The Microphysics of Collisionless Shocks

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Abstract

Shock waves in interplanetary (IP) space are of considerable interest due to their potential to damage ground based electronic systems and their ability to energize charged particles. The energization of charged particles at IP shocks has the obvious extrapolation to supernova shock waves, which are thought to be a candidate for generating the most energetic particles in the universe. The observations and theory behind collisionless shock wave evolution suggest that IP shocks should, for the most part, be stable structures which require energy dissipation. In a regular fluid, like our atmosphere, energy dissipation is accomplished through binary particle collisions transferring the loss of bulk flow kinetic energy to heat. Plasmas are mostly collisionless fluids, thus requiring other means by which to dissipate energy.

The studies herein were performed using wave and particle data primarily from the Wind spacecraft to investigate the microphysics of IP shock energy dissipation mechanisms. Due to their lower Mach numbers, more simplified geometry, and quasi-perpendicular nature, IP shock waves are an excellent laboratory to study wave-particle related dissipation mechanisms. Utilization of multiple data sets from multiple high time resolution instruments on board the Wind spacecraft, we have performed studies on the transition region microphysics of IP shocks.

The work began with a statistical study of high frequency (\( \gtrsim 1 \text{ kHz} \)) waveform capture data during 67 IP shocks with Mach numbers ranging from \( \sim 1-6 \) found ion-acoustic wave amplitudes correlated with the fast mode Mach number and shock strength. The ion-acoustic waves (IAWs) were estimated to produce anomalous resistivities roughly seven orders of magnitude above classical estimates.

Another study was an examination of low frequency waves (0.25 Hz < \( f < 10 \text{ Hz} \)) at five quasi-perpendicular IP shocks found the wave modes to be consistent with oblique precursor whistler waves at four of the events. The strongest event in that study had low frequency waves consistent with shocklets. The shocklets are seen simultaneously with diffuse ion distributions. Both the shocklets and precursor whistlers are seen simultaneously with anisotropic electron distributions unstable to whistler anisotropy and heat
flux instabilities. The IP shock with upstream shocklets showed much stronger electron heating across the shock ramp than the four events without upstream shocklets.

Further investigation of the atypical IP shock found the strong heating to be associated with large amplitude (> 100 mV/m) solitary waves and electron Bernstein waves. The observed heating and waveforms are likely due to instabilities driven by the free energy provided by reflected ions at this supercritical IP shock, not the DC macroscopic fields. The particle heating observed for the event with shocklets was observed to be different from other events with similar shock parameters, suggesting a different dissipation mechanism.

The work presented in this thesis has helped increase the understanding of the microphysics of IP shocks in addition to raising new questions regarding the energy dissipation mechanisms dominating in the ramp regions. The initial work focused on a statistical study of high frequency waveforms in IP shock ramps. The study results suggested a re-evaluation of the relative importance of anomalous resistivity due to wave-particle interactions. This assertion was further strengthened by the atypical particle heating observed in the 04/06/2000 event which we claimed clearly showed a dependence on the observed waveforms. Thus, the nearly ubiquitous observations of large amplitude IAWs in the ramp regions of IP shocks raises doubts about ignoring these high frequency fluctuations. In addition to these findings, we also observed a low frequency wave mode which is only supposed to exist upstream of quasi-parallel shocks with small radii of curvatures.

All of these findings have increased our knowledge of collisionless shock energy dissipation, but they have raised many questions regarding our current theories. We have raised doubts regarding the use of the solar wind electron distributions as one particle population. We have showed evidence to support the energy dependence of wave-particle interactions between low frequency whistler waves and ≤1 keV electrons. Thus, we conclude that in the analysis of IP shocks the microphysics can no longer be disregarded.
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Chapter 1

Introduction

1.1 The Sun and Solar Wind

The closest star to Earth, the sun, is a large ball of ionized gas called a plasma. The sun has intense variable magnetic fields which are the drivers of most solar and geomagnetic activity. The sun’s outer atmosphere, the corona, is hundreds of times hotter than lower visible layer of the solar disk. The apparent violation of the second law of thermodynamics has provoked immense interest in the dynamics of the solar atmosphere. The outer atmosphere is not bound to the sun and its expansion is observed as an ever present flow called the solar wind [Lang, 2000]. The solar wind is composed of electrons, protons, alpha particles, and heavy ions. The particle nature of the solar wind was first measured by a Russian scientist named I. Gringauz [Gringauz et al., 1960]. However, Gringauz et al. [1960] could only determine a relative ion flux and direction of flow. It was not until 1962 that Marcia Neugebauer and Conway W. Snyder used more than 100 days of Mariner 2 data to show that charged particles were continuously emanating from the sun [Neugebauer and Snyder, 1962].

The sun’s magnetic field has been found to flip roughly every 11 years going through periods called solar maximum and solar minimum. The maximum and minimum periods are defined based upon the number of sun spots observed. Sun spots are seen as dark localized regions on the visible disk of the sun and are now known to contain intense magnetic fields (up to $\sim$0.3 Tesla, or $\sim$10,000 times the terrestrial magnetic field strength). These localized regions of intense magnetic field are called active regions.
and are responsible for the solar storms which can excite a myriad of local terrestrial phenomena (e.g. the aurora). The variation of the sun’s magnetic fields have impacts on the expansion of the solar wind into the surrounding heliosphere.

It is now known that the sun expells roughly $10^9$ kg of mass per second at speeds ranging from roughly 300 km/s to over 800 km/s [Marsch, 2006]. The solar wind flow is almost entirely in the radial direction and its speed has strong latitudinal variations during solar minimum as discovered by the Ulysses spacecraft [McComas et al., 1998; Phillips et al., 1995]. Later studies found that after removing radial gradients, the high latitude regions showed little latitudinal variation [McComas et al., 2000]. During solar maximum, the variations in solar parameters are much less well defined.

1.1.1 Collisionless Fluids

A fluid made of positively and negatively charged particles is called a plasma. Plasmas can be considered a neutral gas if their microscopic electric fields are screened out over a distance known as the Debye length given by:

$$\lambda_{Ds} = \sqrt{\frac{e^2 k_B T_s}{n_s e^2}}$$  \hspace{1cm} (1.1a)

$$\lambda_{Ds} = \frac{V_{Ts}}{\omega_{ps}}$$  \hspace{1cm} (1.1b)

where $\varepsilon_o$ is the permittivity of free space, $k_B$ is the Boltzmann constant, $T_s$ is the temperature, $e^2$ is the charge squared, $n_s$ is the number density of particle species $s$, $V_{Ts}$ is the thermal speed of species $s$ ($= \sqrt{k_B T_s/m_s}$), and $\omega_{ps}$ is the plasma frequency of species $s$ defined as:

$$\omega_{ps} \equiv \sqrt{\frac{n_s q_s^2}{m_s \varepsilon_o}}$$  \hspace{1cm} (1.2)

