular speed we have $r_{ce}/r_{ci} = m_e/m_i$, with the consequence that ions penetrate considerably deeper into the shock transition than electrons. A similar condition holds for the inertial lengths, which lead to a less restrictive ratio $\lambda_e/\lambda_i = \sqrt{m_e/m_i}$. The difference in penetration depth generates a positive space charge over a fraction of the distance r_{ci} (or λ_i) and give rise to an upstream directed electric field component which accelerates electrons downstream and inhibits electron reflection.

The second problem is also related to the small electron gyroradius (or inertial length) because it restricts the shock normal angle of electron reflection to near perpendicular. A third difficulty is that for electrons in a stream like the solar wind the kinetic energy of the streaming motion $\frac{1}{2}m_e V^2 \ll T_e$ is usually much less than the electron thermal energy. In fact, for a streaming velocity of $V = 1000 \text{ km s}^{-1}$ the kinetic energy is just $\sim 5 \text{ eV}$, while the solar wind electron temperature is $T_e \simeq (50-100) \text{ eV}$. Even though the shock is supercritical, the dominant electron speed is the electron thermal velocity. Writing $r_{ce}^2/r_{ci}^2 = (m_e/m_i)(T_e/K_i)$, with $K_i = m_i V_i^2/2$ the ion kinetic energy, we have $r_{ce}/r_{ci} =$

$\sqrt{T_e/K_i}(\lambda_e/\lambda_i)$. In the solar wind $K_i \sim 1$ keV. Hence, due to the large electron temperature, $r_{ce}/r_{ci} \sim \sqrt{m_e/10m_i}$ is just a factor 3 smaller than the ratio of the inertial lengths. However, the high electron temperature means that the electrons are about isotropic with negligible flow speed, and flow-dependent electron reflection *per se* as considered in the next section does practically not depend on the flow.

7.5.2 The Sonnerup-Wu Mechanism

A simple reasoning based on a first attempt by *Sonnerup* [1969] was put forward by *Wu* [1984] in order to explain electron reflection from a nearly perpendicular shock. In a purely kinematic picture, the energy gain of a shock reflected particle depends on the de Hoffmann-Teller velocity

$$\mathbf{V}_{\text{dHT}} = \frac{\mathbf{n} \times \mathbf{V}_1 \times \mathbf{B}_1}{B_{1n}}, \quad B_{1n} \equiv \mathbf{n} \cdot \mathbf{B}_1 \quad (7.36)$$

where $\mathbf{n}$ is the shock normal vector. The shock velocity along the upstream magnetic field is $V_{s\parallel} = V_{\text{dHT}} \sin \Theta_{Bn} + V_{1\parallel}$. We recall that the above definition of $\mathbf{V}_{\text{dHT}}$ makes sure that in the de Hoffmann-Teller frame the bulk upstream flow has no component perpendicular to the upstream magnetic field. The complete upstream flow is along $\mathbf{B}_1$ in this frame, i.e. the de Hoffmann-Teller frame translates in the perpendicular direction with the velocity of the magnetic field lines projected to the shock surface. However, at a nearly perpendicular supercritical shock the parallel flow speed $V_{1\parallel}$ is small. In the expression for $V_{s\parallel}$ it is then negligible.

Considering the particle distribution in the de Hoffmann-Teller frame is achieved by subtracting $V_{s\parallel}$ from the parallel upstream particle speed $v_{\parallel}$. Assuming that a particle can be mirror reflected at the shock under conservation of the particle magnetic moment implies that the particle velocity measured in the de Hoffmann-Teller frame has a pitch angle against the magnetic field that is larger than the loss-cone angle $\alpha > \alpha_{lc}$ defined through $\sin^2 \alpha_{lc} = B_1/B_{os}$, where B_{os} is the overshoot (or maximum) magnetic field in the quasi-perpendicular shock transition. This is shown schematically in Figure 7.19. This loss cone angle is completely determined through the compression ratio of the shock. For a compression ratio of 3 the loss cone angle is $\alpha_{lc} = 35^\circ$ covering a large part of the distribution function for large de Hoffmann-Teller speeds corresponding to large $V_{s\parallel}$. The narrowest loss cone would be obtained for the largest (fluid) compression ratio 4, yielding just $\alpha_{lc} = 30^\circ$. These numbers are based on the Rankine-Hugoniot MHD relations and do not account for the overshoot magnetic field, which narrows the loss cone angle a little more.

Since most of the particles are not reflected but pass the shock, these particles are lost from the distribution, which is equivalent to the fact that particles that can possibly be reflected must already have sufficiently high velocity, a conclusion having been drawn already by *Fermi* [1949]. One is thus forced to assume that the electron distribution consists of a thermal core distribution and a superposed halo distribution $F_e(v) = F_{\text{core}} + F_{\text{halo}}$. The number density of reflected particles follows from

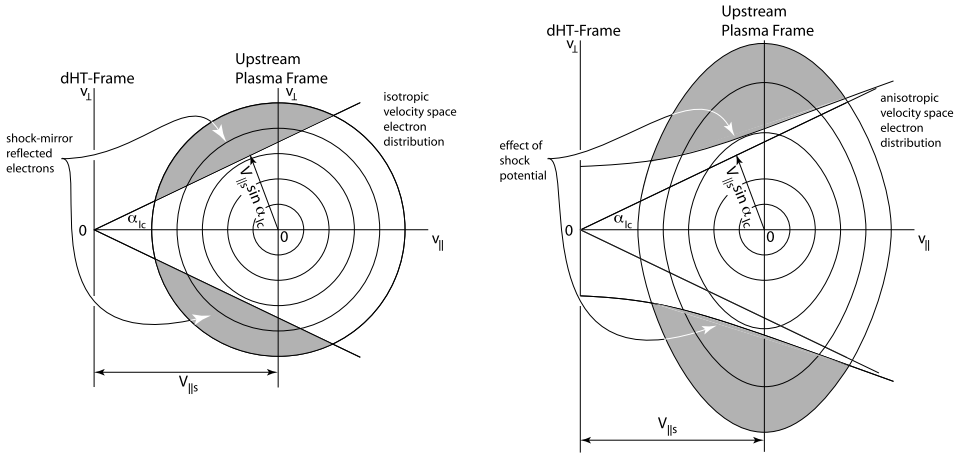


Figure 7.19: Electron phase space distribution in velocity space seen from two frames, the upstream plasma frame and the de Hoffmann-Teller frame, on the *left* the isotropic distribution [after Wu, 1984], on the *right* the anisotropic distribution with $T_{\perp} > T_{\parallel}$. The dHT-frame is obtained translating the parallel flow velocity in the direction along the upstream magnetic field by the parallel shock velocity $V_{\parallel s}$. The loss-cone angle seen from the dHT-frame cuts out the particles of low speed which pass the shock. The higher energy particles (*shaded*) are ‘trapped’ in the upstream flow magnetic field and can become reflected. On the *right* the enlarging effect of the attractive ambipolar electric shock potential on the loss cone is shown schematically. For being mirror reflected, electrons require quite a large temperature or perpendicular anisotropy.

$$N_{\text{refl}} = 2\pi \int_0^{\infty} dv_{\parallel} \int_{v_{\parallel} \tan \alpha_{lc}}^{\infty} v_{\perp} dv_{\perp} F_e(v_{\parallel}, v_{\perp}) \quad (7.37)$$

where the distribution has to be transformed into the de Hoffmann-Teller frame. Clearly, according to the above assumptions, only the halo distribution contributes to this integral. In this formulation the electrostatic ion-reflecting field in the shock ramp is not included. This can however be corrected as has been shown by Goodrich & Scudder [1984]. The effect of the electrostatic shock potential is to accelerate the impacting electrons in the direction downstream of the shock. This widens the electron loss cone and reduces the number of reflected electrons further.

As long as sufficiently high energy electrons are present in the upstream plasma some of them will become reflected due to the action of the ‘shock mirror force’ and will return into the upstream flow where they experience the upstream electric induction field and become accelerated in the direction opposite to the reflected ions thereby increasing the current in the foot of the quasi-perpendicular shock. (Note that this amplification of the shock foot current further increases the vulnerability of the shock foot to the excitation of various current driven instabilities like the Buneman two-stream and modified two-stream instabilities.) This perpendicular acceleration further increases the electron energy and thus increases the chance of the electrons to become a second time reflected. In addition it introduces an anisotropy (see the right part of the figure) into the accelerated electron distribution with higher perpendicular than parallel energy, elongating the electron distri-

bution into the perpendicular direction. Again this is in favour of the reflectivity but also may serve as a source for the whistler instability with the consequence of enhanced pitch angle scattering and progressing isotropisation. Hence, for a small number of electrons multiple reflection is programmed into this mechanism of reflection.

7.5.3 Hoshino's Electron Shock Surfing Mechanism

This magnetic mirror reflection mechanism completely neglects any electric potential drop in the shock transition. There are several contributions to such a potential drop. The first is due to the inertia difference between ions and electrons, which induces the general ambipolar large-scale shock electric field which we have discussed earlier in relation to ion shock surfing. Electrons, however, are attracted by this potential and will not be reflected. (This has been noted above in passing when speaking about the shock electric field effect which diminishes the magnetic mirror effect.) However, the intensification of the current in the shock front by the mirror effect can cause several instabilities to grow and can substantially modify the structure of the electric field such that electrons can under certain conditions become reflected, thus intensifying the mirror reflection process, an idea that has been made use of in the *Hoshino* [2001] electron shock acceleration model and the related simulations by *Hoshino & Shimada* [2002] and *Amano & Hoshino* [2007].

Electron Shock Surfing Mechanism

The fact that the passing electrons – i.e. the non-reflected part of the upstream plasma electron population that crosses the shock for afterwards constituting the downstream plasma electrons – is heated and partially thermalised by the action of micro-instabilities in the shock transition layer when crossing the shock led *Hoshino* [2001] to propose that these micro-instabilities might be capable not only of thermalising the plasma but also of accelerating particles (ions and electrons) to large energies (as sketched on the right in Figure 7.3, and detailed in Figure 7.20 electrons). Since it is the electrons which are most vulnerable to micro-instabilities this acceleration would be of particular interest in electron acceleration. In *Hoshino's* words, these “non-thermal accelerations that occur in the shock transition layer/shock front layer are still controversial, but at least they are believed to serve as ‘seed’ particles which are subject to further acceleration to much higher energies by the diffusive shock acceleration” mechanism. In contrast to diffusive acceleration which dwells on multiple shock crossings and reflections “the acceleration in the shock front region is provided when the particle traverses a thin shock layer, and the acceleration time may be shorter than that of the diffusive shock acceleration” which therefore “may be called (a) ‘fast’ process”. Justification for this assumption *Hoshino* [2001] seeks from the recent observation of large amplitude highly localised electric fields in shock transitions in space [*Matsumoto et al*, 1997; *Bale et al*, 1998], measurements that have been confirmed by later observations [*Bale et al*, 2002; *Behlke et al*, 2004; *Hull et al*, 2006; *Oka et al*, 2006; *Bale & Mozer*, 2007] and have been referred to in previous chapters. According to these observations the shock transition is a region where large amplitude solitary structures or BGK modes are generated.

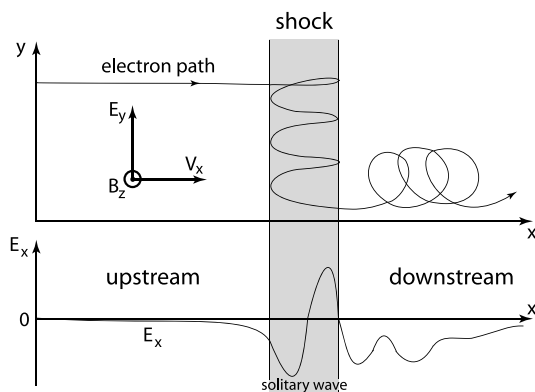


Figure 7.20: Sketch of Hoshino's electron shock surfing mechanism [after *Hoshino, 2001*]. The electron after arriving at the shock is trapped for a while inside the shock potential which in this case is assumed to be a solitary wave. After having become accelerated the electron leaves the shock downstream.

The most recent measurements by the Polar spacecraft [*Bale & Mozer, 2007*] during a bow shock transition suggest parallel electric fields $E_{\parallel} \lesssim 100 \text{ mV/m}$, and perpendicular electric fields $E_{\perp} \lesssim 600 \text{ mV/m}$ on parallel scales comparable to the electron inertial length λ_e . Such fields correspond to localised electrostatic potentials of several 100 V along the magnetic field and $\lesssim 1 \text{ kV}$ perpendicular to the ambient shock magnetic field. Production of these structures requires the inclusion of electron dynamics and refers to nonlinear kinetic plasma theory. It is thus quite natural to assume that these structures are accompanied by electron acceleration and heating, in particular as on these short scales the electron magnetic moment is not conserved anymore and the electrons become effectively non-magnetic and vulnerable to prompt acceleration in the parallel and also in the perpendicular electric fields.

Considering the presence of the small-scale electric field structures, which are solitary structures of the family of BGK modes of spatial scales $\lambda_e \sim (10-100)\lambda_D$ several Debye lengths long, has two additional interesting aspects. The first is that these solitary structures expel electrons from their interiors, i.e. they represent localised negative potentials which act repulsing on the electrons. Hence, while the charge separation in the shock front that has been mentioned previously to cause a problem in electron reflection, accelerates electrons downstream, the localised solitary structures compensate and overcompensate for this effect. Electrons can in this way become effectively reflected from the shock in spite of the shock ramp bearing a larger scale electron-attracting potential. The second is that electrons have a chance of being reflected only when impacting on one of these solitary structures and feeling the solitary wave field. It is just the fraction of upstream electrons that collide with the solitary structures that is reflected. This explains why not all particles of smaller than maximum reflection energy will be reflected from the shock ramp. The particle should in addition collide with the solitary structure. For efficient acceleration it should also become trapped. This is possible in principle in two ways: firstly, by entering

the solitary structure in which it can become trapped because these BGK modes are positively charged and lack electrons inside. The electron for this to happen must overcome the negative wall around the BGK mode (which is in principle a quadrupolar structure built of the bipolar interior and the reflecting negative electron wall around; often only the bipolar structure is seen when the wall electrons are convected away). Secondly, by jumping from one solitary BGK structure to the next being some limited time in space-time coherence with the BGK modes.

As Figure 7.20 suggests, both types of these reflected/trapped electrons experience the upstream motional electric field and become accelerated along the shock front in the direction opposing the direction of acceleration of reflected ions. Again, this increases the reflected particle current along the shock front. This current is thus striated consisting of a distributed shock foot current which contains much stronger current filaments in those narrow regions where the reflected electrons flow along the shock. The energy gained by the electron can be estimated from the electron equation of motion $\dot{p}_x = -eE_x - ev_y B_z$, where $p_x = m_e \gamma v_x$ is the electron momentum, and E_x is the amplitude of the electrostatic BGK mode in the shock transition region. The trapped electron is accelerated as long as the electric force $eE_x > ev_y B_z$. During the non-adiabatic phase the electron experiences, when feeling the E_y motional field, the velocity v_y increases quickly until this inequality inverts and the electron escapes from the BGK mode soliton. The energy it can reach in this case follows from the condition that the two forces have equal magnitude, or $v_y = E_x/B_z$. Estimating this quantity requires knowledge of the solitary amplitude. Since in the Buneman two-stream instability case the BGK modes take their energy from the electron-ion current

$$\epsilon_0 E_x^2 \sim \zeta m_e N v_d^2, \quad v_d \sim 2V_1 \quad (7.38)$$

where v_d is the current drift velocity which is the difference between the reflected ions and the inflowing electrons which is assumed to be responsible for the Buneman instability. This expression follows from equating the electrostatic energy in the BGK mode to the current drift energy (note that the current is carried by the electrons). The factor ζ is the conversion efficiency which from Buneman instability theory is taken as $\zeta \sim \frac{1}{4} \approx (m_e/m_i)^{\frac{1}{3}}$. On defining the upstream motional field $E_u = V_1 B_1$ the above expression yields for the BGK structure amplitude

$$\frac{E_x}{E_u} = \frac{2c}{V_{1A}} \left(\frac{\alpha m_e}{m_i} \right)^{\frac{1}{2}} \simeq \frac{2c}{V_{1A}} \left(\frac{m_e}{m_i} \right)^{\frac{2}{3}} \quad (7.39)$$

In the solar wind the Alfvén velocity is between $50 < V_{1A} < 150 \text{ km/s}$ yielding $15 < E_x/E_u < 45$, which corresponds to solitary wave amplitudes of $E_x < 500 \text{ mV/m}$ well in the range of observations in near Earth space.

The maximum velocity that the electron can attain is then obtained from the equilibrium between the electric and Lorentz forces yielding $v_y^{\max} \lesssim 2 \mathcal{M}_A c (m_e/m_i)^{\frac{2}{3}}$ which corresponds to a maximum energy

$$\frac{\mathcal{E}_e^{\max}}{m_e c^2} \lesssim 2 \mathcal{M}_A^2 \left(\frac{m_e}{m_i} \right)^{\frac{4}{3}} \sim 10^{-4} \mathcal{M}_A^2 \quad (7.40)$$

This estimate shows that for medium large Mach numbers of the order of, say, $\mathcal{M}_A \sim 10$ electrons can gain a maximum energy of a few per cent of their rest energy, which is in the ~ 10 keV range. This should in principle be sufficient for entering the Fermi cycle. On the other hand, for very large Mach numbers the possibility arises that the electron remains trapped for very long time in the BGK solitary structure, in which case the inequality is inverted, and the electron may gain energy quite far above its rest energy. This will happen for Mach numbers $\mathcal{M}_A > 100$, cases that are realised in astrophysical systems.

Shock Surfing Simulations

Several increasingly sophisticated numerical simulations have been performed in order to check the proposal of the electron shock surfing mechanism. The question arises which instability is responsible for the generation of the solitary structures which could reflect the electrons. There are a number of instability candidates which we have discussed in Chapter 5 on quasi-perpendicular shocks. These are the Buneman two-stream instability, the modified two-stream instability, and the ion sound instability.

The Buneman instability requires electron thermal speeds smaller than the electron-ion current drift speed $v_d = v_{di} - v_{de} \approx -v_{de} > v_{eth}$. Taking the reflected ions and incoming electrons it can marginally grow with growth rate $\sim 0.03\omega_{pe}$. It immediately heats the electrons until it interrupts its growth by letting the drift velocity drop below thermal speed. It, however, quickly forms the required BGK modes which reflect and also trap low energy electrons. On the other hand the electrons heated by the Buneman instability may lie outside the de Hoffmann-Teller loss cone in which case they can become mirror reflected. These electrons are further accelerated in the foot thereby increasing the foot current until it drives another instability.

However, *Matsukiyo & Scholer* [2003, 2006] have shown by particularly tailored full particle PIC simulations that the electron current driven modified two stream instability is a stronger instability, growing faster than the ion cyclotron frequency, and in addition has a lower threshold than the Buneman two-stream and ion acoustic instabilities. It heats the electrons parallel to the magnetic field and produces BGK phase space holes. The situation is quite complex switching between the different reflected and inflowing species, their densities and temperatures and the angles to the magnetic and electric fields. Only simulations can help understanding the acceleration of electrons.

Figure 7.21 shows the final electron distribution function (differential energy flux) as the result of a high Mach number ($\mathcal{M}_A = 32$) full particle electron shock surfing PIC simulation for a quasi-perpendicular shock. This high Mach number simulation required that the mass ratio was set to the unrealistic value $m_i/m_e = 20$. The distribution is given on two scales, showing the evolution of a long nonthermal tail. Note that the dotted line in the linear scale is the corresponding Maxwellian. At energies $\gamma/\gamma_0 < 2$ the distribution is exponential (γ_0 is the relativistic energy of the incident flow electrons). At higher γ it deviates becoming non-exponential. The log-log representation shows that the high energy tail is clearly power law $\propto \gamma^{-\alpha}$ with power $\alpha \sim 10$. This is a very steep power law, however. Nevertheless it confirms the acceleration of electrons which *Hoshino & Shimada* [2002] trace this acceleration back to the action of the Buneman instability.

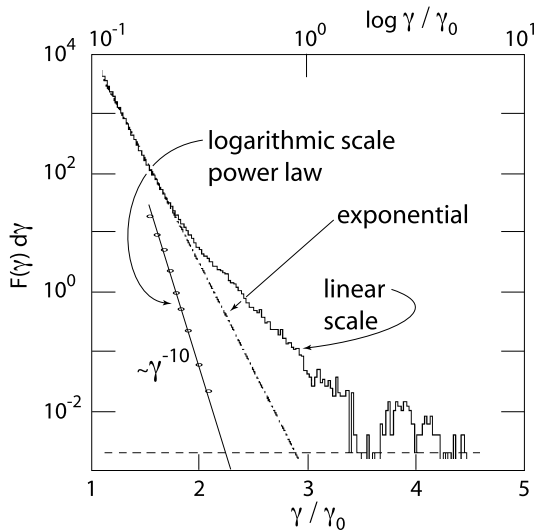


Figure 7.21: A high Mach number ($\mathcal{M}_A = 32$) electron shock surfing PIC simulation for quasi-perpendicular shocks, showing the downstream electron distribution [after *Hoshino & Shimada, 2002*]. The original data are given on linear energy scale (*bottom abscissa*) and rescaled into logarithmic energy scale (*top abscissa*). The *dash-dotted line* in linear scale indicates an exponential. The rescaled logarithmic curve is linear in log-log and therefore represents a power law which is quite steep, $\propto \gamma^{-10}$. The electron spectrum is exponential at low and power law at high energies.

Hoshino & Shimada [2002] in their one-dimensional PIC simulation have shown that in interaction with one single BGK mode the electron cannot gain more energy than $\gamma/\gamma_0 = 1.26$. Hence, the acceleration to higher energies requires a mechanism of shock surfing, i.e. further interaction in the motional electric field and possibly secondary reflections. The BGK modes serve mainly the reflection of electrons including some pre-acceleration. The path of an electron along the shock is in these simulations shown in Figure 7.22. This figure, on the left shows the time evolution of the shock normal electric field component E_x . This field exhibits a quasi-periodic structure in the foot of the shock showing that the field is concentrated in narrow spatial regions which move in a characteristic way first in upstream direction, turn around and propagate in downstream direction until decaying away. Two selected particle orbits have been superimposed on the field, one of the particles not being affected by the presence of the electric field (labelled electron 1), and another particle that interacts with the wave electric field (labelled electron 2).

Both electrons enter the shock with the upstream convective flow velocity along the straight line at the bottom of the figure. Electron 1 does not feel the wave electric field because of some obscure reason when encountering the shock. It passes over the crest of E_x . In the back of the electric field it starts performing an oscillatory motion as it is now stopped becoming a slowly moving member of the downstream flow. The central panel shows that its energy gain is small due to adiabatic heating only, while the right

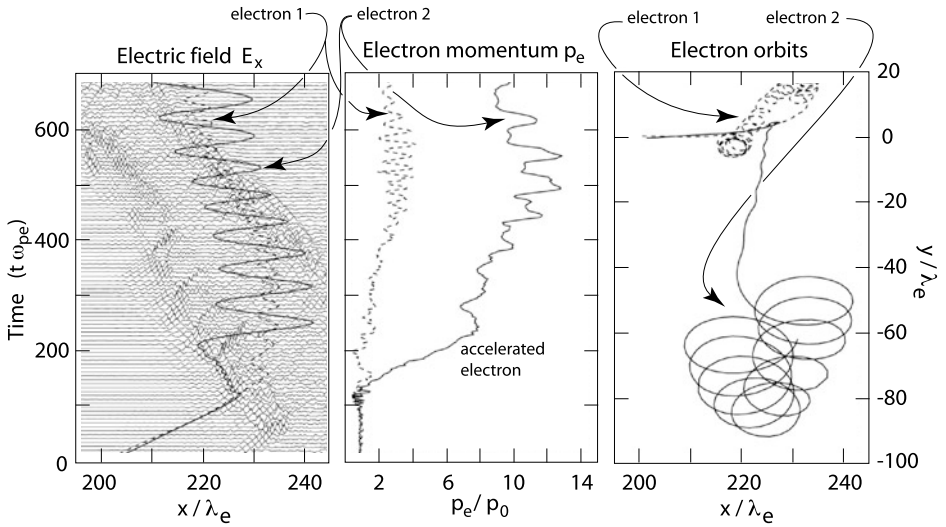


Figure 7.22: High Mach number ($\mathcal{M}_A = 32$) electron shock surfing PIC simulation for quasi-perpendicular shocks, showing *Left*: the time-stacked evolution of the electric wave field E_x in the shock transition with two electron orbits overlaid. The *dashed orbit* belongs to an electron not in resonance with the wave, the *solid line* is the resonant electron. In contrast to electron 1, electron 2 performs large excursions around the shock position before leaving the shock. Electron 1 also moves with the shock but crosses it without gaining energy. *Centre*: Evolution of the total momentum of the two electrons in time. Electron 1 gains very little momentum/energy after entering the shock, while electron 2 initially gains very much momentum and then enters a nonlinear state where the gain is slower. *Right*: The two electron orbits in the (x, y) -plane. Electron 1 moves only a small distance in both x and y , while electron 2 performs a long initial jump in y at constant x after which it becomes trapped in the wave and bounces back and forth with its enlarged gyroradius [after Hoshino & Shimada, 2002].

panel shows that it has been moderately heated in perpendicular direction performing its gyration motion.

Unlike electron 1, electron 2 when encountering the repelling electric field around about time $t\omega_{pe} \sim 100$ is stopped in its downstream directed motion and has become slightly reverted into $-x$ direction. For roughly 100 plasma periods it stays trapped in the electric field and follows its path in upstream direction. The central panel shows that during this time it becomes violently accelerated about eight-fold times. Having gained that much energy it manages to escape the field and pass to downstream performing now large oscillation in x . The two-dimensional representation of its orbit shows that similar to electron 1 it has escaped from the electric field but has shifted a long distance along the shock in y direction. The oscillations in x seen in the left panel unmask as being projections of its gyro-orbit showing that electron 2 has been accelerated non-adiabatically perpendicular to the magnetic field during its contact with the localised electric field. The lack in acceleration of electron 1 can thus be interpreted that electron 2 has encountered the shock at a position y without wave electric field. The electric field must hence have been highly localised. In fact, from the right panel one realises that electron 1 encountered the shock at $y \sim 10\lambda_e$ only. It is also interesting to see from the right panel that both electrons have not

moved far downstream during the time interval of this simulation shown. They are hanging around at a fixed location in x while the shock is moving in the direction upstream. It is only in the shock frame that they have become displaced downstream of the shock.

Considerable progress has been achieved recently in understanding the electron shock surfing process through high resolution one-dimensional fully relativistic full particle PIC simulations by *Amano & Hoshino [2007]* of a quasi-perpendicular ($\Theta_{Bn} = 80^\circ$) high Mach number shock ($\mathcal{M}_A = 15$, Alfvén velocity $V_A = 0.05c$) and for an ion-to-electron mass ratio of $m_i/m_e = 100$, upstream frequency ratio $\omega_{pe}/\omega_{ce} = 20$ fixing the plasma density for a given upstream magnetic field, and plasma beta $\beta_i = \beta_e = 0.08$. Note that for resolution of accelerated electrons a simulation should have to be relativistic. The mass ratio is still unsatisfactory when referring to the investigation of *Matsukiyo & Scholer [2006]* that the Buneman instability ceases to be important at realistic high mass ratios $m_i/m_e = 1836$ in which case it is replaced by the modified two-stream instability. In addition, one-dimensionality of the simulation misses obliquely propagating waves and electrostatic structures and can thus be only approximate. Nevertheless, the new simulation confirms shock surfing as a viable mechanism of electron acceleration and reveals a number of additional properties of shock surfing.

Detailed Electron Dynamics

Figure 7.23 gives an overview of the simulations at a late time $t\omega_{pe} = 12000$ (corresponding to $t\omega_{ci} = 5.5$). From top to bottom the panels show the incident ion phase space ($\gamma_i v_{ix}$), two cuts through the electron phase space ($\gamma_e v_{ez}, x$) and ($\gamma_e v_{ex}, x$), electron kinetic energy $\mathcal{E}_e = (\gamma_e - 1)m_e c^2$, magnetic field B_z , and the shock normal electric field component E_x . All data have been normalised accordingly (see the caption). The uppermost panel shows the incoming, reflected and shock heated/thermalised ions.

The second panel from top shows the gyrating incoming electrons as well as at negative velocities $v_{ez} < 0$ the surfing electrons. The third panel shows the heating of the electron inflow and surfing electrons. The fourth panel is the electron energy showing the enormous electron heating in connection with the flow entering the shock foot region and interacting with the reflected ion component. It is also worth noting that in the three electron panels fluxes of highest energy electrons flowing along the shock and along the magnetic field in $-z$ direction can be identified at a distance far upstream of the shock foot region. The effect of the reflected ions is seen in the formation of the foot in the magnetic field.

However, the most interesting panel is the bottom panel which shows the shock normal electric field component which in the shock ramp evolves from the ambipolar charge separation while in the shock foot is entirely due to the action of micro-instabilities. Large amplitude signatures of solitary structure BGK modes evolve close to the edge of the shock foot. Here the relative drift between reflected ions and inflowing electrons is largest. They are followed deeper in the foot by a multitude of smaller amplitude but still non-linear signatures of electrostatic turbulence. In the present simulation these structures evolve because of the growth of the Buneman two-stream instability, and we note that at a higher ion-to-electron mass ratio this picture would change somewhat because the modified two-stream instability would become more important than the Buneman instability. The other

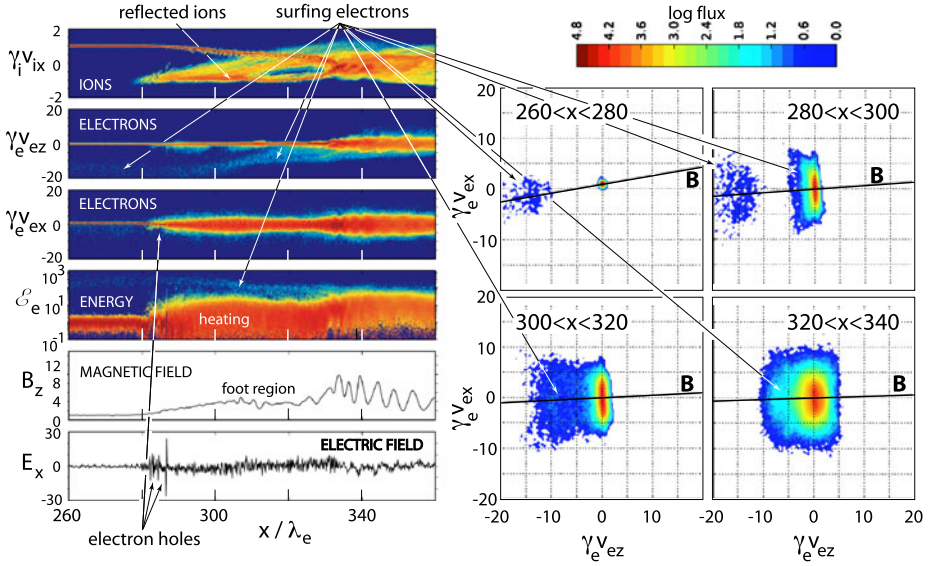


Figure 7.23: One-dimensional electromagnetic full particle PIC simulation of high Mach number ($\mathcal{M}_A = 15$) electron shock surfing for quasi-perpendicular shocks ($\Theta_{Bn} = 80^\circ$, $\beta_i = \beta_e = 0.08$) at time $t\omega_{pe} = 12000$, corresponding to $t\omega_{ci} = 5.5$ [after *Amano & Hoshino, 2007*]. The ion-to-electron mass ratio is $m_i/m_e = 100$, and the ratio of electron plasma to cyclotron frequency is $\omega_{pe}/\omega_{ce} = 20$ implying a very low plasma density. The initial magnetic field $\mathbf{B}_0 = B_{0z} \sin \Theta_{Bn} \mathbf{e}_z$ is in z direction. *Left:* The figure shows from top to bottom the ion phase space (v_{ix}, x), two components of the electron phase space (v_{ez}, v_{ex}, x), electron kinetic energy $\mathcal{E}_e = (\gamma_e - 1)m_e c^2$, with $\gamma_e = (1 - v_e^2/c^2)^{-1/2}$ and v_e the speed of each electron in the simulation frame, magnetic field component B_z and electric field component E_x . Velocities are normalised to the injection speed V_1 , magnetic field to upstream field, electric field to $E_{0y} = V_1 B_{0z}$, kinetic energy to $(\gamma_1 - 1)m_e c^2$, with $\gamma_1 = (1 - V_1^2/c^2)^{-1/2}$. The spatial scales are normalised to the electron inertial length $\lambda_e = c/\omega_{pe}$. *Right:* Electron phase space distributions taken in different space intervals in the *left part* of the figure successively approaching the shock transition. The direction of the local average magnetic field in the respective spatial range is shown as the straight line labelled **B**. The colour coding of the particle counts as function of the velocity components is logarithmic as given in the *bar on top* of the figure. The relevant information about reflected ions, surfing and accelerated electrons has been indicated by *light arrows*. Note that for resolution of the weakly relativistic energetic electrons, the code has to be relativistic. Hence the velocities given here in both phase space representations are actually the components of the 4-velocities $\mathbf{v}_e \rightarrow \gamma_e \mathbf{v}_e$.

instability concurring with the Buneman mode is the ion-acoustic instability. Its growth rate is, however, of the order of $>10\omega_{ci}^{-1}$, too long for becoming important in the distant foot but presumably responsible for some of the heating in the shock transition region.

On the right in Figure 7.23 the electron phase space distribution is plotted in four spatial sectors from the edge of the shock foot right up to the shock ramp. At and outside the foot edge the distribution consists of the electron inflow and a dilute but very energetic upstream electron beam component which escapes at high negative speed along the magnetic field. Closer to the shock in the outer part of the shock foot the inflowing electrons have become heated perpendicular to the magnetic field and show a component of electrons that separates from the bulk electrons wanting to flow along the magnetic field in

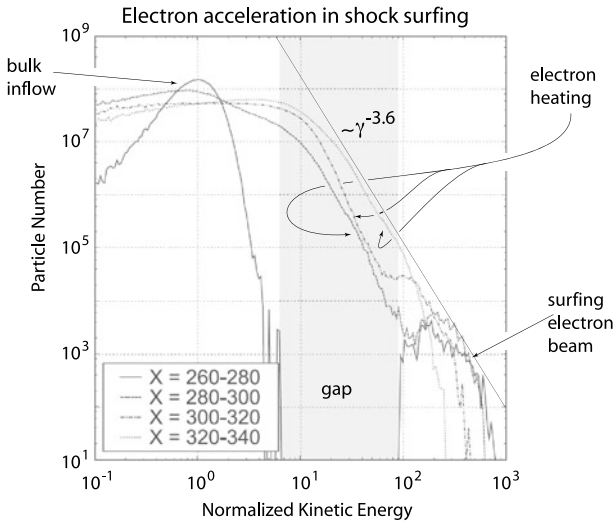


Figure 7.24: Spatial evolution of electron energy spectra for the four space intervals on the right in Figure 7.23 in electron acceleration by shock surfing [after Amano & Hoshino, 2007]. Progressing from the upstream foot boundary into the shock the electron spectrum widens with increasing electron temperature and in the centre of the shock evolves from an exponential into a distribution that over a certain energy interval exhibits a clear power law shape. Note the gap between the surfing beam and the bulk inflow distribution which is gradually filled by heating.

z direction. In addition the surfing electron beam is seen at high but lower speed along $\mathbf{B}_1$ than in the former region at the foot edge. In the inner part of the foot the bulk electrons have become heated even more while a large group of electrons begins to completely separate from the bulk and to flow away along the field. These are the surfing electrons just during their first phase of acceleration. Finally, just at the ramp the electrons are substantially heated, but the distribution consists already of two parts, the bulk and another hot component at negative v_{ez} , the source population of the surfing electrons of which not all will participate in the surfing.

The electron energy distribution functions in these four regions are shown in Figure 7.24. The interesting property of these distributions is that the field aligned surfing beam seen most pronounced at the foot edge does survive through all the four upstream regions. It is forming a bump on the distribution function, while the bulk distribution is heated ever more with approaching the shock. The height and width of this bump varies slightly, but the beam electrons are still identifiable until close to the shock, and neither their energy nor their intensity does not vary strongly. This suggests a very fast acceleration mechanism. Moreover, the bumps occur only on the upstream distribution and are thus identified as shock reflected electrons. Since electrons of such energies have not been present initially in the original inflowing electron distribution, these bumps and the corresponding electrons cannot be the result of mirror reflection by the Sonnerup-Wu mechanism until the distribution has become heated as there have not been any particles outside

the loss cone. For the Sonnerup-Wu mirror reflection mechanism the heating has to come first, and then the electrons can be reflected and accelerated in the upstream motional electric field. This acceleration is in the direction perpendicular to the magnetic field, however, while the electrons leave along the magnetic field. Note, however, that close to the shock transition the surfing electrons are heated almost isotropically (see Figure 7.23 on the right).

Heating of the bulk distribution is clearly seen in Figure 7.24. The distribution widens, and between the beam electrons and the bulk maximum evolves into a region that exhibits an approximate power law shape with variable power law. The final shape in this simulation is found to have power $F_e(\gamma_e) \propto \gamma^{-3.6}$. However, it is clearly the surfing beam electrons that possess high energy of several $10^2(\gamma_1 - 1)m_e c^2$. For instance, in the solar wind with its bulk energy of $\mathcal{E}_{e1} \sim (\gamma_1 - 1)m_e c^2 \sim 100 \text{ eV}$, this corresponds to an energy $\mathcal{E}_e \sim$ of several 10 keV, presumably sufficient for entering into the Fermi cycle for electrons. This can be concluded from a comparison of scales. For the Fermi mechanism the shock must be an infinitely thin surface. Electrons should thus have gyroradii $r_{ce} \gtrsim \Delta_s$. Since the shock width is of the order of $\Delta_s \sim 10^3 \text{ km}$, this implies electron energies of $\mathcal{E}_e \gtrsim 50 \text{ keV}$ in a $B_1 \approx 5 \text{ nT}$ magnetic field. In addition sufficient scattering of these electrons off upstream and downstream turbulence is required for further acceleration.

Further insight into the surfing mechanism is again obtained by following a selected surfing electron path. This is done in Figure 7.25. In contrast to Figure 7.22 this time the left panel shows the particle trajectory superimposed on the stacked magnetic field instead of the wave electric field. This has the advantage to directly see the motion of the shock ramp and formation of shock foot during quasi-perpendicular shock reformation.