The condition for defining a gas as a plasma is defined as:

$$n_s \lambda_{Ds}^3 \gg 1$$  \hspace{1cm} (1.3)

which states that the number of particles in a Debye sphere is large. Most plasmas are considered collisionless, which means that the Coulomb collision time, $\tau_c$, is much
larger than any relevant plasma time scale. In other words, the average mean free path of a particle is often much larger than the gradient scale lengths of waves or density fluctuations. The average mean free path of a charged particle in a fully ionized plasma can be estimated by multiplying the average time between collisions, $\tau_c$, and the average speed of the particles (assume $V_{Ts}$). A particle with speed $V_{Ts}$ undergoes a collision in unit time $\tau_c$ with the particles that lie within a cylinder of volume $\pi r_c^2 V_{Ts}$, where $r_c$ is the limiting impact parameter for which the deflection of the incident particle is sufficiently large to make the change in momentum comparable to the initial momentum of the incident particle. Given that the particle number density is $n_o$, the effective collision frequency is:

$$\nu_c = \frac{1}{\tau_c} \approx \pi r_c^2 V_{Ts} n_o$$

(1.4a)

$$\approx \frac{n_o e^4}{32\pi (m_e \varepsilon_o)^2 V_{Ts}^3}$$

(1.4b)

where we have assumed $r_c$ to be $\approx \delta_o/(m_e V_{Ts}^2)$ and $\delta_o = e^2/(4\pi \varepsilon_o)$. We used $m_e$ because their light mass makes it easier to cause significant deflections of their trajectories, which we define as a collision. A typical proton kinetic energy in the solar wind is roughly 10 eV, which corresponds to a thermal speed $V_{Tp} \sim 31$ km/s. The typical solar wind particle density is roughly $10^7$ particles per meter squared, thus a 10 eV proton experiences roughly $4 \times 10^{-7}$ collisions per second or $\tau_c \sim 2.5 \times 10^6$ seconds. So a typical solar wind proton moving at 31 km/s travels, on average, $7.75 \times 10^{10}$ meters in this time, which is roughly half of an astronomical unit or AU ($\sim 1.5 \times 10^{11}$ m). Thus, it is likely that the particles will interact with something else before colliding with another particle.

Since plasmas can act like a fluid by exhibiting bulk continuous motions and though they are not dominated by collisions, they can produce discontinuities [Kellogg, 1962]. However, in a plasma, the particles are affected by not only the regular fluid parameters, they are also affected by electric and magnetic fields due to their individual charges. The consequences of the added forces result in new types of discontinuities and shock parameters.

Discontinuities are changes in relevant parameters on scale lengths shorter than the relevant communication scale lengths of the medium. For a regular fluid, the relevant scale length is the mean free path for binary collisions. For a plasma, the relevant scale lengths are the particle gyroradius, $\rho_{gs}$, and/or inertial lengths, $c/\omega_{ps}$, where $s$
represents the particle species \(i.e.\) electrons or ions) and \(c\) is the speed of light. The parameter, \(\rho_{gs}\) is defined by:

\[
\rho_{gs} \equiv \frac{V_{\perp}}{\Omega_{cs}} \tag{1.5}
\]

where \(V_{\perp}\) is the velocity of a particle perpendicular to the magnetic field (often assumed to be the thermal speed of the particle), \(\Omega_{cs}\) is the particle gyrofrequency \((= q_s |B|/m_s)\), \(n_s\) is the particle number density, \(q_s\) is the particle charge, and \(m_s\) is the particle mass.

For reference, a typical proton in the solar wind near 1 AU will have a gyrofrequency of roughly 0.1 Hz \((B_0 \sim 6.56 \text{ nT})\), a plasma frequency of roughly 600 Hz \((n_e \sim 8.2 \text{ cm}^{-3})\), a thermal speed of roughly 30 km/s \((T_e \sim 13 \text{ eV})\), a thermal gyroradius of roughly 300 km, and inertial length of roughly 80 km. For future reference, we define microscopic as scales smaller than a proton gyroradius, \(\rho_{gp}\), and macroscopic as scales larger than \(\rho_{gp}\). There is not a sharp distinction between the two other than to say that the phenomena involved in macroscopic effects often can be treated as a fluid while microphysics requires kinetic treatments.

\section{1.2 Basic Hydrodynamics and Shock Physics}

\subsection{1.2.1 Nonlinear Wave Steepening and Discontinuities}

Fluids are a ubiquitous substance found in every region of the universe. Fluids can exist in a collision dominated or collisionless state. One of the characteristic properties of this state of matter is a set of phenomena known as discontinuities. Discontinuities can be described as a mathematical property defining a step-like or abrupt transition between two points along a curve. Discontinuities can occur in fluids in different forms. The simplest discontinuity is a density gradient with no flow across the sharp boundary, known as a contact discontinuity. An example of a contact discontinuity in a regular collisional fluid is a chemical explosion where a new source of hot gas is introduced to a system.

Wave discontinuities can arise from a nonlinear wave, which has steepened to a point where the flow, density, or temperature changes direction and/or magnitude in a
distance shorter than the characteristic scale length of the medium. The characteristic length is defined by the type of information transfer in the medium. In a regular fluid, the characteristic scale length is the mean free path for binary particle collisions. The mean free path of particles in our atmosphere is on the order of a micrometer and there are roughly $10^{19}$ particles per centimeter cubed at standard temperature and pressure. Thus, the number of binary collisions is easily sufficient to dissipate the free energy in regular hydrodynamic discontinuities.

Discontinuities require the nonlinear growth a wave propagating in a medium, called nonlinear wave steepening. The steepening can result from a dependence of the phase velocity on the amplitude of the wave, causing the larger amplitude waves to catch up with and overtake the smaller amplitude waves. For instance, acoustic waves having a finite amplitude of any wave mode must always steepen, if some form of resistance to steepening is ignored \cite{Shu,1992}. The equations of fluid mechanics are fundamentally nonlinear. Almost all nonlinear waves have a tendency to steepen except for some dispersive waves whose different Fourier components propagate at different speeds for small amplitudes. These waves have special finite amplitude solutions, called solitons, maintain their shape as they propagate because dispersion exactly balances steepening.

The process of wave steepening can be thought of as a series of small successive step waves trailing each other. The leading waves alter their respective downstream fluid speeds which define the propagation speeds of the trailing waves, allowing them to overtake the leading waves. A compressional wave, for instance, increases the fluid speed behind the compression which allows any compressional wave trailing the first compressional wave to propagate faster than the first. Thus, the intrinsic nature of compressional waves makes them prone to wave steepening \cite{Kennel et al.,1985}. The waves steepen until either some form of dissipation limits the steepening or the wave reaches what is called a gradient catastrophe and breaks (e.g. water waves seen as white caps on lakes). However, if there exists enough resistance to growth, the steepened wave may reach a balance between steepening and the resistive forces. The resisting component is energy dissipation due to non-conservative forces, much like friction between two surfaces. Energy dissipation is the transfer of energy from one form (e.g. kinetic energy) to another form (e.g. heat) through an irreversible process. The distinction of irreversible is necessary because reversible forms of dissipation cannot, by themselves,
produce specific types of discontinuities like shock waves.

1.2.2 Introduction to Shock Physics

Shock waves are type of propagating wave that occur obliquely throughout the universe as a discontinuous transition, called the ramp, between two regions defined as upstream (unshocked) and downstream (shocked). Shock waves occur in collision dominated media like our atmosphere and in collisionless media like interstellar space. They are seen in our atmosphere around supersonic aircraft and in space as supernovae, galactic/planetary/cometary bow shocks, and transient interplanetary (IP) events.