The physics is however similar to what was concluded by *Hoshino & Shimada [2002]*. The particle moves in into the shock at the upstream convective flow velocity indicated by the straight line at the bottom. When encountering the shock foot edge containing the dense group of reflected gyrating ions the inflow motion of the electron is suddenly truncated, and the electron starts oscillating in x for $\sim 100\omega_{ce}^{-1}$ electron cyclotron periods around an almost stationary position. During this time the shock ramp approaches the electron at the speed at which the shock jumps ahead quasi-periodically during quasi-perpendicular shock reformation while the electron remains in the foot region not being able to cross the shock ramp. At the contrary, it surfs along the shock ramp and is taken over by the next cycle of reflected ions, until the end of the simulation turning together with the shock reflected foot ions into the direction upstream away from the ramp in the formation of the next reformation cycle.

We know of course that the electron is not being at rest at position x . The electron moves in fact in the $-y$ -direction along the shock being subject to acceleration by the upstream motional electric field. The excursions in its path in x seen are the projections of the gyrations of the particle into the (x, t) -plane. Note that the injected electron had so low temperature that the gyrations remained hidden in the convective straight line electron path. The gyration becomes visible now because the electron has been violently accelerated at the encounter with the edge of the shock foot.

This acceleration is seen from the central panel where the kinetic energy $\mathcal{E}_e = (\gamma_e - 1)m_e c^2$ of the electron is plotted as function of simulation time. At the instant of

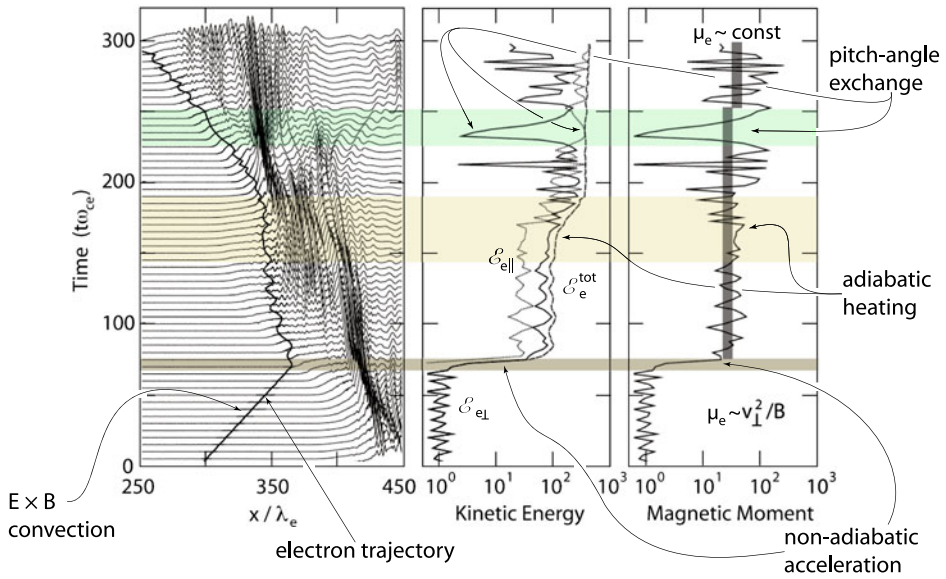


Figure 7.25: Quasi-perpendicular shock electron acceleration by shock surfing [after Amano & Hoshino, 2007] *Left:* The time-stacked evolution of the magnetic field in the shock simulation plotted after $t\omega_{ce} = 300$ ($t\omega_{pe} = 6000$). Two cycles of quasi-perpendicular shock reformation (with period $\sim 2\omega_{ci}^{-1}$) are seen with the shock ramp jumping forward (to the left) against the upstream flow. The surfing electron trajectory (solid line) runs along the outer boundary of the shock foot. The electron is practically enslaved by the group of gyrating foot ions which are in excess of the flow and need to be charge neutralised. *Centre:* Evolution of the electron energy in perpendicular and parallel components and also the total electron energy. The strongest acceleration happens in the first encounter of the electron with the foot ion group. It experiences further acceleration each time the reformation cycle ends. Note the stronger perpendicular acceleration of the electron causing an electron anisotropy $T_{e\perp} > T_{e\parallel}$. *Right:* The same for the electron magnetic moment μ_e showing that the initial energy gain within a time interval of $\Delta t\omega_{ce} \sim 5$ is highly non-adiabatic with the magnetic moment changing drastically. Afterwards the average magnetic moment (grey bar) is constant. Note also the phase of adiabatic heating in the increasing magnetic field which is mainly in the perpendicular energy.

the electron impact on the shock foot edge the electron energy suddenly increases by a factor $\sim 10^2$. This acceleration is non-adiabatic which is indicated by the sudden change of the electron magnetic moment in the outer right panel at this time. The non-conservation of the first electron adiabatic invariant $\mu_e = \mathcal{E}_{e\perp}/B$ signals that the electron is interacting on a scale shorter than the electron gyro-radius being effectively non-magnetic which can happen only when either the magnetic gradient is extraordinarily steep, which is not the case as seen from the left outer panel, or an electrostatic interaction takes place on a scale shorter than the electron gyro-radius. It can thus be attributed to large amplitude solitary BGK modes which are excited at the edge of the shock foot forming small scale electric potentials. This process is the same as in the simulations by Hoshino & Shimada [2002].

From the central panel it is seen that the acceleration is predominantly in the perpendicular direction which is due to the trapping of the electron at the BGK mode where it comes to rest and experiences the motional electric field over the time of roughly $\sim 5\omega_{ce}^{-1}$

gyroperiods. After this time the magnetic moment does not change anymore until the next reformation cycle starts at time $t\omega_{ce} \sim 250$. The further increase in energy the electron experiences between $140 < t\omega_{ce} < 190$ is purely adiabatic caused by the electron moving up into the stronger ramp magnetic field which acts as a mirror and reflects the electron by the mirror force as proposed in the Sonnerup-Wu mechanism. Since the thermal upstream electrons cannot become reflected by no means by the Sonnerup-Wu mechanism, however, one realises that the pre-acceleration of the electron by the upstream convection electric field which is made possible by trapping the electron in the electrostatic BGK wave field, is a necessary condition for shock reflection of electrons and their further acceleration. In between the two acceleration phases, the non-adiabatic and the adiabatic mirror reflection, there are some reversible variations in the magnetic moment without changing the particle energy that are accompanied by reversible exchanges between the parallel and perpendicular energies of the electron. They are induced by changes in the magnetic field direction (probably caused by phase locked whistler waves attached to the shock ramp) and related pitch angle variations which do not interest us in the context of electron acceleration.

Applied to the solar wind the electron energy gained in the non-adiabatic phase is roughly ~ 10 keV, corresponding to a $\Delta\phi \sim 10$ kV electric potential drop. In a $B_1 \sim 5$ nT field this yields a large acceleration rate of $d\mathcal{E}_e/dt \sim 1.5$ MeV/s. If the acceleration is due to the motional solar wind electric field of $E_{sw} \sim 2.5$ mV/m, the acceleration length is of the order of $\ell_{acc} \sim 4 \times 10^3$ km along the shock. Over this distance the electron remains to be in close contact with the BGK solitary wave. After leaving the BGK mode the accelerated electron has a large pitch angle against the magnetic field. It can now enter the Sonnerup-Wu mechanism for reflection at the shock since it will be found outside the loss cone which for a shock compression ratio of $B_2/B_1 \sim 3$ is $\alpha_{lc} \approx 35^\circ$.

Digression on Quasi-parallel Shock Surfing

From the above discussion of theory and simulation of electron surfing and acceleration at quasi-perpendicular shocks we learn that it is the combined effect of the interaction between the reflected ion and incoming electron components, the generation of localised electrostatic solitary structures of the BGK mode family, electron heating and acceleration in the motional electric field, and the shock mirror effect that are responsible for the generation of high energy electrons at quasi-perpendicular shocks. Even though this combination is very complicated, it works only for quasi-perpendicular shocks, and it works the better the closer the shock is to being perpendicular as then the number of reflected ions at a given supercritical Mach number is highest and the mirror effect on the electrons is largest. This raises the question whether or not quasi-parallel shocks could as well reflect and/or accelerate electrons.

The question is difficult to answer. It has, however, been realised from hybrid and full particle simulations that quasi-parallel shocks reform due to the interaction of large amplitude magnetic pulsation. These pulsations can give rise to the excitation of high frequency electrostatic waves and broadband electric noise which can be understood only as the signature of many solitary like structures of the BGK type family. Observations by CLUSTER have demonstrated the presence of such waves (see Chapter 6 on quasi-parallel shocks).

Structures of this kind, as we know, are generated only when the electron dynamics is taken into account. Quasi-parallel shocks therefore seem to host a highly active electron dynamics.

Moreover, during the reformation of quasi-parallel shocks the direction of the shock normal angle fluctuates considerably, identifying the quasi-parallel shock in many cases as a wave dominated locally quasi-perpendicular shock on scales of the order of the tangential extension of the large amplitude pulsations or SLAMS. This evokes the conclusion that on those scales, locally, the same processes of electron reflection may take place as in the quasi-perpendicular case. Mirror reflection of electrons will then become possible.

Ion reflection and foot formation happen along the front of a pulsation when the pulsation (SLAMS) approaches the shock, though on a smaller scale than in the extended quasi-perpendicular case. This might nevertheless be sufficient for electron surfing on the pulsation and subsequent acceleration. The whole effect could be extended over a much larger spatial region than in the quasi-perpendicular case where it is restricted solely to the shock foot region. In the quasi-parallel shock case it could be filling the entire region of presence of large amplitude pulsations in front of and at the ramp of the quasi-parallel foreshock. The related problems have not yet been properly considered and are open for future investigation.

7.6 Conclusions

Particle acceleration by shocks is just one aspect of collisionless shock physics even though for a large number of predominantly astrophysical applications it occupies the place of the most important property of a shock. Surprisingly, even after six decades of intensive research on this subject the original proposal made by Fermi still survives in both of its versions, first and second order Fermi acceleration. However, considering the shock acceleration literature one cannot escape the impression that the whole field is stagnating for the last twenty years. The progress since the mid-eighties has been miniscule even though the number of papers has multiplied enormously. That shocks accelerate particles, in the first place ions, is highly probable. It is supported by observations. Test particle simulations do support it as well but cannot elucidate the real acceleration mechanism. Reasons for obtaining power law particle spectra have been put forward, but the power laws obtained depend on the kind of theory and on the settings of the simulations. Theoretical reasons have been given for marginally flat power laws which in some cases seem to be in agreement with observation, but the field still looks erratic.

While the basic acceleration mechanism for energetic ions by some kind of diffusive first-order Fermi acceleration will probably ultimately cover the problem, electron acceleration is still not understood sufficiently well and, what concerns the acceleration of Cosmic Ray electrons, will presumably be solved only when considering highly relativistic shocks. In addition, ion acceleration still suffers from many open problems of which the injection problem is one of the most important and probably the most tantalising. Analytical theory has so far been unable – and will probably remain to be unable – to provide a satisfactory or even just convincing answer to the question where the high energy ions originate

that the diffusive Fermi acceleration mechanism needs to push them up the ladder to the observed high energies and the measured power law spectra. Bluntly spoken, the injection problem is too complex to be treated analytically. Lee's heroic self-consistent solution has a very limited range of applicability; test particle investigations ignore the effect of the shock itself and are valid only for very small numbers of accelerated particles which already possess high energies. In all of its versions even the most sophisticated of the test particle simulations, those which include different shock geometries, turbulence or radiation, cover just the generation of high energy tails on the distribution function and do not contribute to the understanding of the injection problem.

Investigation of quasi-parallel shock physics has in recent years shown that quasi-parallel shocks evolve in a way that is intrinsically non-stationary. The reform due to the interplay between the shock-accelerated diffusive ion component and the generation of large amplitude upstream waves, so-called pulsations. Both processes are closely related. The quasi-parallel shock in the collisionless supercritical state does not exist without the diffuse ion component which generates upstream pulsations which on their own reform the shock when being advected downstream and arriving at the shock. Hence ion acceleration into the diffuse component cannot be understood without including the generation of pulsations. This suggests that diffuse upstream ions are generated in a process which is based on the trapping of quasi-parallel shock-reflected ions between the old and new (pulsation) shock ramps where the ions experience kind of a Fermi acceleration process by bouncing between two mirrors while at the same time feel the upstream convection electric field – which differs from Fermi acceleration – which accelerates them along the shock ramp thereby producing near energetic isotropy of the ion component. The latter becomes possible because the arrival of the pulsation turns the shock normal into quasi-perpendicular direction.

On the other hand, so far self-consistent simulations of particle acceleration to high energies are still in their adolescence and could not simulate this process. The simulation boxes and times are insufficiently small and short in order to cover sufficiently large spatial ranges and sufficiently long acceleration times; realistic mass ratios which have turned out to be of crucial importance must compete with requirements on resolution, box sizes, dimensionality and simulation time which all are restricted. Nevertheless, the injection problem will be solved only on the way of such self-consistent simulations. Leakage of ions from downstream to upstream for providing the seed population is out of question. There is light over the horizon, however. For it seems that quasi-parallel shocks can accelerate ions until they overcome the injection threshold for entering into the Fermi acceleration cycle. These ions are self-consistently produced by the shock by accelerating them out of the bulk thermal distribution. The mechanism is still unclear being obscured by the assumptions on geometry, insufficient particle numbers, and low dimension. So far the acceleration of the required high energy ions has been seen only in one-dimensional hybrid simulations applying the tricks of energetic particle splitting, while electron dynamics has been completely ignored. From quasi-perpendicular shock dynamics it has, however, become transparent that the structure, formation and space-time dynamics of the shock will be understood only when considering the realistic high ion to electron mass ratio. There is no reason to believe that this would be different for quasi-parallel shocks.

Therefore what is required are full particle PIC simulations at realistic mass ratios in three dimensions, large boxes and many particles. In the present chapter we have discussed several different aspects of the particle acceleration in one or the other approximation. Indeed deepest though quite preliminary insight into some of the problems has been obtained on the way of simulations that were specially tailored to illuminate a particular problem like, for instance, the problem of perpendicular diffusion of energetic particles that have been accelerated under shock conditions. In this case it was found that the diffusion is weak while at the same time being time dependent and much faster than classical diffusion, a finding that in analytical theory has not yet been taken into account. Hybrid simulations applying energetic particle splitting, on the other hand, could demonstrate that there are good reasons to believe that quasi-parallel shocks are capable of pushing some particles over the Fermi injection threshold. Still, these simulations could not answer why this is so and what mechanism is responsible for this effect. Possibly it is the arrival of large amplitude upstream pulsations at the shock which cause a mirror effect on trapped ions when reforming the shock forcing those ions into a surfing orbit on the shock. Since electron dynamics is not considered in this case, nothing can be said about the role of higher frequency waves, BGK modes, and solitary structures which have been identified in observations to exist in the shock transition region.

The latter structures have turned out to be of crucial importance in quasi-perpendicular shock acceleration of electrons. Full particle simulations, still applying very low mass ratios – and therefore probably favouring the wrong instability – and being performed in one dimension only, have convincingly demonstrated that it is the combined ion and electron dynamics that is responsible for acceleration of electrons by the shock surfing mechanism. Since it is known from quasi-perpendicular shock reformation that higher dimensions are of crucial importance and that full mass ratios modify the picture completely in many ways, not least by replacing the Buneman two-stream instability with the modified-two stream instability, even these simulations do probably not yet provide a realistic view. Nevertheless, in which however way the BGK modes in the shock foot are generated, by ion-electron or by electron-electron interaction, their presence is crucial for electron shock surfing and electron acceleration to high energies, electron beam formation at the shock and electron heating.

At the quasi-parallel shock it is expected that similar mechanisms are at work simply because the upstream wave spectrum transforms the quasi-parallel shock locally on the small scale into a quasi-perpendicular shock, opening up the possibility of a similar process of electron acceleration also at quasi-parallel shocks. It is completely unknown what the effect would be on the ions, in particular as in the quasi-parallel foreshock the acceleration could proceed over a substantial part of the foreshock upstream of the shock where the large amplitude pulsations and shocklets are formed. Self-consistent simulations combined with high resolution observations in situ the accessible shocks in space are the only way to answer those questions. Finally, the non-relativistic shocks treated here can contribute only to the understanding of the shock formation process, shock structure and the particle injection process. The system in the solar wind is not large enough and the upstream and downstream turbulence required to provide the scattering centres is not strong enough to shock-accelerate particles to the high cosmic ray energies.

7.7 Update – 2012

The following is a very brief update referring to the most relevant literature on particle acceleration. We do not go into depth at this place. We only point the reader to some of the literature which has been published since this chapter has been written, i.e. since early spring 2008. The number of papers published on shock acceleration during this period has been huge. However, there is no need to review them at this place because only very few of these papers are of some relevance for non-relativistic shocks. Most papers are devoted to the currently very fashionable discussion of relativistic shocks which turn up in astrophysical problems as supernova remnants, gamma ray bursts and relativistic jets [for a contemporary review of most of the relevant facts on relativistic shocks cf., e.g., *Bykov & Treumann, 2011*]. Of the relevant papers, as well, very few contribute anything really substantial to the acceleration process of electrons or ions out of the thermal background at non-relativistic shocks. Most of them add a little to the diffuse shock acceleration problem discussed in length in this chapter while the new ideas are left to the relativistic domain. Because of this reason we can restrict to only a few notices on papers which we find of some interest. Otherwise the reader who is interested in the most recent, mostly numerical investigations of adding to diffusive shock acceleration is referred to the literature.

In a very recent paper *Hamza & Meziane [2011]* tried to extend the theory of the diffusion coefficient to include resonance broadening. Such broadening may be caused by electrostatic turbulence as a nonlinear effect in wave-particle interaction different from the usual quasilinear approach. This interaction might be located in the shock foreshock and in the shock transition as well, when referring to Bernstein-Green-Kruskal modes, i.e. including a possibly large number of phase-space holes. Whether this is important or not to the problem of particle acceleration by shocks remains unclear, however, even though one can, of course, heuristically define diffusion coefficients also in this case if only a sufficiently large number of holes exist in the foreshock-shock region allowing for a statistical treatment and the calculation of transport coefficients. This is probably the case as we have discussed in this book, at least under conditions of strong shocks with large cross-shock potential and/or large amounts of reflected particles.

Amano & Hoshino [2010] attacked the important problem of how electrons can be accelerated in shocks. In principle these authors pick up the idea proposed earlier that electrons can be reflected from shocks when the Mach number is large enough by the flow-loss cone mechanism. Then the electrons perform a shock drift for some time along the shock. It is argued that the reflected electrons excite whistlers upstream in the usual way thereby inducing quasilinear pitch angle diffusion and scattering and causing further shock acceleration. In this mechanism the shock-diffusive acceleration mechanism is applied as well to the electrons in this second step. The authors suggest that this mechanism might work at supernova remnant shocks whose Mach numbers are sufficiently high, in fact relativistic, for this mechanism to work. These authors derive for the critical injection Mach number the approximate formula $\mathcal{M}_A^{crit} = \frac{1}{2} \sqrt{\mu \beta_e} \cos \Theta_{Bn}$, where $\mu = m_i/m_e$ is the mass ratio and β_e the electron plasma-beta. The latter is assumed to be close to unity. This expression is a function of the shock normal angle and thus vanishes for perpendicular shocks. The theory thus applies best to quasi-parallel shocks. Indeed, quasi-parallel shocks

are believed to be strong particle accelerators in various ways. The present suggestions applies a conventional mechanism to the badly understood acceleration of electrons.

Considerable progress has, on the other hand, been achieved in the theoretical understanding of the generation of high-energy excess tails on the electron distribution function. More than a decade ago it had been proposed that in collisionless plasmas electron phase space distributions exhibiting might excessively long and flat high non-exponential energy tails of algebraic power (so-called frequently observed power law tails) may represent quasi-stationary thermodynamic equilibria [Treumann, 1999a, b; Treumann & Jaroschek, 2008]. In this case the observed power would be related to the selfconsistent turbulent level of wave turbulence. Recently Yoon [2012a, b] and Yoon *et al* [2012] demonstrated that, indeed, for electrons in interaction with self-excited Langmuir turbulence the nonlinear interaction reproduces the proposed form of the distribution function and allows for the determination of the wanted relation between the power law coefficient and the level of wave turbulence. This theory is essentially analytical [Yoon, 2012a, b] but is heavily backed by numerical simulations [Yoon *et al*, 2012]. Interestingly, it requires the solution of the full wave kinetic and turbulent Vlasov equations including spontaneous emission and absorption of the waves. its importance for electron acceleration awaits exploitation. Moreover, one may expect that in magnetised plasmas similar interaction will be at work as well as interactions which cause the acceleration of ion. Possibly, the solution of particle acceleration must be looked for in this direction.

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Final Remarks

Coming to an end always means that one is making a cut, leaving many question not asked and others not answered. We are very well aware that the present volume is incomplete in many one respect. It has been started under the premise to provide a frame for the understanding of the collisionless shocks in near Earth space at the current state of our knowledge at the time of writing (Part I and Applications Chapters 10 and 11 in spring 2008). In finishing writing one, however, realises that this goal has only partially been achieved. We have just touched some of the more interesting problems which concern the physics of shock waves in the collisionless plasmas of our surrounding space environment, giving a rough overview of coronal shocks, planetary bow shocks including bow shocks of comets, a recollection of interplanetary shocks and finally a compilation of our knowledge about the heliospheric Termination Shock. None of these presentations can be said that it would be complete. Many if not most of the questions have been left unanswered. It is our hope that we have at least formulated some of them making them available for further thought and further investigation.

Collisionless shocks pose one of the most interesting problems in classical many particle physics. First there is the confusing term collisionless. How can shocks be collisionless? We hope we have provided sufficient material for that this question can be settled even though the mechanisms of how a collisionless shock becomes a shock could only be illuminated tentatively. Neither has the mechanism become transparent of how anomalous dissipation is generated in subcritical shocks, nor could we pin down a definite mechanism of how supercritical shocks achieve a state where they are in some kind of dynamical quasi-equilibrium far from the thermal and thermodynamic equilibria. The solutions we have provided hold just under the premise that we are dealing with supercritical though still relatively low-Mach number collisionless shocks.

The philosophy implied is that these shocks, in some way, do in fact manage to reject part of the momentum and energy that is supplied by the upstream flow from passing across the shock. It has not become clear, however, why this should be so. Why should a shock *want* to avoid breaking if breaking would be the easiest way to dissipate excess flow momentum and flow energy and maintaining a thermodynamic equilibrium state? Neither have nor could we answer this question.

It seems that in the range of the shocks that are accessible to us in interplanetary space shock formation is still advantageous over breaking. But we do not know why. Possibly, because in the accessible parameter range (nonrelativistic Mach numbers, plasma β , density, magnetic field strength, temperatures...) breaking is thermodynamically not favoured as a dissipation mechanism? Possibly the magnetic field is to be blamed as it cannot easily be warped sufficiently strongly? We do not know, and the available simulations have not

been designed to investigate the competition between breaking the flow and maintaining a shock. We have assumed that it would be so and that shocks maintain a (thermodynamically unpleasant) non-equilibrium state for long, practically infinite time, at least that long as the upstream flow exists. Maybe this is a question of Mach number that at sufficiently high Mach numbers no such state can be maintained anymore and a shock will always break. At high Mach numbers on the other hand the shock becomes relativistic, electrons become heavy, radiation comes into play and, finally, particles can be created and annihilated. Relativistic shocks have been excluded from consideration from the very beginning simply because all shocks in the heliosphere are non-relativistic. It is the domain of astrophysics where relativistic shocks reign.

Rejecting inflow momentum and energy in the non-relativistic regime implies rejecting mass. Supercritical collisionless shocks reflect part of the inflow. This we have discussed in great detail for quasi-perpendicular and quasi-parallel shocks pointing out the similarities and differences which are known from observation and simulation. Only in recent years simulation techniques have developed far enough and computer power has increased sufficiently to step into this interesting field. Because of this reason we could only hint on the direction in which the physics of collisionless shocks will evolve in the near future: the investigation of the effects of realistic electron-to-ion mass ratios on shock formation and maintenance, inclusion of much larger particle numbers, design of very large simulation boxes, three spatial dimensions, high spatial and temporal resolution, embedding of local high-resolution simulations into global lower resolution simulations, adaptive grids and so on.

The use of several of these (still intermediate) steps in combination or alone will during the next decade provide us with a wealth of new insight into the physics of collisionless shocks, insight which we can only vaguely imagine at the present time. The evolution will go in the direction of investigating the instabilities that are responsible for the rejection of momentum and energy, reflection of ions and acceleration of ions and electrons under quasi-perpendicular and quasi-parallel conditions. Some questions of interest are related to the resolution of the particle state in the very shock ramp at the place where the rejection of momentum and energy and the reflection of particles takes place. What is happening there? Which are the mechanisms that stop the flow and heat the plasma while throwing a fraction of the flow back upstream? Which instabilities are involved? The investigation of the microphysics of this region will bring much new insight into the very physics of shocks. Are there microscopic electric fields involved? What is their generation mechanism? What is the role of the magnetic field on these short scales? What are the differences between the mechanisms in quasi-perpendicular and quasi-parallel shocks on these scales?

Of quasi-perpendicular shocks we know by now that some instabilities in the foot and ramp regions, like the modified-two stream instability, are of vital interest from the point of view of shock formation and maintenance. These instabilities also accelerate particles, but it is not yet known in which way and under what conditions they contribute to the production of the energetic shock particle component. Their nonlinear state poses a big question because it forces one to include the strong inhomogeneities of the plasma and fields. Acceleration of the reflected particles takes place in the foot and ramp and is dif-

ferent for electrons and ions. It is not yet known how it changes with increasing Mach number and changing plasma parameters.

Of quasi-parallel shocks we know by now that they are in continuous reformation from pulsations which are the result of the evolution of the upstream wave population which interacts with the diffuse particle component. So far the role of electrons in this processes has not been elucidated sufficiently. Nevertheless, we know that there is an interplay between the diffuse ions and the pulsations. The ions are accelerated when surfing between the old and new shock ramps and experiencing the tangential electric field. Electrons are accelerated in a different shock-surfing process which makes use of the modified-two stream instability or other still unknown or overlooked instabilities. Exchange of pulsations in quasi-parallel shocks makes the quasi-parallel shock locally quasi-perpendicular such that the convection electric field acceleration process can work and electrons can be scattered as well.

This continuous exchange of pulsations is also the generator of large amplitude downstream turbulence thereby providing the scattering centres downstream of the shock which are necessary for the Fermi mechanism to become effective. The observed compressive component is probably just the surviving pulsations which formerly constituted the shock ramp and became convected downstream after quasi-parallel shock reformation. This is the first step to an understanding of downstream shock turbulence which, we believe, will not reach a state of fully developed turbulence exhibiting power law spectra and being composed of a continuum of waves in place of a collection of single modes unless the downstream region is very broad. This latter is the case possibly only at the giant planets and behind the Termination Shock in the Heliosheath. Only there does the size of the downstream region exceed the microscopic shock width sufficiently much that the shock boundary condition does not anymore affect the state of the plasma turbulence.

Particle acceleration in shocks is the other only marginally understood problem. It seems that quasi-parallel shocks, which live from the existence of the suprathermal diffuse particle component that surrounds them and interacts with the upstream pulsations, are capable of injecting a sufficient number of particles into the Fermi mechanism. It is not yet known how this proceeds in the shock, by which mechanism and with what efficiency, but it is probable that it is by the interaction of the approaching pulsation with the trapped particle component which is involved self-consistently into pulsation generation and generation of energetic diffuse ions. In any case it seems that the shock can manage to inject enough particles to sufficiently high energies for first order Fermi acceleration to take over.

However, microscopic physics may also work inside the shock to produce some very high energy particles, in particular high energy electrons. And, at an extended shock, particles can become trapped inside magnetic loops upstream as well as downstream of the shock which have both feet on the shock surface. The trapped particles experience the oscillation of the footpoints and, by bouncing back and forth inside the loop, can in principle reach high energy. Such mechanisms may combine to accelerate particles to cosmic ray energies. This seems to be the case at the Termination Shock and could be extrapolated to astrophysics as well.

Another tantalising question is the generation of shock radiation. Astrophysical shocks radiate. Planetary bow shocks do also radiate but only in a weak sense. They emit plasma

radiation at harmonics of the local plasma frequency. Other radiation has not been detected from them, no optical radiation, no x-rays. The emission measure is simply too small for producing a susceptible radiation. Coronal shocks on the other hand radiate also in plasma harmonics. They, however, show in addition another kind of radiation which is believed to originate in the shock ramp: the so-called backbone radiation. Why do planetary bow shocks not show this kind of radiation? This question has not been answered yet while being of crucial importance for astrophysical application. Are planetary bow shocks, including the Termination Shock which is definitely not a radiator, incapable of radiating because their density is too low, their Mach numbers are too low, their magnetic fields are too low? Currently, the question remains to be unanswered.

This implies the other question of what do we see when we speak of astrophysical (non-relativistic) shocks as these are accessible to us only through the radiation they emit? At what Mach number or density or magnetic field, in short at what set of plasma parameters starts a shock to become a radiator if it remains, say, still to be non-relativistic? A simple answer is that it will start to radiate whenever there is no other way of rejecting the incoming momentum and energy densities than through the generation of radiation. In this case radiation might become the only way of getting rid of the excess momentum and energies. But when does this happen? And why? Why does the shock want to avoid breaking also in these cases? Or is breaking involved and we simply have not discovered it yet? Some of the numerical simulations show indications of shock wave breaking in the high resolution ion and electron phase space densities where some small groups of particles overturn the bulk flow and disappear, just as it is known from the bubbling cascades in breaking waves. However, we do not really know. Nor do we understand the complicated results of the simulations. Nor are the simulations informative enough. Proper design of simulations (such that their results are reliable) and proper interpretation of the simulation results (with the help of particularly tailored analytical theory) is kind of high-level art that only few simulationists are capable of.

Possibly the shock at this stage helps itself in using the excess momentum and energy densities to generate a different plasma state: it might start self-consistently exciting magnetic fields via a non-dynamo mechanism based on the Weibel instability. If the secondary Weibel-magnetic field becomes strong enough it might build up a magnetic barrier, braking the flow, diluting the plasma and providing the conditions that are necessary to ignite the electron-cyclotron maser, the strongest electromagnetic radiation process we know that works in the weakly relativistic to relativistic plasma regimes. Questions like this one cannot easily be investigated. They await their numerical treatment.

At the current state of the art we are still far away from a complete theory and understanding of all the processes that are involved into collisionless supercritical shocks even as long as they are non-relativistic. Such are the shocks in the heliosphere. Astrophysical shocks require, in most cases, a relativistic treatment which has been outside the scope of this volume. It would involve not only upstream velocities approaching the velocity of light which sometimes is mistakenly assumed. It would require the modification of the shock by its self-generated energetic particle component, inclusion of the effect of radiation on shock structure, and particle creation, thereby violating the assumption of particle number conservation. Processes of this kind require a field-theoretical approach.

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Part II

Applications:

Two Kinds of Collisionless Shocks in the Heliosphere

Introduction

After finishing the theoretical Part I of this treatise we realised that it would make sense to add a Part II on Applications. This part should contain more extended observations and measurements of the real shocks which occur in the Heliosphere and should compare these observations with theory. There are basically four types of such shocks in the Heliosphere, presumably all being non-relativistic: shocks in the solar corona, interplanetary shocks, planetary and cometary bow shocks, and the celebrated Heliospheric Termination shock. It would have made sense (and we would have liked) to describe all four types of these shocks in detail. In previous publications [e.g. the AGU monographs 34–35 on “Collisionless Shocks in the Heliosphere”, edited by R. G. Stone and B. T. Tsurutani, AGU, Washington D.C. 1985, and the less systematic Cospar Conference Proceedings “Physics of Collisionless Shocks”, Adv. Space Res. 15 (8/9) Elsevier-Pergamon 1995, edited by C. T. Russell, and AIP Conference Proceedings Vol. 781, edited by G. Li, G. P. Zank and C. T. Russell, AIP, Melville N.Y. 2005] this has been done to the extent of the information which, at those times, had been available. Though it has originally been our intention to properly and as possible completely review the material and knowledge which during the past two decades has been accumulated on shocks in the heliosphere, in this second part on Applications we just cursory deal with only two of the heliospheric shock families: planetary bow shocks and the Termination Shock. The reason for this restriction is twofold.

Firstly, the accumulated observational material is still in fast evolution, in particular what concerns solar coronal and interplanetary shock waves, causing the field to change quickly. We therefore just refer to the knowledge which, from our conservative point of view, can be considered as established. As for an example, recent claims attributing the generation of coronal shocks to something called solar tornados below and in the lower solar atmosphere, indicate that this field evolves rapidly and is subject to sudden changes. The notion that such a kind of *Solar Meteorology* might be responsible for the generation of shocks and, more generally, causes the observationally well established CMEs, mass ejections from the solar corona, is still debatable. People tend to quickly coin new paradigms, models, the validity of which is disputed but in this case, if confirmed, might substantially change our view of shock wave generation in the corona. Moreover, the lower corona is by no means collisionless. Hence the generation and properties of shocks emanating from it has little in common with the really collisionless bow shocks, interplanetary shocks, and the Heliospheric Termination shock. Coronal shocks are, in addition, strong radiators in the radio and also in x rays. This distinguishes them in another way from the collisionless shocks in question even though bow shocks are also very weak radio emitters, as we have described in Part I. These facts have led us to exclude coronal shocks from this review.

Secondly, limited time inhibits a complete survey of the observations of the heliospheric shocks. For this reason we have omitted the wide field of interplanetary shocks from inclusion into the present Part II as well. Interplanetary shocks may have several reasons for being produced in the heliosphere. The simplest one is that they are related to coronal shocks which manage to propagate up from the corona into interplanetary space. Another reason is the ejection of large amounts of matter from the corona into interplanetary space like CME's which, being sufficiently fast with respect to the solar wind, drive blast shock waves in front of the ejected and expanding matter. Since the matter behind behaves like a material piston of complicated ionic composition, these shocks behave differently from bow shocks. They require a particular treatment. Moreover, interplanetary shocks may also evolve from instabilities in the collisionless solar wind, in which case they grow nonlinearly from fluctuations, steepen and become either weak or, in particular cases, strong shocks. The majority of these shocks are subcritical in this case but can also become super-critical if growing and steepening from waves which in the solar wind frame propagate upstream. In a stationary observer's frame they would lag behind the solar wind in this case. Such shocks can well be understood on the basis of the theory in Part I. In general they fit quite well into the picture developed, sometimes with the additional even simplifying property that they are large scale and thus close to being planar shocks, i.e. they just are examples of one particular kind of shocks, depending on their shock normal angle Θ_{Bn} .

In the following Chapter 10 we briefly describe the properties of the known planetary and cometary bow shocks. This chapter has been finished in Spring 2008. For reasons of time we didn't update this chapter. Too large a number of observational papers would have to be included. However, in essence, most of them only illuminate very particular questions on the shock structure which do not go far beyond the knowledge reviewed here. In particular, no paradigm change is in sight anywhere, such that the cursory review given in Chapter 10 can be taken as a grossly still valid state of our current knowledge which can and will in the near or more distant future be completed by new *in situ* observations.

Chapter 11 contains the knowledge on the Heliospheric Termination Shock as of Spring 2008 when the Voyager spacecraft were crossing it in a number of most spectacular and internationally celebrated events. We, at that time, included the most recent observations and discussed them in the light of the then available theoretical knowledge. Since progress has been made since, based on the Voyager observations during the Termination Shock crossings, after the crossing in the Heliosheath, and also in data analysis and interpretation, we have added a short update to this chapter without changing the philosophy of our initial writing. The material given can be taken as background for further development when referring to more recent or future observations and theoretical calculations. Among the most interesting new conclusions is that the magnetic field, i.e. the interstellar magnetic field and the compressed magnetic field of the heliospheric current sheet in the Heliosheath seem to play a much more important role than was believed so far. Claims have been published that the compressed current sheet magnetic field in the Heliosheath near the nose of the Heliosphere undergoes violent reconnection, independent of the Termination Shock, which should lead to isotropisation of the field and particle distribution, as well as to plasma heating. It is also asserted that, possibly and contrary to original

belief, the Heliosphere does not cause a Helio-bow shock in the surrounding interstellar medium. The relative velocity between interstellar gas and Heliosphere seems to be substantially less than previously assumed, while the Heliosheath plasma has become more thermalised and about stationary. In this respect one may look forward to further spacecraft observations and numerical simulations including composition, magnetic fields and reconnection which will provide deeper insight into the real conditions at the outer bounds of the Heliosphere.

Planetary Bow Shocks

Abstract. Our present knowledge of the properties of the various planetary bow shocks is briefly reviewed. We do not follow the astronomical ordering of the planets. We rather distinguish between magnetised and unmagnetised planets which groups Mercury and Earth with the outer giant planets of the solar system, Mars and Moon in a separate group lacking magnetic fields and dense atmospheres, and Venus together with the comets as the atmospheric celestial objects exposed to the solar wind. Asteroids would, in this classification, fall into the group together with the Moon and should behave similarly though being much smaller. Extra-solar planets are not considered as we have only remote information about their behaviour. The presentation is brief in the sense that our *in situ* knowledge is rather sporadic yet, depending on just a countable number of bow shock crossings from which just some basic conclusions can be drawn about size, stationarity, shape and nature of the respective shock. The only bow shock of which we have sufficient information to deal in sufficient depth with its physics is Earth's bow shock. This has been reviewed in other places in this volume and therefore is mentioned here only as the bow-shock-paradigm in passing.

10.1 Introduction

Among the shocks in the heliosphere planetary bow shocks occupy an exclusive position due to their comparably easy accessibility. It is in particular the Earth's bow shock that has provided the deepest insight into the structure and formation of collisionless shocks. Much of the physics that was learnt from its continuous observation by spacecraft crossing it multiply has been dealt with in the chapters of Part I in this volume. In the present chapter we provide a brief review of some of the new developments in planetary bow shock physics that have not been mentioned yet and go beyond the existing excellent review paper of *Russell* [1985] and the short and more specialised papers contained in Section 6 on Planetary Shocks in *Russell* [1995].

When speaking about planetary bow shocks we have in mind all the different kinds of shocks which are generated when a heavenly solid obstacle is put into the high speed flow of the solar wind. Since the dominant such obstacles are the planets, planetary bow shocks also dominate this selection even though they are not just the most abundant in the solar system and heliosphere. The more abundant ones are the bow shocks around the many solid asteroids which occupy some parts of interplanetary space sometimes passing the vicinity of Earth in the inner region of the solar system, each of them creating its own bow shock.