Across a shock wave, there is a discontinuous change in flow speed that results in a bulk change in kinetic energy across the ramp. The finite difference in kinetic energy is a source of free energy in the shock wave which must be dissipated in some way to conserve energy. It is well known that energy dissipation can be achieved easily in a regular collision dominated fluid, like our atmosphere, through binary particle collisions. The binary collisions transfer the free energy by randomly increasing their average kinetic energy, also known as heating. The loss of information through the randomization of particle kinetic energies produces an increase in entropy, thus the shock transition is irreversible. Shock waves manage to heat particles in relatively short distances defined by a characteristic scale size in the medium. The mean free path for binary particle collisions of particles in our atmosphere is on the order of a micrometer and there are roughly $10^{19}$ particles per centimeter cubed at standard temperature and pressure. However, as we discussed in Section 1.1.1, the average mean free path for a proton in the solar wind is on the order of an AU and the typical number density is $10 \text{ cm}^{-3}$. Thus, shock waves which exist in the interplanetary or interstellar medium cannot rely upon binary particle collisions to dissipate their energy.

Collisionless shock waves are capable of efficiently heating and/or accelerating charged particles. For instance, the collisionless shocks involved in supernovae or binary star collisions are thought to produce some of the highest energy cosmic rays \cite{Blandford and Eichler, 1987}. Cosmic rays are also produced locally, in our solar system by transient IP events or at the heliospheric termination shock. These particles can damage and/or destroy spacecraft or injure/kill astronauts. The transient IP events...
can also affect planetary magnetic fields by compressing the field, after resulting in strong aurora that induce currents in long wires, pipelines, or power grids. The induced DC currents are capable of destroying power stations [Pirjola 1989; Pirjola et al. 2000].

1.2.3 Shock Parameter Definitions

A shock is usually a relatively stable discontinuity, shown qualitatively in Figure 1.1. Here, stable refers to the temporal structure of the discontinuity. This is in contrast to the breaking water waves preferred by surfers, which are not stable. They are a nonlinearly steepened wave which does not have enough energy dissipation to balance the steepening. Shock waves often do have enough energy dissipation to balance the steepening, thus they are seen as a time-stationary discontinuity.

In many plasma regimes, the magnetic field is the direction of reference, whereas in a shock, the shock normal vector becomes the important direction. If we let \( \hat{n} \equiv \text{shock normal vector}, \theta_{Bn} \equiv \text{shock normal angle or angle between} \hat{n} \text{ and upstream (subscript 1) magnetic field, } B_1, \) then we can define the tangential and normal components of any given vector as:

\[
Q_n \equiv Q \cdot \hat{n} \\
Q_t \equiv (\hat{n} \times Q) \times \hat{n} \\
\equiv Q - (Q \cdot \hat{n}) \hat{n}
\]

where \( Q_n(t) \) refers to an arbitrary normal(tangential) vector with respect to the shock normal vector.

Figure 1.1 is an illustrative cartoon used to show the change in various plasma parameters one might measure in space across the shock. The parameters are shown in the shock reference frame. In the figure, each panel corresponds to a different plasma parameter(s) magnitude versus time. The panels are, from top to bottom: 1) the magnitude of the normal, \( U_n \), and tangential, \( U_t \), flow speeds (in the shock frame), 2) the magnitude, \( |B| \), and tangential component, \( B_t \), of the magnetic field, 3) the particle pressure, \( P (= n_o k_B T) \), 4) the magnitude of the average temperature, \( T \), and 5) the particle number density, \( n_o \). This figure is for illustrative purposes; real collisionless shock crossings are much more complicated and far less laminar.
Figure 1.1: An illustrative example of how the parameters change across a shock in the shock reference frame. The top panel shows the normal, $U_n$, and tangential, $U_t$, flow speeds, the second panel is the same but for the magnetic field. The third through fifth panels show the pressure, temperature, and density change across the shock. The image is specifically for a fast mode shock in a collisionless plasma. The subscripts, $n$ and $t$, refer to normal and tangential directions respectively with respect to the shock normal vector. For a regular fluid, simply ignore the panel that refers to the magnetic field.

For future purposes and definitions, an illustrative example of the shock geometry can be seen in Figure 1.2. The blue line represents the shock normal vector, $\mathbf{n}$, the red $\mathbf{B}_1$($\mathbf{B}_2$) represent the upstream(downstream) magnetic field vectors, $\mathbf{B}_{n1}$($\mathbf{B}_{n2}$) represent the tangential component of the upstream(downstream) magnetic field vectors, and $\theta_{Bn}$ is the shock normal angle. The shock normal angle is defined by:

$$\theta_{Bn} \equiv \tan \left| \frac{\mathbf{B}_n}{B_n} \right|$$

$$\quad = \tan \left| \frac{(\hat{n} \times \mathbf{B}) \times \hat{n}}{\mathbf{B} \cdot \hat{n}} \right|$$

where $\hat{n}$ is the shock normal vector as defined above.
Figure 1.2: An illustrative example of the often assumed geometry of a collisionless shock wave. The shock is assumed to have a planar surface with a well defined normal vector, $\mathbf{n}$. The angle between the incident magnetic field, $\mathbf{B}_1$, and $\mathbf{n}$ is $\theta_{\mathbf{n}}$. 
1.2.4 Conservation Relations

When observing collisionless discontinuities in space with satellites, it is necessary to show that time series measurements of specific quantities satisfy a sequence of conditions in order to verify the discontinuity as a shock wave [Vinas and Scudder, 1986]. In regular fluid dynamics, a discontinuous transition between an upstream and downstream state must conserve mass, momentum, and energy when considering a one-dimensional time-stationary discontinuity. The conservation of mass results from integrating the mass continuity equation over some volume which contains the shock transition. The continuity equation is a solution to a moment of the Boltzmann equation given by:

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \nabla f + \frac{\mathbf{F}}{m} \cdot \nabla \mathbf{v} f = \frac{\delta_c f}{\delta t}$$

(1.8)

where $f$ is the distribution function, $f = f(\mathbf{v}, \mathbf{x}, t)$, $\mathbf{F}$ is some external force, $\nabla_{\mathbf{v}}$ is the gradient operator in velocity space, and $\delta_c / \delta t$ is a collision operator. When $\delta_c / \delta t \to 0$, as in a collisionless plasma, the Boltzmann equation becomes the Vlasov equation.

We can define the number of particles, $dN$, in a phase space element, $d^3x \, d^3v$, in the following way:

$$dN = f(\mathbf{v}, \mathbf{x}, t) \, d^3x \, d^3v$$

(1.9)

where one assumes that the number of particles in $d^3x \, d^3v$ is sufficiently large such that $f(\mathbf{v}, \mathbf{x}, t)$ can be regarded as a continuous function of $\mathbf{v}$ and $\mathbf{x}$. The total number of particles can be found by integrating over all of phase space. The following quantities will be used later, so we define:

$$n_s = \int_V f_s(\mathbf{v}, \mathbf{x}, t) \, d^3v$$

(1.10a)

$$\mathbf{U}_s = \frac{1}{n_s} \int_V \mathbf{v} f_s(\mathbf{v}, \mathbf{x}, t) \, d^3v$$

(1.10b)

$$W_s = \frac{m_s}{2} \int_V v^2 f_s(\mathbf{v}, \mathbf{x}, t) \, d^3v$$

(1.10c)

$$\mathbf{P}_s = m_s \int_V (\mathbf{v} - \mathbf{U}_s) \, (\mathbf{v} - \mathbf{U}_s) f_s(\mathbf{v}, \mathbf{x}, t) \, d^3v$$

(1.10d)

where $n_s$ is the number density, $\mathbf{U}_s$ is the average velocity, $W_s$ is the kinetic energy density, $m_s$ is the mass, and $\mathbf{P}_s$ is the pressure tensor of species $s$ [Gurnett and Bhattacharjee, 2005].
Moment equations are calculated by multiplying the Vlasov equation by powers of the velocity then integrating over velocity space. As mentioned before, the continuity equation is a solution to a moment equation, specifically the zeroth moment given by:

$$\frac{\partial \rho_m}{\partial t} + \nabla \cdot (\rho_m \mathbf{U}) = 0 \quad (1.11)$$

where $\rho_m$ is the mass density, given by Equation 1.10a multiplied by the particle mass, and $\mathbf{U}$ is the bulk flow velocity, given by Equation 1.10b. The right-hand side of this equation is zero because no sources/losses of mass inside the discontinuity are assumed to be present; mass is neither created nor destroyed in the discontinuity. Momentum conservation requires a little more work because the Vlasov equation is multiplied by $v$ before integration is performed. After a great deal of arithmetic and a few assumptions, the momentum equation can be shown to be:

$$m_s n_s \left[ \frac{\partial \mathbf{U}_s}{\partial t} + (\mathbf{U}_s \cdot \nabla) \mathbf{U}_s \right] = n_s \mathbf{F} - \nabla \cdot \mathbf{P}_s \quad (1.12)$$

where $\mathbf{F} = q_s (\mathbf{E} + \mathbf{v} \times \mathbf{B})$, is the Lorentz force here.

In a magnetized fluid, known as magnetohydrodynamics or MHD, the conservation of each moment is altered slightly to accommodate the electric and magnetic forces involved. The conservation of momentum and energy can be written as:

$$\frac{\partial}{\partial t} (\rho_m \mathbf{U}_s) + \nabla \cdot \left[ \rho_m \mathbf{U}_s \mathbf{U}_s + \left( P + \frac{B^2}{2\mu_0} \right) \mathbf{I} - \frac{\mathbf{BB}}{\mu_0} \right] = 0 \quad (1.13a)$$

$$\frac{\partial}{\partial t} \left( \frac{1}{2} \rho_m U_s^2 + \frac{P_s}{\gamma - 1} + \frac{B^2}{2\mu_0} \right) + \nabla \cdot \left[ \frac{1}{2} \rho_m U_s^2 \mathbf{U}_s + \frac{\gamma P_s}{\gamma - 1} \mathbf{U}_s + \frac{\mathbf{E} \times \mathbf{B}}{\mu_0} \right] = 0 \quad (1.13b)$$

where $P$ is some scalar isotropic pressure, $\gamma$ is the statistical mechanics parameter which depends on the number of degrees of freedom, $\gamma = (d + 2)/d$ with $d$ being the degrees of freedom (also the ratio of specific heats). For closure, the pressure is typically defined by the adiabatic equation of state, $PV^\gamma = \text{constant}$. For future reference, a one-dimensional (degree of freedom) gas has $\gamma = 3$, an isothermal gas has $\gamma = 1$, a three dimensional gas has $\gamma = 5/3$, and an incompressible gas has $\gamma = \infty$ [Gurnett and Bhattacharjee, 2005].

If we let $\{\}$ define the difference between downstream and upstream quantities across a discontinuity, then the set of relationships defined by these conservation laws are
known as the Rankine-Hugoniot relations, given by Garnett and Bhattacharjee [2005]:

\[
\{\rho_m U_n\} = 0 \quad (1.14a)
\]

\[
\left\{ (\rho_m U_n) U_t - \frac{B_n B_t}{\mu_0} \right\} = 0 \quad (1.14b)
\]

\[
\left\{ \rho_m U_n^2 + P + \frac{B_t^2}{2\mu_0} \right\} = 0 \quad (1.14c)
\]

\[
\left\{ \left( \frac{1}{2}\rho_m U^2 + \frac{\gamma P}{\gamma - 1} + \frac{B^2}{\mu_0} \right) U_n - (\mathbf{U} \cdot \mathbf{B}) \frac{B_n}{\mu_0} \right\} = 0 \quad (1.14d)
\]

\[
\{ U_n (\mathbf{n} \times \mathbf{B}_t) + (\mathbf{U_t} \times \mathbf{n}) B_n \} = 0 \quad (1.14e)
\]

\[
\{ \mathbf{n} \cdot \mathbf{B} \} = 0 \quad (1.14f)
\]

where Equation 1.14a corresponds to mass continuity, Equation 1.14b is the tangential momentum, Equation 1.14c is the normal momentum, and Equation 1.14d is energy, while the last two equations are specific to MHD and Maxwell’s equations. Equation 1.14e results from an assumption made in MHD where the plasma has an infinite conductivity. Using a simplified version of Ohm’s law and assuming an ideal MHD plasma, then \( \mathbf{E} = -\mathbf{U} \times \mathbf{B} \). Using this assumption, we assume the discontinuity is a sharp boundary and so there can be no sudden time variation of \( \mathbf{B} \) within the layer. This results in Faraday’s law changing to \( \nabla \times \{ \mathbf{E} \} = 0 \), which after linearizing (discussed in the next chapter) goes to \( \mathbf{n} \times \{ \mathbf{E} \} = 0 \), where \( \mathbf{n} \) is the shock normal vector. Equation 1.14f results from linearizing Maxwell’s equation which claims no existence of magnetic charges (monopoles). Using this set of equations, one can determine the relevant shock parameters as described by Vinas and Scudder [1986]. Using the adiabatic equation of state we can replace the scalar pressure by recalling that the speed of sound in a fluid, \( C_s \), is given by \( \partial P/\partial \rho_m \) and using \( P \rho_m^{-\gamma} = \text{constant} \), we see:

\[
C_s^2 = \frac{\partial P}{\partial \rho_m} \quad (1.15a)
\]

\[
\frac{\partial P}{\partial \rho_m} = (\text{constant}) \gamma \rho_m^{\gamma - 1} \quad (1.15b)
\]

\[
= \frac{\gamma P}{\rho_m} \quad (1.15c)
\]
which results in Equations 1.14c and 1.14d going to:

\[
\left\{ \begin{array}{l}
\rho_m \left( U_n^2 + \frac{C_s^2}{\gamma} \right) + \frac{B_t^2}{2\mu_o} = 0 \\
\rho_m \left( \frac{1}{2} U^2 + \frac{\rho_mC_s^2}{\gamma - 1} + \frac{B^2}{\rho_m\mu_o} \right) U_n - \left( \mathbf{U} \cdot \mathbf{B} \right) \frac{B_n}{\mu_o} = 0.
\end{array} \right.
\]

(1.16a)

(1.16b)

(1.16c)

As mentioned previously, when an object/wave propagates through a medium faster than the relevant speed of information transfer, a shock wave may form. In a collisionless plasma, the relevant speed is known as the fast mode or magnetosonic speed. A magnetosonic wave is a compressive mode where the density fluctuations are in phase with the magnetic fluctuations leading to a wave whose phase speed depends on both the sound speed and Alfvén speed, given by:

\[
2C_f^2 = (C_s^2 + V_A^2) + \sqrt{(C_s^2 - V_A^2)^2 + 4C_s^2V_A^2 \sin^2 \theta_Bn} = (C_s^2 + V_A^2) + \sqrt{(C_s^2 + V_A^2)^2 - 4C_s^2V_A^2 \cos^2 \theta_Bn}
\]

(1.17a)

(1.17b)

where \( V_A = B/\sqrt{\mu_o \rho_m} \) is the Alfvén speed. There are three different relevant speeds in a plasma, \( C_f, C_{\text{slow}}, \) and \( V_A \) corresponding to the three MHD wave modes. Therefore, one might expect there to be three different Mach numbers given by:

\[
M_f = \frac{U_n}{C_f} \\
M_{\text{slow}} = \frac{U_n}{C_{\text{slow}}} \\
M_A = \frac{U_n}{C_A}.
\]

(1.18a)

(1.18b)

(1.18c)

Though the transition across the shock is an abrupt change, allowing little time for heat to flow, it would be incorrect to use only the adiabatic equation of state to describe the shock. The reason is that the irreversible nature of the shock implies that the fluid is not in a succession of equilibrium states [Gurnett and Bhattacharjee, 2005; Shu, 1992]. The Rankine-Hugoniot relations only refer to the asymptotic values far from the shock transition region, not within the dissipative scale lengths [Kennel et al., 1985]. As discussed previously, the bulk kinetic energy across the shock is not conserved, thus it is transferred to some other form in the shock transition region. The transformation
of the energy into some other form (i.e. heat) results in the irreversibility of the shock. One can show that the entropy increases across the shock, proving that the transition is irreversible, but we will not discuss that here \cite{Gurnett_and_Bhattacharjee,2005;Shu,1992}.