These asteroid bow shocks are, however, of small sizes and thus do not attract our main attention. Since the asteroids do not have atmospheres and at least to our knowledge no magnetic fields either, their bow shocks are of similar nature as the bow shocks of other

non-magnetised solid bodies in the solar system like most planetary satellites, in particular like the moon, and can thus be considered to be smaller relatives of the lunar bow shock. There is, however, one distinction which is indeed related to the very size of these objects. The physics of their bow shocks changes completely relative to those like the lunar bow shock if the size of the asteroid becomes comparable to the ion inertial length. In this case – possibly – no bow shock can evolve at all as from the plasma point of view the obstacle (asteroid) becomes a very heavy, possibly charged and (due to the solar wind induction effect) magnetised very large “particle”.

Other heavenly bodies which pass the solar system are comets. These are non-magnetic but gaseous objects. As a consequence the bow shocks developing around them differ from those of the magnetised planets but have much in common with non-magnetic planets possessing dense atmospheres like Venus with the main difference of the lacking gravity, the completely different composition of the atmosphere, and the size as comets are bodies that are as small as asteroids. Nevertheless, from the plasma point of view they cannot be considered to behave like particles just because of the out-gassing of their atmospheres which forms an atmospheric gaseous cloud around the cometary body that is several orders of magnitude larger in extent than the cometary body itself.

In the following we will, therefore, consider three classes of planetary bow shocks: bow shocks around magnetised planets, cometary bow shocks including the Venusian bow shock, and the lunar bow shock (as a paradigm for other non-magnetised non-atmospheric bow shocks).

We are interested in the physics of shocks here. In the present case this means that we are interested in the physics of bow shocks under the various differing external plasma and planetary atmospheric conditions. Their shapes and macroscopic properties have been reviewed elsewhere. While the available information about the known planetary bow shocks decreases steeply with distance from Earth, the presentation will naturally be rather short, sticking just to the few conclusions that can be drawn from the sporadic measurements that are available to us from the few spacecraft crossings *in situ* of the more remote planetary bow shocks.

Earth’s bow shock is the only bow shock that has been sufficiently monitored to conclude about the formation and internal physics of a bow shock. This has been made use of in various different places in this volume (and in the earlier publications). Here, we are treating Earth’s bow shock just as the paradigm of a magnetised planetary bow shock, and this we will do only briefly. We are less interested in the shapes and gasdynamic properties of bow shocks as they do not tell us very much about the structure, internal dynamics and dissipation processes acting inside a bow shock. Shapes and gasdynamic properties of bow shocks have been reviewed in depth in other papers [see, for instance, *Slavin et al*, 1985; *Russell*, 1985, 1995; *Spreiter & Stahara*, 1995, among others], based on similarities and scaling to Earth’s bow shock and the (relatively meagre) statistics available from the small number of planetary bow shock crossings. Figure 10.1 provides an overview of some of the planetary bow shock shapes properly scaled to their stand-off distances from the centres of their respective planets as obtained from the assumption of similarity to Earth’s bow shock and the available knowledge from the locations of the sporadic spacecraft bow shock crossings.

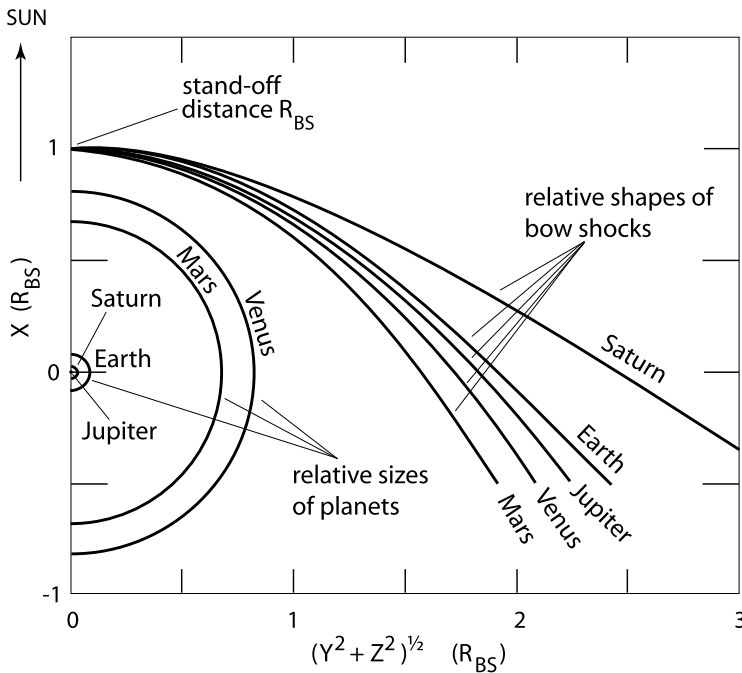


Figure 10.1: The semi-theoretical shapes of some planetary bow shocks. Length scales have been normalised to the stand-off nose distances R_{BS} of the corresponding planetary bow shocks. In this representation the sizes of the planetary bodies are inverted. Jupiter shrinks to a point at the origin while the small planets Mars and Venus become grossly oversized. Mercury is not yet included here. Of the outer giant planets only Saturn is shown. [data taken from *Slavin et al, 1985*].

10.2 Terrestrial Type Bow Shocks

The terrestrial planets are Earth, Mercury, Venus, Mars, and – though not a planet – a body of similar consistence: the Moon. This classification applies to their interior structure, composition and densities. However, from the bow shock viewpoint the classification does not follow this scheme. Here one rather distinguishes the magnetised and non-magnetic planets which leads to group together Earth with Mercury and the giant magnetised planets, Jupiter, Saturn, Uranus, and Neptune. Surprisingly one of the Jovian moons, Ganymede, has an own magnetic field and could also be included. However, like the other satellites of the giant planets including Titan, Ganymede never leaves Jupiter's magnetosphere to enter the solar wind. It possesses a magnetosphere which in the Jovian magnetospheric plasma flow stands up as an obstacle. Therefore a bow shock he would possess would be of purely Jovian origin and required super-Alfvénic flow in the Jovian magnetospheric plasma. Since this could only occasionally be the case we do not include Ganymede or any other close planetary satellite (with the exception of the Moon) into this collection.

10.2.1 Earth's Bow Shock

The paradigm of a planetary bow shock embedded into a super-magnetosonic stellar wind flow is the Earth's bow shock wave. Its continuous observation has fertilised the understanding of collisionless shock physics like no other shock neither in space nor in the laboratory. The fortunate properties of the bow shock – we are simply speaking of the bow shock if speaking of Earth's bow shock – are multitude: it is a supercritical shock, it is a magnetised shock, it is a curved shock, and because of this latter property it at one and the same time possesses regions of perpendicular, quasi-perpendicular, oblique, and quasi-parallel nature permitting to study all the properties that these different shocks develop separately and in their relation. In addition it is a sufficiently extended shock of tangential diameter $\gg \Delta_s, \lambda_i, r_{ci}$ that by far exceeds its width Δ_s , the ion inertial length λ_i and the ion gyro-radius r_{ci} .

Macroscopically the bow shock reacts sensitively to the variations in the upstream flow, the upstream turbulence, direction of the interplanetary magnetic field and to macroscopic disturbances in the upstream flow thereby permitting to study these dependencies and variations and the different states of a collisionless shock that are caused by the variations in the upstream plasma parameters like the respective speeds, temperatures and densities of the flow, the composition of the upstream plasma, and the upstream state of turbulence. It also allows to study the downstream properties of the shocked plasma behind the different quasi-perpendicular and quasi-parallel parts of the shock, the development and properties of the downstream turbulence, particle dynamics and the downstream mixing state. This is due to the blunt nature of the obstacle that is responsible for the existence of the bow shock, Earth's magnetosphere with its comparably strong geomagnetic surface magnetic field of $B_E(1 R_E) \sim 4.3 \times 10^4$ nT that is much larger than the average solar wind magnetic field $B_{SW}(1 \text{ AU}) \sim 10$ nT at Earth's orbit and therefore exerts a strong resistance to the solar wind flow, digging a large blunt-nosed cavity into the solar wind flow and forcing it to deviate from its original radial heliospheric direction to turn around this cavity, the magnetosphere. This bluntness of the magnetosphere gives sufficient space in upstream direction between the magnetosphere and bow shock for the accumulation of shock disturbed plasma.

The bow shock that is created in this interaction between the solar wind flow and the geomagnetic field is located at an upstream distance that follows from hydrodynamic theory as the envelope of the characteristics of waves that can propagate upstream against the flow in the time the flow needs to pass from the shock to the magnetosphere. For the bow shock this distance at the nose of the bow shock in solar direction is in the average found at a geocentric radius of $R_{BS} \sim (12-14) R_E$. It is of approximately hyperbolic shape widening with distance from the Earth-Sun line in such a way that its radius at the terminator becomes approximately $\sim 27 R_E$.

The relatively large range of variation in distance is due to the variation in the solar wind conditions. A simple formula that approximately describes the bow shock nose location distance $\Delta_d = R_{BS} - R_{MP}$ from the nose location of the magnetosphere R_{MP} , i.e. the magnetopause distance along the Earth-Sun line at local noon, as function of the shock density compression ratio N_2/N_1 has early been given by *Spreiter et al* [1966] from their

gasdynamic calculations as

$$\Delta_d \approx 1.1 \frac{N_1}{N_2} R_{\text{MP}}, \quad \frac{N_1}{N_2} \approx \frac{\gamma+1}{\gamma-1} - 2 \frac{\gamma+1}{\mathcal{M}^2(\gamma-1)^2} \quad (10.1)$$

where γ is the adiabatic index, and $\mathcal{M}$ the (magnetosonic) upstream Mach number, which in the hydrodynamic approximation used is the ordinary gasdynamic Mach number. The upstream magnetic field is not included here since the magnetic pressure in the upstream solar wind is small and in this early theory not much was known about the quasi-parallel and quasi-perpendicular differences between shocks. However, what concerns the nose distance the magnetic effect on the global form is not so important. It is usually incorporated into a constant factor $K \sim O(1)$ in the definition of the magnetopause nose distance R_{MP} which is defined in terms of the planetary radius R_{P} for a planetary dipole magnetic field of surface strength B_{P} through the magnetic and dynamic pressure balance as

$$R_{\text{MP}} = \left(\frac{KB_{\text{P}}^2}{m_i \mu_0 N_i V_1^2} \right)^{\frac{1}{6}} R_{\text{P}} \quad (10.2)$$

The factor in parentheses is the square of the Alfvén speed based on the planetary surface magnetic field, the upstream solar wind density N_i and the upstream velocity V_1 . The sixth root makes this distance only weakly dependent on the factor K . The relative variation of the distance Δ_d between the magnetopause and bow shock is given by

$$\frac{1}{\Delta_d} \frac{d\Delta_d}{dP_{\text{SW}}} = - \frac{1}{6P_{\text{SW}}} \quad (10.3)$$

where $P_{\text{SW}} = m_i N_i V_1^2$ is the ram pressure of the solar wind.

The above formulae suggest the intuitively reasonable result that an increase in the compression ratio will push the bow shock closer to Earth. For the largest MHD compression ratio of $N_2/N_1 = 4$ this distance is $\sim 0.275 R_{\text{MP}}$. At the Earth with $R_{\text{MP}} \approx 10 R_{\text{E}}$ the bow shock nose is at a geocentric distance of $R_{\text{BS}} \approx 12.75 R_{\text{E}}$. Because of the assumption of maximum compression ratio this is somewhat at the lower edge of the average distance noted above. But the above handy formula is not solely restricted to Earth's bow shock; it can be applied to any other planetary bow shock as long as the magnetic effects on the shape and dynamics of the shock can be neglected which is the case as long as $\beta_1 \ll 1$, where $\beta_1 = \mu_0 m_i N_1 V_1^2 / B_1^2$ is the non-relativistic upstream plasma β that is based on the kinetic energy of the flow. Figure 10.1 shows the semi-theoretical relative shapes of the bow shocks for some planets where all lengths have been normalised to the theoretical stand-off distances of the corresponding planetary bow shocks. In this representation the sizes of the various planets are inverted. Jupiter shrinks to a point at the origin while Mars and Venus become oversized.

These expressions can be applied to any planet immersed into the solar wind stream under the condition that its magnetic field is to good approximation a dipole field. The solar wind velocity V_1 is a constant throughout the solar planetary system. However, the solar

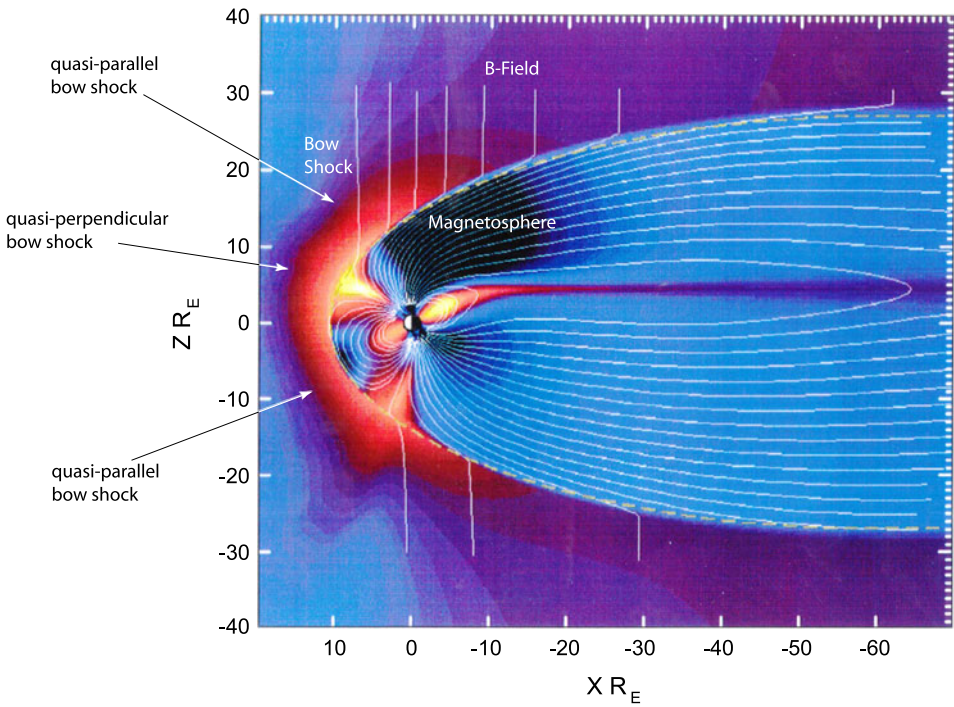


Figure 10.2: View of the whole magnetosphere and bow shock system embedded in the solar wind stream. The figure is obtained by application of the empirical magnetospheric Tsyganenko-model that is based on observations. Shown is the cross section that includes the Sun-Earth line, terminator and rotation axis of Earth for summer on the northern hemisphere and for a North-South directed magnetic field vertical to the ecliptic being completely reconnected to the magnetospheric magnetic field. The Sun is on the far left. The bow shock is the sunward edge of the *red region* on the left in front of the magnetosphere. Colouring indicates plasma temperature/thermal pressure with *red* highest and *dark* lowest.

wind density changes about like $N_i \propto R^{-2}$ such that the solar wind dynamical pressure decreases with heliocentric distance. In application to other planets it is therefore convenient to measure R in units of 1 AU and to use the average solar wind density at Earth's orbit as reference. This density is in the average $N_i = 5 \times 10^6 \text{ m}^{-3}$.

Figure 10.2 gives an impression of the bow shock-magnetosphere system as obtained from the empirical data based Tsyganenko-model for the particular case of northern hemispheric summer and strictly north-south directed solar wind magnetic field such that the field is completely reconnected to the geomagnetic field inside the magnetosphere. The solar wind is streaming from the left against the geomagnetic field creating the magnetosphere and causing the bow shock in front of the magnetosphere.

The figure shows the asymmetry of the bow shock and magnetosphere that is caused by the inclination of Earth's rotation axis against the ecliptic. Slowing down of the solar wind flow and down-tail re-acceleration of the plasma is recognisable from the relative distance between adjacent magnetic field lines. Behind the bow shock the distance between

the field lines is small. When the downstream plasma couples newly to the solar wind the distances between the field lines increase along the magnetospheric tail in anti-solar direction. Colouring indicates the plasma temperature and thermal pressure. It shows the transition from the cool solar wind flow across the shock to the hot and dense plasma behind the shock in the magnetosheath, with the bow shock being the boundary between both plasmas, itself being responsible for the heating and compression of the plasma and magnetic field.

The figure also shows that even in this case when the magnetic field is exactly perpendicular to the ecliptic plane the bow shock is bent in a way that two large regions of quasi-parallel shocks surround the nose region which is quasi-perpendicular. This property of the bow shock and its practically continuous and uninterrupted presence make it so attractive for the investigation of the properties of supercritical collisionless quasi-perpendicular as well as quasi-parallel shocks. This is why most of our knowledge of the structure and physics of collisionless supercritical shocks has been stimulated by investigations of the bow shock. In the following we briefly describe the global conditions prevalent at the terrestrial bow shock before turning to the discussion of the differences and similarities between the Earth's bow shock and the bow shocks encountered near the other heavenly bodies in our solar system.

Solar Wind Conditions

Compared to Earth's bow shock the information obtained on the remaining planetary bow shocks is still sparse. In order to be able to classify the global properties of the true planetary bow shocks, i.e. the bow shocks of the planets and their satellites, one needs to know the respective upstream solar wind conditions. In other words one needs to know the variation of the main properties of the solar wind with heliocentric distance. These are given in Table 10.1 for the radial heliocentric distances R of the planets measured in AU. Mercury, the planet with orbit closest to the Sun, has two entries because of its large orbital eccentricity. The decrease in solar wind ram pressure P_{SW} with heliocentric distance is due to the $N \sim R^{-2}$ solar wind dependence on radius. $\mathcal{M}_s$ and $\mathcal{M}_A$ are the respective sonic and Alfvénic Mach numbers, $\beta = 2\mu_0 NT/B^2$ is the ratio of thermal to magnetic pressures. Since $B \sim R^{-1}$, the radial dependence of β is mainly determined by the radial variation of the solar wind temperature $T = T_i + T_e \approx T_e$; outside the Mars orbit the solar wind cools about adiabatically. Finally, the spiral angle of the interplanetary magnetic field $\mathbf{B}$ against the radial direction of the solar wind flow in the ecliptic plane increases gradually to become about perpendicular to the solar wind already at the orbits of the outer giant planets. This table is built on both theory and observation by the Helios and Mariner 10 spacecraft.

In addition the solar wind exhibits a susceptible level of turbulence, fluctuations in the magnetic field, density and flow velocity, dependence on the solar cycle and the well known magnetic sector structure. Occasionally, when the planet passes a sector boundary the magnetic field direction may change abruptly and the planetary bow shock may be subjected to the passage of a current sheet which affects the structure of the planetary bow shock and creates effects like the formation of High Flow Anomalies (HFA). Moreover,

Table 10.1: Solar wind conditions upstream of the planets

Planet	R (AU)	Spin $T_p^\dagger$ (hrs)	$10^{-9}P_{\text{SW}}$ (J m $^{-3}$)	$\mathcal{M}_s^*$	$\mathcal{M}_A^*$	β	$\mathbf{B}_{\text{SW}}$ -angle ($\alpha_{\text{Parker}}^\circ$)
Mercury	0.31	1407.50	26.5000	5.5	3.9	0.5	17
	0.47		11.0000	6.1	5.7	0.9	25
Venus	0.72	−5832.50	5.0000	6.6	7.9	1.4	36
Earth	1.00	23.93	2.5000	7.2	9.4	1.7	45
Moon		655.73					
Mars	1.52	24.62	1.1000	7.9	11.1	2.0	57
Jupiter	5.20	9.93	0.0920	10.2	13.0	1.6	79
Saturn	9.60	10.49	0.0270	11.6	13.3	1.3	84
Uranus	19.10	−17.24	0.0069	13.3	13.3	1.0	87
Neptune	30.20	16.11	0.0027	14.6	13.3	0.8	88

* Respective sonic (subscript s) and Alfvénic (subscript A) Mach numbers.

† Spin period of planet in terrestrial hours. The sign indicates direction of rotation.

Solar wind data taken from *Slavin & Holzer [1981]*.

Coronal Mass Ejections (CMEs) cause interplanetary shocks to pass over the planet and interact with the bow shock. These effects are transient, however, temporarily changing the character, type, shape and structure of the bow shock.

Properties

The Earth's bow shock has served as the paradigm of a supercritical shocks in the previous chapters of this volume. We may briefly summarise its various properties as follows:

The Earth's bow shock is a permanent feature of the solar wind-magnetosphere system, standing in front of the magnetosphere at an average geocentric nose distance of $R_{\text{BS}} = (12-14) R_E$, a distance that may change substantially depending on the variations of the solar wind pressure, i.e. depending on the compressional state of the Earth's magnetosphere. Occasionally the bow shock can be found at a distance as close as $\lesssim 10 R_E$ or also much farther out at nose distances $\gtrsim 15 R_E$. Its shape is about hyperbolic forming a shield in front of the magnetosphere that opens up to the flanks and poles. Since its shape is convex around the magnetosphere when looked at it from the Sun, it has at least one spot on its surface where the solar wind magnetic field is tangential and the shock is perpendicular, respectively quasi-perpendicular. The location of this perpendicular spot and the quasi-perpendicular part around it depends on the spiral angle of the magnetic field. Large parts of the bow shock towards its flanks are quasi-parallel. The bow shock is highly supercritical with Mach number $6 \lesssim \mathcal{M} \lesssim 12$, occasionally even larger. It is a strong shock in the sense that the shock compression ratio is close to what is known to be the maximum of the compression ratio of a supercritical MHD shock $2.5 < N_2/N_1 \sim B_2/B_1 < 4$. This ratio is based on the values of the fields well in front and well behind the shock. When using the overshoot fields the compression ratio can be higher than the MHD maximum. This is attributed to kinetic effects that take place in the shock ramp and transition region on

scales $< \lambda_i$, shorter than the ion inertial length where electron kinetics comes into play. The width of the shock transition is of the order of $1 \lesssim \Delta_s / \lambda_i \lesssim (3-4)$, depending on whether one is talking of the quasi-perpendicular or quasi-parallel bow shock. The width of the quasi-perpendicular shock ramp is usually at the lower end of this interval, $\Delta_{\text{ramp}} \sim \lambda_i$ or less.

The quasi-perpendicular bow shock reflects a substantial amount of ions back upstream into the solar wind and already retards the solar wind by a fraction of $\sim (10-30)\%$ in velocity. It possesses a foot region which has been measured to be slightly broader than one ion gyro-radius based on the upstream field, $\Delta_{\text{foot}} \sim r_{ci}$. This foot is caused by the current carried by the reflected particles in the foot which are subjected to acceleration in the upstream motional electric field that is locally tangential to the shock surface. Observationally it has not yet been identified to what extent the reflected, heated and accelerated electrons contribute to this current. It is also not precisely known which instability is responsible for the shock reformation. Full particle simulations with realistic mass ratio suggest – without experimental proof – that the instability is dominated by the electrons, being of the type of modified two-stream instability. Reformation might then be driven by electrons which for high Mach numbers can become reflected from localised electric fields in combination with the magnetic mirror force.

Such localised fields have been observed only recently to be of the order of a few 100 mV/m yielding localised potential wells $\Delta\phi \sim \text{kV}$ in the ramp and overshoot region. They are capable of reflecting and accelerating charged particles. A substantial fraction of such localised potential drops of the order of few 10 mV/m is found already in the foot region, but the main place where they exist is in the shock ramp and overshoot. From simulations, the overshoot is known to be the region where the strongest ion reflection occurs and the upstream ion flow mixes into the downstream hot ions while in the ramp ion phase space vortices develop where ions are locally trapped and which are related to oscillations in the magnetic field with the magnetic maxima occurring at the places of the vertexes of the vortices. So far the particle observations have not been able to resolve similar structures. These are however expected in the regions of the localised electric fields which in the plasma wave spectra appear as broadband noise reaching from low frequencies near the ion plasma frequency up to well above the electron plasma frequency.

The quasi-parallel bow shock possesses an extended electron foreshock and a broad ion foreshock which are parts of the shock transition.

The ion foreshock contains upstream ULF waves which at large distance from the shock are quasi-periodic but increase in amplitude when approaching the shock transition. Several types of such waves have been observed. The quasi-periodic waves are mainly Alfvén-ion cyclotron waves having upstream directed phase velocity, magnetosonic waves evolve into shocklets, and some of the waves evolve into large amplitude pulsations of irregular shape implying a broadband magnetic spectrum. They are generated when the upstream waves are convected downstream toward the shock and start resonantly interacting with the upstream diffuse ion component that is produced at the shock. Approaching the shock the density of diffuse super-thermal ions increases about exponentially causing the interaction to readily become nonlinear, causing the pulsation wave amplitudes to grow and steepen during the downstream convection toward the shock ramp. In fact each

pulsation when becoming sufficiently steep and large amplitude starts behaving like a part of the shock, reflecting some ions and generating other waves, e.g. standing whistlers, in front of the pulsation.

The large amplitude pulsations have been found of being the generators of the shock in quasi-parallel shock reformation. Quasi-parallel shock reformation works by replacing the shock front with newly arriving large amplitude pulsations. This is a local process since each pulsation is a three-dimensional structure of relatively irregular shape, limited width in the direction normal to the shock and limited extension in both directions along the shock surface. This makes the quasi-parallel bow shock highly irregular and variable in time and space.

Production of the diffuse super-thermal energetic particles and the very mechanism of generating the pulsations is not yet understood satisfactorily. The parallel shock is capable of accelerating ions nonadiabatically, injecting them into the shock environment where they may become the seed population for further acceleration by the first order Fermi process. This might be done by a shock surfing mechanism which requires trapping of the ion component at the shock transition. At the bow shock the threshold energy for ions for stepping into the Fermi acceleration is in the range of $\sim(50-100)$ keV/nucleon. The particle trapping and acceleration probably occurs by trapping of ions between the old shock ramp and the new pulsation. Acceleration of particles takes then place when the trapped ions feel the motional tangential electric field which is continuous across the shock for the time of trapping during the approaching pulsations. Probably electrons are accelerated in a similar way since it has been found that the upstream waves and in particular the pulsations with their transverse magnetic fields cause a local turning of the shock normal in the perpendicular direction. The quasi-parallel bow shock probably behaves locally like a quasi-perpendicular shock on distances $< \lambda_i$ ahead of the nominal shock ramp. In addition electrons are accelerated in the shock ramps by microscopic non-adiabatic processes leading to electron hole formation and electron heating.

The electron foreshock on the other hand contains electron beams that are reflected from the tangential part of the shock. These electron beams excite Langmuir waves and cause electromagnetic radiation from the shock. Both has been observed from remote as well as in situ of the bow shock. The Langmuir waves are restricted to the magnetic flux tube that is tangential to the shock surface. They thus indicate that the region of observation is magnetically connected to the shock. The radiation, on the other hand, can be observed from remote. It is usually emitted at the harmonic $\omega \approx 2\omega_{pe}$ of the local plasma frequency at the location of the excitation of Langmuir waves. It thus carries information about the plasma density upstream of the shock.

Many of the processes noted here have not yet been fully understood even though the Earth's bow shock is the best investigated shock in the entire solar system. Some of the claims made above are even speculative and are based on the combination of insufficient observations combined with insufficient (one-dimensional or hybrid) simulations which do not yet describe the complex and interwoven processes which act in forming the bow shock.

10.2.2 Mercury's Bow Shock

Mercury, the smallest and closest to Sun planet, is in several respects similar to Earth and different from the other planets. Relative to its size Mercury has the largest iron core and the thinnest mantle of all terrestrial planets. Because of this reason we treat its bow shock at this place first after Earth's bow shock in order to demonstrate the similarities and differences.

Mercury might have developed into another Earth if it would not have been so close to the Sun where it was exposed to a continuous bombardment by other heavy celestial bodies that were attracted by the Sun. Its surface speaks of the many collisions it had to survive. However, being on a stable orbit, it did quite well despite that in these collisions Mercury has been skinned losing its crust and part of the mantle such that practically only its iron nucleus surrounded by a comparably thin mantle shell was left over.

Its surface also speaks of the strains Mercury experienced during its evolution and still experiences on its extremely eccentric orbit around the Sun which the Sun exerts on it, driving tectonic activity on Mercury. Still, Mercury managed to maintain a weak interior magnetic field the origin of which is not yet clear and gives rise to speculations about a surviving dynamo with contributions that are driven by induction from the outside. The magnetic moment of Mercury is about $<1\%$ of that of the Earth. Mercury's dipole moment is $M_{\text{Merc}} = 4 \times 10^{-4} M_{\text{E}} = (4.9 \pm 0.2) \times 10^{21} \text{ nT m}^3$ [Ness, 1978]. These and other questions have been discussed in a recent monograph on Mercury [Balogh *et al*, 2008]. The internal dynamo Mercury possesses may be the result of surviving tectonic activity or is kept alive by the strains.

Mercury's magnetic field though weak surrounds the planet and interacts with the solar wind which is quite strong at the Mercury orbit (see Table 10.1). Thus Mercury behaves like a magnetic planet and necessarily develops a bow shock in front of its magnetosphere. So far there have been only three spacecraft encounters with Mercury that provided magnetic measurements near Mercury, two of them encounters of Mariner 10, and one recent encounter of MESSENGER. Mariner 10 found the magnetopause solar wind stagnation point (the nose distance of the Mercury magnetosphere) at a planetocentric distance of $R_{\text{MP,Merc}} = (1.85 \pm 0.15) R_{\text{Merc}}$ (the Mercury radius being $R_{\text{Merc}} = 2439 \pm 1 \text{ km}$), which is very close to the planetary surface and leaves very little space between the bow shock and magnetopause distances at the nose. Applying the canonical formula Eq. (10.1) gives for this distance just $\Delta_{d,\text{Merc}} = 1240 \text{ km}$. (Note also that the determination of the magnetopause and shock distance is quite uncertain, in Table 10.2 the values given are smaller even though they are based on the same data of Mariner 10. This is caused by a high uncertainty of the Mercury magnetic field, the variability of the solar wind, Mercury's orbital eccentricity and the location where the spacecraft crosses the bow shock which implies the assumption of a magnetospheric model.) Compared to the solar wind ion gyro-radius of $r_{ci} \sim 300 \text{ km}$ at Mercury this Mercury magnetosheath transition region is only $4r_{ci}$ wide. Mariner 10 passed the bow shock at the flanks where the distance between bow shock and magnetopause is large enough to distinguish the two boundaries in the magnetic field recordings. From the behaviour of the magnetic field at bow shock encounter

Table 10.2: Relevant planetary properties

Planet	R_P (km)	M_P ($T m^3$)	M_P/M_E	$ \mathbf{B}(R_P) ^{\ddagger}$ (μT)	$\mathbf{B}$ -Tilt (θ°)	R_{MP} (R_P)	R_{BS} (R_P)
Mercury	2440	$3.6 \cdot 10^{12}$	$\sim 4 \cdot 10^{-5}$	~ 0.26		$< 1.2^\dagger$ $1.37^\dagger$	~ 1.5
Venus	6051	$1.26 \cdot 10^{11}$	$1.3 \cdot 10^{-4}$	0.001		0.56^*	
Earth	6378	$9.9 \cdot 10^{16}$	1	31	+11.3	~ 11	
Moon	1738	$1.6 \cdot 10^9$	$1.6 \cdot 10^{-8}$	< 0.0003		0.23^*	
Mars	3397	$2.6 \cdot 10^{12}$	$2.7 \cdot 10^{-5}$	< 0.0005			$\sim 1.5^{**}$
Jupiter	71492	$1.56 \cdot 10^{20}$	20000	428	-9.6	50–100	
Saturn	60268	$4.6 \cdot 10^{18}$	580	22	0	16–22	
Uranus	25559	$3.9 \cdot 10^{18}$	50	23	-59	~ 18 –25	$\lesssim 33$
Neptune	24764	$2.2 \cdot 10^{17}$	27	1.42	-47	23–26	$\lesssim 34$

† Variation due to the large Mercury eccentricity.

‡ Surface field strength.

* No global magnetopause; global magnetic data are upper limits.

** Mars Global Surveyor observations [Acuña *et al.*, 2001].

Giant planet data from Bagenal [1992].

it was concluded that one of the bow shock passages was at a quasi-perpendicular bow shock, the other happened to cross a more quasi-parallel shock [Ogilvie *et al.*, 1974; Ness, 1978].

The recent MESSENGER passage [Solomon *et al.*, 2008; Anderson *et al.*, 2008; Slavin *et al.*, 2008] gave a similar result. Figure 10.3 shows the magnetic field along the MESSENGER path across the Mercury bow shock-magnetosphere system. As before Mariner 10, the MESSENGER spacecraft detected a magnetosphere around Mercury. It crossed the bow shock and magnetosheath twice at different locations. This is illustrated in the figure by the use of a fluid model simulation [Kabin *et al.*, 2000]. On its inbound path the spacecraft entered from a quiet solar wind across a nearly perpendicular flank bow shock which exhibited a small foot region in front of a weak shock of compression ratio < 2 , afterwards staying in a broad flank magnetosheath. The sheath region adjacent to the shock was quite calm, large amplitude oscillations in magnetic field and direction occurred only close to the magnetopause encounter and are probably standing mirror modes known from Earth's magnetosheath. The fluctuations are seen in the polar and azimuthal angles of the magnetic field. The amplitude of low frequency fluctuations < 10 Hz in the magnetic field are plotted in the lowest panel. Interestingly the shock crossing does not occur as a spectacular event in these fluctuations.

The outbound bow shock crossing exhibits a highly fluctuating magnetosheath and another sharp much stronger shock crossing of compression factor ~ 3 and a large overshoot. There are strong fluctuations in front of this shock crossing and a more disturbed solar wind following the crossing. However, the nature of this crossing is interpreted as being also more quasi-perpendicular than quasi-parallel with an extended active shock foot region in front of the shock ramp. (Possibly, however, this is not a foot – as the

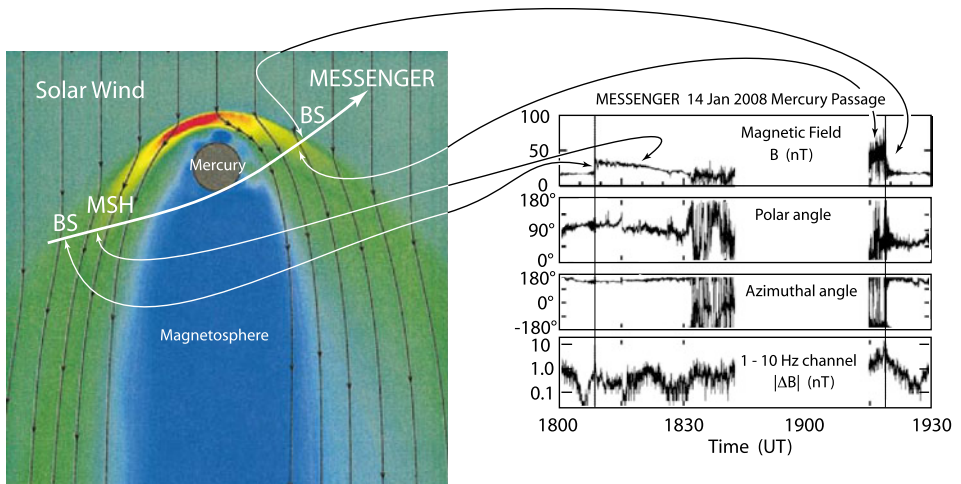


Figure 10.3: The MESSENGER magnetic field observations [after *Anderson et al, 2008*] during the Mercury bow shock crossings. Shown is the spacecraft orbit near Mercury overlaid on an MHD simulated Mercury magnetosphere [simulation from *Kabin et al, 2000*]. Only the shock and magnetosheath crossings are shown. The time from 1845 UT to 1915 UT the spacecraft was in the Mercury magnetosphere. The two bow shock crossings are quite different of different compression ratio and dynamics. Note the small foot at the first crossing and extended magnetosheath behind the bow shock. In the second crossing the magnetosheath is narrow and very turbulent, the shock has a large overshoot and a much broader foot.

extension is considerably larger than an ion gyro-radius – but indication of a foreshock, in which case the shock would be quasi-parallel while being locally quasi-perpendicular.) The angular variation in the magnetic field shows that the shock angle has varied very strongly during this crossing (suggesting that the shock might have been quasi-parallel, indeed, rather than quasi-perpendicular). The bow shock at this location seemed very variable. Particle observations indicate that during this crossing energetic ions were present in the immediate shock environment [*Zurbuchen et al, 2008*]. Taken these observations together with the Mariner 10 observations one concludes that Mercury possesses a supercritical bow shock that is continuously present and is a moderately strong shock. It develops a foot/foreshock region and a highly turbulent downstream magnetosheath and has regions of different strength as well as possesses diffuse ions. Even though this MESSENGER passage gave no confirmed indication for a genuine parallel shock region, Mariner 10 had indeed crossed a quasi-parallel Mercury bow shock.

All this is very similar to Earth's bow shock though on a smaller scale and in a different environment. However, a difference between Earth's and Mercury's bow shocks lies in the fact that Mercury has been found to be surrounded by a comet-like cloud of ions. These ions are of planetary origin, probably being created by sputtering from the Mercury surface [*Slavin et al, 2008*]. Since it is known that such sputtered ions when entering the solar wind become pickup ions and are accelerated, they could considerably affect the nature of the Mercury bow shock by reflection and acceleration. Thus Mercury's bow shock is in some sense a hybrid bow shock somewhere in between a comet and a magnetised planet. The

implications of this finding for the structure and dynamics of the Mercury bow shock are not yet known.

10.2.3 Jupiter's Bow Shock

Jupiter is not only the largest planet, it has also the strongest magnetic field believed to be generated by a huge internal dynamo that is driven by differential rotation in the planet's interior. Indeed Jupiter's rotation period is short, less than half as long as a terrestrial day. Table 10.2 shows the surface strength of this field to be more than an order of magnitude larger than the Earth's surface field. Together with the decreased kinetic solar wind pressure at the 5.2 AU heliocentric distance of Jupiter this strong field promises Jupiter to be surrounded by a huge magnetosphere that forms a giant obstacle in the solar wind and causes an extended bow shock to stand in front of the magnetosphere.

Simple scaling from Earth yields a stand-off distance for the magnetospheric nose of 4.2 times the Earth's magnetopause distance measured in Jovian radii. This yields a theoretical value of $R_{MP,J} = 46 R_J$, a value that underestimates the observed distance which varies between $50 R_J$ and $100 R_J$. This variation is due to the internal magnetospheric dynamics of the planet which is determined by the fast planetary rotation which exerts enormous centrifugal forces acting on the cold plasma in the inner magnetosphere, the Jovian plasmasphere. These forces cause a strong deviation of the plasmasphere from dipolar symmetry stretching the plasma radially out to become a fat sheet that extends out to the magnetopause, carrying the Jovian magnetic field with it and increase both the plasma and magnetic pressures at the contact with the solar wind. Since this happens in the symmetry plane of the Jovian rotation the nose of the Jovian magnetosphere deviates from bluntness and becomes more cusp-like than the Earth's magnetosphere.