In this thesis, I utilize shock parameters determined by \cite{Kasper,2007} using the conservation relations discussed in this section.

1.2.5 Energy Dissipation Mechanisms in Collisionless Shocks

Collisionless shock waves are known to be efficient mechanisms by which charged particles can be heated and/or accelerated. The collisionless nature of these shocks prevents binary particle collisions from dominating the transfer of bulk kinetic energy to thermal energy. Thus, energy dissipation mechanisms have been a focus for research \cite{Mellott,1984} since the prediction of collisionless shock waves \cite{Kellogg,1962}. Three possible energy dissipation mechanisms discussed here are wave dispersion \cite{Mellott_and_Greenstadt,1984}, anomalous resistivity due to wave-particle interactions \cite{Gary,1981}, and particle reflection \cite{Edmiston_and_Kennel,1984;Kennel,1987}.

The type of energy dissipation important in a collisionless shock wave depends strongly on the magnetosonic or fast mode Mach number, $M_f$, shock normal angle, $\theta_{Bn}$, and the ratio of particle to magnetic pressures called the plasma beta, $\beta$ \cite{Mellott,1984}. At low Mach number shocks, the energy dissipation has been shown to include wave dispersion, thermal conduction, and anomalous resistivity \cite{Mellott_and_Greenstadt,1984;Mellott,1985;Wilson_III_et_al.,2007}. As the Mach number increases, one can show that there exists a Mach number where resistivity cannot limit the wave steepening. For a true shock transition, the form of energy dissipation, within a few wavelengths, must act to increase upstream perturbations (i.e. whistler waves) approaching the shock and decrease downstream perturbations leaving the shock \cite{Kennel_et_al.,1985}. However, when the downstream flow speed exceeds the downstream sound speed, stationary point analysis shows that perturbations downstream will grow. This implies that the shock is unstable. Typically when one includes dispersion, conduction, resistivity, or viscosity into the Rankine-Hugoniot relations, the mechanisms producing these terms are not considered. In other words, the inclusion of these terms into the conservation relations across a discontinuity are devoid of any physically significant content. Since the analysis
is done on the asymptotic values conserved across the ramp, the specifics of these effects can be ignored and treated simply as sources or losses of energy. The type of dissipation also affects the scale length of the shock transition region. In the following paragraphs, we will discuss the types of energy dissipation focused on in this thesis.

There are three basic scale lengths in a two-fluid model of a plasma, $c/\omega_{pe}$, $c/\omega_{pi}$, and $\lambda_{De}$. These scale lengths serve as deterministic limiting factors for the time scales over which currents or charge densities can change. So as a nonlinear compressional wave steepens to dispersive scale lengths (larger than resistive scales), the dispersion can limit steepening. The dissipation is accomplished by the radiation of a nonlinear wave which carries the short-wavelength energy in the dispersive range away from the compressive nonlinearly steepened wave. The nonlinear wave, originally radiated by the steepened front, is later dissipated by wave damping. Whether these short-wavelength waves propagate faster or slower than the steepened front determines whether they lead or trail the front itself. If the front forms into a shock and dispersion is limiting the steepening of the shock, then the radiated nonlinear waves must damp out asymptotically away from the shock. In the simplest case, the entire wave train is time-stationary in the shock frame, which implies that the nonlinear waves are phase standing with respect to the shock. In other words, their wave vector is aligned with the shock wave vector and their phase speeds match. To carry energy away from the shock, the wave group velocity must be greater than (less than) the upstream (downstream) flow speed if it phase stands upstream (downstream) [Kennel et al., 1985].

At short wavelengths, the fast mode converts to an elliptically polarized whistler wave whose phase speed exceeds the fast mode phase speed. As the wavelength continues to decrease, the fast mode phase speed increases which implies a leading nonlinear wave train for shocks propagating obliquely to the magnetic field. It should be noted that only low Mach number shocks exhibit a leading wave train and that all dispersive scale lengths are larger than resistive dissipative scale lengths. Wave dispersion has been well investigated at low Mach number (≲2) collisionless shock waves by Mellott and Greenstadt [1984] and Mellott [1984, 1985].

Ion-acoustic waves (IAWs) have long been believed to provide anomalous resistivity in collisionless shock waves [Sagdeev, 1966] because the current threshold is low when $T_e > T_i$ and the Debye length is small compared to electron inertial scales. In the
solar wind, $\omega_{pe} \gg \Omega_e$, which causes the IAWs to become decoupled from the magnetic field. This allows the waves to interact with the bulk of the electron distribution, which can lead to energy dissipation through particle heating \cite{Dum_1974, Gary_1981}. In the case of collisionless shock waves, IAWs are thought to be driven by a relative drift between electrons and ions across the magnetic field, or a cross-field current. Anomalous resistance results from radiated waves, due to a relative drift between electrons and ions, scattering the drifting particles in such a way as to reduce the drift between the two species. Thus, the waves act to mediate a transfer of momentum from the electrons to the ions \cite{Kennel_1985}. It is important to note that there are number of possible wave modes/instabilities which have been suggested to provide anomalous resistivity at quasi-perpendicular shock waves besides current-driven IAWs \cite{Gary_Sanderson_1970, Gary_1970a, 1981, Lemons_Gary_1978}, including lower-hybrid waves \cite{Lemons_Gary_1978}, lower-hybrid drift instability \cite{Davidson_Gladd_1975}, two-stream or modified two-stream instabilities \cite{Kellogg_1965, Lemons_Gary_1978}, electron cyclotron drift instability \cite{Forslund_1970, 1971, 1972, Lampe_1971a, 1971b, Wong_1970}, and the Buneman instability \cite{Buneman_1959, Kellogg_1964}.

Only a few studies have been done on resistive dissipation \cite{Wilson_III_2007}, and its significance with respect to the other possible mechanisms is not well understood. It is thought that wave-particle interactions produce an anomalous resistivity that limits the cross-field currents in the shock ramp that produce the abrupt change in magnetic field \cite{Gary_1981}. Though some studies suggest that wave-particle interactions are not necessary to heat the downstream plasma \cite{Hull_1998, 2000, Hull_Scudder_2000, Hull_2001}, recent observations \cite{Wilson_III_2007} and simulation studies using realistic mass ratios \cite{Petkakiet_al_2006} suggest otherwise.

In classical fluid theory, there is a Mach number called the critical Mach number, $M_{cr}$, above which resistive energy dissipation can no longer limit wave steepening \cite{Edmiston_Kennel_1984, Kennel_1985, Kennel_1987}. To determine this parameter, one assumes that the downstream normal flow speed is equal to the downstream sound speed. Thus, any acoustic waves which would normally propagate with the shock in the ramp would be convected into the downstream making them no longer
Figure 1.3: An example of a supercritical quasi-perpendicular IP shock observed by Wind on 10/21/2001 illustrating the different regions of a collisionless shock. The foot of the shock is highlighted by the blue region, the ramp by green and overshoot by red.
a source of energy dissipation in the ramp. The often quoted $M_{cr} \approx 2.7$ is valid only for a perpendicular shock (shock normal angle, $\theta_{Bn} = 90^\circ$) propagating into a cold plasma. For more realistic ranges of $\theta_{Bn}$ and temperature, typical solar wind conditions will actually yield $M_{cr} \sim 1-2$ [Edmiston and Kennel, 1984; Kennel, 1987], values often observed at interplanetary (IP) shocks which suggested particle reflection may occur even at low Mach number shocks.

In the simplest treatment one can ignore the specifics of the reflection process and treat the reflected particle population as an extra energy loss term in the Rankine-Hugoniot relations. However, the consequences of the microphysics behind ion reflection leads to a number of instabilities which help dissipate energy. For instance, if ions reflect and stream upstream of the shock they can excite ion-ion instabilities [Gary et al., 1981]. The reflected ions can also interact with the incident electrons producing a modified two-stream instability, which can act to slow the incident plasma [Matsukiyo and Scholer, 2003, 2006a]. The deceleration of the incident flow compresses the plasma producing a shock foot. The relative fraction of reflected ions could potentially be reduced by cross-field current driven instabilities in the shock ramp [Forslund et al., 1984].

Figure 1.3 highlights the different regions of a quasi-perpendicular shock observed by the Wind spacecraft on 10/21/2001. The structure of collisionless shocks is often defined by the magnitude of the magnetic field. The shock itself requires a discontinuous transition from a high speed upstream to a low speed thermalized downstream, referred to as the ramp (see green region of Figure 1.3). The ramp of a magnetized shock can be defined as the duration from the point of lowest magnetic field immediately preceding the discontinuous transition to the point of highest magnetic field immediately following the discontinuous transition [Farris et al., 1993].

Particle reflection/energization at collisionless shocks has been extensively studied [Paschmann et al., 1980, 1981, 1982; Meziane and D’Uston, 1998; Meziane et al., 2004a,b]. Furthermore, ISEE 1 and 2 observations of ion reflection at a subcritical terrestrial bow shock have found evidence to suggest that the transition between subcritical and supercritical may not be abrupt [Greenstadt and Mellott, 1987]. Thus, ion reflection is likely an ever-present dissipation mechanism for quasi-perpendicular shocks, but it does not become a dominant mechanism until higher Mach numbers.

The magnetic foot, a region immediately preceeding the ramp with thickness of
roughly an ion inertial length \((c/\omega_{pe})\), is created by reflected gyrating ions (blue region of Figure 1.3) \cite{Leroy1982, Livesey1984, Paschmann1982}. The gyrating ions are produced by nearly specular reflection, where roughly 20% of the incident ions are reflected \cite{Leroy1982, Paschmann1980, Thomsen1985a}. The foot region itself supports a large fraction of the energy dissipation in supercritical shock waves by slowing and pre-heating the incident plasma \cite{Leroy1982}. The ions skip along the shock front but the electrons remain constrained to the magnetic field and the relative drift between the two produces a current. It is this current that causes the small increase in magnetic field strength which defines the foot. Such gyrating ions are sources of free energy for multiple instabilities \cite{Matsukiyo2006a} and they can produce instabilities upstream of the shock which alter the Rankine-Hugoniot conditions \cite{Scholer1971}.

Giacalone et al. \cite{1991} examined the effect of the shock overshoot on reflection coefficients of ions at collisionless shocks. They found that an increase in the overshoot amplitude increased the number of reflected particles and the loss of energy by the transmitted particles. Leroy \cite{1983} found that a portion of the incident ion population is directly transmitted and accumulates in the immediate post-ramp region developing an overshoot as a natural consequence of the Rankine-Hugoniot conservation relations. Mellott and Livesey \cite{1987} found that magnetic overshoots resulted from a combination of reflected ions and turbulence in supercritical shocks. Paschmann et al. \cite{1982} observed gyrating ions at a nearly perpendicular bow shock crossing. The gyrating ions were clearly distinct from the specularly reflected ions. Such gyrating ions are thought to influence if not produce the shock overshoot.

Figure 1.3 is four IP shock examples of 2.5 minutes magnetic field data taken from the Wind spacecraft centered on the shock ramps. The two panels on the left are examples of quasi-parallel shocks and the two on the right are quasi-perpendicular. The top left panel is an approximately laminar quasi-parallel shock and the turbulent counterpart (directly below) shows the chaotic transition region of a supercritical quasi-parallel shock. However, for both quasi-perpendicular examples, the fields are much more laminar when viewed on these time scales. The figure is useful in illustrating the range of turbulence one can observe in IP shock waves.
Figure 1.4: Four examples of a super and subcritical quasi-perpendicular and quasi-parallel IP shocks observed by Wind. The two panels on the left are examples of sub and supercritical quasi-parallel IP shocks. The two panels on the right are the same for quasi-perpendicular. Notice that the term laminar is a relative term for each type of shock wave.

1.3 Shock Types

1.3.1 Bow Shocks

Bow shocks form when an object obstructs a supersonic flow in some medium. Examples would be a planet or comet in the solar wind. The solar wind is highly supersonic, thus an object roughly at rest with respect to the sun would appear to the solar wind as a ballistic object moving at supersonic speeds. They were predicted to exist upstream of planetary magnetospheres by Kellogg [1962] and first measured by Sonett et al. [1964].

An illustrative cartoon of the Earth’s foreshock, Figure 1.5, defines the different regions of a bow shock. There are two main types of collisionless shock geometries, quasi-parallel and quasi-perpendicular. Quasi-parallel shocks occur when the angle between the shock normal vector (the blue \( \hat{n} \)) and the interplanetary magnetic field (IMF), called the shock normal angle or \( \theta_Bn \), is \( \leq 45^\circ \). A quasi-perpendicular shock
Figure 1.5: This figure is an illustrative example of the Earth’s foreshock, bow shock, and magnetopause. The different regions are labeled appropriately along with illustrative examples of quasi-parallel versus quasi-perpendicular shock definitions. The yellow dots are supposed to represent electrons while the red protons. Due to the higher speeds of electrons, on average, one can see the electron foreshock extends further upstream than the ion foreshock. Also note, the quasi-parallel section of the bow shock appears much more turbulent than the quasi-perpendicular due to the reflected ions. This image is adapted from Plate 1 in *Tsurutani and Rodriguez*, [1981].
occurs when $\theta_{Bn} > 45^\circ$. One of the more important differences between a collisionless shock and a shock wave in a regular fluid is the capacity for collisionless shock waves to communicate with the upstream medium. This is accomplished by reflecting particles off the shock front that propagate away from the shock \cite{Edmiston1984, Greenstadt1987, Kennel1987}. The region upstream of a shock that can communicate with the shock is called the foreshock region. As a consequence of their much smaller mass, electrons can often travel further upstream than ions. This can be seen as the yellow region illustrated in Figure 1.5, called the electron foreshock. The ion foreshock corresponds to the red region of Figure 1.5.