At the Jovian orbit the Mach numbers are large, $\mathcal{M} > 10$ even under undisturbed conditions. The huge obstacle of Jovian magnetosphere thus causes a huge bow shock to stand in front of it. This bow shock is highly supercritical and should have a high compression ratio $N_2/N_1 \lesssim 4$. Straightforward application of Eq. (10.3) gives for the distance between the Jovian nose magnetopause and the Jovian bow shock the interval $13 R_J \lesssim \Delta_{dJ} = 0.275 R_{MP,J} \lesssim 28 R_J$. This is a substantial distance which, even in terms of numbers of planetary radii, is larger than the geocentric distance of Earth's bow shock from Earth! Since the solar wind magnetic field at the Jupiter orbit is of the order of $B \sim 1$ nT, the 1 keV ion gyro-radius becomes $r_{ci} \approx 750$ km, vanishingly small compared with the distance between the bow shock and magnetosphere. Moreover, the solar wind number density at the Jovian orbit has decreased from that at Earth's orbit by a factor of 27. The ion inertial length becomes $\lambda_i(5.2 \text{ AU}) \approx 530$ km. The Jovian bow shock is thus extremely thin, measured in Jovian radii. Everything happens on the plasma scales and not on the scales of the obstacle which only determine the gross scales like stand-off distance and curvature radius.

One therefore expects that, up to the lesser plasma densities and magnetic fields, the Jovian bow shock behaves very similarly to Earth's bow shock. No large surprises are expected in crossing it except for the slightly different plasma properties like lower densi-

ties, temperatures and magnetic fields as well as higher kinetic β corresponding to higher Mach numbers. This implies a somewhat stronger shock, stronger ion reflection, larger two stream and modified two stream instability growth rates and, possibly, a better efficiency in particle acceleration. In addition, due to its considerably larger curvature radius, the quasi-perpendicular and quasi-parallel surface areas on the Jovian bow shock are much larger than in the case of Earth's bow shock. This larger area combined with the much wider magnetosheath region will support particle acceleration simply because the fore-shock region is more extended and the downstream magnetosheath is broad enough for letting the turbulence evolve into a state closer to the state of well developed turbulence than in the case of Earth's magnetosheath.

The above values are the theoretical expectations which are based on the knowledge of Earth's bow shock. Jupiter's bow shock has been crossed by a small number of spacecraft, Pioneer 10, Pioneer 11, the two Voyager spacecraft, the Galileo mission, and Ulysses. Already the two Pioneer spacecraft confirmed the existence of a bow shock at Jupiter. Pioneer 10 observed upstream waves in association with bursts of relativistic electrons from Jupiter. Pioneer 11 crossed the Jupiter bow shock six times at planeto-centric distances in the interval $78 \leq R_{BS,J} \leq 110 R_J$. The average shock velocity in the solar wind frame was found to be of the order of $V_{BS} \sim 100$ km/s. A turbulent magnetosheath was detected where the average magnetic field was draped around the Jupiter magnetosphere in a way suggesting a blunt magnetopause shape. Moreover, large amplitude waves were reported to be observed in the magnetosheath [Smith *et al.*, 1975].

In 1979 the two Voyager spacecraft passed Jupiter on their way out into the outer heliosphere and interstellar space. During this passage which preceded the passages of the other giant planets of the outer solar planetary system they also crossed Jupiter's bow shock when entering close to the planet in a swing by manoeuvre. Figure 10.4 shows the path of Voyager 11 near Jupiter in February 1979. Again the bow shock was crossed six times (indicated by circles on the spacecraft path). Each time it was crossed multiply, indicating the very high variability of the Jovian bow shock location. The magnetic recording of three of the crossings (on days 59 and 60, 1979) is plotted in the upper right panel of the figure. All three crossings belong to the class of quasi-perpendicular shocks exhibiting steep compressions of the magnetic field and some indication of a foot in front of the shock ramp. The compression is about a factor 3, identifying the Jovian bow shock as a fairly strong but not a super-strong shock. Two minutes of plasma wave spectrogram during the shock crossing [see Scarf *et al.*, 1979a, b] are given on the lower right. The spectrogram covers the range of frequencies < 10 kHz which includes the nominal plasma frequency based on the expanding solar wind model. Indeed, prior to the shock crossing an intense plasma frequency line at ~ 6 kHz was detected that is seen here as a red emission band which disappears roughly 5 s before the spacecraft enters the shock. Approaching the shock the frequency increases gradually thereby mapping the increase in plasma density. Taking the nominal shock velocity of 100 km/s the 5 s between the end of the plasma line and the shock crossing correspond to a distance of 500 km upstream of the shock, roughly comparable to the upstream ion gyro-radius and the width of the shock foot. Indeed, some weak broadband noise begins at this time filling the gap from below. Referring to our knowledge about the processes taking place in the quasi-perpendicular shock foot we may

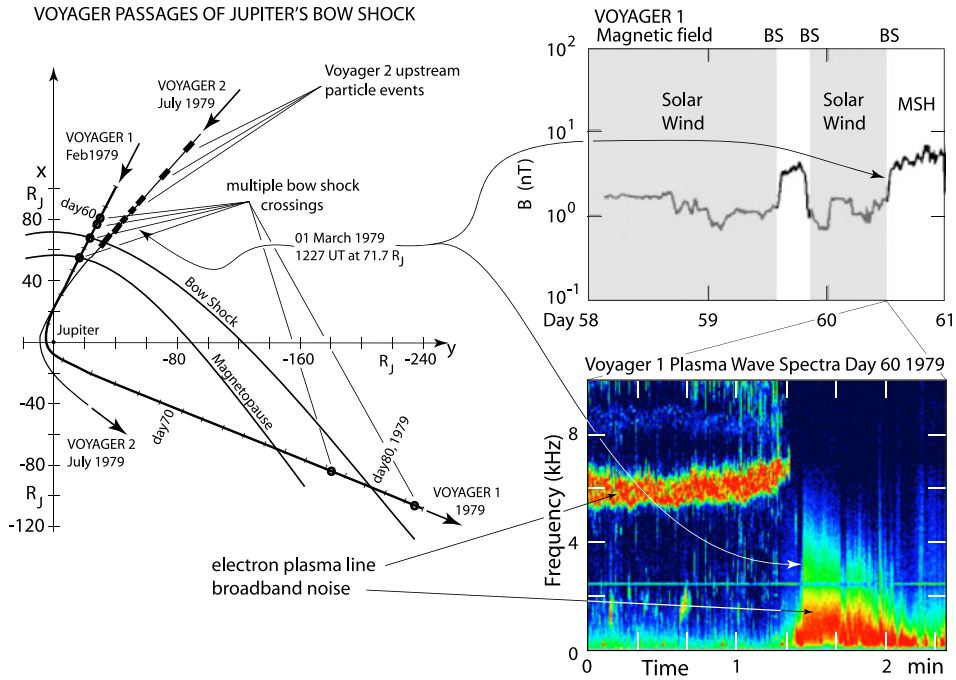


Figure 10.4: Summary of Voyager 1 crossings of the Jovian bow shock. *Left:* Projected Voyager 1 orbit in 1979 around Jupiter and the nominal bow shock and magnetopause locations in the (x, y) -plane with x pointing towards the Sun. The numbers along the Jupiter path give the days of the year 1979. The *dots* along the orbit indicate the multiple crossing of the Jupiter bow shock showing its high variability [after *Ness et al., 1979*]. Also shown is part of the Voyager 2 orbit in July 1979. The *black bars* on this orbit indicate the orbital segments when energetic particle events have been detected upstream of the bow shock [after *Krimigis et al., 1985*]. *Right:* Magnetic field (*top*) and plasma wave spectra (*bottom*) during selected crossings. The magnetic field shows the increase (decrease) when Voyager passes the shock inbound (outbound). Two minutes of plasma wave spectra are shown below for the third crossing in the *top* figure (*red* is highest, *black* lowest wave intensity). Voyager is in the electron foreshock before crossing as seen from the intense plasma frequency at 6 kHz. Passage of the shock occurs as a sharp cut-off of the plasma frequency and the start of broadband low frequency noise (courtesy by D. A. Gurnett, U Iowa).

conclude that these are the waves which are generated by two-stream and modified two-stream instabilities in the shock foot. Finally, when entering the shock ramp the broadband noise suddenly intensifies at the low frequencies and the bandwidth increases above the pre-shock plasma frequency. This is the typical signature of a quasi-perpendicular shock ramp crossing. For a shock of compression factor 3 the ramp plasma frequency is above 10 kHz and thus not captured in this spectrogram. Similarly, any radiation that is possibly emitted at the harmonic of the plasma frequency in front of the shock is at ~ 12 kHz outside the spectrogram. There is indication of a weak emission line at 8–9 kHz the origin of which is not known. More than a decade later, in February 1992 the *Ulysses* spacecraft on its swing-by path around Jupiter observed similar plasma wave emissions from the Jovian bow shock [Stone et al., 1992].

Similar to Voyager 1, Voyager 2 who crossed Jupiter's bow shock in July 1979 observed sporadic energetic electron bursts related to electron plasma oscillations and ion acoustic waves for several weeks prior to crossing the shock [Gurnett *et al*, 1979]. In analogy to Earth's bow shock these observations suggested magnetic connection to Jupiter's bow shock. When approaching the planet, energetic ion bursts were detected within the energy interval $30\text{keV} \lesssim \mathcal{E}_i \lesssim 5\text{MeV}$ [Krimigis *et al*, 1985]. The intervals when this happened are shown in Figure 10.4 as black bars along the Voyager 2 orbit. The energy spectra of these particles in the energy range $0.1\text{MeV} \lesssim \mathcal{E}_i \lesssim 2\text{MeV}$ followed a power law shape in energy. Above 300keV the spectra were void of protons being dominated by Oxygen and Sulphur, which indicates that Jovian magnetospheric ions were leaking out and were accelerated by the Jovian bow shock.

In toto Voyager 2 experienced five multiple crossings of Jupiter's bow shock between $99R_J$ and $62R_J$, in good agreement with the Pioneer 11 observations of the location of Jupiter's bow shock. On 5 July 1979 Voyager 2 reported a multiple crossing at planetocentric distances between $R \sim 69R_J$ and $R \sim 64R_J$ in a time interval ~ 5 hrs long, witnessing very high variability of the location of the bow shock. Since the Jupiter day is only 9.9 hrs long, the planet performed half a full rotation during this time such that part of the bow shock variability was caused by the wobbling of Jove's magnetosphere and the flapping of its dayside plasmaspheric nose. One may conclude from this that Jupiters bow shock position and form are highly variable being affected not only by the conditions in the solar wind but also by the Jovian magnetosphere in a way that is much stronger than Earth's bow shock is affected by Earth's magnetosphere.

The Jovian bow shock has also been crossed by the Ulysses spacecraft in February 1992 when it passed Jupiter. During its Jovian flyby on 2 February 1992 Ulysses confirmed the findings of the previous four missions to Jupiter and also explored the high-latitude dusk environment of Jupiter [Balogh *et al*, 1992]. The only inbound shock crossing seen by Ulysses occurred at 1733 UT on the planetocentric distance of $113R_J$ at a shock normal angle $\Theta_{Bn} = 36^\circ$ identifying this shock crossing as being quasi-parallel. In the energetic plasma measurements the inbound bow shock passage occurred as a sudden steep increase in the energetic particle fluxes [Lanzerotti *et al*, 1992] indicating ion heating. High energy particles had already been detected a few days before the encounter of the shock but were probably escaping energetic ions from the Jovian plasma environment. On the Ulysses outbound pass on 14 February 1992 three bow shock crossing were observed at planetocentric distances of $109R_J$, $112R_J$, and on 16 February at $149R_J$. These crossing were found under shock normal angles $\Theta_{Bn} = 80^\circ$ and 82° , respectively and were thus identified as quasi-perpendicular shocks. All these crossings found the shock at amazingly large distances from Jupiter suggesting that Jupiter's outer environment is strongly affected by the Jovian magnetospheric plasma content and its very fast rotation. We will return to this point below. Ulysses also detected low frequency waves in the foreshock and in the magnetosheath, both transverse and compressional and found evidence of magnetic mirror modes downstream of the Jovian bow shock.

In the late nineties the Galileo spacecraft crossed it several times, and Cassini, during its Jupiter flyby on the way to Saturn skimmed the Jovian bow shock at the time when Galileo was also in orbit. Achilleos *et al* [2004] analysed the magnetic field of 30 passages

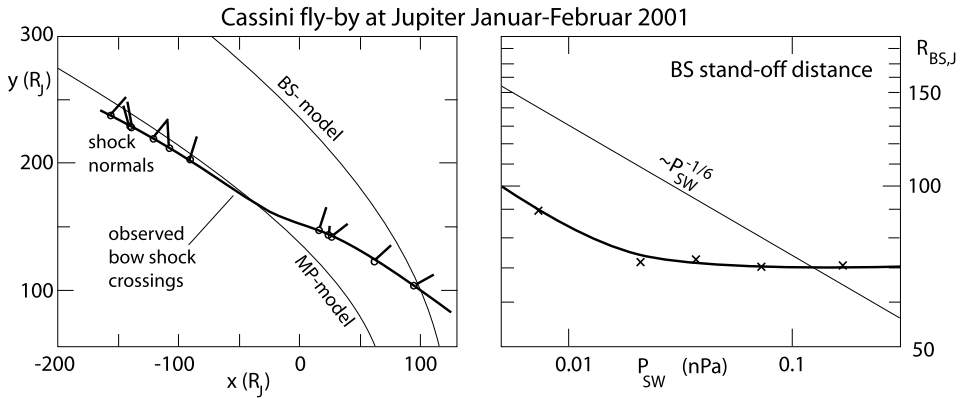


Figure 10.5: Jupiter's bow shock location as inferred from crossings of the Cassini spacecraft during its fly-by at Jupiter on its way to Saturn. *Left:* The model positions of the Jovian bow shock and magnetopause and the observed locations of the bow shock showing including the shock normal directions. One notes the variations in the shock normal which indicate the non-stationary conditions at the bow shock. *Right:* The inferred dependence of the stand-off distance of the Jovian bow shock as function of the solar wind pressure. The straight line shows the predicted theoretical dependence. The observation indicates independence of the stand-off distance from solar wind pressure for pressures exceeding 0.02 nPa [data taken from *Achilleos et al, 2004*, courtesy American Geophysical Union].

of the more than 40 Cassini spacecraft crossings of Jupiter's bow shock between October 2000 and April 2001. Their important finding is that the stand-off distance of the Jovian bow shock is relatively independent of solar wind pressure for pressures above $P_{SW} \gtrsim 0.02$ nPa. At these pressures the average shock nose distance was around $\sim 70 R_J$. At lower solar wind pressures the stand-off distance of the shock increased. This is interesting as it suggests that the Jovian magnetosphere becomes incompressible when it is compressed to a distance of about $\sim 70 R_J$, an effect which may be due to the presence of the extended internal magnetospheric corotating Jovian plasma disk which adds to the resistance of the magnetosphere against further compression. At solar wind pressures $P_{SW} \lesssim 0.02$ nPa the magnetosphere is sufficiently expanded to follow the theoretical law $R_{BS,J} \propto P_{SW}^{-1/6}$ [cf., e.g. *Slavin et al, 1985*]. This is also concluded from the determined shape of the bow shock on the left in Figure 10.5 which shows that the stand-off distance of the bow shock in the ecliptic is farther out than at higher latitudes. At higher latitudes the effect of the plasma disk vanishes and the shock position is supported only by the Jovian magnetic field in the magnetosphere. However, these single spacecraft measurements are ambiguous leaving it open whether this ecliptic expansion was not a temporal effect caused by a change in solar wind pressure. Such a change has indeed been observed simultaneously by Cassini and Galileo [*Kurth et al, 2002*] and has been interpreted as deformation of the magnetopause.

The second interesting result of this investigation is that the shock normal exhibits substantial variations. The Jovian bow shock is highly variable in time, which is not an unexpected result seen in Figure 10.5. The farthest Cassini bow shock crossing where plasma wave observations were available occurred at planetocentric distance of $140.2 R_J$ again

confirming the existence of Langmuir waves and ion acoustic waves in the foreshock, and an intense burst of broadband emissions during the shock ramp passage [Kurth *et al*, 2002]. 12 of these crossings were well-defined quasi-perpendicular shocks, 5 were quasi-parallel, the rest were not well defined oblique shocks. Cassini found the bow shock to extend over $700 R_J$ down the flank of Jupiter's magnetosphere [Szego *et al*, 2003]. Plasma observations during these crossings showed that the shock crossings were characterised by the expected strong ion thermalisation which should coincide with reflection of ions and the appearance of an electrostatic shock potential. Both effects gradually weaken with increasing distance from the bow shock nose. In particular the shock potential softens rapidly, indicating a weaker and more variable shock at increasing downstream distance. Krupp *et al* [2004] refer to energetic particle observations during Cassini bow shock crossings of the far downstream bow shock which show the occurrence of energetic ions >36 keV and in a few cases also relativistic electrons. Though the ions are clearly related to the existence of the shock probably being of diffusive origin, the relativistic electrons are not attributed by these authors to acceleration by the Jovian bow shock; rather they believe they leak out from the Jovian magnetosphere due to reconnection. Thus the picture is rather confused as electron acceleration can also be achieved by trapping in magnetic flux tubes that are connected to the shock, which is not unreasonable in view of the huge extension of Jupiter's bow shock.

10.2.4 Saturn, Uranus, Neptune

The remaining three giant outer planets are all strongly magnetised (see Table 10.2). As a consequence they possess magnetospheres. Since the solar wind dynamic pressure is relatively low at their locations, these magnetospheres are very large, and strong bow shocks at large stand-off distances from the planets evolve in front of them. Moreover, their relatively fast spin periods (Table 10.1) together with their large radii suggest that their magnetospheres are also subject to strong corotation of a presumable inner magnetospheric plasma component. Whether such a component exists is a question as the solar UV radiation turns out to be rather weak at the distances of these planets. Nevertheless, cosmic ray produced ionospheres due to ionisation of their atmospheres will release sufficient plasma along the magnetic field to populate the inner corotating regions of the magnetospheres. In the absence of such an ionosphere the planetary magnetic fields are still strong enough for producing a magnetosphere and its shielding bow shock.

The Saturnian Bow Shock

The bow shock of the next closest of the outer giant planets, Saturn, was first encountered by the Pioneer 11 spacecraft in 1979 [Smith *et al*, 1980] at a distance of $23.7 R_S$. When Pioneer 11 left the planet, having crossed its magnetosphere, monitoring the Saturnian magnetic field dipole and confirming the existence of a corotating plasma between $4 R_S \lesssim R \lesssim 16 R_S$, it again experienced multiple bow shock crossings between distances of $49.4 R_S \leq R \leq 65 R_S$. Similar to Jupiter, the corotating plasma component (called Saturn's plasma torus) plays some role in the dynamics of the Saturnian magnetosphere. Its

particle energies are in the range $0.1 \lesssim \mathcal{E} \lesssim 8$ keV, and plasma densities are of the order of $N \sim 5 \times 10^5 \text{ m}^{-3}$, and temperatures $T \sim 10^6$ K. Thus the plasma- $\beta \sim 1$ up to an inward distance of $\sim 6.5 R_S$. Probably, this plasma is caused by photodissociation of water frost of Saturn's rings [Frank *et al*, 1980].

One year later Voyager 1 and then Voyager 2 passed Saturn in November 1980 and August 1981, respectively, crossing its bow shock and magnetosphere and confirming each others observations. Altogether the observations of these spacecraft identified Saturn's bow shock to be very similar to Jupiter's bow shock. The stand-off distances have already been given in Table 10.2 showing that, again as in the case of Jupiter, the Saturnian bow shock is far upstream (in absolute units of planetary radii) and thus very thin being of plasma scale width and leaving a broad (in terms of plasma scales) magnetosheath downstream behind it. At Saturn's orbit the solar wind is even colder ($T_i \sim 1.4$ eV, $T_e \sim 2$ eV) and more dilute than at Jupiter with $N \sim 1.1 \times 10^5 \text{ m}^{-3}$, the Alfvén and magnetosonic speed are low, $V_A \sim 20$ km/s and $V_{ms} \sim 28$ km/s, respectively [Scarf *et al*, 1981]. Thus the Mach numbers are high, of the order of $\mathcal{M} \sim 15\text{--}20$, and the shock is strong. From Voyager 1 an inbound dayside high Mach number quasi-perpendicular shock crossing at planetocentric distance of $32.4 R_S$ and an outbound morning flank quasi-parallel shock crossing at $78 R_S$ was reported, the latter possessing an extended turbulent foreshock region containing much low-frequency upstream turbulence [Ness *et al*, 1981].

Plasma wave observations from both Voyager spacecraft upstream of the shock confirmed occasional magnetic connectivity from Langmuir and ion acoustic waves and suggesting the presence of fast electron beams. Just prior to crossing the bow shock strong Langmuir wave bursts were observed, and the shock manifested itself in broadband low frequency noise as was expected [Gurnett *et al*, 1981; Scarf *et al*, 1982]. One surprising observation of Voyager 1 was that downstream of the shock the plasma wave intensities were rather low, much less than in Earth's magnetosheath, an observation that still awaits explanation. Scarf *et al* [1981], from comparison of Earth, Jupiter and Saturn plasma wave spectra, speculates that there is a relation between Mach number and wave excitation. This is indeed the case as at high Mach numbers other instabilities are excited. However, the downstream absence of high wave levels is more difficult to understand. Rather it has to do with the enormous downstream extension of the giant planet magnetosheaths where the high levels of turbulence are confined to the downstream region adjacent to the bow shock with thickness of the order of a couple of energetic passing ion gyroradii. Farther downstream outside the thermalisation distance the plasma probably forgets for the presence of the shock and normalises at a sub-magnetosonic flow state of enhanced temperature and density. On the other hand, Voyager 2 observed very low plasma densities and plasma frequencies at the bow shock crossings. From bow shock models these low values seemed unusual, and it was suggested that the wake of Jupiter's extended magnetotail swept over Saturn during the passage of Voyager 2.

Two decades later it was the Cassini spacecraft on its way to Saturn's moon Titan which crossed the bow shock of Saturn 17 times. Cassini's first contact with Saturn's bow shock was on 22 March 2004, when it was hit by electron beams arriving from the direction of Saturn along the interplanetary spiral magnetic field. These electrons indicated magnetic connectivity to the bow shock at an upstream distance of $825 R_S$ and made themselves

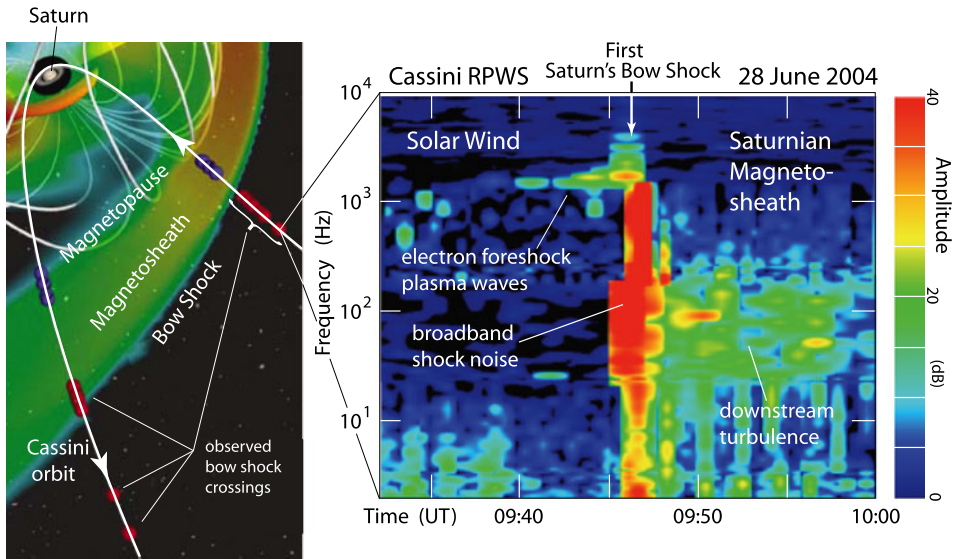


Figure 10.6: The first Cassini spacecraft crossing of Saturn's bow shock at a distance of $49.2 R_S$ on 28 June 2004 as seen in plasma wave spectra (courtesy by D. A. Gurnett, U Iowa) on the *right*. The crossing occurs at about 0945 UT. About 5 min prior to the shock the spacecraft detects electron plasma waves which are an indication of magnetic connection between spacecraft and shock. As expected, the shock is seen as a sudden spectral broadening from low frequencies to frequencies well above the plasma frequency. The crossing lasts for ~ 5 min. Afterwards the spacecraft is in the turbulent Saturnian magnetosheath. The picture on the *left* shows a 3d-MHD simulation of Saturn's magnetosphere using nominal solar wind values at 9.6 AU [section taken from Gombosi & Hansen, 2005]. Overlaid is the inbound-outbound Cassini orbit around Saturn. The *red dots* indicate observed bow shock crossings, to the first of which belong the plasma wave spectra on the *right*. *Blue dots* are magnetopause crossings. Note the large inflation of Saturn's bow shock-magnetosphere system compared with the nominal simulation.

visible by exciting a burst of Langmuir waves [Gurnett *et al*, 2005]. Such bursts have been detected in the following during the Cassini approach more frequently until the first crossing of the Saturnian bow shock on 28 June 2004 at 0945 UT which happened at radial planetocentric distance $49.2 R_S$. Figure 10.6 on the right shows the dynamic plasma wave spectrum of this first Saturnian bow shock crossing in the RPWS recordings between 3 Hz and 10 kHz. In the Cassini plasma instrument RPWS data the crossing appears as the abrupt transition from the Saturnian electron foreshock plasma oscillations at frequency $f \sim 1.78 \text{ kHz}$ to a burst of low frequency broadband electrostatic noise. The Langmuir electron plasma oscillations correspond to a plasma density $N \sim 4 \times 10^4 \text{ m}^{-3}$ in the dilute upstream solar wind plasma, in very good agreement with the r^{-2} -decrease of the solar wind density from 1 AU to 9.6 AU. In the time following this crossing the shock position oscillated six times back and forth across the Cassini spacecraft. The left part of the figure shows the Cassini orbit around Saturn overlaid on a 3d-MHD simulation study of the solar wind-Saturnian magnetosphere interaction [Gombosi & Hansen, 2005] using the nominal solar wind data at the Saturn orbit. The red dots along the path indicate the various bow

shock crossings observed by Cassini. Blue dots indicate magnetopause crossings. It is interesting to note that the simulation grossly underestimates the inflation of both the Saturnian magnetosphere and bow shock.

In the plasma wave spectra the Saturnian bow shock is preceded by ~ 5 min of plasma oscillations, the trace of which broadens in both directions to higher and lower frequencies just before shock crossing, an effect known from Earth's bow shock and probably related to the presence of the shock foot and the excitation of either electron acoustic, Buneman or modified two-stream instabilities. This is followed by the intense broadband noise burst of the shock transition with maximum intensity around $\sim (0.1-1)$ kHz, one minute later extending up to ~ 5 kHz. The whole transition lasts for just $\lesssim 3$ min. After this time Cassini is in the Saturnian magnetosheath which in the plasma wave spectrum appears as broadband ion acoustic noise $\lesssim 30$ Hz that is modulated by the spacecraft spin, and by about isotropic electric noise between 30 Hz and 300 Hz, indicating that Cassini entered the region of downstream magnetosheath plasma turbulence. In the plasma waves this behaviour is typical for a quasi-perpendicular shock crossings as is known from crossings of Earth's bow shock wave. Taking the view that the shock is strong and of maximum compression, the density and magnetic fields should be compressed by at most a factor 4, implying a shock plasma frequency $\lesssim 4$ kHz. Clearly the broadband noise detected during the shock crossing extends beyond this value up to ~ 6 kHz or more. Hence, it cannot be produced by a plasma instability. Rather it is the result of crossing spatially localised electrostatic structures of central frequency around 100–1000 Hz.

Little is known about the shock velocity. The electron and ion inertial lengths are in the respective intervals $1.5 \text{ km} \lesssim \lambda_e \lesssim 3 \text{ km}$ and $700 \text{ km} \lesssim \lambda_i \lesssim 1300 \text{ km}$. If we assume that the width of the shock transition is of the order of the ion inertial length, then the 2–3 min crossing time of the shock correspond to the spacecraft velocity, suggesting that the shock was about stationary during this first crossing, possibly standing at its maximum possible inflation distance. Moreover, assuming that the scales of the localised small scale structures responsible for the broadband noise are of the order of the electron inertial length, then the bandwidth of the noise suggests that these structures move at velocities of the order of $\sim 1000 \text{ km/s}$. This speed should be less than the electron thermal velocity, suggesting that the electron temperature in the shock is at least of the order of $T_e \gtrsim 50 \text{ eV}$. When referring to the upstream low electron temperature in the solar wind which, for adiabatic expansion, is of the order of $T_e \sim (0.5-1) \text{ eV}$, this implies strong electron heating in the shock transition.

The larger than expected inflation of the Saturnian magnetosphere and bow shock points on a more Jupiter-like than Earth-like magnetosphere of Saturn. This is confirmed by a comparative statistical study [Hendricks *et al*, 2005] using data from Pioneer 11, the two Voyager spacecraft, and Ulysses.

Another question concerns the nature of the shock. As in the case of Jupiter, the bow shock of Saturn shows, in the overwhelming number of cases, a sharp transition from solar wind to the magnetosheath which is typical for quasi-perpendicular shocks. This is seen in Figure 10.7 for the first 7 Cassini bow shock crossings. The absolute changes in the shock ramp are not very high; they are, however, fairly steep with well expressed overshoots. They also possess extended shock feet region which contain ULF wave activity.

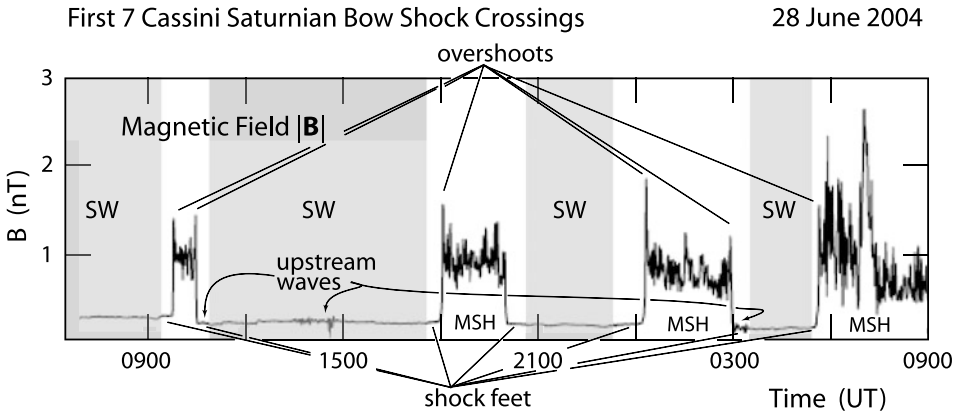


Figure 10.7: The first 7 Cassini spacecraft crossing of Saturn’s bow shock [data taken from *Achilleos et al, 2006*] in the magnitude of the magnetic field. All of these crossings are quasi-perpendicular shock crossings with steep ramps, overshoots, feet and indications of ULF upstream waves in their feet. The magnetosheaths are fairly disturbed in the ULF wave range. In one case ULF upstream waves not related to a shock crossing are observed which might either indicate switching into the ion foreshock of a quasi-parallel part of the bow shock or approach to a shock without touching it. The first of these Cassini inbound crossings at 0945 UT is the one shown in the plasma wave spectra in Figure 10.6.

The very first of these crossings is the same as that shown in Figure 10.6 in the high frequency plasma waves. Interesting enough, these spectra indicate that the shock foot is a region of reduced ion acoustic wave activity, which may be in favour of the modified two-stream instability which has lower frequency and is oblique such that convective Doppler shift affects it very little. Since the electron cyclotron frequency is < 10 Hz, these waves which propagate on the whistler branch are not seen in the spectra and do not occur either in the rough magnetic measurements. *Achilleos et al [2006]* have determined the shock strengths. These turn out to be considerably weaker than expected, mostly with compression ratios $B_2/B_1 \lesssim 2$. Only in one case this ratio is closer to 3. On the other hand, the measured overshoots have been found in the interval $2 < B_{\text{over}}/B_1 < 3.5$. Shock ramp widths have been determined to fall between $(0.13\text{--}1.11)\lambda_i$ with the sixth crossing making an exception of having a ramp width of $\sim 2.75\lambda_i$. Neglecting this crossing, the average ramp width turns out to be $0.6\lambda_i$ which is in agreement of what is known from Earth’s bow shock. Shock velocities have been found from passing the shock ramp to be in the average $V_S \sim 124 \pm 108$ km/s. There is a discrepancy in the case of the first crossing where the velocity found is 363 km/s, much larger than the shock velocity determined from the plasma wave spectrum which was between 10–20 km/s. This might result from the fact that the plasma wave spectra in the Figure 10.6 do not resolve the shock ramp and thus cover a wider spatial range that includes the overshoot and therefore underestimates the velocity.

The more perpendicular than parallel behaviour of the outer planetary bow shocks is probably related to the more circular nature of the spiral interplanetary magnetic field with increasing heliocentric distance. In the ecliptic plane, the plane of the Cassini orbit, the

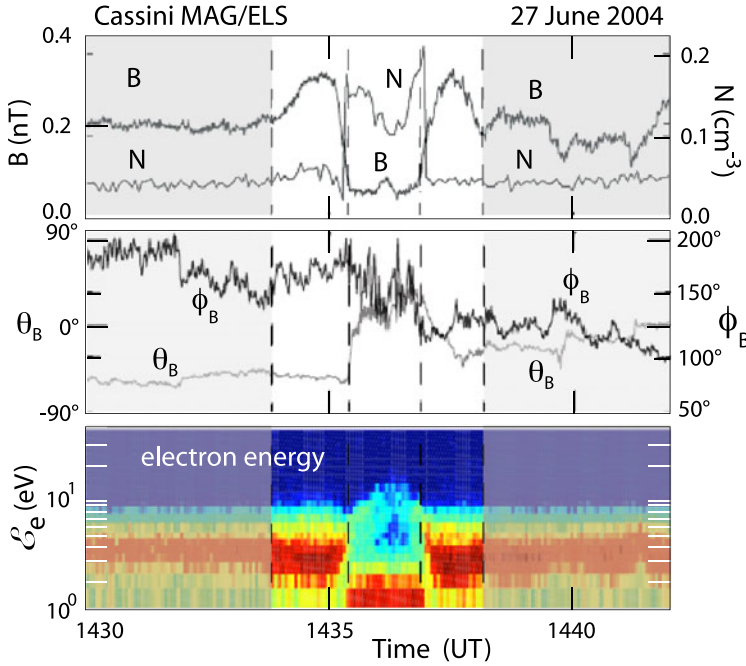


Figure 10.8: Passage of Cassini across an Hot Flow Anomaly near the Saturnian bow shock but before getting into contact with the shock [data taken from *Masters et al., 2008*]. Shown is the variation in the magnetic field B and density N . The magnetic field develops compressions at the edges of the HFA prior to dropping to very low values inside. Simultaneously the plasma density increases steeply inside the magnetic cavity. The turn in the direction of the magnetic field vector across the HFA is seen in the two magnetic angles in the *middle panel*. The *bottom panel* shows the electron energy colour coded from the solar wind across the HFA back into the solar wind. The electrons are retarded in the boundaries and heated. Inside the HFA the electrons do not move anymore while assuming a high velocity spread, i.e. a high temperature roughly a factor 1.5 higher than outside.

interplanetary magnetic field at the orbit of Saturn is almost azimuthal. Moreover, since the Mach numbers increase with heliocentric distance and the dependence of the velocity-shock normal angle $\theta_{Vn} < \cos^{-1}(\mathcal{M}^{-1})$, the opening angle of the bow shocks for larger Mach numbers increases thus favouring quasi-perpendicular shock formation. However, this might not always be the case as one particular recent observation shows [*Masters et al., 2008*] where Saturn obviously was crossed by current sheets of the character of tangential discontinuities in the solar wind, across that the magnetic field direction rotates by an arbitrary angle possibly also changing magnitude. As a consequence the convection electric field normal to the discontinuity surface may also change sign, depending on the angle of rotation of the magnetic field. If this happens in a way that the electric field to both sides points into the tangential discontinuity, a Hot Flow Anomaly (HFA) develops at contact with the shock because the shock reflected ions are accelerated into the discontinuity where they accumulate. This obviously happened at Saturn as shown in Figure 10.8, and the Saturnian bow shock reacted in about the same way as the terrestrial bow shock

by evolving Hot Flow Anomalies. These events were observed by Cassini during its first two orbits around Saturn. The Saturnian HFA is somewhat different from those observed at Earth as it has two magnetic walls at its boundaries while the near-Earth HFAs always possessed strong shocks which confined the current and plasma in its interior. Nevertheless, the observation of this structure shows that the Saturnian bow shock behaves in some respect similar to the terrestrial bow shock.

Uranus' Bow Shock

The planet Uranus is particular among the outer giant planets because its magnetic axis is inclined towards the ecliptic by an angle of 59° against its rotation axis. Rotating faster than Earth and being of four times larger diameter, the magnetosphere is in a fast wobbling motion, flapping up and down with respect to the ecliptic. Moreover the magnetosphere has a strange configuration with the poles being so close to the ecliptic, which makes the cusps of the magnetosphere much more vulnerable to direct solar wind inflow and to widening than on Earth. Probably these conditions do not have very strong effects on the Uranian bow shock, however, as the conditions at the bow shock are determined mainly by the gross properties of the obstacle. The distance between the shock and the magnetospheric obstacle is very large, comparable to Jupiter and Saturn (see Table 10.2), such that one does not expect that the behaviour of the magnetosphere directly transfers to the bow shock. The Uranus bow shock is probably controlled by the solar wind. One, thus, does not expect to encounter large differences between the bow shock of Uranus and those of Saturn and Jupiter, except for the fact that the solar wind is more dilute, colder than at the latter planets, and the magnetic field is weaker, which implies a lower plasma β and larger Mach number.