1.3.2 Coronal Mass Ejections

Coronal Mass Ejections (CMEs) are large transient eruptions of plasma from the sun’s surface. Once a CME leaves the Sun, it is usually called an interplanetary CME (ICME). They are the result of the largest explosions in the solar system releasing up to $\sim 10^{34}$ ergs (or $\sim 10^{27}$ joules) This is roughly the equivalent of the energy released by detonating 20 million 100 Megatons nuclear bombs \cite{Lang2000}. CMEs were first discovered using coronagraphs on the OSO-7 and Skylab spacecraft during the early 1970’s \cite{Kahler1992}. The CME is observed in white light as a bright ‘umbrella-like’ structure on the leading edge of a dark cavity. The light observed is caused by Thomson scattering of photons by free electrons. Thus the intensity depends on the density of electrons, not their temperature. The bright leading edge is a high density plasma known as the sheath while the dark trailing cavity is a low particle density high magnetic field region, called a magnetic cloud. Upwards of 5-50 billion tons of plasma can be ejected at speeds over 2000 kilometers per second in some of the strongest events \cite{Lang2000}. The eruptions can take anywhere from a few minutes to hours to leave the solar surface \cite{Schwenn2006}.

Figure 1.6 is an example of coronagraph images of two different erupting CMEs taken from the High Altitude Observatory \cite{MacQueen1980}. The images are white light images taken on April 14th, 1980 and October 24, 1989. The leading edge of the CME, seen as a bright light bulb-shaped emission, is the sheath region of the shock wave formed by the erupting CMEs.
CMEs are typically assumed to result from large scale changes in the magnetic field topology of the sun. Solar flares were previously thought to cause many of the phenomena now known to be associated with CMEs [Bougeret, 1985; Gosling et al., 1968; Schwenn, 2006]. CMEs can cause large geomagnetic storms either by compressing the magnetosphere due to a high ram pressure or by having a large and continuously southward magnetic field component allowing for reconnection at the Earth’s magnetopause [Luhmann et al., 1998; Schwenn, 2006]. CMEs are also known to accelerate high energy particles towards Earth which can, in extreme cases, cause lethal radiation doses to spacecraft and astronauts [Schwenn, 2006].

Figure 1.7 is a cartoon which attempts to show a possible 2D magnetic field topology and various regions in and around a CME. The upstream region is the area farther away from the sun than the red bow-shaped line (marking the shock ramp) while the downstream is closer to the sun. The black(green) lines correspond to open sunward(antisunward) magnetic field lines while the blue lines correspond to closed field lines. As with the bow shock example, the quasi-parallel section of the shock is accompanied by low
frequency electromagnetic waves which are less easily excited at the quasi-perpendicular sections.

As a CME leaves the surface of the sun, it accelerates and compresses the plasma on the leading edge of a magnetic cloud (seen in Figure 1.7). If the CME accelerates at a sufficient rate and reaches a high enough speed, the wave forming on the bright leading edge (seen in Figure 1.6) steepens. Since compressional waves have the fundamental property that they steepen, the compressed plasma at the leading of the CME steepens unless sufficient energy is removed from the wave. When a CME reaches a speed faster than the local magnetosonic speed ($C_f$ or fast mode speed), a shock wave can form. The magnetic cloud of a CME, or ejecta, is often separated from the leading plasma. The

Figure 1.7: A cartoon of a CME with magnetic cloud propagating away from the sun. The black(green) lines correspond to open sunward(anti-sunward) magnetic field lines while the blue lines correspond to closed field lines. The cartoon illustrates the different regions of the shock.
shock front is on the leading edge of a plasma sheath which results from compression, deflection, and heating of the solar wind. The ejecta are separated from the sheath by a tangential discontinuity [Schwenn, 2006]. A tangential discontinuity is a special type of contact discontinuity where no mass or magnetic flux cross the boundary. An example of a tangential discontinuity is the Earth’s magnetopause in regions where magnetic reconnection is not occurring. CMEs are known to expand nearly as fast as they propagate radially outward [Siscoe and Schwenn, 2006; Siscoe et al., 2006]. Thus, by the time they reach 1 AU, they are often in excess of 1 AU in diameter. Since IP shocks are extremely large compared to the Earth’s bow shock, one can typically assume a planar geometry. This assumption simplifies the shock normal vector calculation rather significantly. CME-driven shock waves are also more likely to be quasi-perpendicular than quasi-parallel at 1 AU which has implications on dissipation mechanisms [Gary, 1981]. In this thesis, I will focus on observations of waves and particle distributions at IP shocks.

1.4 Particle Distribution Observations in Solar Wind and at Shocks

1.4.1 Distribution Function Moments

Particle distributions functions are useful because they are solutions to the equations of motion described by the Vlasov and/or Boltzmann equations, thus they allow for a statistical description of a fluid.

For instance, the heat flux is an important quantity because it describes the flow of thermal energy into or out of some volume. In the solar wind, electrons carry the flow of energy due to their lighter mass than that of the much heavier ions. The electron heat flux is due to a relative drift between the higher energy and lower energy electrons. Since the electron heat flux is responsible for the expansion of the solar corona [Feldman et al., 1975], it is important to study and understand how the heat flux is generated, maintained, and altered in the solar wind.

The heat flux, also called the kinetic energy flux in the solar wind reference frame, is derived from the third moment of the distribution function. Though in its general
mathematical form the heat flux is simply a third moment integral resulting in an asymmetric rank-3 tensor. Typically the heat flux is seen as a tensor in the second moment integral. However, one usually assumes certain symmetries to reduce the calculation to a rank-2 tensor, which can be treated as a matrix. In practical applications, such as calculating the heat flux for a real particle distribution function, the heat flux is derived from the third moment. In general, the third moment is known as the skewness, or the measure of the degree of asymmetry of a distribution function. Further assumptions can reduce the calculation to a rank-1 tensor, which can also be a simple vector. The vector form of the heat flux can be shown as:

\[ \mathbf{q} = \frac{m_s}{2} \int_V d^3v \mathbf{v}^2 f_s(x, \mathbf{v}, t). \] (1.19)

The heat flux vector, \( \mathbf{q} \), is the form most often used in plasma physics because it is more physically tangible and easier to work with.

In this thesis, I examine evidence of wave-particle interactions observed as scattering of the electron heat flux vector in Chapter 4.

### 1.4.2 Previous Observations and Examples of Electron Distributions

Electron velocity distributions have been measured in the tenuous solar wind for over 30 years [Feldman et al., 1975]. A number of very distinct properties were quickly distinguished, starting with the observation that the solar wind electrons are often composed of a cool dense Maxwellian core (typically \(< 60 \text{ eV}\)) and a more tenuous hot halo component consistent with a power-law [Feldman et al., 1975]. The energy at which the core component no longer dominated the distribution was defined as the break energy. Later a very prominent field-aligned signature was noticed, now called the strahl (German for bright or beam), in the high energy halo electrons [Pilipp et al., 1987a] and lastly a high energy super halo [Lin et al., 1996]. The differences in functional form of the two electron components is important because they can be used to infer their progenitor/source [Lin et al., 1996]. For instance, any distribution can be shown to relax to a Maxwellian if it is in a collision dominated medium [Gurnett and Bhattacharjee, 2005] and a power-law can illuminate the possible mechanisms which led to the energization of these particles [Lin et al., 1996].
Due to their much lighter mass than hydrogen and helium, electrons carry the bulk of thermal energy flow away from the sun. The electron heat flux, or more appropriately, the kinetic energy flux in the plasma rest frame, is the dominant component of energy flux in the solar wind. This heat flow is also thought to play a critical role in the expansion of the solar corona, thus studying electron moments in the solar wind have been motivated by many factors \cite{Feldman1975}.