Uranus was passed by Voyager 2 in January 1986 (see Figure 10.9). A first contact with its environment was dated on 22 January 1986 at distance $140 R_U$ when Voyager 2 detected a burst or 28-keV ions coming from the planet which lasted for ~ 12 hours [Krimigis *et al*, 1986]. Two days later, just before bow shock encounter, these fluxes increased by two orders of magnitude, briefly decreasing after shock passage and increasing to 10^3 times the interplanetary flux level after entering the Uranus magnetosphere. The bow shock crossing itself was not spectacular in the particle fluxes, however. Instead, the bow shock crossing was identified in magnetic field [Ness *et al*, 1986] and plasma wave observations [Gurnett *et al*, 1986]. In the latter the shock crossing occurred 0728 spacecraft time on 24 January 1986 at planetocentric distance of $23.5 R_U$ as a 4 min long broadband electric wave burst that was preceded hours before by electron plasma oscillations in the 1.78 kHz channel in Langmuir waves, signalling magnetic connection and the emission of electron beams from the shock and fixing the pre-shock solar wind plasma density to a surprisingly high value of $N \sim 4 \times 10^4 \text{ m}^{-3}$, three times higher than expected at the Uranian distance of 19 AU. Somewhat surprising was the low plasma wave activity in the shock foot which Moses *et al* [1989] attribute to the unusually high plasma- β during this upstream crossing comparably near to the Uranian bow shock nose. Again the strong broadband burst is an indication of the two stream instabilities acting in the shock ramp transition and overshoot which produces ion heating and plasma scale localised wave structures. Similarly to the

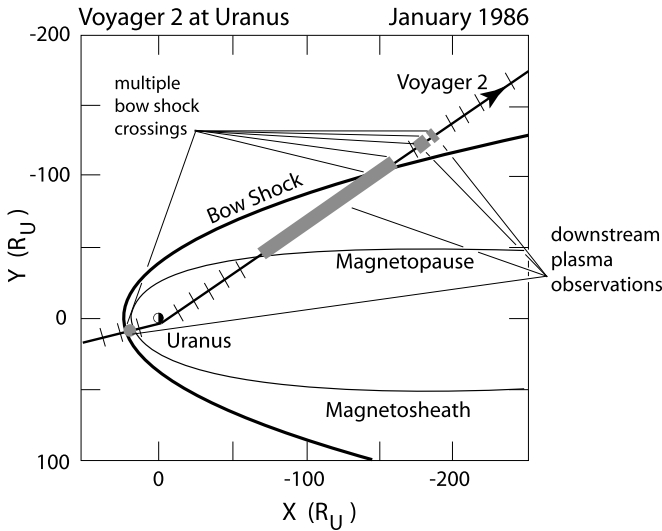


Figure 10.9: The path of Voyager 2 around Uranus in January 1986. Indicated are the multiple bow shock crossings as seen in the plasma instrument [Bridge *et al.*, 1986]. The grey bars indicate the observation of thermalised magnetosheath plasma downstream of the bow shock which is bound by the shock and gives a rough indication of their geometry. Note the flapping of the bow shock at the flank crossings at January 27–29.

other giant planets passed by Voyager, the wave activity in the adjacent downstream magnetosheath was very low. This first crossing has been identified as a quasi-perpendicular shock [Bridge *et al.*, 1986] in a high Mach number plasma with plasma $\beta \sim 1$.

Neptune's Bow Shock

The outermost planet, also a giant, is Neptune. He was the last to be visited by manmade spacecraft, when on 24 August 1989 Voyager 2 reached it on its way out of the solar system. As suspected, Neptune was found to possess a strong and complex intrinsic magnetic field, a magnetosphere in interaction with the solar wind, and a bow shock standing at an upstream distance of $34.9 R_N$ [Ness *et al.*, 1989; Belcher *et al.*, 1989]. This bow shock was left at the tail flanks where its signature was still detectable. The sunward compression factor of the shock was $B_2/B_1 \sim 2\text{--}3$ indicating a fairly strong shock. It also possessed a susceptible overshoot, a foot region and some magnetic fluctuations in front of it. Thus the shock was quasi-perpendicular. Not surprisingly, the plasma wave instrument aboard Voyager 2 observed all the features that are common during a supersonic (Mach number ~ 6) collisionless bow shock crossing: upstream plasma oscillations indicating magnetic connectivity to the shock and electron beams emitted from the shock, broadband electrostatic turbulence in the bow shock crossing and transition region, and a drop to the instrument noise level in the adjacent magnetosheath [Gurnett *et al.*, 1989], similar to what had been seen in the other giants, Saturn and Uranus. The plasma oscillations were in the 562-Hz

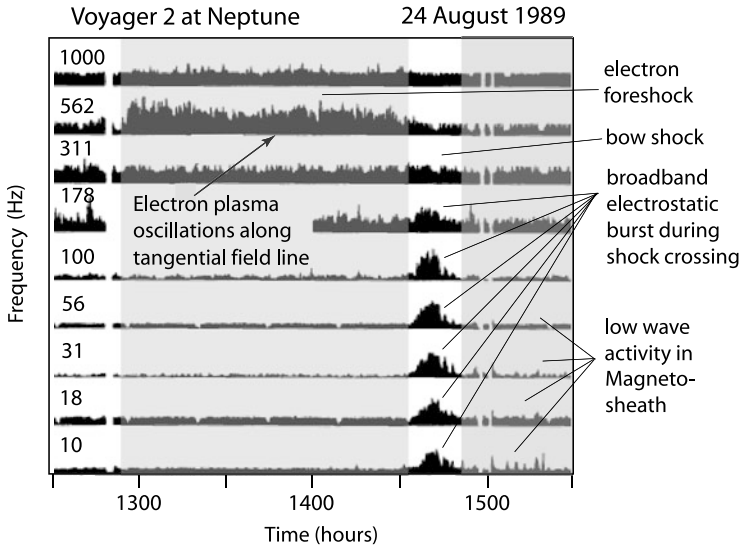


Figure 10.10: The Voyager 2 encounter of the Neptunian bow shock on 24 August 1989 as seen in the plasma wave recordings [after Gurnett *et al.*, 1989]. Shown are the electric wave field intensities in *black shadowing*. The intensity scale in each of the frequency bands shown is three orders of magnitude from 10^{-6} mV/m to 10^{-3} mV/m. From 1300 Spacecraft time on until hitting the bow shock Voyager 2 was in the electron foreshock as indicated by the intense electron plasma oscillations emitted along the shock tangential magnetic field line. The shock encounter around 1520 (*left white* in the figure) appears as the intense broadband noise up to 311 Hz. Afterwards the spacecraft is in the surprisingly quiet Neptunian magnetosheath.

channel which corresponds to an upstream plasma density of $N \sim 3.9 \times 10^3 \text{ m}^{-3}$. Figure 10.10 shows the plasma wave recording during this first encounter of the Neptunian bow shock.

10.3 Mars and Moon

Mars and Moon are the two celestial bodies in our nearer spatial environment which either have no or a very small magnetic field, no or a very thin atmosphere, which come into direct contact with the solar wind and are large enough (see Table 10.2) to distort the solar wind flow as their diameters are larger than the gyro-radius of a solar wind ion.

10.3.1 Lunar Mini-Bow Shocks

The typical ion solar wind ion gyro-radius at Moon is of the order of $r_{ci}(1\text{AU}) \sim (400\text{--}1000)$ km depending on the magnetic field strength. Hence, if the Moon possessed an own magnetic field or if the passing solar wind would induce a magnetic field in the body of the Moon, this lunar field would deflect the solar wind, cause a magnetosphere around Moon and a detached bow shock standing in front of its magnetosphere. The size of

the magnetosphere and stand-off distance of this bow shock would depend on the strength of the interior or induced fields. As long as nothing is known about the proper interior field, nothing can be said about the presence of such a magnetosphere and/or bow shock [Sonett & Colburn, 1967]. However, an induced field would at most be of the same order as the solar wind field depending on the lunar conductivity.

Thus the first observation of the lunar magnetic field was expected with great curiosity. These observations were provided by the Lunar Orbiter Explorer 35 in 1967 with the result that the Moon does indeed disturb the solar wind generating a plasma diluted wake behind its body while the Moon does not show any indication of a magnetosphere and definitely does not possess a bow shock [Colburn *et al*, 1967].

This finding is of interest as it tells that the Moon does also not possess any significant lunar magnetic field, whether intrinsic or induced. The absence of any remarkable induced field, in addition, tells something about the value of the bulk lunar resistivity which must be larger than 10^6 Ohm m for an interplanetary magnetic field strength of 5 nT [Sonett & Colburn, 1967]. The Moon absorbs part of the impacting solar wind, part of it passes the Moon like a dielectric medium, and the distortions that the solar wind and its magnetic field experience are entirely due to the deflection of the frozen-in solar wind magnetic field and particles around the Moon. The Moon creates a shadow in the solar wind screening its backside from the solar wind particles which have not enough thermal spread in order to fill the gap behind the Moon. The minimum length of this shadow in lunar radii is $\sim 6R_L$. This particle shadow, which appears as a lunar wake is bound to its sides by current sheets of width the order of an ion gyro-radius. Their task is to cancel the magnetic field in the region behind the Moon that is void of plasma. They are dynamically active and cause distortions of the magnetic field and solar wind plasma from current drift instabilities behind the Moon, but no bow shock wave is generated in front of the Moon.

Transverse fluctuations in the magnetic field that have been detected upstream and downstream near the Moon [Ness *et al*, 1968; Ness & Schatten, 1969] with frequencies in the range $0 < f \lesssim 1$ Hz on magnetic field line that cross the Moon's wake have been shown to result from just these instabilities. Electrons that are accelerated in the unstable wave fields escape along the magnetic field to upstream of the Moon, where they cause ballistic instabilities [Krall & Tidman, 1969] and excite electron plasma oscillations which mistakenly can be taken as the signature of a lunar bow shock in similarity to such signatures upstream of the bow shocks of the magnetised planets.

Apart from the absence of any large-scale intrinsic lunar magnetic field, its consequences for the internal dynamics of the Moon, and the limitations on a global lunar conductivity [Sonett, 1975] that are set by the absence of a lunar bow shock wave, it cannot be concluded that the Moon does not have any local magnetic fields which result from remnant magnetisation of its rocky material or otherwise from highly conducting materials in which magnetic fields of solar wind strength could be induced when they become exposed to the solar wind magnetic field passing the Moon. Such localised internal magnetic fields may well cause mini-magnetospheres which are screened against the solar wind by mini-bow shocks that exist only in the region where the lunar magnetisation occurs. Situations of this kind have recently been investigated by numerical simulation [Harnett & Winglee, 2000] after measurements performed during the Apollo missions suggested that such local

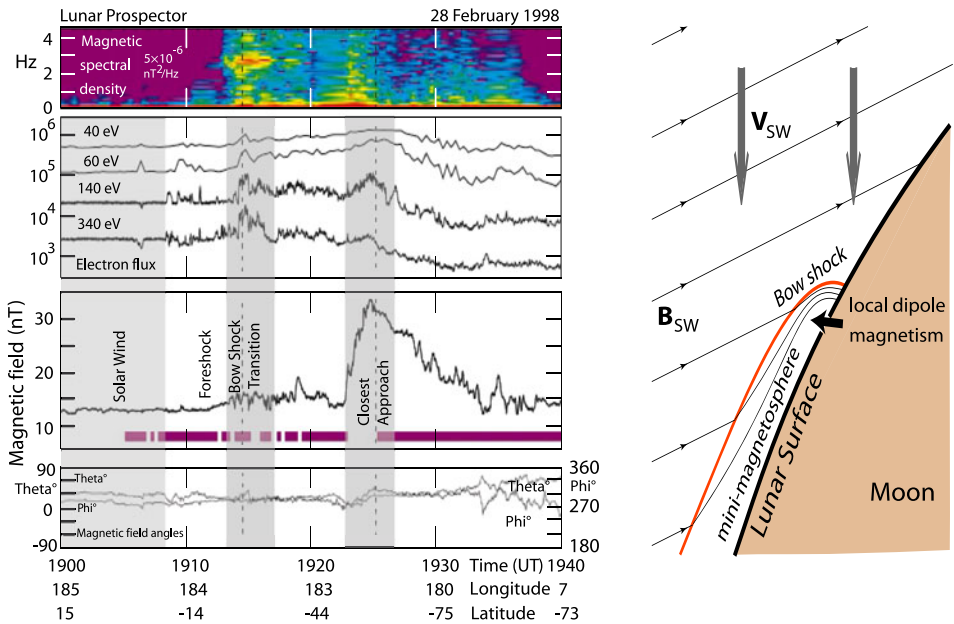


Figure 10.11: The passage across a mini-magnetosphere close to the lunar surface by the Lunar Prospector on 28 February 1998 [after Lin *et al.*, 1998]. The spacecraft passes from the solar wind into a region similar to the electron foreshock with energetic electron beams emanating from the lunar direction along the magnetic field to the spacecraft. Afterwards a shock-like discontinuity is crossed which has the character of a detached bow shock which shields a strong magnetic field from a local crustal magnetism on the Moon. The low frequency wave spectrum shows the magnetic fluctuations that are related to the foreshock and bow shock. Four energy channels of electrons are shown. The two bottom panels give the magnitude and two angles of the magnetic field. The field reaches 30 nT. On the right the model of the mini-magnetosphere with its mini-bow shock is sketched.

fields up to strengths of 10^3 nT and extensions >100 km are indeed present on the Moon [Sharp *et al.*, 1973; Dyal *et al.*, 1974; Hood & Schubert, 1980; Hood *et al.*, 1981; Lin *et al.*, 1988] and the Lunar Prospector mapped the lunar surface magnetic field and found surface fields $\lesssim 300$ nT over extensions of $\lesssim 1200$ km [Lin *et al.*, 1998]. The condition for such magnetosphere/bow shock systems to exist is related to the ion gyro-radius, i.e. the extension of the magnetised region must be substantially larger than the solar wind ion gyro-radius.

Figure 10.11 shows a case of the Lunar Prospector observations with indication of a mini-magnetosphere and bow shock. The spacecraft moves from the solar wind into a region which resembles an electron foreshock with enhanced fluxes of electrons coming from the Moon along magnetic field lines. Then it crosses a signature that resembles a shock transition followed by a disturbed field zone and a strong increase of the magnetic field at closest approach to Moon.

Harnett & Winglee [2000] found that such mini-magnetospheres can indeed be formed under the condition that the magnetic anomaly field strength on the Moon is >10 nT at an altitude of 100 km above the lunar surface, corresponding to a surface field of ~ 300 nT

whenever the solar wind density is $<4 \times 10^7 \text{ m}^{-3}$. Such lunar mini-magnetospheres have a horizontal extension of the order of 100 km and are extremely vulnerable with respect to variations in solar wind pressure and magnetic field direction. They would appear and disappear, shrink and broaden and wander around depending on solar wind conditions making the lunar surface a magnetically living animal with magnetospheres and bow shocks arising and dying away in regions of high magnetisation or high conductivities. Average observed local lunar surface magnetic field strengths at 100 km altitude are of the order of $\sim 2 \text{ nT}$ and should not lead to the generation of mini-magnetospheres and mini-bow shocks.

These considerations of the Moon and its localised magnetic fields can be extended also to the larger asteroids of sizes well above the solar wind ion gyro-radius. If these objects are magnetised or consist of highly conducting material, mini-magnetospheres might form around them either being supported by the internal magnetism or by the magnetic fields that can be induced in their conducting interiors. In both cases mini-bow shocks might exist around them having similar properties as the proposed lunar mini-bow shocks, deflecting the solar wind and contributing to its very localised heating.

10.3.2 Mars – A Pile-Up Induced Bow Shock

It has been believed for long that Mars possesses an internal magnetic field though weak but strong enough to form a magnetosphere around the Red Planet that has been named after the Roman *Deus Belli*. Mars was always the beloved darling of man onto that he has dreamt away to build his cities and bunkers under the Martian crust after having escaped from Earth. Previous spacecraft like Phobos and Mars Pathfinder could not definitely exclude the existence of a Martian magnetic dipole field. The proof that Mars has no global internal field was left to the Mars Global Surveyor [Acuña *et al*, 1998, 2001]. The Phobos spacecraft had detected substantial ultra low frequency wave activity in the magnetic field during the Mars approach [Russell *et al*, 1990] that was attributed to a Martian bow shock wave respectively the existence of a Martian global field.

The Mars Global Surveyor set an upper limit of $B < 0.5 \text{ nT}$ [Acuña *et al*, 2001] on a global surface field on Mars. To the great surprise, when mapping the magnetic surface field of Mars, the Mars Global Surveyor found that Mars possesses several cratereous regions of fairly strong local remanent magnetic fields of magnitude $B \lesssim 200 \text{ nT}$, indicating that in early times Mars was magnetised, had a working dynamo and left some signatures of these times in his crustal minerals. Thus, had Mars no atmosphere it would today behave very similar to the Moon, lacking a global magnetosphere and a global bow shock but instead being covered with regions of mini-magnetospheres and mini-bow shocks which reacted highly dynamically to the fluctuations in solar wind pressure and magnetic field.

However, Mars has an own thin atmosphere, the existence of that was confirmed. Similar to Earth's atmosphere, the Martian atmosphere becomes ionised by the solar ultraviolet radiation and forms an ionosphere on the illuminated side of the planet which prevents the solar wind from impacting the Martian surface. Mars, in this sense, behaves similar to Venus with the main difference that its atmosphere and hence its ionosphere as well are thin. Nevertheless pressure balance between the solar wind and the ionosphere creates

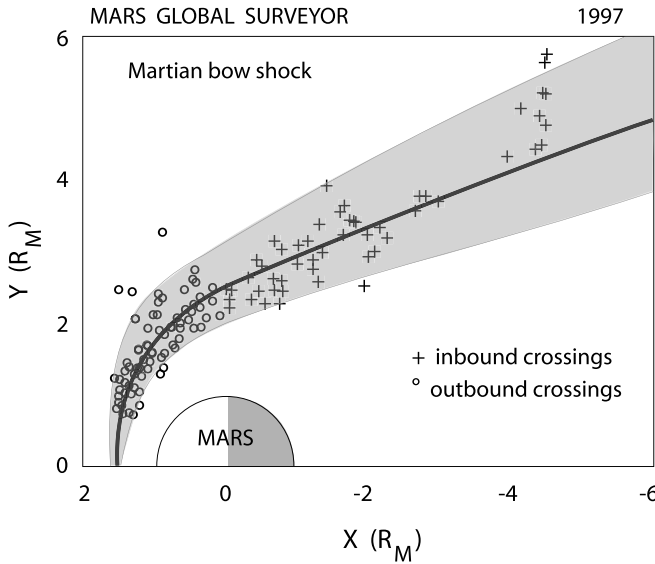


Figure 10.12: The Mars Global Surveyor passages across the Martian bow shock. Circles indicate the inbound passages which were all near the stagnation point. The outbound crossings were at the morning flank of the bow shock. Large flappings of the shock are seen by the scatter in the data. The range of variation of the location of the Martian bow shock has been indicated by shading [data taken from *Acuña et al., 1998*].

an ionopause and a blunt detached bow shock that stands in front of the ionopause at a planetocentric stand-off distance of $\sim 1.5 R_M$.

The Mars Global Surveyor spacecraft crossed the Martian bow shock many times on its outbound and inbound passes before settling on a circular mapping orbit around Mars at altitude of ~ 400 km. Figure 10.12 shows a collection of the identified locations of the bow shock in the projection onto the (x, y) -ecliptic plane [data taken from *Acuña et al., 1998*]. The inbound crossings are exceptionally at the flank-side bow shock and are at distances farther away from the planet than the outbound crossings which were located on the morning side of the planet. It is interesting to see that the shape of the Martian bow shock is highly non-stationary, obviously flapping around an average position with quite large amplitude. Moreover, the bow shock seems to be more blunt on the dayside than other planetary bow shocks and might inflate even more at larger tailward distances. There are different reasons for this flapping. It may be caused by fluctuations in the solar wind pressure as well as by fluctuations in the Martian ionospheric ionisation that is affected by the atmospheric conditions on Mars. Part of it may also be affected by the average local crustal magnetic field strength over one Martian rotation which contributes to the resistance of Mars to the solar wind, but the main reasons can probably be seen in the interplay between the atmospheric and solar wind conditions.

Apart from this global shape of the Martian bow shock, the microscopic signature is physically of more interest. Figure 10.13 shows some of the Mars Global Surveyor measurements during two of the first crossings of the Martian bow shock. The top two panels give

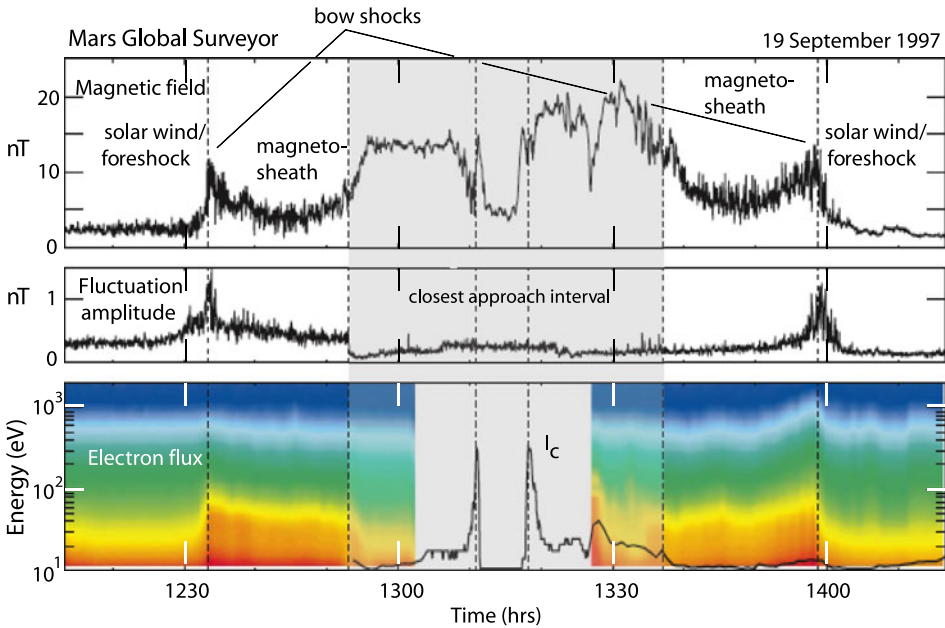


Figure 10.13: The Mars Global Surveyor passage near the red planet on 19 September 1997 as seen in the magnetic field, ultra low frequency magnetic fluctuations and electron flux measurements [data taken from *Acuña et al., 1998*]. Contact with the bow shock is preceded by increased magnetic fluctuations and increasing electron energy. Two shock crossings (inbound and outbound) are seen as steep magnetic transitions, peaked fluctuation levels and high electron thermal spreads, followed by a broad magnetosheath. The obstacle here is not the magnetic field but the thin Martian atmosphere. The *black line* in the electron spectrogram is the Langmuir probe current which is a measure of electron densities and temperatures close to the planet.

the magnetic recordings: magnetic field strength and rms fluctuation amplitude versus time on 19 September 1997. The bottom panel gives the colour coded spectrogram of the electron energy fluxes (in electrons/s cm² sr eV) in the energy range $10\text{ eV} \lesssim \mathcal{F}_e(\mathcal{E}) \lesssim 2\text{ keV}$. Highest fluxes are shown in red, lowest fluxes in blue. The absolute scale of the fluxes is not given here. The shaded region in the figure is the inner part of the spacecraft passage at closer approach to the planet which does not interest us in this connection except for the strong increase in the magnetic field before reaching the ionopause.

The spacecraft crosses the shock on its inbound path at about 1235 Spacecraft Time. Prior to this crossing it is in the solar wind detecting enhanced fluctuation levels in the low frequencies with increasing level the closer it comes to the shock. There is a well expressed shock foot region with high and nearly constant fluctuation amplitudes from 1230–1234 Spacecraft Time. The shock crossing exhibits a steep ramp and overshoot, both of them are also seen in the peaked fluctuation levels. These are followed by a broad turbulent magnetosheath region with enhanced but lower levels until spacecraft entrance into the increasing magnetic field which is a magnetic pile-up boundary standing in front of the ionopause at the start of the shaded area. This pile-up is created by the draping and

compression of the interplanetary magnetic field around the highly conducting ionosphere that inductively screens the planet from the interplanetary field. It is this magnetic pile-up that acts as the obstacle for the solar wind and forces the detached bow shock to develop. Would it not exist, would the bow shock probably be attached to the atmosphere due to the unbroken penetration of the solar wind into the atmosphere.

The electron energy fluxes also show an increase in energy with entrance into the bow shock foot region, corresponding to electron heating, and a broad intense burst of electrons related to the shock crossing and showing its signature up to ~ 1 keV. Clearly the shock is a region of plasma heating also in this case. However, in the lowest energy channels the electron fluxes have become very strong indicating that more electrons participate than are present in the solar wind, which implies that the plasma is compressed. Indeed, the compression in the magnetic field from the ~ 2 nT solar wind field to the $\sim (5-6)$ nT shocked magnetosheath field indicates a moderately strong shock compression factor of about 2–3. Note also the decrease in electron flux and cooling at the magnetic pile-up boundary.

The outbound shock crossing at 1358 is similar but with a much longer and better developed foot region, compression of roughly a factor 3, higher overshoot exceeding 10 nT and a broad downstream magnetosheath region which, however, is less turbulent than the inbound magnetosheath as may be inferred from the rms amplitude of the magnetic field. The fluctuations concentrate solely on the shock transition and parallel the stronger electron heating which now extends to electrons > 1 keV. We should remember that the first inbound crossing was a tail bow shock crossing while the outbound crossing is a dayside crossing that is closer to the planet in the stronger compressed magnetic field and plasma. In both cases, however, the shocks are still similar to the known shocks of the other planets, both being quasi-perpendicular. This must be so in the case of Mars because the Martian ionosphere almost screens the solar wind magnetic field from the planet. This Mars has in common with his sister planet Venus.

10.4 Venus and the Comets

In contrast to the other planets and the Moon, Venus and the comets have in common that they are non-magnetic and both have dense and extended atmospheres even though their atmospheres are of completely different origin. Thus the interaction of Venus and the comets with the solar wind is basically different from that of the other planets. One, however, expects that Venus gives rise to the formation of a strong shock wave that stands in front of the planetary atmosphere either as a detached or atmospherically attached wave. The reason is that Venus is large enough to deflect the super-magnetosonic supercritical solar wind around its body and dense extended atmosphere while the velocity of compressional waves that are needed to maintain the deflection of the solar wind is too small to compete with the solar wind flow. Whether comets behave the same way like Venus or not, depends, however, on their particular properties. Before coming to discuss some of the observations of the cometary interaction with the solar wind in the inner heliosphere we turn to Venus and what has been learned about its interaction with the solar wind from the various observations *in situ* its environment.

10.4.1 The Venusian Bow Shock

A Venus-like interaction is defined as an interaction that is associated with a planet that has no intrinsic magnetic field but a substantial atmosphere. Indeed, the upper limit for the intrinsic magnetic field of Venus (set by the Pioneer Venus Orbiter after initial measurements by the Venera satellites) is $\sim 10^{-5}$ times that of Earth. Venus also does not have any remanent crustal magnetic field left over from an early active magnetic dynamo period, because the temperatures in Venus' crust are expected to lie above the Curie point implying complete destruction of any remanent magnetism. The solar wind reaches directly to the top of the ionosphere, assuming that an ionosphere exists and must not be created by solar wind impact. This has been the case for Mars as we have discussed above. Venus, however, with its denser atmosphere is, in addition, much closer to the Sun (cf. Table 10.1). Being at about half the distance of Mars the solar radiation is four times more intense than at Mars and 1.6 times more intense than at Earth. The dense Venus atmosphere absorbs this optical and extreme ultraviolet radiation and produces a highly conducting ionosphere at altitudes below 300–800 km with maximum ionisation between 125–170 km altitude. Venus could possess an induced magnetic field if its planetary conductivity would be high enough. However, this is prevented by the presence of the ionosphere that screens the lower atmospheric layer and the planet from the solar wind and interplanetary field. The interaction between the solar wind and Venus is thus directly with the Venusian ionosphere [for reviews see, e.g., Luhmann, 1986; Luhmann & Cravens, 1991] which forms the obstacle to the solar wind. The surface of pressure balance between solar wind ram pressure and ionospheric thermal plasma is the ionopause. Farther out the solar wind magnetic field piles up, a blunt bow shock [see, e.g., Wolfe *et al*, 1979; Scarf *et al*, 1979c] is generated at a distance upstream leaving a magnetosheath like region between the shock and the ionopause. The stand-off distance of the bow shock is found at a planetocentric distance of $\sim 1.5 R_V$ at solar maximum and depends on solar activity as this regulates the ionospheric electron content. In the magnetosheath the solar wind is deflected around the ionopause in a way similar to the deflection around a magnetopause. Small-scale interplanetary magnetic flux tubes can sometimes break through the ionopause being found deep in the atmosphere where they form so-called flux ropes which remain to be connected to the solar wind magnetic field and are convected around the planet.

The solar wind density at Venus is four times higher than at Mars, and the solar wind dynamic pressure is twice that at Earth's orbit. The magnetosonic Mach number at Venus is $\mathcal{M}_{ms} \approx 5.1$ while the solar wind β is slightly less than at 1 AU. In the average the solar wind at Venus is super-critical but not of vastly high Mach number. On the other hand, the spiral solar wind magnetic field is in the average more inclined at Venus than at the outer planets such that one expects that its prospective bow shock possesses a fairly large quasi-parallel area, depending on the strength of the magnetic deflection in the magnetosheath. Such an interaction does not produce a magnetotail like on Earth because there is no planetary magnetic field which could be stretched out into space. An important point is that the magnetic field shows a strong shock of compression ratio ~ 3 in both cases with sharp overshoots. Such shocks should reflect particles upstream along the tangential field and have foot regions. The latter are indeed seen in these crossing, and the inbound cross-

ing indicates also the presence of a foreshock as can be judged from the relatively large amplitude magnetic fluctuations detected ~ 4 –5 minutes before shock entrance.

Backstreaming ions from the Venus bow shock along the magnetic field have been reported in earlier [Frank *et al.*, 1991; Williams *et al.*, 1991] passages of other spacecraft, in particular from the Galileo spacecraft on 10 February 1990. Such ions have been observed in the keV energy range [Moore *et al.*, 1989] and might be the seed population for the more energetic ions in the 120–280 keV energy range seen by Galileo [Williams *et al.*, 1991] after having been accelerated at the Venusian bow shock. Whether this is possible is not known as the energetic ions have gyroradii larger than the curvature radius of the small Venusian bow shock. It is more probable that these ions have been accelerated in a different way, either directly in the Venusian ionosphere, or by Venus lightning, or by bouncing around in the piled-up magnetic field of Venus, in which case the Venusian bow shock would not be responsible for their existence and would produce only the low energy keV ions.

On the other hand, the search for pick-up ions leaking out from the Venus atmosphere and ionosphere into the solar wind gave a low upper limit of $N_{pui} \lesssim 10^3 \text{ m}^{-3}$ for the pick-up ion density in front of the bow shock. This is still consistent with the scale height of the Venus atmosphere. Unfortunately, the measurements could be performed only outside the magnetosheath. Earlier observations [Nagy *et al.*, 1981] by Pioneer Venus in the magnetosheath had given indication for the presence of a substantial number of pick-up ions behind the bow shock which are accelerated in the convection field as well as the strongly fluctuating electromagnetic field in the magnetosheath.

The electron observations by Galileo, on the other hand, confirmed that the Venus bow shock is quite strong exhibiting density increases of a factor ~ 3 , the same as concluded from the magnetic observations. However, only little electron heating $\lesssim 20\%$ was observed in the Galileo shock crossings, which contrasts with the bow shocks of magnetised planets and is also not in agreement with the observations at Mars dealt with above where substantial electron heating by a factor of 2–3 has been found.

Venus has been passed since the 1960s by a substantial number of spacecraft. The most recent mission to Venus was Venus Express. Figure 10.14 shows an example of the high resolution magnetic field recordings by Venus Express on 27 June 2006 with two bow shock crossings [Zhang *et al.*, 2008]. The first crossing being preceded by a short encounter with the Venusian foreshock, followed by the passage of the shock foot and the ramp and overshoot of the bow shock. Afterwards the spacecraft was in the magnetosheath until passing the magnetic pile-up boundary, pile-up region and ionopause followed by an outbound repetition of the same sequence. In both cases the shocks were quasi-perpendicular as indicated by the sharp shock fronts, the overshoot and the foot regions. This passage is very similar to that seen above at Mars showing the close relationship between the bow shocks on these two otherwise so different planets. Similar observations were made by Galileo in February 1990 during the spacecraft encounter with Venus [Kivelson *et al.*, 1991]. Galileo crossed the shock and mapped its width to be shorter along the interplanetary field than transverse to it, confirming the near perpendicular shock structure. Moreover, Galileo also found large amplitude ultra-low frequency magnetic waves in front of the bow shock the amplitudes of which peaked in the magnetosheath and were attributed to the presence of pick-up ions. It was also speculated that some of the structures found might be interme-

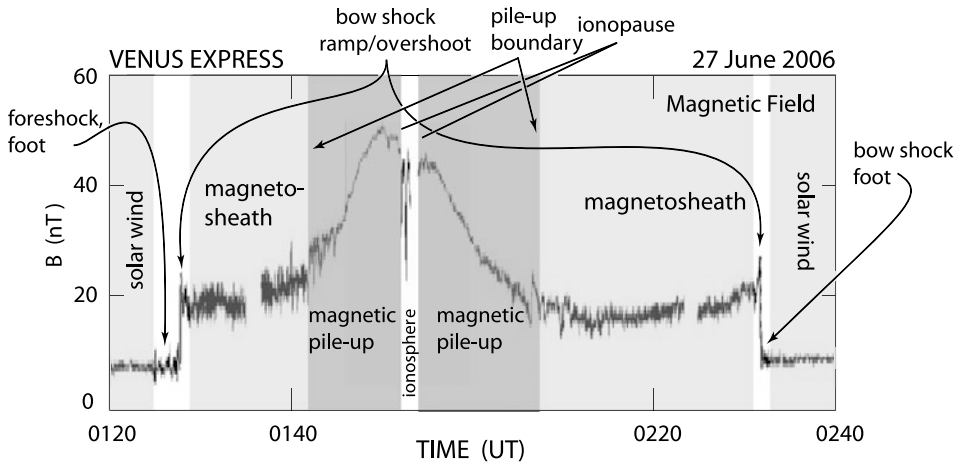


Figure 10.14: Recent magnetic field recordings during a Venus Express passage of the Venusan bow shock [data taken from [Zhang et al., 2008](#)], showing the various different regions and magnetic boundaries in the environment of Venus. During inbound the spacecraft detected a weak foreshock exhibiting magnetic fluctuations, a short foot region, a quasi-perpendicular bow shock ramp and overshoot, entered a disturbed magnetosheath and passed the magnetic pile-up boundary until reaching the ionopause. On the outbound path no foreshock was detected except for a short shock foot.

diated shock waves which could, however, not be proved definitely and thus is highly speculative.

A compilation of long-term magnetic observations of the Venusan bow shock by the Pioneer Venus Orbiter is shown colour coded in [Figure 10.15](#). This figure shows the spatial extension of the region between the bow shock and the ionopause of Venus for solar maximum and solar minimum in the local amplitude of the magnetic field reflecting the control of the Venusan bow shock, external magnetic field and ionosphere by solar activity. At solar minimum the bow shock is found at planetocentric distance substantially farther out than during solar maximum. On the other hand, during solar maximum the magnetosheath and magnetic pile-up regions are stronger compressed and more variable. The reasons for these differences can be found in the fact that during solar maximum the solar wind fluctuates more strongly and has higher Mach number and dynamic pressure causing a stronger compression of the ionopause and a stronger pile-up of the magnetic field than during solar minimum when the solar wind is relatively quiet. Moreover, during solar minimum the Sun is hotter which causes a stronger ionisation of the Venusan atmosphere and thus an increased plasma density and thermal pressure in the ionosphere which lets the ionosphere expand farther out from the planet.

We finally turn to the observation of high frequency waves during the Galileo Venus bow shock encounters [[Gurnett et al., 1991](#); [Hospodarsky et al., 1994](#)]. The plasma wave instrument on the Galileo spacecraft detected the shock in both the electric and magnetic instruments as broadband impulsive noise similar to what is observed in other shock crossings, followed by enhanced plasma wave turbulence in the adjacent magnetosheath. Close

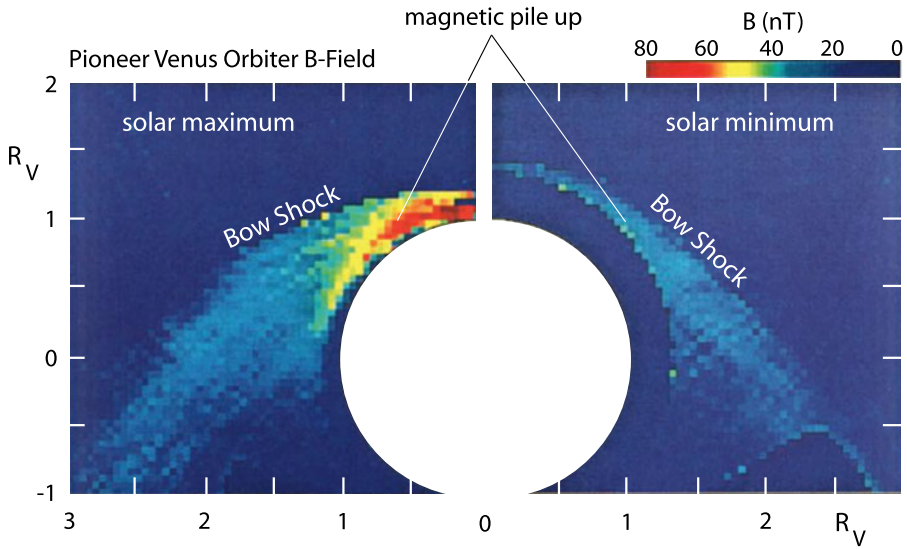


Figure 10.15: The Pioneer Venus Orbiter magnetic field measurements near Venus during solar maximum and solar minimum, showing the control of the magnetic pile up region and bow shock position with solar activity [data taken from *Kallio et al., 1998*, courtesy American Geophysical Union]. At high solar activity the pile up is much stronger and the bow shock stand-off distance smaller than at low solar activity.

upstream to the shock an intense electromagnetic emission band between 5–50 Hz was detected. The nature of this band is consistent with the assumption of electromagnetic whistler or ion cyclotron waves. Since no supra-thermal particle component could be identified from the plasma instruments during this time these waves are probably caused at the shock from where they propagate to spacecraft position in the solar wind. The geometry of how this can happen is not clear, however, as the solar wind is super-Alfvénic.

Electron plasma waves at the higher frequencies of 10–50 kHz in the Langmuir wave band have been observed upstream of the shock in complete analogy to other bow shocks. An example from the Galileo observations is shown in Figure 10.16 exhibiting a very close similarity of these oscillations to those observed in front of Earth’s bow shock and also at the other magnetised planets. These waves had been seen before by Pioneer Venus in May 1979 [*Scarf et al., 1980*]. They consist of two distinct types, the Langmuir waves which map the plasma frequency at roughly ~ 45 kHz and intensify in the presence of a fast electron beam, and a sporadic upstream downshifted emission which signals the presence of supra-thermal electrons and disappears abruptly when the electron foreshock boundary is crossed into the undisturbed solar wind. The electrons come from the bow shock where they must have been accelerated. Our knowledge of the reflection and acceleration mechanism tells that the Venusian bow shock contains highly localised electric field structures in order to produce such intense supra-thermal electrons. It thus behaves very similarly to the bow shocks of the magnetised planets. *Hospodarsky et al [1994]* investigated these oscillations in more detail finding the expected close similarity to the plasma oscillations

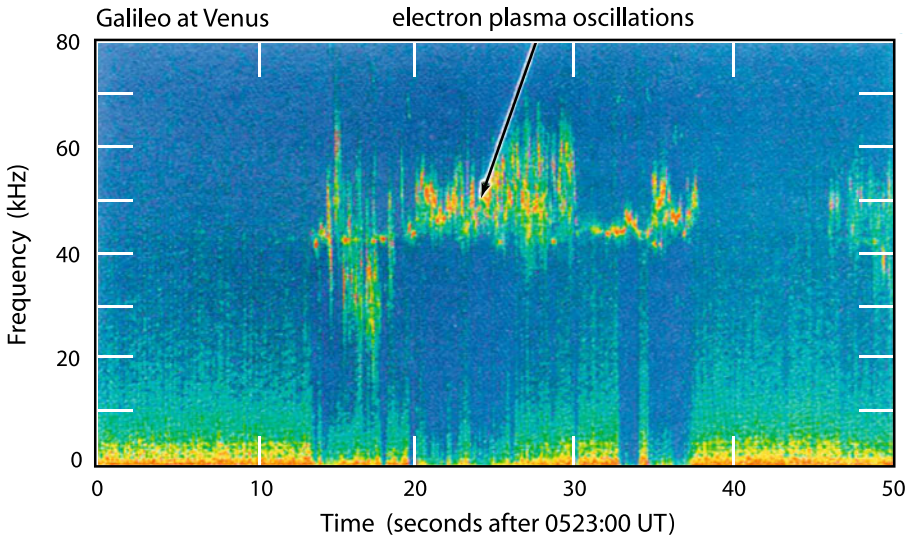


Figure 10.16: Galileo observations of electron plasma waves in the Langmuir mode as are excited by electron beams emitted from the Venusian bow shock into the electron foreshock. Intensification of the waves occurs when the spacecraft crosses the tangential magnetic field line [after *Gurnett et al, 1991*]. Note the high time variability which indicates wave localisation. The disappearance of the low frequency ion acoustic waves in response to the Langmuir waves is also of interest.

and electron plasma waves in the Earth's electron foreshock. These waves are strongly localised and form structures of Debye length scales, in this case the order of $\sim 60\text{--}70$ m, which are generated by nonlinear interaction between the supra-thermal shock-accelerated electrons and the plasma wave component.

We may thus conclude, that under normal solar wind conditions at the Venus orbit (i.e. high speed solar wind, nominal plasma densities and magnetic field strengths) the bow shock behaves similar to the supercritical bow shocks of magnetised planets, i.e. the upstream flow does not care what kind the obstacle is as long as the obstacle consists of a conducting plasma layer which in this case (and the case of Mars) is the ionosphere. This layer forces the magnetic field to pile-up upstream of it thereby inducing a magnetic obstacle which, however, is itself of purely upstream nature, while the ionosphere becomes a screening current layer. This obstacle deflects the solar wind and generates a strong supercritical bow shock wave that is detached from the Venusian ionosphere.

However, occasionally the normal conditions at Venus become violated when slow dense plasma streams encounter the planet. Such cases have recently been reported from Venus Express, when the spacecraft crossed the Venus bow shock several times under sub-critical conditions with Mach number $\mathcal{M} < 1.5$ at shock normal angles $\Theta_{Bn} > 70^\circ$ when no ion reflection takes place [*Balikhin et al, 2008*]. Such cases have not been described in this book as in the heliosphere they are found very rarely and are expected to occur only in its innermost part and under extreme dense slow solar wind conditions. The upstream

solar wind at the times under question was indeed unusually dense. Nonetheless, the conditions at the bow shock were still collisionless such that the shock was probably not maintained by resistive dissipation. The magnetic shock profile was of a kind similar to a Korteweg-deVries shock with either upstream or downstream attached spatially damped oscillations. Since β was small, such a shock should be maintained ballistically with the complete cold solar wind ion beam gyrating in bulk with negligible thermal spread in the shocked magnetic field. We do not go further into this kind of observation as it represents a rare, though interesting, case of subcritical shocks that will be realised only in very dense and comparably slow plasma streams which happen to exist in proximity to the Sun (cf. Chapter 4). Such shocks might exist more frequently in the solar corona or outside our cosmic environment near exoplanets. Also Mercury may sometimes develop this kind of a bow shock when immersed into a dense and slow solar wind stream as generation of such subcritical shocks does not depend on the existence of a neutral atmosphere but only requires the presence of an obstacle like Venus' ionosphere with induced magnetic pile-up or Mercury's magnetosphere.

10.4.2 Cometary Bow Shocks

The first claim of a possible crossing of a cometary bow shock by the Pioneer Venus spacecraft goes back to *Russell et al* [1984]. Comets need not necessarily have bow shocks. Comets consist of a relatively small icy nucleus of the size of an asteroid. This nucleus is covered by a comparably thin mantle of mostly carbon atoms mixed with dust. Where this nucleus is damaged it evaporates and out-gasses attributing an extended gaseous cloud to the comet. However, far away from the Sun comets out-gas very weakly if at all. There is no or very little photo-ionisation. Ionisation is mainly due to charge exchange. Thus the mass added to the solar wind by the comet is negligible. The solar wind that flows around the comet might well induce a magnetic field in the cometary body or in its highly conducting crust layer, but because of the extremely small size of the nucleus the distortion of the solar wind at those distances is miniscule and might not even exist if the particle gyroradii exceed its diameter. However, as the comet approaches the Sun, solar radiation intensity increases and heats up the illuminated part of the comet's outer shells such that the shell starts melting in several places, the ice liquefies and ultimately evaporates, causing cometary outgassing at a susceptible rate until the comet is surrounded by a huge gas cloud many times larger than the nucleus. Part of this gas envelope becomes photo-ionised, charge exchange between solar wind protons and the gas atoms or molecules removes gas particles from it and adds them as pick-up ions to the solar wind. Both processes lead to mass and momentum exchange between the solar wind and the gas envelope of the comet until the rate of this exchange becomes large enough to decelerate and deflect the solar wind in the region where the interaction takes place. The processes involved have been investigated in detail by *Biermann et al* [1967], *Cloutier et al* [1969], *Wallis* [1973a, b] and others, and it has been shown that at large addition rates of mass to the solar wind a cometary bow shock must evolve in order to maintain the deflection of the solar wind in the interaction with the cometary cloud. The reason is again very simply that the slow

compressional waves in the cometary environment are insufficient to deflect the super-magnetosonic solar wind flow around the comet. Hence, a shock must be generated in front of the comet in order to break the flow and divert it around the cometary obstacle.

Comet Giacobini-Zinner

Investigating the nature of comets and their interaction with the solar wind has for long time been one of the major fields of interest in space and planetary physics because of many reasons not the least one the proposal that possibly comets have been responsible for transporting germs in the solar system between the planets and contribute to the spread of life. The first mission to pass by a comet was the International Cometary Explorer ICE. In 1985 it approached the comet Giacobini-Zinner passing it at a distance of ~ 15000 km.

Apart from its optical appearance on the sky, *in situ* information from the comet was obtained first by the ICE plasma wave instrument [Scarf *et al*, 1986] which detected bursts of strong ion acoustic waves almost continuously at times when ICE was closer than 4×10^6 km from the Giacobini-Zinner nucleus. It also observed oscillations at the electron plasma frequency and whistler waves, all being believed to be related to the presence of heavy cometary pick-up ions as well as hot electrons within this region surrounding the comet. In particular the presence of electron plasma waves could suggest magnetic connection to the magnetic field penetrating the cometary coma and allowing supra-thermal electrons to escape from it to reach the spacecraft. However, from magnetic field observations there was no indication for such a connection. Hence the plasma waves were also attributed to the presence of pick-up ions, which might indeed be the case as these should locally excite lower hybrid waves which are known to accelerate electrons into beams along the magnetic field. This whole region was indeed filled with energetic cometary ions (O^+ and OH^+) of energy exceeding 30 keV [Hynds *et al*, 1986; Ipavich *et al*, 1986]. Closer to comet Giacobini-Zinner within a few 100 000 km distance from the nucleus ICE measured increased wave levels and broadband signals. In fact, the signal intensity was reported to increase by more than nine orders of magnitude [Scarf *et al*, 1986], causing saturation of the instrument in some of the channels. At the boundaries of this inner region impulsive broadband noise was measured, preceded upstream with enhanced electron plasma waves (see Figure 10.17) which strongly reminded of shock crossings that are accompanied by electron heating. The electron plasma oscillations indicate magnetic connectivity and generation of electron beams, which is typical for bow shocks. However, only for the second outbound crossing on this day at 1219 UT the magnetic signature of a possible shock could be identified [Smith *et al*, 1986], making this crossing more conclusive of a bow shock pass even though Smith *et al* [1986] prefer to speak of a bow wave also in this case. Because of the perpendicular direction of the interplanetary field at the time of passage one would have expected a quasi-perpendicular shock in this case such that the observed lack of a steep shock signature in the magnetic field can hardly be attributed to the passage of a parallel shock. Moreover, the plasma instrument [Bame *et al*, 1986] did not show the expected signature of a conventional shock transition in neither of the

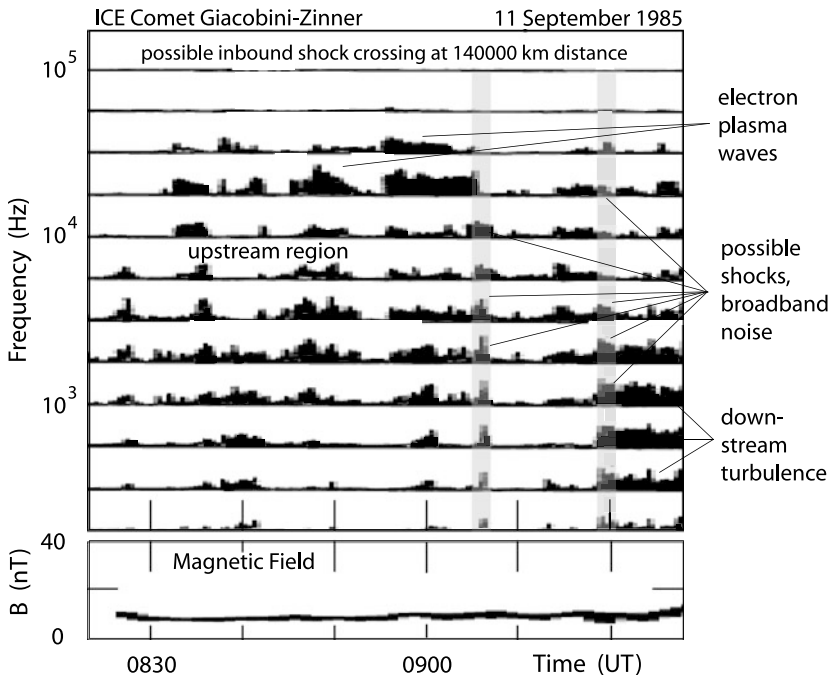


Figure 10.17: The ICE approach to comet Giacobini-Zinner on 11 September 1985 as seen in the plasma waves [after Scarf *et al.*, 1986]. Before the onset of the broadband noise on the inbound passage which is typical for a shock crossing, the also typical electron plasma waves upstream of a shock were detected. Below these broader low frequency waves are seen. After the broadband impulsive shock noise the spacecraft entered the downstream magnetosheath region with intense low frequency turbulence. However, these wave observations are still inconclusive of the presence of a shock as no shock signatures were seen in the magnetic field [taken from Smith *et al.*, 1986] and in the other instruments. The gradual pile-up in the magnetic field starts at about 0940 UT after the observation of the broadband noise and electron plasma waves and reaches up to 20 nT on the inbound pass.

inbound or outbound crossings tempting the authors to speak of a slow transition region instead of a shock. Bame *et al.* [1986] speak of plasma signatures distinctly different from a conventional magnetosonic shock transition. On the other hand a broad magnetic pile-up region was passed exhibiting large-amplitude nonlinear magnetic fluctuations that resemble the upstream pulsations of quasi-parallel shocks with steep fronts and attached high frequency oscillations. On the inner edge of the pile-up region the magnetic field showed signatures of an ionopause thus resembling the crossings of the Venusian shock-pile-up region. Hence, the cometary transition region is by far structured more complicated than the transition regions of planets.

Thus the ICE observations of comet Giacobini-Zinner were inconclusive about the possible existence of a bow shock at a comet in the inner heliosphere. A conclusion was expected from observation of comet Halley.

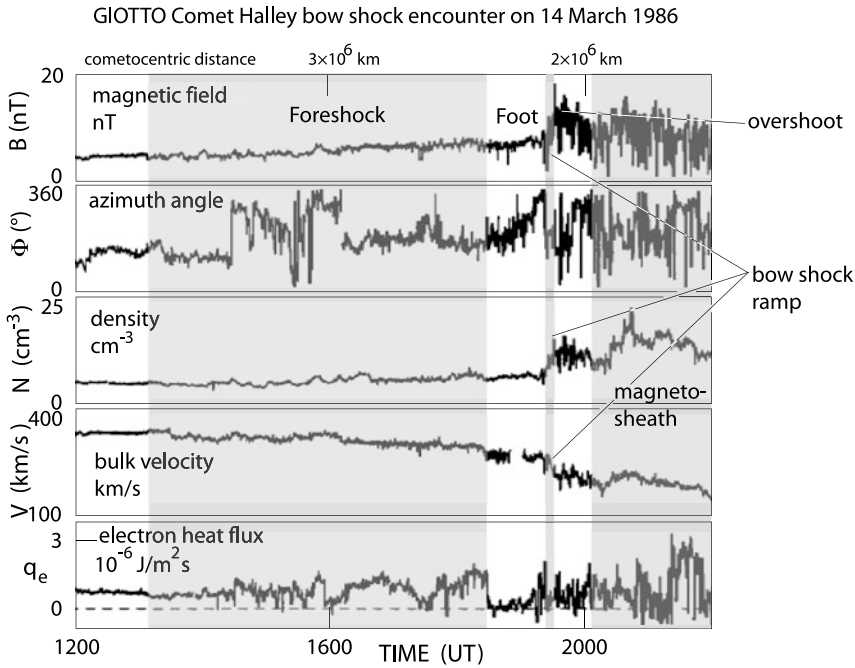


Figure 10.18: Giotto passage across the Halley bow shock on 14 March 1986 [data taken from *Larson et al, 1992*, courtesy American Geophysical Union]. Proving the existence of a cometary bow shock. The spacecraft entered the foreshock at about 1400 UT just before an interplanetary discontinuity passed over it which is seen in the turning of the magnetic field. After 1600 UT it was again in the foreshock, entering the foot region with the continuously increasing magnetic field and crossed the shock ramp at about ~ 1930 UT. The shock compression factor was about 2, and the shock was oblique, neither parallel nor perpendicular.

Comet Halley

The two most spectacular cometary events near the end of the past century were the destruction of comet Shoemaker-Lévy when it hit Jupiter causing one of the greatest nuclear explosions in the heliosphere that could be observed from remote, and the human borne spacecraft visits to comet Halley.

Halley was visited by three spacecraft, the two Russian spacecraft Vega 1 and Vega 2, and by Giotto. The Vega satellites reported the possible observation of a stand-off shock near Halley [*Galeev et al, 1986*]. The existence of a bow shock at comet Halley was confirmed by the precise measurements of Giotto [*Neubauer et al, 1986*]. This shock was crossed two times; on the inbound path, shown in Figure 10.18, the shock was found at cometo-centric distance ~ 1.2 Million km and was preceded by upstream waves. On the outbound crossing the shock was seen as a broad turbulent quasi-parallel shock. Figure 10.18 shows a collection of the observations on 14 March 1986. At around ~ 1930 UT Giotto crossed the shock ramp and entered the overshoot and magnetic pile-up magnetosheath region. It has before that time been in the foreshock where intense electron fluxes

were detected flowing along the interplanetary magnetic field. Deceleration of the solar wind indicated entrance into the shock foot where the magnetic field started increasing and turning in direction while the electron heat flux decreased, an observation that coincides with what is known from planetary bow shocks.

The shock ramp is seen in the steep increase in the magnetic field, strong deceleration of the flow, localised heat fluxes and violent changes in the direction of the magnetic field. The shock is oblique, probably more quasi-parallel than quasi-perpendicular. One of its peculiarities is that the transition takes quite a long time, much longer than on any of the magnetised planet shock crossing where the shock ramp is steep. Moreover, the determination of all these regions is not as clear as in the case of planetary bow shocks or even in the cases of Mars and Venus. The plasma and field quantities in the comet passage are more disturbed, take longer time to adjust to the downstream values, and the compression ratio is relatively small such that the shock is less strong than at a planet. Nevertheless, the data show definitely that large comets like Halley do develop a bow shock in interaction between their gaseous environment and the solar wind when entering the inner heliosphere. However, the properties of the cometary shocks differ to a certain degree from the known planetary bow shocks.

Coates et al [1990] have undertaken a thorough check of the various shock plasma parameters to confirm that the shock that was crossed by Giotto was indeed a shock. They found that the magnetosonic Mach numbers for the Halley shock were quite low, the inbound crossing during passage gave a Mach number of $1.1 < \mathcal{M}_{ms} < 1.8$, the Mach number of the outbound crossing was more stable $1.6 < \mathcal{M}_{ms} < 1.7$. Thus the Halley bow shock was quite weak and just marginally supercritical if the shock normal angle was small enough. The reason for this weakness may be found in the large number of heavy pick-up ions that have been loaded into the solar wind at and around shock location. Deceleration of the solar wind by this heavy mass loading clearly reduces the Mach number. This is shown in the lower panel of Figure 10.19 where the solar wind velocity and the number density of cometary O^+ -ions is plotted during shock approach and crossing. The anti-correlation of both quantities is obvious. The solar wind speed decreases gradually with increasing heavy cometary ion concentration during approach of the comet. The upper part of the figure shows the measured proton energy distribution for those protons which come from the direction of the comet. Before crossing the shock foot and ramp (shaded vertical bars) reflected or escaping protons are measured arriving from Halley. Note that these protons have low energy spread but relatively high energy ~ 1 keV, part of them reaching energies of several keV indicating acceleration in the shock or shock foot region. (Heavy ion acceleration up to < 250 keV at Halley has also been reported from downstream observations [Kecskeméty & Cravens, 1992]. However, this acceleration is not necessarily due to the presence of the Halley bow shock.) After crossing the shock a broad thermal proton distribution of lower energy < 1 keV is found that is typical for the downstream particle distribution. *Gombosi et al* [1991] determined the width D of the region of heavy ions (O^+) outside the bow shock along the Giotto trajectory of being in the spatial interval between $2 \times 10^6 \text{ km} > D > 1.2 \times 10^6 \text{ km}$. This thickness is determined by the ability of the particles to diffuse away from the comet.

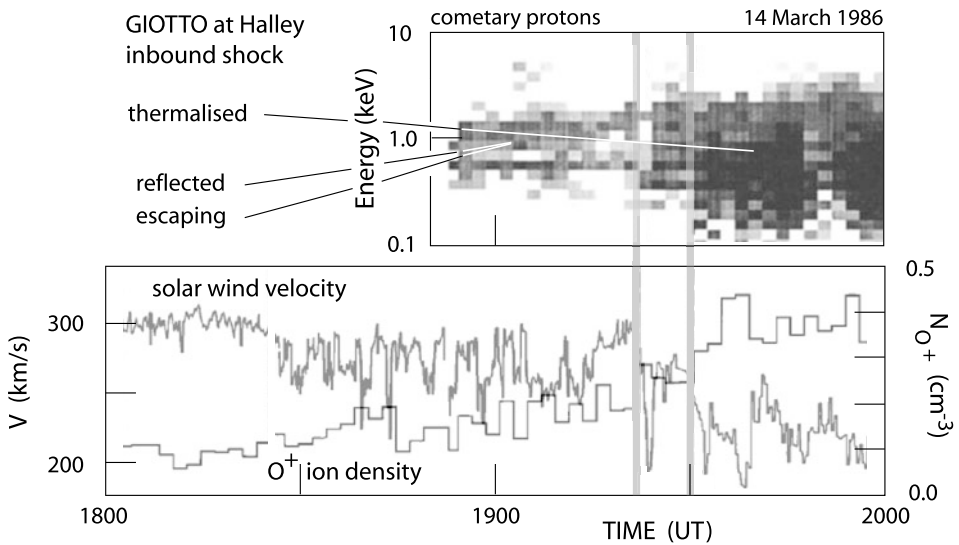


Figure 10.19: Giotto passage across the Halley bow shock on 14 March 1986 [data taken from *Coates et al., 1990*, courtesy American Geophysical Union]. The *lower panel* shows the anti-correlation between the solar wind velocity and water ion content during approach of comet Halley and inbound shock crossing suggesting that mass loading by heavy cometary ions is responsible for the slowing down of the solar wind. The *upper panel* shows protons arriving at Giotto from the direction of Halley during approach of the comet. Cold reflected/escaping protons of energy ~ 1 keV are detected prior to entering the foot and shock ramps, some of the ions have energies up to 5–6 keV, indicating acceleration either in the shock or foot region. Downstream the mean proton energy has decreased to ~ 0.3 keV corresponding to the low mean bulk speed in the *lower panel*. At the same time the proton distribution has been thermalised in the shock transition as see from the ~ 1 keV spread of the distribution.

Giacobini-Zinner and Halley have so far been the only comets from which *in situ* data have become available. Much theoretical and simulational work has been done to investigate the formation of shocks in front of a comet entering the inner heliosphere, work that was mostly based on magneto-hydrodynamic considerations including cometary ions. Even the formation of a shock during Shoemaker-Lévy 9 approach of Jupiter and its interaction with the Jovian bow shock has been simulated [*Lipatov & Sharma, 1994*] in order to show the deformation of the comet, ion heating and particle acceleration. Still, the formation of a shock around a comet has not yet been clarified satisfactorily. The observations suggest that sometimes, predominantly at larger comets, shocks develop. However their structure is different from planetary bow shocks because of mass loading, decrease in Mach number and the differences between the gaseous shells around a comet and the atmosphere of a planet like Venus.

10.5 Conclusions

This chapter has focussed on the discussion of the various planetary bow shocks that can form around the celestial bodies in the solar system and heliosphere. By the nature of

the question this discussion was divided into a review of earth-like bow shocks, non-magnetised planetary bow shocks, atmospheric bow shocks, and finally cometary bow shocks. This division of the celestial objects in the solar system grouped the earth together with the tiny innermost planet, Mercury, and the giant outer planets, while the terrestrial planets Venus and Mars had to be attributed to the remaining groups.

Interestingly enough, all planets and even the larger comets are surrounded by detached bow shocks. These bow shocks, for the magnetised planets, become ever stronger with increasing heliocentric distance from the Sun, not only because of the increase in the internal magnetic fields from the inner planets Mercury and Earth to the giant outer planets but also because of the increasingly cooler and more dilute solar wind and the weaker magnetic field with distance from the Sun. As a consequence the Mach number increases with distance, while the interplanetary magnetic field becomes more tangential to the planetary magnetic field. Thus the shocks become increasingly high Mach number quasi-perpendicular. In spite of this systematic, surprising differences have been found between the canonical paradigm of a planetary bow shock, Earth's bow shock, and the bow shocks of the outer giant planets. At Jupiter some of these differences could be attributed to the presence of the Jovian heavy plasmasphere and fast rotation leading to large centrifugal forces acting on the Jovian plasma and possibly affecting the shape and properties of Jupiter's bow shock. However, a general and interesting observation is that with increasing distance from the Sun the bow shock obviously becomes narrower and steeper having a large overshoot and causing less plasma turbulence in the adjacent magnetosheath than expected. Also, the shock shape is more inflated the farther out from the Sun we are. Finally, particle acceleration at the outer bow shocks seems to be more efficient than at Earth. This might be a result of the much broader magnetosheath at the giant planets which allows the turbulence to develop until quasi-stationarity.

Non-magnetised planets also possess bow shocks depending on whether they have atmospheres or not. The Moon's bow shock does not exist, however, as a global object since no induced or global lunar magnetic field has been found. However, locally, remanent magnetisation in the lunar crust sometimes leads to the formation of highly dynamic mini-bow shocks and magnetospheres. This is a very interesting finding that can be extrapolated to the larger asteroids like Ceres but does, surprisingly, not apply to Mars. The reason is that Mars possesses a dilute atmosphere that is still dense enough to become ionised. Its ionosphere acts as a shield because the solar wind induces a magnetic pile-up of the magnetic field in front of it which then acts in a similar way as a planetary magnetic field even though being of purely interplanetary origin. Its existence is the reason for the formation of a detached Martian bow shock.

Venus with its very dense atmosphere, high surface temperature and proximity to the Sun behaves even differently. Venus is non-magnetic with even no remanent magnetism possible in its crust as its surface temperature is above the Curie point. However, its atmosphere becomes ionised, creates a dense ionosphere and leads to a pile-up of an induced magnetic barrier in a way similar to Mars. In front of this barrier and detached from it a bow shock stands close to Venus with similar properties as Earth's bow shock but the difference of being much smaller in size and being affected by atmospheric pick-up ions from Venus.

Finally, comets can possess bow shocks as the example of comet Halley has shown. However, for this they need to enter the inner heliosphere for being sufficiently illuminated by the Sun, out-gas matter from their interior and generate a gas cloud that surrounds them. Such a cloud behaves similar to though not exactly like an atmosphere of a non-magnetic planet. It becomes ionised by illumination and charge exchange with the solar wind, creates an induced magnetic pile-up and thus a barrier which forces a bow shock to evolve if the comet is large enough. However, the mass loading of the solar wind substantially decreases the cometary Mach number such that the shock becomes weak and may even become subcritical not reflecting any particles. If this mass loading takes place over large distances from the comet, the solar wind might become gradually slowed down when approaching the comet and no shock will evolve at all even though a magnetic barrier forms around the comet. The solar wind then becomes sub-magnetosonic and has time enough to adjust for the presence of the comet and to flow around. This situation might have been observed at Giacobini-Zinner. In general, however, even large comets of the kind of Halley will develop a bow shock only inside a certain heliocentric distance from the Sun. Farther out, even though the Mach numbers are large, the absence of a cometary atmosphere prevents its formation. The only possibility for a shock then is that the fast solar wind stream induces an inductive field in the cometary body and conducting cover layer. This however requires sufficiently large cometary nuclei of size substantially exceeding the local proton gyro-radius in the solar wind stream.

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The Heliospheric Termination Shock

Abstract. The heliospheric Termination Shock is the largest (by dimension) shock in the heliosphere. It is believed that it is also the strongest shock and is responsible for the generation of the Anomalous Cosmic Ray component in the heliosphere. This chapter reviews the gross properties and observations of the Termination Shock. It is structured as follows: 1. The heliosphere, providing the heliospheric stage for Termination Shock formation, 2. Arguments for a heliospheric Termination Shock, 3. The global heliospheric system, 4. Termination Shock properties, 5. Observations: the Voyager passages, radio observations, plasma waves and electron beams, traces of plasma and magnetic field, energetic particles, galactic cosmic rays, Termination Shock particles, the Anomalous Cosmic Ray component, 6. Conclusions.

11.1 The Outer Heliosphere

On its way out from the Sun into the near interstellar environment of our solar system the super-magnetosonic solar wind is not remarkably retarded by its interaction with the celestial bodies in the solar system. These bodies are not very frequent and, even though large by human measures, retarding the solar wind locally, diverting it locally from its straight radial path and forcing it to flow around their magnetospheres or atmospheres, are too small to affect the solar wind on the global scale. The diverted wind, consisting of charged particles that are slower than the solar wind while being heated by the interaction with the bow shocks, behave like pick-up ions which in their proper system feel the solar wind electric field, become accelerated and couple to the solar wind. In this way the celestial bodies in the solar system create hot plasma striations in the solar wind on the scales of the magnetospheric cross sections which interact with the solar wind by plasma instabilities and gradually heat it when moving down the way to the boundaries of the heliosphere. By this they contribute to the downstream solar wind turbulence and at the same time a little retard the adiabatic cooling of the solar wind which should be the consequence of its three-dimensional collisionless outward expansion.

It is only in the outer heliosphere that the solar wind becomes retarded mainly by its interaction with the interstellar gas (conventionally called the Local Interstellar Medium LISM) which it encounters with its full bulk kinetic energy density $\mathcal{E}_{\text{SW}}(r) = \frac{1}{2}m_i N(r)V_1^2$. Since the solar wind density decreases approximately proportional to the surface of a sphere, the bulk flow energy density is roughly given by $\mathcal{E}_{\text{SW}}(r) = \mathcal{E}_{\text{SW},1}R^{-2}$, where $R = r/(1 \text{ AU})$ is the heliocentric radial distance measured in AU, and $\mathcal{E}_{\text{SW},1} \approx 10^7 \text{ keV/m}^3 \approx 10^{-9} \text{ J/m}^3$ is the average bulk solar wind energy density at the Earth's orbit. The interstellar gas in the Local Interstellar Cloud (LIC) wherein our heliospheric system is embedded has an average density of $N_{\text{LIC}} \approx 3 \times 10^5 \text{ m}^{-3}$ and a relative velocity $V_{\text{LIC}} \approx 26 \text{ km/s}$ with

Table 11.1: Properties of the local interstellar medium in the local interstellar cloud*

Parameter	Observation/estimate	Method
Sun/LIC relative velocity (km/s)	26.3 ± 0.4 25.7	He atoms ¹ , absorption lines, Doppler effect ²
Temperature (K)	6300 ± 340 6700	He atoms ¹ , absorption line broadening ²
N_H (atomic hydrogen, m^{-3})	$(2 \pm 0.5) \times 10^5$	pick up ions ³
N_i (protons, m^{-3})	$(0.3-1) \times 10^{-5}$	pick up ions ³
B_{IS} (magnetic field, nT)	0.2–0.4	direction ⁴ : $30^\circ < \alpha < 60^\circ$
CR-pressure P_{CR} (MeV/m^3)	~ 0.2	

* After *Izmodenov & Baranov [2006]*, $\alpha = \angle(\mathbf{B}_{\text{IS}}, \mathbf{V}_{\text{IS}})$.

¹ Witte [2004], Möbius et al [2004], ² Lallement [1996], ³ Gloeckler [1996], Gloeckler et al [1997], Geiss et al [2006], ⁴ Lallement et al [1995, 2005].

respect to the solar system. Its kinetic energy density is thus $\mathcal{E}_{\text{LIC}} \approx 2 \times 10^{-13} \text{ J/m}^3$. Equating $\mathcal{E}_{\text{SW}}(R) = \mathcal{E}_{\text{LIC}}$ one immediately finds that the upwind boundary of the region which is dominated by the solar wind, the heliopause, is to be expected around a heliocentric distance of $R_{\text{HP}} \approx 130 \text{ AU}$. Downwind the heliopause will be at a distance somewhat farther away with the heliosphere probably forming an extended tail.

There are many uncertainties in this estimate but for the solar wind velocity. One of them is that the LISM does not consist just of neutral atoms; it is partially ionised by the UV radiation of nearby stars. In addition it contains a small amount of energetic Cosmic Rays [a recent collection of the composition of the LISM can be found in *Izmodenov & Kallenbach, 2006*, see also Table 11.1]. Moreover, the LISM is weakly magnetised as well. These components and the presence of the weak interstellar magnetic field affect the interaction between the solar wind and the LISM. In particular the ionic component turns out to be important in this interaction.

Nevertheless, from the above estimate one obtains an impression on the order of magnitude of the global size of the region that is dominated by the solar wind. One of the uncertainties is the thermal pressure of the solar wind. With increasing heliocentric distance the number density of neutral gas atoms which leak in across the heliopause increases. These neutral gas particles, predominantly atomic hydrogen, undergo multiple charge exchanges with the fast solar wind protons whereby they transform into cold pickup ions which couple to the solar wind and are heated up to four times the solar wind kinetic energy, a mechanism that has been described by *Winske et al [1985]* as the result of the interaction of the freshly injected and accelerated pickup ion ring distribution with electromagnetic instabilities. They heat the solar wind and at the same time retard it. By this they produce fast streaming neutral hydrogen atoms flowing outward at solar wind speed. This process can repeat itself generating a relatively dense component of such pickup ions in the outer heliosphere which carry a substantial fraction of energy. When interacting with the Termination Shock some of these pickup ions are believed to be accelerated to comparably high

Table 11.2: Densities and pressures of solar wind components at 80 AU

Component	$N \text{ (m}^{-3}\text{)}^*$	$P \text{ (MeV/m}^3\text{)}^*$	$\beta = 2\mu_0 P/B^2$
SW protons	$(7-15) \times 10^2$	$10^{-3} - 10^{-4}$ (thermal)	7×10^{-3}
		$\sim 0.5 - 1.0$ (dynamic)	50–100
Pickup ions	$\sim 2 \times 10^2$	~ 0.15	15–20
ACR		0.01–0.1	1–10

* After *Izmodenov & Baranov [2006]*.

energies [*Fisk et al, 1974; Pesses et al, 1981*] in the >10 MeV range forming the so-called Anomalous Cosmic Ray (ACR) component of the solar wind. In the outer heliosphere the retardation and heating effect on the solar wind and the pressure that is attributed by the pickup ions and the Anomalous Cosmic Rays is believed to be quite substantial. It modifies the properties of the outer heliospheric solar wind.

Another uncertainty is found in the behaviour of the interplanetary field the field lines of which follow geometrically the Parker spiral. At 1 AU the angle of the Parker spiral is roughly 45° against the radial direction, rapidly increasing with radial heliocentric distance. In the outer heliosphere it is believed to be close to 90° . In the ecliptic plane the magnitude of the interplanetary magnetic field decreases $B \propto R^{-1}$. Thus the ratio of the kinetic to magnetic pressure $\beta_{kin} = 2\mu_0 \mathcal{E}_{SW}/B^2 = \text{const}$ is roughly constant. Since at 1 AU $\beta_{kin} > 1$, the solar wind remains high β_{kin} plasma also in the outer heliosphere (see Table 11.2) such that the magnetic field becomes unimportant for the processes at the heliopause. This holds, of course, only as long as the solar wind is not substantially retarded by the presence of a strong Anomalous Cosmic Ray component. Table 11.2 also gives the thermal, pickup ion and Anomalous Cosmic Ray pressures and kinetic to magnetic pressure ratios. In pressure pickup ions are comparable to the dynamic solar wind pressure thus completely dominating the ion-plasma β in the comoving plasma frame. It is only the pickup ion β that has to be taken into account in all the internal plasma processes like instabilities, wave excitation and wave-particle interactions, with a contribution of the Anomalous Cosmic Ray component.

In addition the interplanetary magnetic field exhibits a sector structure with oppositely directed magnetic fields. Since the spiral fields have roughly 90° angle when going along the radial direction outward, one alternately passes through adjacent regions of alternating field directions or, in other words, flows across a washing board formed by the undulating interplanetary current sheet.

The magnetic field at 100 AU, on the other hand, is a factor 100 less than at 1 AU, being about $B_{HP} \sim 5 \times 10^{-2}$ nT in magnitude, a value that is just below the value of the interstellar magnetic field (cf. Table 11.1). This situation is similar to that near Earth where the solar wind also carries a weaker magnetic field than the geomagnetic field in the magnetosphere, the obstacle with that the solar wind interacts at Earth. However, when taking into account a presumable compression factor $B_2/B_1 \sim (2-4)$ at the Termination Shock, the heliospheric magnetic field near the heliopause becomes the same order as the interstellar magnetic field B_{IS} . Since at the heliopause both fields are inclined, they can in principle

reconnect and thus can lead to a much closer coupling between the solar wind and the interstellar medium than is naively believed, causing asymmetries in the geometrical form of the heliopause, leakage of plasma, generation of plasma jets and other effects.

11.2 Arguments for a Heliospheric Termination Shock

Given these conditions and uncertainties one of the most interesting effects in the outer heliosphere is the formation of a shock wave that, on the global scale, surrounds the heliosphere, the heliospheric Terminal Shock (or Termination Shock). It is believed that it surrounds the entire heliosphere because the interstellar gas can theoretically be considered as being at rest compared to the fast solar wind flow.

Its formation is understandable if one considers the super-magnetosonic nature of the solar wind. The constancy of the dynamical $\beta_{kin} \equiv \mathcal{M}_A^2 = \text{const} \gg 1$ implies that the solar wind in the outer heliosphere remains to be a super-Alfvénic and supercritical plasma flow. The violent interaction of the supercritical solar wind flow with the interstellar gas necessarily leads to the formation of a shock inside of the heliopause where the solar wind stream is retarded to sub-magnetosonic speed and is heated.

The location of this Termination Shock can be estimated from the formula for the stand-off distance of a bow shock in front of the blunt obstacle. Clearly the heliopause is a blunt though concave obstacle, a circumstance that yields an inward moving (in the solar wind frame) Termination Shock of as well concave shape. The distance Δ_{TS} of this Termination Shock from the heliopause is given by the empirical formula that holds for bow shocks

$$\frac{\Delta_{TS}}{R_{HP}} \approx 1.1 \frac{N_1}{N_2} \quad (11.1)$$

where N_1 is the total solar wind density inside the heliosphere, and N_2 is the total plasma density outside the Termination Shock in the turbulent heliosheath. The Termination Shock is the bow shock of the local interstellar gas obstacle in the solar wind.

Again, for a strong shock with $N_2/N_1 = 4$ the distance of the Termination Shock from the heliopause becomes approximately $\Delta_{TS} \approx 36$ AU. Expressed as the distance from the Sun this yields a heliocentric distance of $R_{TS} \approx 94$ AU. On the other hand, if the shock becomes weaker – which is the case when N_1 increases due to the addition of a substantial fraction of pickup ions in the outer heliosphere – then the nominal location of the Termination Shock should shift inward into the heliosphere to a position substantially closer to the Sun. Detecting the location of the heliospheric Termination Shock is therefore an indication of the strength of the shock respectively of the state of it being mediated by pickup ions and the pressure of Anomalous Cosmic Rays. The estimated solar wind parameters for the four dominant solar wind components, thermal and dynamic solar wind protons, interstellar pickup ions, and Anomalous Cosmic Ray particles are given at a distance of 80 AU in Table 11.2. The most intriguing observation is that the pickup ion density and pressure are comparable to the solar wind density and dynamic pressure, while the ACR pressure may also contribute up to 20%.

The region outside of the Termination Shock, i.e. the region between the Termination Shock and the Heliopause, is the Heliosheath. By its nature as being the transition region between obstacle and shock it should be a highly turbulent medium of compressed density and magnetic field, of increased temperature and in the ideal case zero normal flow velocity as the simple model described here assumes that there is practically no solar wind flow across the Heliopause into the interstellar gas. The solar wind should be transformed into a flow that is tangential to the heliopause surrounding the solar system in the heliosheath boundary layer while slowly mixing into the surrounding local interstellar medium. This mixing is provided by the weak collisions and charge exchange processes between the solar wind and the interstellar atoms and, if the magnetic field becomes important, is also caused by reconnection at the heliopause.

11.3 The Global Heliospheric System

Since the heliosphere has a radius (upwind heliopause distance) of $R_{\text{HP}} \sim 100$ AU its size is huge. The collisional mean free path is only a few AU (the dominant collisional process being charge exchange with neutrals from the LISM with cross section $\sigma_{\text{ex}} \sim 10^{-18} \text{ m}^2$), still large compared to terrestrial dimensions but short compared with the dimensions of the heliosphere. It is thus believed that it can be described sufficiently well within a hydrodynamic or magneto-hydrodynamic approach. This is certainly true what concerns its gross dimension and structure but does not apply to the interesting local properties of the Termination Shock which – as we know from the previous chapters – has dimension of the ion inertial scale $\lambda_i = c/\omega_{pi}$ with substructure on the electron inertial scale λ_e where the electron physical processes are taking place. From Table 11.2 we find that $\lambda_i \sim 170 \text{ km} \approx 10^{-6} \text{ AU}$. This is much less than the collisional mean free path which is of the order $\sim (50-100) \text{ AU}$, identifying the Termination Shock as a truly collisionless supercritical and very thin shock transition. Nevertheless, its gross shape can be described by fluid theory as it also holds for the planetary bow shocks and there is no good reason to not extrapolate to the Termination Shock as well.

Figure 11.1 taken from Zank & Mueller [2003] shows the results of multi-fluid simulations of the outer heliosphere in the local interstellar medium of the local interstellar cloud (so-called kinetic simulations as they include the multi-component nature of the solar wind and LISM coupling them via collisional interaction including charge exchange). In the left part of the figure the magnetic field is neglected. Thus the system is symmetric, and only a half plane is given space on top for the temperature (T in K), below for the density ($N \rightarrow n$, in cm^{-3}). The colour code means that red is high, blue is low. Correspondingly, the high kinetic energy in the solar wind transforms into high temperature.

The upper part of the figure on the left very clearly shows the formation of the heliosphere in the interaction to which the solar wind is confined. The inner part of the heliosphere is bound to the greenish-yellow region lying inside the Terminal Shock (called here Termination Shock) which has an approximately ellipsoidal shape with the Sun in one of its foci. Outside the Termination Shock in the region of the broad heliosheath the medium is hot until reaching at the heliopause boundary where the flow (blue lines) is diverted

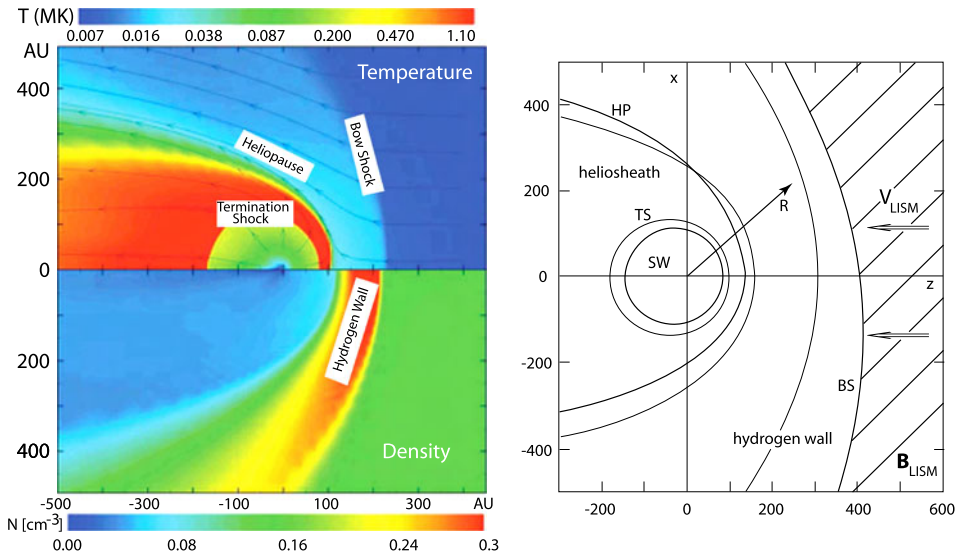


Figure 11.1: The heliosphere. *Left:* Global kinetic simulations [from Zank & Mueller, 2003, courtesy American Geophysical Union] of the interaction of the heliosphere and the local interstellar medium. The *upper half* of the figure shows the kinetic temperature distribution in the heliosphere and local interstellar medium, the *lower part* the density distribution. The ion component of the local interstellar medium and leaking solar wind ions accumulate in the dense hydrogen ion wall. The boundary is the heliopause. A bow shock in the interstellar medium develops when the medium is braked by the heliosphere obstacle. Inside the heliosphere the Termination Shock evolves when the solar wind is stopped. This simulation is without interstellar magnetic field and is thus symmetric. *Right:* The effect of the interstellar magnetic field on the heliosphere is shown as the *dark lines* compared to the *light lines* without magnetic field [after Izmodenov & Baranov, 2006]. No solar wind magnetic field is assumed. The magnetic field in the LISM deforms the heliosphere system and compresses the heliopause and Termination Shock. The values used in this simulation are: $N_{i,LISM} = 6 \times 10^4 \text{ m}^{-3}$, $N_{H,LISM} = 1.8 \times 10^4 \text{ m}^{-3}$, $B_{IS} = 0.25 \text{ nT}$.

towards the flanks and where it mixes with the LISM and cools. This mixing region widens at the flanks of the heliosphere.

On the other hand, outside the heliosphere the LISM is also heated over a broad range of distances of the order of ~ 100 AU width with the flow lines of the LISM being bent around the heliosphere. Note that inside the Termination Shock the yellow colour indicates that the solar wind in its outer parts is hot which is due to the presence of the pickup ions. It is also worth noting – though not of particular importance for the physics of the Termination Shock – that in these simulations a bow shock forms in the LISM far in front of the heliopause. A hydrogen wall develops as the downstream transition layer from the LISM to the heliosphere in this interaction. The reason for this bow shock (the existence of which is still hypothetical and the physics of which is different from magnetised planetary bow shocks being more similar to the Venusian or cometary bow shocks) is that the LISM in this model simulation is supersonic and the interaction with the heliosphere is thus shock-like.

The lower part of the figure shows the density which is low inside the heliosphere but shows a huge enhancement in the upwind pile-up region at the nose of the heliosphere

where the LISM ions are compressed and accumulate. This upwind pile-up is called the Hydrogen Wall. Leakage of the LISM into the heliosphere is indicated by the erosions on the outer heliosphere.

The right part of the figure, redrawn after *Izmodenov & Baranov [2006]*, shows the effect of the presence of an interstellar magnetic field on the shape of the heliosphere and Termination Shock. The light lines correspond to the simulations on the left without magnetic field, the heavy lines include the LISM magnetic field effect. The main effect in this simulation – which does not include a solar wind magnetic field and thus also no reconnection effects – is the deformation by the additional pressure of the LISM magnetic field. This deformation, on the scale shown, is very small at the Termination Shock but strong at the heliopause and heliospheric bow shock wave.

11.4 Termination Shock Properties: Predictions

Calculations like those in the above simulations can only provide the gross properties of the Termination Shock, its presumable location in the heliosphere and its presumable shape. This information is even questionable as long as no solar wind magnetic field has been included. The latter is subsequently imposed, behaving passively, non-dynamically, and having the Parkerian (Archimedean) spiral form which implies that it is about tangential to the nominal shape of the Termination Shock such as it follows from the simulations. Hence the Termination Shock is, in terms of these simulations, a supercritical, almost perpendicular, collisionless shock. It might, however, be modified to some extent by the presence of the energetic and comparably dense pickup ion component. Thus the Termination Shock should be a huge shock surface of a quite strong collisionless shock, by far the most extended shock wave in the entire heliosphere.

All these claims are rather uncertain as long as they cannot either be supported (or falsified) by direct observations (as will be discussed in the next section) or made plausible with the help of suitable numerical simulations on scales which are appropriate for the study of the Termination Shock microphysics. In particular, the claim that the Termination Shock is strictly perpendicular might be mistaken as it is based on information that is extracted from global scales of the order of ~ 0.3 AU only, the presumable average distance between the solar wind corotating magnetic sector structures. Even at the Earth's orbit at 1 AU the direction of the solar wind is quite unstable and is only in the long term and on the large scale average aligned along the Parker spiral. This variability is usually attributed to the proximity of the Sun and the variability of the inner heliospheric solar wind. There is, however, little reason to be certain about the quietness of the outer heliospheric solar wind and the greater stability of its local magnetic field. In fact, Pioneer 11 observations at 35 AU *Smith [1993]* suggest a variation in the magnetic field direction around the Parker spiral of $\pm 25^\circ$.

The presence of the dense pickup ion shell distributions in the outer heliosphere should cause a high level of Alfvénic turbulence of long wavelengths electromagnetic ion-cyclotron waves which in the collisionless cold outer solar wind plasma with its long scales are very weakly damped only and are not hindered to steepen until reaching large amplitudes.

Their transverse character implies that the assumptions on the outer heliospheric solar wind magnetic field of being perpendicular to the direction of the solar wind flow ceases to hold. At the contrary, in the solar wind frame of reference the wave magnetic field is transverse to the magnetic field and might thus have a strong component in the direction perpendicular to the spiral. In this case, the Termination Shock will – on the local scale – behave much more like a quasi-parallel shock than a quasi-perpendicular or even strictly perpendicular shock. Even when the magnetic field is of spiral form is there no reason to believe that the Termination Shock is tangential to the field. The field lines might cross the shock at its one flank and re-enter it on the other flank if the shapes of the magnetic field spiral and Termination Shock do not align which will be the case when the magnetic field behaves completely passively and the interaction between the solar wind and the LISM is due to the flow as is assumed in the above global simulations.

It has been concluded by *Liewer et al* [1993] from hybrid simulations including 10% pickup protons at quasi-perpendicular shocks that below a shock normal of $\Theta_{Bn} = 60^\circ$ pickup protons are very efficiently back-scattered from the shock and generate upstream waves which cause pitch angle diffusion on the particles. The scattering ceases at larger angles, however. Although the Termination Shock is expected to be dominantly perpendicular, *Kucharek & Scholer* [1995] investigated pickup ion reflection for the case that the Termination Shock would temporarily have quasi-parallel character. In contrast to the *Liewer et al* [1993] simulations they included heavier ions suspecting that they are back-scattered even stronger than protons. Their simulation settings are one-dimensional hybrid including mass ratio weighted H^+ and He^+ ions and 5% He^{2+} , Mach number $\mathcal{M}_A \approx 8$, investigating the two cases of $\Theta_{Bn} = 50^\circ$ and 70° thereby confirming the above *Liewer et al* [1993] result of backscattering only for angles $\lesssim 60^\circ$. The initial pickup distributions are accelerated spherical shells. The heavier ions have been treated as test particles in a non-self-consistent way. The result of this simulation is that with increasing mass the reflection of pickups decreases. This can be simply explained by the larger energy of the particles which can overcome the shock ramp potential. Applied to the Termination Shock this implies that for oblique shock normal angles the Termination Shock can well act as a good reflector of pickup ions thus providing a seed population for their further acceleration when they return to the shock.

This has been further investigated by *Scholer & Kucharek* [1999] who right away assumed a quasi-parallel shock in their one-dimensional hybrid simulations of $\Theta_{Bn} = 5^\circ$, $\mathcal{M}_A = 5.3$ and $\mathcal{M}_A = 10.4$ shocks in view of the generation of Anomalous Cosmic Rays (see Table 11.3). These authors adopt the picture that pickup ions are in a first step predominantly accelerated by travelling interplanetary shocks in the region farther inside the heliosphere, and when transported out to the Termination Shock become accelerated there to the high Anomalous Cosmic Ray energies by efficient reflection at the stationary Termination Shock and the action of assumed non-linear trapping and non-adiabatic acceleration. Again the heavier ions than pickup protons are treated as test particles. The main results of these simulations are shown in Figure 11.2. The particle distributions (differential energy fluxes) are shown on the left. The particle reflection coefficients are shown on the right as function of the Mach number. The authors find reflection ratios $R = N_{ref}/N_{in}$ for pickup protons of 42%, for He^+ of 25% compared to the respective shock generated

Table 11.3: Ion reflection ratios $R = N_{ref}/N_{in}$ at quasi-parallel shocks*

$\mathcal{M}_A$	Stationary shock				Travelling shock	
Ion	5.3	10.4	10.4	10.0		
	Pickup (%)	Diffuse (%)	Pickup (%)	Diffuse (%)	Pickup (%)	Diffuse (%)
H ⁺	42	0.9	~29		10.3	
He ⁺	25	0.6	~31		6.2	<0.6

* ($\Theta_{Bn} = 5^\circ$, $\beta_i = \beta_e = 0.5$, $N_{H^+}/N = 10^{-3}$, 5% He²⁺).

Data taken from Scholer & Kucharek [1999].

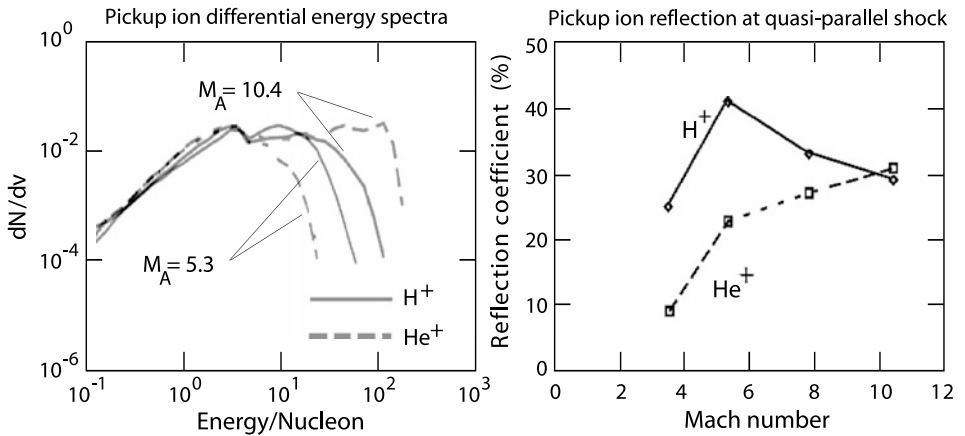


Figure 11.2: Hybrid simulations of pickup ion reflection at quasi-parallel shocks ($\Theta_{Bn} = 5^\circ$) in view of acceleration of Anomalous Cosmic Rays at the Termination Shock [data from Scholer & Kucharek, 1999, courtesy American Geophysical Union]. *Left:* The differential energy spectra of the pickup ions for two different Mach numbers of $\mathcal{M}_A = 5.3$ and 10.4 (in units of the upstream kinetic energy). The low energy spectra are identical but the higher Mach numbers yield higher heavy pickup ion energies. At the low Mach number proton pickups reach higher energy. At the high Mach number the He^+ ions are accelerated stronger than protons. *Right:* The reflection coefficients as function of Mach number. Protons reach maximum reflexivity around medium Mach number while the reflection coefficient for He^+ pickups still increases and overturns the proton coefficient at about Mach number 10. Pickup reflection coefficients are in the range of tens of percent, two orders higher than the shock reflexivity of diffuse ions.

0.9% and 0.6% diffuse proton and He^+ ions in the foreshock. High Mach numbers imply higher final energies of the heavier ions (see the left part of the figure). This can clearly be explained only when the non-adiabatic behaviour of the pickup ions at the shock dominates over the shock potential effect at the higher Mach numbers. This can be understood because the amplitude of the shock produced waves increases with Mach number, as had been shown earlier [Scholer et al, 1997].

It follows that shocks can indeed inject substantial amounts of pickup ions into the Fermi cycle. For the pickups this fraction is much higher than the fraction of shock produced upstream particles because of the very low velocities of the pickup ions in the shock normal direction, since they have merely been accelerated into the direction of the tangen-

tial upstream motional electric field perpendicular to the shock normal. Thus their energy is low enough for being reflected from the shock potential and being trapped in the field close to the shock for a while where they become accelerated non-adiabatically to high energies, a process that is known from shock surfing. The distribution of these reflected and accelerated pickup ions has the shape of a pancake with more energy in the perpendicular direction. Such a component excites two types of waves: anisotropy-driven whistlers and ion-ion instabilities when back entering the upstream solar wind.

Applied to the Termination Shock which is located in the midst of the highly increased fraction of interstellar pickup ions this implies that the Termination Shock could be an excellent place to reflect and further accelerate those pickup ions – protons and heavier less abundant elements – to become Anomalous Cosmic Rays. The condition for this to happen is that the Termination Shock is not an ideal quasi-perpendicular shock but has regions where locally the shock normal angle to the interplanetary magnetic field is substantially smaller than $\Theta_{Bn} \lesssim 60^\circ$. Such regions are expected to exist possibly at its flanks where in the idealised picture the spiral interplanetary heliospheric magnetic field should cut through the Termination Shock at some angle. However, they can also be generated by surface instabilities of quasi-perpendicular shocks which are driven by the gyrating reflected ions in their accelerated motion in the upstream motional electric field along the shock surface [Lembège & Savoini, 1992; Lowe & Burgess, 2003; Burgess & Scholer, 2007]. Such an instability resembles a flow instability of the Kelvin-Helmholtz family but is entirely due to the kinetic effect of the reflected ions. It requires a supercritical shock and sufficiently low β_i . When it sets on, the surface of the quasi-perpendicular shock ceases to be smooth but develops “ripples” which locally affect the shock normal direction.

More recent full particle PIC simulations performed for quasi-perpendicular shocks [Scholer *et al.*, 2003] with real mass ratios $m_i/m_e = 1836$ have investigated the β_i dependence of shock reformation, ion reflection, and the cross-shock potential and have found that the shock potential divides between the foot and the ramp and is restricted to scales of $\sim 4\lambda_e$. At the low thermal ion- β_i that is expected at the quasi-perpendicular Termination Shock coherent ion trapping and acceleration of pickup ions, as is required in shock surfing mechanisms, is thus improbable. If this turns out to be valid then the acceleration of pickup ions at the Termination Shock will mainly be due to acceleration in those regions where the character of the Termination Shock is quasi-parallel and should thus lead to anisotropies in the spatial distribution of the Anomalous Cosmic Rays and pancake forms in the energetic ion velocity distributions. There are still a number of caveats in this conclusion which is based on the one hand on hybrid simulations, on the other hand on just one-dimensional PIC simulations still leaving much space for interpretations, speculations and also surprises.

On the other hand the same simulations have also shown that quasi-perpendicular shocks like the expected Termination Shock cannot be stationary but undergo reformation at low $\beta_i < 0.4$ at the simulation Mach number $\mathcal{M} = 4.5$ unless the pickup ions in the outer heliosphere contribute to such an increase in β_i that reformation ceases to occur. This effect might be partially compensated by a low plasma-to-cyclotron frequency ratio $\omega_{pe}/\omega_{ce} \sim 1$. From Table 11.2 we have that in the foot region of the Termination Shock $\omega_{pe}/\omega_{ce} \sim 1.6$, even if all plasma components, including pickup ions, are taken

into account (remembering that the ratio N/B^2 is independent of radius in the expanding solar wind). This favours reformation even though the high pickup β in Table 11.2 tends to suppress it when the pickup ions are involved in the formation of the shock. The large mass ratio simulations also show that real reformation of quasi-perpendicular shocks does not so much depend on the number of gyrating ions than on the electromagnetic modified two-stream instability in the foot of the shock. In this case the β_i dependence is not sufficiently well known and the shock might reform anyway.

It is thus probable that the Termination Shock is rather a non-stationary shock than a stationary quasi-perpendicular shock. Locally it will be oscillating back and forth around its nominal position, a process that is driven by the variations in the interplanetary solar wind medium and its magnetic field direction and sector structure. It necessarily leads to surface waves running along the shock. The Termination Shock can thus be expected to exhibit quasi-periodic or also irregular variations in its shock normal angle Θ_{Bn} , the direction of its normal with respect to the magnetic field. Because of the high pickup ion density near the Termination Shock one might even speculate that the capability of the Termination Shock of reflecting a fraction of hot pickup ions might cause shock reformation on the scale of a pickup ion gyro radius that is produced by accelerated gyrating reflected pickup ions and involves an instability between the pickup ions and the upstream electrons of the main flow. The related questions have not been investigated yet. In the following we will discuss the more recent observations in the vicinity of the Termination Shock as they have been provided by the Voyager spacecraft.

11.5 Observations: The Voyager Passages

There is no way to observe the Termination Shock from Earth. With one exception, even from near Earth space the Termination Shock cannot be observed in any radiation. The exception is the Anomalous Cosmic Ray component. It is called anomalous because it has little to do with the real Cosmic Rays from the Galaxy. Though it has still not been approved, it is believed that these energetic particles come from the Termination Shock where they have been accelerated, forming the energetic tail of the outer heliospheric pickup ions. They are thus generated inside the heliosphere having their source in the outer heliosphere. The process by which they are produced maps the composition of the LISM, not only containing hydrogen ions but also other elements. The highest abundance has been found for helium, oxygen, and neon. Anomalous Cosmic Rays therefore tell more about the LISM than the Termination Shock. Not much information can be extracted about the Termination Shock from observing them from remote because they are distributed about homogeneously in angle throughout the heliosphere. Below in the section on energetic particles we discuss the *in situ* evidence of their generation by the Termination Shock that is based on the Voyager passage.

Both Voyager 1 and Voyager 2 spacecraft on their ways out into interstellar space (see Figure 11.3) ultimately crossed the Termination Shock. Voyager 1 passed it in December 2004 at radial heliocentric distance ~ 94 AU, the distance predicted by simple scaling arguments. Voyager 2 crossed it end of August 2007 at ~ 84 AU and at higher helio-latitude than

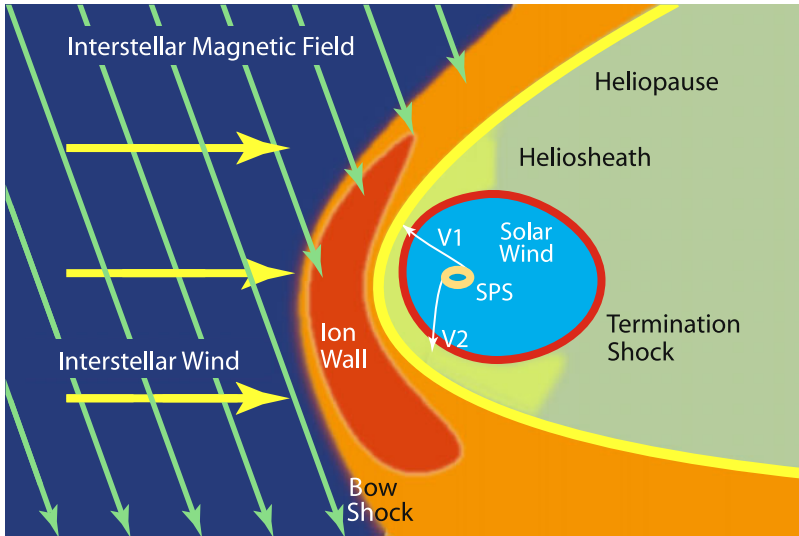


Figure 11.3: Sketch of the upwind heliosphere in the magnetised interstellar wind blowing against the solar wind, creating a Bow Shock and Ion Wall in the LISM, and the Heliopause. SPS is the solar planetary system showing the two orbits of the Voyager 1 (V1) and Voyager 2 (V2) spacecraft. V1 is shown to cross the Termination Shock (TS) at larger northern distance than V2 in the south. The Termination Shock is the *red ring* that confines the solar wind. The (exaggerated) asymmetry of the TS and location of SPS in the solar wind is caused by the interstellar magnetic field. Outside of the Termination Shock and the heliopause is the Heliosheath. Colour symbolises density, *dark blue* the low LISM density up to *blue* the low solar wind density. Density gradients are not included except for the compression of the upwind heliosheath which is drawn in *lighter green*. The pickup ion component in the outer upwind solar wind is not made visible.

Voyager 1, much closer in than expected from scaling but not completely unexpected. The asymmetry between North and South is attributed to the action of the interstellar magnetic field and to the variations in the solar wind with solar cycle. An asymmetry of the heliosphere had been predicted already earlier from Lyman α observations [Lallement, 1993] and from three-dimensional MHD simulations including an interstellar magnetic field [Ratkiewicz *et al*, 1998]. It was recently confirmed again by direct upwind/downwind Lyman α asymmetry measurements and was indeed attributed to the interstellar magnetic field direction [Lallement *et al*, 2005].

Before coming to the Voyager measurements we briefly elaborate on radio observations from remote and in situ Langmuir wave measurements.

11.5.1 Radio Observations

Radiation emitted from the Termination Shock was expected in plasma waves and in radio waves. Intense radio emission from the outer heliosphere has indeed been observed by the Voyager spacecraft [Kurth *et al*, 1984]. But it was soon recognised [Gurnett *et al*, 1993] that this radiation could not have been generated at the Termination Shock

unless the latter would have had very unusual properties that are not expected of any shock. Depending on the heliocentric distance of the Termination Shock between 80 AU and 100 AU the local plasma density is between $5 \times 10^2 \text{ m}^{-3} < N < 8 \times 10^2 \text{ m}^{-3}$. From Table 11.2 the total plasma density is of the order of $N(80 \text{ AU}) \lesssim 10^3 \text{ m}^{-3}$. The local plasma frequency at the Termination Shock cannot be higher than $\omega_{pe} = 2\pi \times 280 \text{ Hz}$, roughly an order of magnitude below the observed radiation frequency which was in the range $1.4 \text{ kHz} < f < 3.6 \text{ kHz}$. In order to be emitted in the second harmonic at $f \sim 2\omega_{pe}/2\pi$ from the Termination Shock the density at the Termination Shock should be 36 times higher than expected. The strongest shocks permit only a factor 4 density compression. Hence, the observed radio emission cannot come from the Termination Shock. *Gurnett et al [1993]* therefore rightly attributed the radiation to come from the outer heliospheric boundary instead. The mechanism of its generation is still unclear, but *Gurnett et al [1993]* and *Gurnett et al [2003]* could convincingly show that its emission was related to the interaction of occasional fast solar wind streams and travelling interplanetary shocks with the heliopause [see also the extended discussion in *Treumann & Pottellette, 2006*].

That the observed radiation has not been produced at the Termination Shock is less a surprise than the now established fact that no radiation at all has been observed that is coming from the Termination Shock. The radio emission detected by the Voyager spacecraft from the outer heliosphere on its way out were at frequency less than the local plasma frequency in the inner heliosphere and thus being completely screened. Radiation at 300–600 Hz could not have been observed until Voyager would have passed a distance of at least 35 AU. However when Voyager was beyond this distance no emission was detected.

The Termination Shock remained silent in the radio band even until the Voyager spacecraft had passed it after 2004. Any strong shock that is supercritical and is capable of reflecting electrons should emit radio waves, however. This holds for the Earth's bow shock, other planetary bow shocks, for solar and interplanetary Type II bursts which emanate from the corresponding coronal and interplanetary shocks; it was true even for the AMPTE IRM artificial plasma cloud injections into the solar wind [*Gurnett et al, 1985*]. It is thus very strange that the Termination Shock which has been believed to be the strongest shock in the heliosphere is actually a radio-silent shock. Thus finding the reasons for this quiescence should illuminate the plasma conditions that prevail at the Termination Shock. It should also contribute to the understanding of the conditions under which a shock is capable of producing radio waves.

Current radiation theories require the presence of electron beams that propagate along the ambient magnetic field and excite Langmuir waves at the local plasma frequency. Subsequent wave-wave interactions between these Langmuir waves are believed to cause emission of radio waves at harmonics of ω_{pe} . These mechanisms require either counter-streaming Langmuir waves in order to be able to satisfy the wave momentum conservation condition, or they require scattering of Langmuir waves off either thermal ions or ion acoustic waves. Some less explored theories prefer instead electron acoustic waves. But the basic condition is the presence of electron beams. These are naturally expected to be injected upstream from a sufficiently strong supercritical quasi-perpendicular shock.

Below we will argue that the Voyager observations indeed have identified both the presence of electron beams as also the Langmuir waves that are excited by these beams which makes the absence of radiation even more mysterious. This absence implies that the simple merging process of either two counter streaming Langmuir or one Langmuir and one ion acoustic wave is set out of action. The first process requires generation of back-scattered Langmuir waves. If the required ion acoustic fluctuations are lacking no back-scattered Langmuir waves will be generated. This happens when the electron temperature is too low for ion acoustic waves to be excited by the solar wind heat flux instability, i.e. when ion Landau damping is strong. Scattering off thermal ions has weak efficiency and in the cold solar wind plasma might be suppressed. However, it remains to be unclear whether the hot and sufficiently dense pickup ions do not contribute. Possibly their densities are insufficient. Electron acoustic waves will also not be excited upstream because the electron temperatures are too low. Therefore the absence of intense radio emission from a Termination Foreshock becomes plausible.

11.5.2 Plasma Waves and Electron Beams

The real question is: Why is the Termination Shock radio silent? This question is even more stringent as the Langmuir waves required for generating the expected shock radiation have indeed been observed, and that long before the Voyager spacecraft have crossed the Termination Shock, already from 11 February 2004 on during six time intervals, the last of them on 15 December 2004 [Gurnett & Kurth, 2005] at radial heliocentric distance 94.1 AU, shortly before the Voyager 1 verified and spectacular crossing of the Termination Shock in December 2004 took place.

Figure 11.4 on its right shows Gurnett and Kurth's dynamical spectrum of the electrostatic plasma waves detected with Voyager 1 on 15 December 2004 between 1800 UT and 21 UT in the frequency range from 178–300 Hz, in excellent agreement with the estimated local plasma frequency upstream of the Termination Shock. Langmuir wave amplitudes of the order of $(2-3) \times 10^{-6}$ V/m have been measured, one order of magnitude above background noise. The top of the left part of the figure shows the Voyager 1 magnetic field trace with the Termination Shock crossing happening several hours after the observation of the electron plasma waves. On the day of 15 December 2004 an intense spike in the electron flux was also detected.

It is intriguing to see the sporadicity of the electron plasma waves. Since their intensities are sufficiently far above thermal noise they can have been excited only by electron beams that flew along the outer heliospheric magnetic field and were connected to the Termination Shock. Electron beams propagating sunward along the spiral interplanetary magnetic field having energies >26 keV have been reported by Decker *et al* [2005] to have been observed by Voyager 1 throughout the year 2004. Since this has been so for almost the full year Voyager 1 was occasionally magnetically connected to the Termination Shock from a distance of 91.0 AU over a range of ~ 3 AU. On the day of 15 December 2004 a particularly intense spike in the electron flux was also detected. (There are indications that the spacecraft was in the foreshock of the Termination Shock already from 85 AU on which implies occasional magnetic connection over 9 AU.) It is not known to which part of the

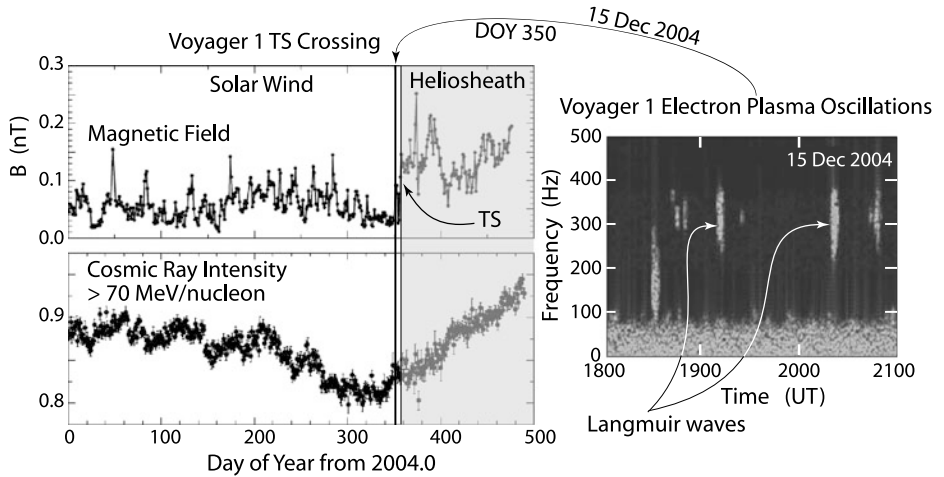


Figure 11.4: *Right:* Electron plasma wave dynamic spectrogram [data taken from [Gurnett & Kurth, 2005](#)] on 15 December 2004 shortly before the Voyager 1 crossing of the Termination Shock. The signals occur sporadically indicating magnetic connection of the spacecraft to the Termination Shock and the arrival of electron beams of a few keV energy. The local plasma density determined from the observation is of the order of $N \sim 10^3 \text{ m}^{-3}$. *Left:* Heliospheric magnetic field magnitude showing the crossing of the Termination Shock (*top*) and Cosmic Ray Intensity showing the unexpected continuous increase with entry into the heliosheath [data taken from [Burlaga et al, 2005](#)]. The vertical black line indicates the time observation of Langmuir wave spectra on the *right*.

Termination Shock the magnetic field lines were attached, initially probably more to the flanks of the Termination Shock as was suggested by [Gurnett & Kurth \[2005\]](#) referring to the spiral geometry of the interplanetary magnetic field.

The observation of the intermittence of the plasma waves also suggests that the Termination Shock cannot be strictly perpendicular, neither can it be stationary. Its shock normal must exhibit oscillations in its direction with periods of the order of hours. [Gurnett & Kurth \[2005\]](#) suggested surface waves to propagate along the Termination Shock and causing quasi-periodic undulations. In addition the Termination Shock must be strong enough to accelerate electron beams in the direction upstream of the Termination Shock. Since the electrons are known to be very cold at the heliocentric distance of the Termination Shock the generation of these beams could hardly have been by the Sonnerup-Wu mechanism, for the electron distribution would be completely inside the loss cone. Thus if there is no mechanism that generates a hot electron halo on the outer heliospheric electron distribution function, one expects that the mechanism of electron acceleration can only be due to intrinsic properties of the Termination Shock like those required by Hoshino's electron shock surfing mechanism. This implies strongly nonlinear processes to take place in the shock like the excitation of BGK modes and solitary waves.

In this respect the most recent passage of Voyager 2 through the Termination Shock in the southern heliosphere is of particular importance. Voyager 2 also observed electron plasma oscillations well before contacting the Termination Shock itself, i.e. 30 days ahead of the first crossing [[Gurnett & Kurth, 2008](#)] thus confirming the former Voyager

1 observations. From these observations an upstream (electron) plasma density of $N_e \approx 1.2 \times 10^3 \text{ m}^{-3}$ was inferred. As before the presence of the Langmuir waves is convincing evidence for the capability of the Termination Shock to eject low density electron beams of keV energy along the heliospheric magnetic field upstream. Those beams were undetectable for the Voyager 2 plasma instrumentation qualifying these Langmuir wave observations as of high diagnostic value.

Much more important than the observation of Langmuir waves is the detection of plasma waves during crossing of the shock ramp. Examples of the dynamic spectrum of plasma waves in such crossings have been given in Figures 6.17 and 6.20 for quasi-perpendicular shocks and in Figure 6.31 for a quasi-parallel shock. In every case the shock appears in the electric wave field as a broadband signal extending from the lowest frequencies up to frequencies beyond the local electron plasma frequency. The intensity of these signals is highest at the lowest frequencies. The difference between quasi-parallel and quasi-perpendicular shocks is mainly in the more turbulent behaviour of the quasi-parallel shocks where the broadband noise signals are distributed over a larger interval in space. Such signals are interpreted as the frequency-space signature of a number of spatially localised electrostatic structures which cross over the spacecraft leaving their trace in a broad Fourier spectrum that is independent on the cut-off frequencies which otherwise govern the dispersion curve traces of plasma waves. Because of the probably large number of solitary structures or BGK modes present in the case of a shock passage, these broadband traces are much more intense than when a single BGK mode is crossed.

Gurnett & Kurth [2008] reported the observation of just such broadband electric signals in coincidence with the crossing of the Termination Shock. In fact, the Termination Shock could have been identified from these traces alone as in the outer heliosphere there would be no reason for the generation of broadband signals of this kind before entering the heliopause region. Nevertheless, the occurrence of the intense broadband noise signals coincided with plasma and magnetic field observations of the shock crossing thereby confirming their interpretation as shock crossings. Of the five Termination Shock crossings identified in the magnetic and plasma data, two were also seen in the plasma wave observations, two fell into plasma wave data gaps, and one did not show a typical shock signature [*Gurnett & Kurth, 2008*].

The event preferred by *Gurnett & Kurth [2008]* is the third crossing TS-3 (their Event B). Figure 11.5 shows the broadband signal detected during this crossing which is typical for a supercritical shock transition as described above. The scale of the rms electric wave field amplitude is given on the right. It holds for all 7 wave bands of the instrument. Note that the 17.8 Hz channel suffers from enhanced instrumental noise, and the 311 Hz channel is polluted by a gyro interference. *Gurnett & Kurth [2008]* scaled the Termination Shock spectra to similar spectra observed at the Neptune bow shock crossing finding very close similarity between both objects.

The shock crossing occurs in the non-shaded time interval lasting for roughly three minutes. Prior to this the spacecraft was in the solar wind, afterwards it is in the heliosheath. Broadband waves with amplitudes of several 10^{-6} V/m were detected during the shock crossing. The highest amplitudes are in the 56.2 channel well below the plasma frequency

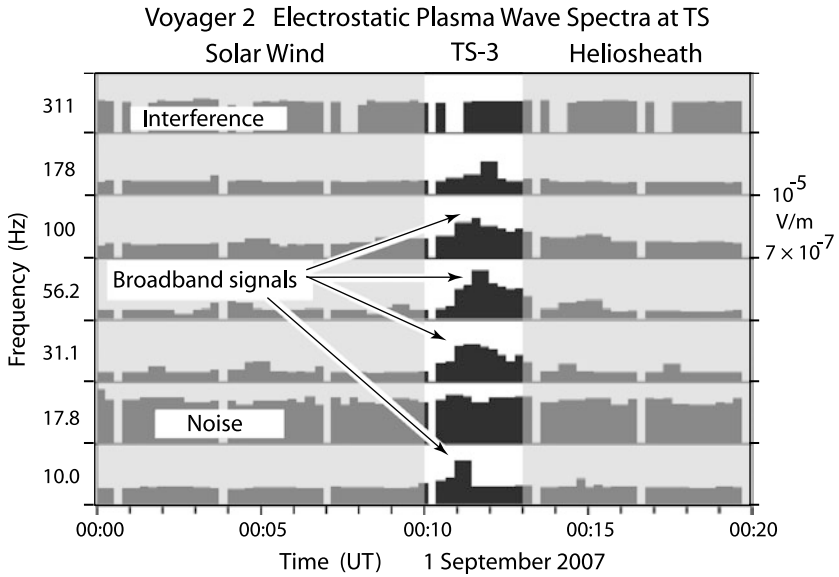


Figure 11.5: Broadband noise signals observed during the Voyager 2 crossing TS-3 of the Termination Shock on 1 September 2007 from the solar wind to the heliosheath [after Gurnett & Kurth, 2008]. This signal is typical for a shock ramp crossing resembling all bow shock crossings or interplanetary travelling shocks. It is an indication of the presence of small scale electrostatic structures in the shock transition region. The scale on the right holds for each of the 7 frequency bands. The height of the black region gives the rms electric wave field amplitude as function of time in units V/m.

and well above the electron cyclotron frequency (which was around $f_{ce} = 7.3$ Hz, embedded in the 10.0 Hz channel). From the magnetic traces and plasma observations which we discuss below it is found that the entire shock transition occurs on a scale of the order of $\Delta_{TS} \sim 1\lambda_i$, which is very thin but in agreement with the expectations on the width of a supercritical collisionless shock. On this short scale it is expected that violent nonlinear processes will take place mainly involving magnetised electrons and unmagnetised ions. It is thus clear that the observed broadband emissions are caused by nonlinear interactions in the shock ramp.

Since the electron temperature is very low and the electron velocity in the 100 eV to keV range is well above the electron thermal speed, the responsible instability is expected to be the Buneman two stream mode which leads to BGK structures. This resembles very closely the well known behaviour of planetary bow shocks. The nonlinear wave structures are most important for the internal dynamics of the shock as well as for the trapping and acceleration of particles. Thus their identification is of crucial importance for the understanding of the behaviour of the Termination Shock which as such serves as the paradigm of a non-relativistic high Mach number supercritical shock generated in the interaction between a magnetised stellar wind and the magnetised interstellar environment.

11.5.3 Traces of Plasma and Magnetic Field

In situ information on plasma and magnetic field properties from the Termination Shock has been obtained recently when the Voyager spacecraft crossed the Termination Shock from the inner heliosphere into the heliosheath. Voyager 1 which travels at a higher speed than Voyager 2 crossed it in December 2004 at heliocentric distance of 94 AU and heliographic latitude 34° N after having been in the Termination Shock foreshock from 85 AU on. The slower Voyager 2 crossed the Termination Shock between 30 August and 1 September 2007 at heliocentric radius 84 AU and helio-latitude 25° S after having been in the foreshock from 75 AU on, surprisingly finding the Termination Shock substantially farther in than expected from the Voyager 1 crossing. The distance between the spacecraft was 110 AU and 45° in longitude. On Voyager 1 the plasma instrument failed. Hence only magnetic field information was available to identify the shock crossing. This is shown in the Voyager 1 magnetic field trace on the left in Figure 11.4. During the non-shaded period the Voyager 1 spacecraft was in the solar wind, crossed the Termination Shock when the magnetic field suddenly increased by a factor of roughly 3 in mid December and remained high afterwards when Voyager 1 had entered the shock transition and the heliosheath. At the time of writing it is still in the heliosheath at high magnetic field levels heading for the heliopause.

Voyager 2 in the southern heliosphere crossed the Termination Shock with working instrumentation. Magnetic field and plasma measurements across the Termination Shock were published by *Burlaga et al* [2008] and *Richardson et al* [2008] respectively. Figure 11.6 shows three minutes of the Termination Shock crossing TS-3 [Burlaga et al, 2008] in 48 s time resolution around the Termination Shock ramp. The thin horizontal line in the upper panel is the value of the solar wind magnetic field prior to the crossing. Shown here are the shock foot, ramp and overshoot region. Afterwards the magnetic field dropped to the heliosheath value which is a factor of $B_2/B_1 \approx 1.7$ higher than the solar wind magnetic field identifying the shock as a magnetically relatively weak compression. (Note that the overshoot is roughly a factor 5 higher, which is due to currents flowing in the overshoot region!)

In the average the solar wind magnetic field was at a Parker spiral direction of $\sim 270^\circ$ azimuthal angle. Entering the foot and shock ramp this angle deviates up to $\sim 30^\circ$ from this direction. The inclination angle in the solar wind fluctuated around -60° , in the foot it is more around 30° and exhibits oscillations between 0° and $\sim 50^\circ$. Similar oscillations occur in the ramp. The angular trace shows that the shock already in the overshoot returns to the average spiral direction. It is the inclination which catches the change in shock direction becoming close to 0° in the overshoot and heliosheath. Altogether the changes in direction are relatively small. The shock crossing TS-3 is hence about quasi-perpendicular and is also supercritical as it develops a well pronounced foot. The Mach number was about $\mathcal{M} \approx 5$.

There are strong and irregular fluctuations in the magnitude of the foot magnetic field which must be related to the reflected gyrating particle distribution which carries the foot currents and excites electromagnetic waves. The fluctuations in the ramp are of very large amplitude. From the full particle simulations described in Chapter 5 we know that these

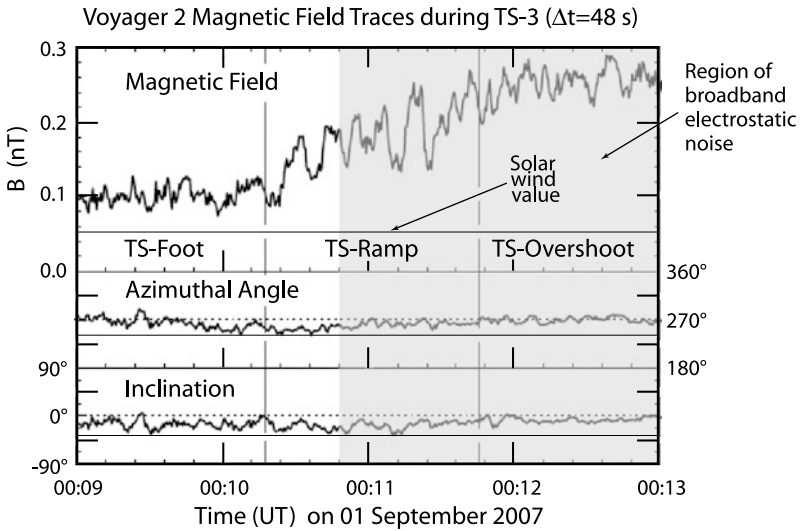


Figure 11.6: The Voyager 2 Termination Shock crossing TS-3 in the magnetic trace [after *Burlaga et al, 2008*]. Shown is the foot region where the magnetic field is about twice the undisturbed solar wind field, the ramp regions with its strong coherent oscillations. These are from simulations known to be caused by ion trapping with the maxima in the magnetic field coinciding with the vertexes of the ion phase space rings. The transition to the overshoot is seen as incoherent fluctuations in the magnetic field. The shaded region is the time interval when the plasma wave instrument detected the broadband electrostatic noise [*Gurnett & Kurth, 2008*] which becomes strongest after 00:11 UT in the transition from ramp to overshoot.

fluctuations are directly related to the trapping of ions in the ramp region with the magnetic maxima coinciding with the vertexes in the ion velocity distribution, as shown in Figure 5.31. They are thus indications of the ion dynamics, formation of vortices, ion heating and acceleration. In the overshoot these fluctuations become shorter wavelength which is also known from the simulation results and is related to the smearing out of the ion phase space vortices and the strong ion heating in the overshoot region.

When comparing these observations with the plasma wave observations by *Gurnett & Kurth [2008]*, the most interesting observation is made that the entire broadband noise spectrum observed in the plasma waves is restricted to the time interval 0010:45-0013 UT which is the transition from the ramp into the overshoot. This is the region where the overshoot currents flow. In fact, this was not unexpected and is in agreement with the AMPTE IRM observations of the broadband noise spectra. From the numerical simulations of a quasi-perpendicular shock [*Matsukiyo & Scholer, 2003, 2006*] it is also known that in this region the ion flow is stopped, and the ions are most strongly heated with the trapped ion vortices losing identity.

As noted above, the magnetic field observations during the Voyager 2 passage across the Termination Shock identified five Termination Shock crossings within roughly 12 hours starting with an inbound shock crossing and ending with an inbound crossing. This is possible only for the Termination Shock being highly dynamical and oscillating back and

forth. This oscillation can have different reasons, the most probable reasons are that the shock is unstable with respect to waves that travel along its extended surface causing ripples on the shock surface. Such waves can be assumed to travel at a velocity that is of the order of the local Alfvén speed, which is of the order of $V_A \sim 150$ km/s. The times between the crossing TS-1 to TS-5 were 3 hours, 4 hours, and 2 hours, respectively. Thus the average wavelength of those ripples or surface waves will be of the order of ~ 30000 km. Since the ion inertial length is $\lambda_i \sim 6000$ km, this is not more than one surface wave wavelength of $\lambda_{\text{sfw}} \sim 5\lambda_i$.

The plasma data [Richardson *et al*, 2008] of the Termination Shock crossings did, in principle, not provide much surprises. They confirm that the shock causes a steep decrease in the flow velocity and a comparable increase in plasma density, which is in agreement with the notion of the Termination Shock being a quasi-perpendicular shock. Except for the determination of the bulk plasma parameters the main conclusion is that the solar wind plasma is *not heated strongly* when crossing the Termination Shock. The bulk of the energy lost by the solar wind stream goes into the pickup ions. Only when taking into account this effect, Richardson *et al* [2008] were able to satisfy the necessary shock condition that the Termination Shock decelerates the solar wind flow to Mach number $\mathcal{M} \lesssim 1$. This is what could have been expected from the particle simulations [Kucharek & Scholer, 1995; Scholer & Kucharek, 1999] however, which showed that pickup ions benefit most in energy by being reflected from the shock. They can be accelerated up to energies several ten times the upstream flow energy. While this effect is strongest at quasi-parallel shocks, there is some discrepancy left between the interpretation of the observations and the predictions from the simulations. Possibly quasi-perpendicular shocks are better reflectors of pick-up ions than was expected. Or otherwise the shock, for the pick-up ions behaves more like a quasi-parallel shock.

11.5.4 Energetic Particles

As we have noted, the most vivid interest in the Termination Shock is from the side of the acceleration of ions (and possibly also electrons) to high energies into the range of Cosmic Ray energies. Theory has predicted that the Termination Shock would serve as the source region of the Anomalous Cosmic Ray component that populates the heliosphere. This component is clearly accelerated only in the heliosphere or in its boundary regions and consists of high energy interstellar ions which are produced when entering the heliosphere. Anomalous Cosmic Rays are thus a subset of the interstellar pickup ions in the heliosphere. The Voyager 1 and Voyager 2 Termination Shock crossings have therefore been looked forward with the greatest expectations. Would the Termination Shock be identified as the ultimate source of the Anomalous Cosmic Ray component?

As noted above, the plasma observations at the low energy end of the particle distribution have clearly shown that it is not the plasma that picks up the bulk energy loss of the flow across the Termination Shock. Neither is the plasma heated sufficiently, nor is it decelerated sufficiently. The deceleration could be accounted for by including the pickup ion component [Richardson *et al*, 2008]. In this section we now analyse the Voyager obser-

uations of both, the energetic particle observations [Decker *et al*, 2005, 2008] and the measurements of the Anomalous Cosmic Ray component [Stone *et al*, 2005, 2008].

Figure 11.4 in its lower panel on its right shows the big surprise provided by the energetic particle observations during the first Termination Shock crossing by Voyager 1. All expectations culminated in finding the maximum in the energetic ions (and possibly even electrons) right at the Termination Shock crossing. Surprisingly this was not the case. Instead, about 50 days before the shock crossing the energetic particles >70 MeV/nucleon dropped gradually to a minimum that relaxed during approach to the shock. However, the fluxes of energetic particles increased almost continuously with progressing distance of the Voyager 1 spacecraft from the Termination Shock when moving into the heliosheath. This observational result was not only most surprising, it also was taken as killing the assumption that the Anomalous Cosmic Ray component is caused at the Termination Shock by a shock surfing/first-order Fermi acceleration process acting on the interstellar pickup ions in the outer heliosphere. Indeed, when the maximum of the Anomalous Cosmic Ray flux is not found at the location of the Termination Shock but much farther out in the heliosheath or possibly even at the heliopause or in the ion wall transition region or the interstellar bow shock it is difficult to accept that the Termination Shock accelerates these particles to the high energies they possess in the inner heliosphere.

Galactic Cosmic Rays

Now, after Voyager 2 has crossed the Termination Shock the particle spectra at the two spacecraft can be compared. This has been done in Figure 11.7 which shows the differential energy fluxes of the energetic protons as function of energy per nucleon. The plot is divided into the three particle components Termination Shock protons TSP of energy $<(3-5)$ MeV, Anomalous Cosmic Rays in the range $(6-60)$ MeV, and Galactic Cosmic Rays >90 MeV.

There is quite a number of interesting results in this figure. The first is that the GCR spectra for the two Termination Shock crossings at the different locations in the heliosphere do about agree. The Voyager 1 crossing being farther out saw a slightly more intense GCR flux. However the difference is not large suggesting that at the two locations the heliospheric / heliosheath screening effect on the GCR component is comparable. In other words, even though the crossing of the Termination Shock at Voyager 2 was roughly 10 AU farther inside the heliosphere than the Voyager 1 crossing, the boundary of the heliosphere cannot have been much farther out than at Voyager 1. This agrees with the claim that the entire heliosphere is deformed being smaller in its southern than in its northern part. Since the Lyman α observations of Lallement *et al* [2005] it is known that this deformation is the result of the presence of the comparably strong interstellar magnetic field that penetrates the LISM in the Local Interstellar Cloud (see Figure 11.3). With further progress into the heliosheath the GCR intensities on both spacecraft have increased at the same rate and are practically (within the error of observation) equal, again reflecting the heliospheric filter effect which cause the positive radial gradient in the GCR flux. We note than, in principle, from these observations the heliospheric cosmic ray diffusion coefficient could be determined.

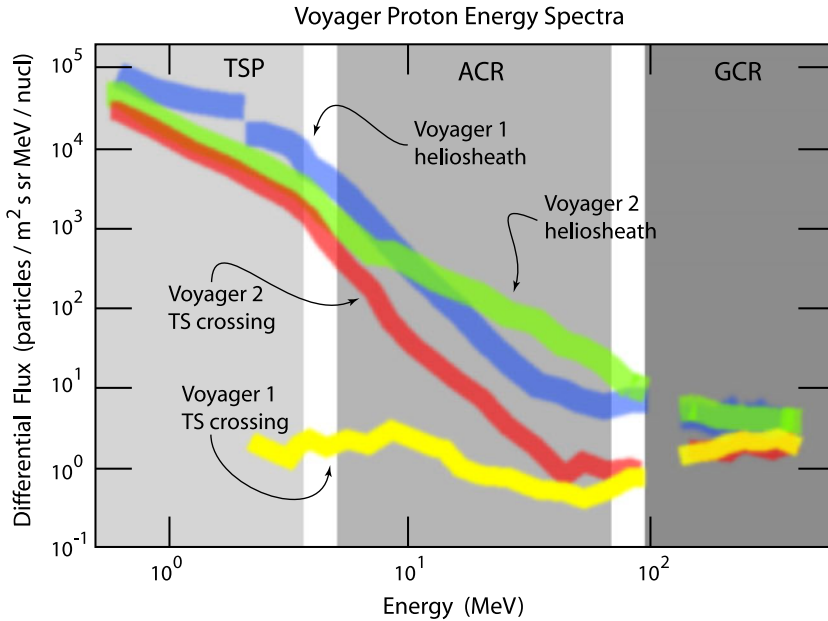


Figure 11.7: The Voyager high energy proton spectra at Termination Shock crossing and in the heliosheath [data taken from *Stone et al., 2008*]. Shown are the proton differential energy fluxes per nucleon (in the case of the protons this is per particle) as function of energy. The width of the lines corresponds about to the error or the measurement. The energy range is divided into the three different populations: Termination Shock protons $< (3-5)$ MeV, Anomalous Cosmic Rays $(6-60)$ MeV, Galactic Cosmic Rays > 90 MeV.

Termination Shock Particles

These particles are essentially the pickup ions that are generated in the charge exchange process between the interstellar gas that penetrates the heliosphere plus a small fraction of upstream protons from the solar wind flow that have been reflected and accelerated at the Termination Shock. The acceleration of pickup ions proceeds in two steps. First they become accelerated in the upstream convection electric field $\mathbf{E}_u = -\mathbf{V}_u \times \mathbf{B}_u$ until they merge into the main flow. In this process they perform cycloid orbits in the plane perpendicular to the magnetic field and reach velocities $v = 2V_u$ twice as high as the streaming velocity or energies per nucleon $\mathcal{E}_{pu}/\text{nucleon} = 4\mathcal{E}_u$ four times the upstream streaming energy. In a second step, then, they become reflected from the shock and are accelerated further. Remember that because of their higher energy and gyroradii pickup ions are much better reflected from a shock than the main flow particles and should thus become much stronger accelerated in first order Fermi acceleration than upstream protons. We should, however, keep in mind that this is the expectation from essentially one-dimensional numerical hybrid, test particle and also PIC simulations of quasi-perpendicular respectively quasi-parallel shocks, which is to be proved by observations.

It is interesting to compare the conditions at 1 AU and at the location of the Termination Shock at 80–100 AU. Since the nominal solar wind speed $V_u \sim 500$ km/s is assumed to be constant throughout the heliosphere, the motional electric field E_u decreases with distance as the interplanetary magnetic field. At 1 AU its value is $E_u \sim 2.5$ mV/m. At the location of the Termination Shock it is roughly two orders of magnitude less in the interval $0.025 < E_u < 0.03$ mV/m. One might thus suspect that its role is diminished. However, what counts are the scales. Because of the weaker magnetic field in the outer heliosphere the gyro-radius and thus the radius of the pickup particle cycloid orbit is two orders of magnitude larger such that the pickup ion becomes accelerated at a hundred times slower rate but reaches the same energy as at 1 AU. On the other hand, the energy of a particle that is reflected from the shock, experiences the upstream motional electric field, and becomes accelerated along the shock surface, depends on the scale tangential to the shock surface. At 1 AU this scale is typically of the order of an Earth radius $1 R_E$ for acceleration at the Earth's bow shock. The shock reflected particle thus gains an energy $\Delta \mathcal{E}_{BS} \sim 16$ keV in one reflection. The typical scale in the outer heliosphere is 1 AU, however. Hence, the typical energy gain of a shock reflected particle at the Termination Shock is in the range $3.75 < \Delta \mathcal{E}_{TS} < 4$ MeV, which shows the different orders of magnitude between the effects of a small planetary bow shock and the huge heliospheric Termination Shock.

Energy spectra of the Termination Shock particles (TSP) are shown on the left in Figure 11.7. Voyager 1 detected a low and about constant pickup ion flux in the TSP range, of the order of ~ 2 particles/(m² s sr MeV/nucleon). This contrasts with the 3–4 orders of magnitude higher TSP fluxes detected in the southern heliosphere at the Voyager 2 Termination Shock crossing. As expected, in the heliosheath the Voyager 2 fluxes increase weakly, gradually becoming comparable to the Voyager 1 heliosheath fluxes in the TSP range. These have gradually increased by 3–4 orders of magnitude when Voyager moved into the heliosheath a fact that is not understood yet. Note that the latter three spectra break at about ~ 3 MeV, suggesting that the TSP proton energy range ends here, while the ACR proton energy range starts from here up. An indication of this break is also seen in the Voyager 1 shock crossing curve where the values above 3–4 MeV can be interpreted as already being ACR particles.

The Anomalous Cosmic Ray Component

The measurement of the ACR component is shown in the central part of Figure 11.7. Voyager 1 observed very low ACR fluxes when crossing the Termination Shock. In fact, when interpreting the spectral break at roughly (3–5) MeV as the transition from pickup ions to ACR then the maximum around ~ 10 MeV is the increase due to the residual ACR seen by Voyager 1 at the Termination Shock crossing. These ACR cannot have come from the TSP pickup ions at the Voyager 1 position, however. Their origin must be from remote. The spectral increase above ~ 50 MeV is the low energy extension of the Galactic Cosmic Ray spectrum. The latter is also seen in the Voyager 2 Termination Shock crossing.

The Voyager 2 Termination Shock crossing finds a much higher ACR flux which can well be interpreted as being the high energy tail of the TSP particle spectrum starting at the spectral break at $\sim (3\text{--}5)$ MeV and merging with the GCR spectrum at ~ 50 MeV.

In the heliosheath Voyager 1 observes the same ACR spectrum however a factor of 2–3 more intense thus suggesting that Voyager 1 is approaching the source of the ACR in the heliosheath while, obviously, the Termination Shock crossing of Voyager 2 is still not in the ACR acceleration region. But the entrance of Voyager 2 into the heliosheath brings a new surprise. The ACR spectrum follows only in its low energy part at $\sim(3\text{--}8)\text{ MeV}$ the Voyager 2 Termination Shock and Voyager 1 heliosheath ACR spectrum. From $\sim 8\text{ MeV}$ on to roughly $\sim 100\text{ MeV}$ it substantially exceeds the Voyager 1 heliosheath fluxes until it merges into the GCR spectrum which takes over at the higher energies.

Comparing the spectral slopes we find that the Voyager 2 Termination Shock crossing and the Voyager 1 heliosheath ACR spectra (after the spectral break) have approximate power law shapes with power $\alpha \sim -3.2$, while the Voyager 2 heliosheath spectrum has power law $\alpha \sim -1.5$. Voyager 2 thus is in a region where the acceleration process of the Anomalous Cosmic Ray component is most active and probably caused by turbulent acceleration [Fisk *et al*, 1974; Fisk & Gloeckler, 2006]. Thermodynamic arguments for turbulent power laws have been given by Fisk *et al* [2007], Treumann [1999a, b], and Treumann & Jaroschek [2008]. In this respect it is of interest that the Voyager 1 TSP proton fluxes below $\lesssim 7\text{ MeV}$ have since remained about constant indicating that Voyager 1 has entered the heliosheath region where the flux has reached maximum and remains to be homogeneous while the ACR fluxes in the range from 12 MeV/nucleon to 22 MeV/nucleon on both spacecraft, Voyager 1 and Voyager 2, continue to increase in intensity at about similar rates [Stone *et al*, 2008]. At higher energies $\gtrsim 60\text{ MeV/nucleon}$ the increases are identical and slow and the absolute values of the fluxes are the same, indicating the gradual contribution of the GCR component with approaching the heliopause.

The tantalising question that comes up when looking at these measurements is what the role of the Termination Shock is in the acceleration of the pickup ions and generation of the Anomalous Cosmic Ray component in the heliosphere. The observations contradict the simple idea of shock acceleration by multiple reflection in which case the maximum of ACR intensity should be found at and around the location of the Termination Shock. Since this is not the case for the Voyager 1 passage, the Voyager 1 orbit did not cross the region of particle acceleration at the shock. On the other hand, the Voyager 2 shock crossing would not be in disagreement with the shock acceleration role, would not the fluxes at the ACR energies for Voyager 2 as well increase farther out in the heliosheath causing the spectrum to be extraordinarily (though not impermissibly) flat in the ACR energy range. Obviously the example of the Termination Shock shows that the mechanism of shock acceleration is not very well understood yet.

A possible though very simple geometric interpretation has been put forward by McComas & Schwadron [2006]. This model is shown in Figure 11.8. It proposes a blunt Termination Shock with the spiral magnetic field lines of the interplanetary field cutting through the Termination Shock at the flanks of the heliosphere. Similar geometries had already been used by Gurnett & Kurth [2005] in their explanation of the observation of upstream Langmuir waves through magnetic connectivity to the Termination Shock. The assumption of McComas & Schwadron [2006] that has been repeated and elaborated on in follow up papers [Schwadron & McComas, 2007; Schwadron *et al*, 2008] is that the ACR particles are accelerated along the flanks of the Termination Shock where they have sufficient time

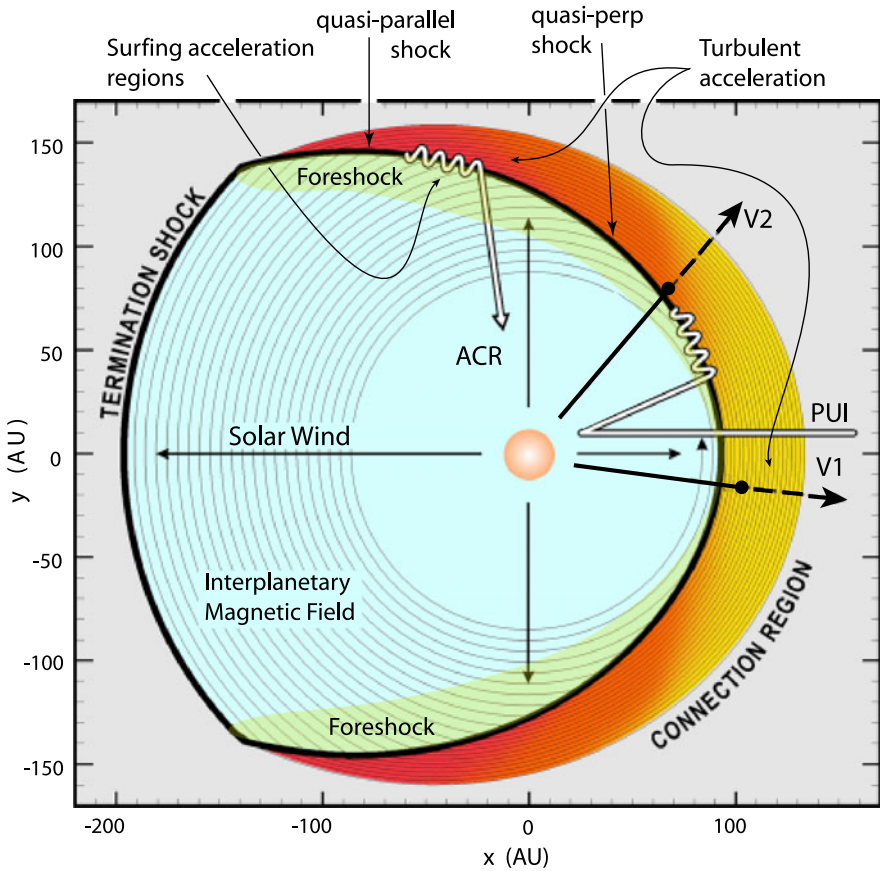


Figure 11.8: The ecliptic projection of the orbits of the Voyager 1 and Voyager 2 spacecraft, a blunt nose model of the Termination Shock, and the interplanetary magnetic field spiral [basic drawing taken from *McComas & Schwadron, 2006*, courtesy American Geophysical Union]. In this model the spiral interplanetary magnetic field lines cross the Termination Shock at the flanks of the Termination Shock where the shock is effectively quasi-parallel and possesses and extended inside foreshock, suggesting that the Anomalous Cosmic Ray component is accelerated at the TS flanks. Maximum fluxes of Anomalous Cosmic Rays are thus not seen in the Voyager 1 TS crossing close to heliospheric nose location where the Termination Shock is quasi-perpendicular. Voyager 1 should also see lower fluxes of accelerated pickup ions PUI. Voyager 2 should observe higher accelerated ACR fluxes. According to this picture Voyager 2 would still be in the quasi-perpendicular shock. Two types of acceleration shown are surfing acceleration of particles along the shock and turbulent acceleration in the extended heliosheath. In addition particles could be accelerated in a bouncing motion between the footpoints of a spiral field flux tube on the Termination Shock through the turbulent heliosheath.

to surf along the Termination Shock and become accelerated by crossing the shock back and forth. However, when flowing away from the shock along the field lines they will be found at the heliosphere nose position to occur farther out in the heliosheath. Thus the Voyager 1 observations are lacking these high fluxes which is in agreement with just seeing a slight enhancement around 10 MeV at Voyager 1 during the Termination Shock crossing

and the spectral increase of ACR at Voyager 1 in the deep heliosheath. The 10 MeV increase in this model is just the particles of the pickup ion distribution that are accelerated close to Voyager 1 at the shock. The prediction for Voyager 2 to observe much higher fluxes is also in agreement with observation and does well apply to the Voyager 2 crossing of the Termination Shock and the power law tail extending the pickup ion spectrum into the ACR range. The flux minimum upstream of the Termination Shock during approach of Voyager 1 shown in Figure 11.4 is also in agreement with such a model if only Voyager 1 has not crossed the source region. The higher fluxes observed earlier away from the Termination Shock are then due to energetic particles flowing from the Termination Shock flanks inward along the spiral magnetic fields.

The two observations of the very low energetic pickup ions at the Voyager 1 crossing and the enormous increase in the ACR spectra above 10 MeV far over what Voyager 1 observed is not really explained by this purely geometrical model, however. Obviously there are much less pickup protons at the Voyager 1 location. Close to the nose 10 MeV pickup ions must have been accelerated from roughly 4 keV/nucleon initial energy by a factor of 2×10^4 implying considerably more than one reflection from the Termination Shock, which implies a large area on the Termination Shock surface being involved in ion reflection. On the other hand, the flat spectrum of the ACR seen by Voyager 2 in the heliosheath suggests the strong involvement of heliosheath turbulence in the acceleration process. This is very similar to acceleration of energetic particles in ordinary solar wind turbulence. The heliosheath seems to host rather high turbulence levels that are capable of accelerating ACR.

One would thus tentatively conclude that the Termination Shock *alone* is not the main source of the Anomalous Cosmic Rays; the Termination Shock is, however, necessary for ACR in order to produce the high turbulence in the heliosheath as well as the seed population of pre-accelerated pickup ions which are accelerated further outside the Termination Shock by the heliosheath turbulence while still being confined to the heliosheath magnetic field. Voyager 2 being at the foot points of the average heliosheath magnetic field at the flank of the Termination Shock does observe the accelerated ACR particles while Voyager 1 is about to enter the heliosheath sufficiently deep for observing them. At the flank location of Voyager 2 the Termination Shock is much closer to being quasi-parallel than quasi-perpendicular. It is thus reasonable to assume that here the acceleration of ions into a diffuse ion component takes place and the surfing mechanism evoked is the kind of surfing we have discussed in the quasi-parallel shock theory: trapping of ions between approaching upstream and downstream pulsations and acceleration in the tangential convection electric field which produces the diffuse energetic ion component. The spiral interplanetary magnetic field crossing the Termination Shock and extending into the heliosheath until crossing the Termination Shock another time on the other side of the heliosphere provides a large magnetic bottle. Further acceleration of the ions can then proceed on the way of trapping the energetic particles in such spiral magnetic heliosheath flux tubes which have their foot-points on the Termination Shock. This was suggested long ago to happen at interplanetary shocks [Balogh & Erdős, 1985, 1991]. Oscillations of these foot points caused by fluctuations of the Termination Shock combined with the bouncing back and forth of the trapped particles are capable of pushing the energy of the trapped particles further up. Since, nat-

usually the acceleration efficiency depends on the length of the flux tube the maximum of the fluxes will be located deeper in the heliosheath and not at the Termination Shock. Thus the combination of quasi-parallel shock acceleration at the Termination Shock flanks, turbulent acceleration in the heliosheath and trapped acceleration in heliosheath flux tubes whose footpoints are located on the Termination Shock may explain the observations and provide a mechanism of generation of high energy Cosmic Ray particles. The *McComas & Schwadron* [2006] model and these conclusions are all quite speculative and have to be confirmed by future continuing observations provided by the two Voyager spacecraft.

11.6 Conclusions

The Termination Shock of the heliosphere is the remotest shock in the entire universe that has been subject to *in situ* observation by two spacecraft, Voyager 1 and Voyager 2. The next shock that is expected to be crossed by them, the heliospheric Bow Shock in the interstellar medium lies decades ahead. The Voyager crossings have not simply confirmed the theoretical predictions about the structure and nature of the largest – in the sense of dimensions – shock in the heliosphere. They have also added new problems and new knowledge which will form the base of future thought, modelling and investigation based on the available data.

It has been believed that the Termination Shock is a completely perpendicular or quasi-perpendicular shock. The observations have confirmed that – in the average – the Termination Shock is indeed quasi-perpendicular in the regions crossed by the two spacecraft. On the other hand it has also been found that it is variable on various time scales. Over large distances of several AU it has been occasionally connected to the Voyager spacecraft long before they ultimately entered and crossed the shock. This connection was identified from typical shock reflected particle fluxes, both ions and electrons, in sunward direction and from the observation of the Langmuir waves that are excited by electron beams along the spiral interplanetary magnetic field. One of the great surprises is the absence of any radio radiation that is emitted from the Termination Shock which for other shocks in the solar system and heliosphere is one of the main shock identifiers from remote. The conditions why the Termination Shock is not radiating may lie in the Termination Shock environment. Possibly the particles are too cool. But no firm explanation has been given yet.

In plasma waves the Termination Shock has turned out to behave like a strong collisionless supercritical shock possessing a narrow $\sim 1\lambda_i$ wide region of very large amplitude broadband plasma wave turbulence covering the shock ramp and overshoot transition the source of which must be sought in the electron dynamics inside the shock. This turbulence will sign responsible for particle trapping, surfing and heating, i.e. for transport and acceleration processes in the shock. In the magnetic field the Termination Shock shows the typical signs of an extended foot, ramp and overshoot, signatures of ion trapping and formation of ion phase space vortices with maximum magnetic field strength at their vertices.

Our current knowledge of shock dynamics lets us conclude that the Termination Shock is capable of reflecting ions which are accelerated in the foot, and the shock will be subject

to reformation that is driven by the reflected gyrating ion fluxes and by current instabilities like the modified two-stream instability. These foot currents and the quasi-periodic reformation modulate the shock along its surface, causing ripples and variations in the shock normal direction, sometimes providing it a quasi-parallel character which explains the occasional remote magnetic connectivity. However, the Termination Shock seems not to be such a strong shock as had been expected. Though being supercritical, it has compression ratio $\lesssim 2$ which might be an indication of mediation of the shock by the high density of pickup ions that are present in its environment.

A peculiarity of the Termination Shock is that it is embedded into the spiral solar wind magnetic field which is divided into sector structures roughly 0.3 AU apart in radial direction. These field lines may cut the Termination Shock at two positions on the flanks forming extended loops in the heliosphere with foot points on the Termination Shock. This kind of geometry may affect the property of the Termination Shock to accelerate pickup ions to the high energies of ACR when confining the shock surfing particles downstream in the loop configuration. Fluctuations of the foot points on the Termination Shock of these downstream spiral loops which can be caused by shock surface waves and oscillations of the Termination Shock and breathing of the heliosheath magnetic field in combination with the heliosheath turbulence may be responsible for heating and acceleration.

We may state in summary that the Termination Shock acceleration mechanism of charged particles in the energy range of the Anomalous Cosmic Rays is not yet understood, and the Termination Shock has not contributed yet sufficiently to its ultimately clarification, even though this clarification was expected to be achieved from the Voyager crossings. At the contrary, it seems that acceleration of particles into ACR is the result of a complicated interplay between the Termination Shock, the local geometry of the spiral interplanetary magnetic field at its contact with the Termination Shock, and the Heliosheath turbulence. Acceleration at the Termination Shock is not a process that is homogeneously distributed over the shock surface. It is regional, apparently being restricted to certain regions on the Termination Shock surface, and is otherwise affected by the conditions and the state of turbulence in the Heliosheath.

11.7 Update – 2012

Since the writing of this article three years have passed. During this period there has been some activity concerning the interpretation of the Voyager observations of the Heliospheric Termination shock which require a brief update of this chapter without claiming completeness.

The few most interesting findings which go beyond the above results concern the asymmetry of the termination shock, its weakness, and the particle distributions. The former, as discussed above, is a consequence of the magnetic field configuration, the weakness results from the very large contribution of the newly ionised interplanetary neutrals which as pick-ups [Wu *et al*, 2008; Müller *et al*, 2008; Zank *et al*, 2010] contribute strongly to the mass and energy of the termination shock environment. The particle populations are affected by the pick-up and secondary effects caused at the weak shock with the solar

wind simply passing through the termination shock being little accelerated and thus ‘cold’ downstream of the termination shock [Burrows *et al*, 2010].

In order to interpret the observations of the reduced effect of the solar wind and the dominance of pick-up ions [Intriligator *et al*, 2010], Scholer & Matsukiyo [2010] added one-dimensional numerical simulation at a quasi-perpendicular shock in the presence of a large fraction of pick-up protons as expected under heliospheric termination shock conditions. As expected they find that >50% of the pick-up proton population should be reflected from the termination shock while being split into two populations. The first results from reflection in the electric foot-field of the termination shock, the second population comes from in or even behind the ramp. This suggests that though the shock is probably not strong, the shock potential is large enough to reflect protons while the major potential drop occurs over the shock foot being caused by the large contribution of the energetic pick-up protons to the reflection and foot formation. In the quite realistic quasi-perpendicular case investigated the velocity difference between solar wind and pick-up protons in the foot excites the modified two-stream instability thus leading a large fraction of the heating to occur already in the pick-up shock foot while in the now quite extended shock transition the colder solar wind ions behave adiabatically thus being heated only very little of the solar wind in the termination shock [which has also been addressed by Burrows *et al*, 2010].

Particle acceleration in the light of these finding has been reviewed by Jokipii *et al* [2007]. What concerns the distribution function, so Chen *et al* [2008] suggested that it might change during crossing the termination shock from log-normal to normal. This might not be surprising, however, in the presence of even weak heating.

Finally, Webber *et al* [2007] have observed and followed the interaction of a large interplanetary shock with the heliospheric termination shock with both Voyager 1 and Voyager 2 spacecraft, finding substantial decreases in energetic particles after crossing from in to out accompanying the decrease in interplanetary shock speed by a factor of two from $\sim 600 \text{ km s}^{-1}$ to $210\text{--}270 \text{ km s}^{-1}$, and a 30–50% decreases in anomalous He.

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