The halo electrons were originally named due to their often isotropic shape in phase space, much like a halo, and nonthermal properties \cite{Feldman1975}. The halo electrons were also observed to drift relative to the cold dense thermal core of the distribution, producing a net kinetic energy flux. This portion of the distribution drifting relative to the cold core is the part responsible for the finite electron heat flux observed in the solar wind streaming away from the sun along the magnetic field \cite{Feldman1973}. It was later observed that the electron heat flux in the solar wind would occasionally "disappear." The effect is referred to as a heat flux dropout (HFD), which is currently assumed to result from a magnetic disconnection from the sun \cite{Crooker2003, Pagel2005, deKoning2006, deKoning2007}. Examining the width of the suprathermal electron pitch-angle distributions (PADs) during solar electron burst events with the ACE spacecraft. Simulations and theory have been pursued to explain the discrepancies between the expected heat flux values and the observed values \cite{Vocks2003, Vocks2005}. A number of instability mechanisms driven by electron heat flux in the solar wind have been theoretically examined in detail by \cite{Gary1975, Gary1994, Gary1999}.

Many different distribution types observed in the terrestrial magnetosphere have been observed near collisionless shocks or in the solar wind. For instance, \cite{Larson1996} examined loss-cone like distributions at the bow shock with the Wind spacecraft, \cite{Gosling1989} examined suprathermal electrons in the bow shock and magnetosheath with ISEE 1 and 2, and temperature anisotropies were examined in the solar wind between 0.3 and 0.8 AU with Helios 1 and 2 \cite{Pilipp1987} and at IP shocks \cite{Wilson2009} with the Wind spacecraft.

Figure 1.8 shows an example of an electron distribution seen in the solar wind observed by the Wind spacecraft. This figure is representative of many of the particle distribution function figures shown later in this thesis. The top panel is a two dimensional
Figure 1.8: An example of an electron particle distribution. The top panel, the contour plot, is 2D projection of the particle distribution function where the horizontal axis is the velocity parallel(anti-parallel) to the magnetic field and the vertical axis is perpendicular to the magnetic field. Four vectors are projected onto the contour plot: heat flux ($Q_f$), solar wind velocity ($V_{sw}$), shock normal vector ($n_{sh}$), and the sun direction. The bottom panel is a plot of the parallel and perpendicular cuts of the distribution function.
contour projection of the three dimensional distribution function where the horizontal axis is the parallel/anti-parallel (with respect to the magnetic field) velocity direction and the vertical axis is perpendicular velocity. The axes range from -20,000 km/s to +20,000 km/s. The distribution function was constructed assuming gyrotropy, or azimuthal symmetry with respect to the magnetic field. Although this is not a necessary process, it often makes analysis easier. On the top panel, there are four straight lines corresponding to the direction of the electron heat flux (red line), solar wind velocity (black line), shock normal vector (blue line), and sun direction (purple line). Each line is projected into the same coordinate system as the distribution contours. The bottom panel shows the parallel (black line) and perpendicular (blue line) cuts of the distribution function shown in the contour plot. The vertical axis corresponds to the magnitude of the phase space density ($s^3 km^{-3} cm^{-3}$), or units of the distribution function, while the horizontal axis is velocity (with the same limits as the contour plot).

The transfer of energy from a nonthermal source to heat is a topic of considerable interest in collisionless plasmas because energy transfers to heat tend to occur through binary collisions in a regular fluid or free-free Bremsstrahlung interactions in the solar atmosphere [Aschwanden, 2005]. The lack of sufficient binary collisions in a plasma requires other mechanisms by which to transfer energy [Gary et al., 1975]. Electron heating has been examined at the bow shock using ISEE 2 data [Feldman et al., 1982], compared to ion heating at the bow shock using ISEE 1 and 2 data [Thomsen et al., 1985a], compared to the de Hoffmann-Teller frame potential across IP and bow shocks with ISEE spacecraft [Schwartz et al., 1988], used as a proxy for the de Hoffmann-Teller frame potential across collisionless shocks [Hull et al., 2000; Hull and Scudder, 2000], and in association with atypical wave modes at IP shocks [Wilson III et al., 2009].

Electrons have been theorized to drive and/or damp electromagnetic fluctuations in many regions of space. Early predictions of electron-ion energy/momentum exchanges, thermal conductivities, and heat flux originally suggested electron(ion) temperatures which were higher(lower) than the observed values in the solar wind [Gary et al., 1975; Spitzer and Härm, 1953]. Thermal conductivities, produced by binary particle collisions, were originally predicted to produce an electron heat flux proportional to a constant multiplied by the electron temperature gradient [Gary et al., 1975]. The proposed solutions included a variety of plausible explanations, but the most likely is heat flux
generated instabilities affecting the solar wind plasma [Gary et al., 1975, 1994, 1999; Gurnett et al., 1979a; Saito and Gary, 2007; Wilson III et al., 2007, 2009].

Electron beams are known to be a source of free energy for instabilities and have been examined with ISEE 1 data [Fitzenreiter et al., 1984], in the ionosphere with sounding rockets [Kellogg et al., 1986], in the foreshock with the Wind spacecraft [Fitzenreiter et al., 1996], in association with Type III radio bursts and impulsive electron events using Wind [Ergun et al., 1998a], and to map the structure of the source region for Type II radio bursts in IP shocks observed by Wind [Pulupa and Bale, 2008].

In this thesis, I utilize electron distribution measurements to examine particle heating signatures at IP shocks and to look for evidence of wave-particle interactions.

1.4.3 Previous Observations and Examples of Ion Distributions

The irregular turbulence seen in magnetic field measurements upstream of planetary bow shocks for over 40 years [Fairfield, 1969] often had frequencies in the range of the ion cyclotron frequency. Thus, the examination of ion particle distributions in association with the waves seemed like the natural thing to do. Three ion populations, reflected, intermediate, and diffuse, are commonly found in planetary foreshocks. The distinction between different distributions became important when it was discovered there was a strong correlation between the distribution type and location within the foreshock [Hoppe et al., 1981; Paschmann et al., 1979, 1981]. Later studies found there to be what appeared to be boundaries between different regions of the foreshock that strongly depended on the angle between the interplanetary magnetic field and the X-GSE direction [Meziane and D’Uston, 1998]. Since, the first observations of ions upstream of collisionless shocks, they have been a topic of intense interest due to their capacity to carry energy away from the shock and interact with waves.

Reflected Ions

Reflected ions have a beam-like feature in their distribution with bulk speeds on the order of 1-5 times the solar wind speed. They occur predominantly near shocks with shock normal angles, \( \theta_{bn} \), between 30\(^\circ\)-75\(^\circ\) (see Figure 1.9) [Bonifazi and Moreno, 1981, Paschmann et al., 1981; Schwartz et al., 1983]. They have broad thermal spreads and are often anisotropic with \( T_{\perp b}/T_{\parallel b} > 1 \) [Paschmann et al., 1981]. The beams typically
have \( n_b/n_o \sim 0.3-13\% \), where \( n_b \) is the number density of the beam and \( n_o \) is the density of the solar wind \cite{Bonifazi1981, Paschmann1981}. Reflected ions carry most of their energy in the form of kinetic energy and not thermal energy \cite{Bonifazi1981}.

Figure 1.9 is an example of a reflected ion distribution function observed by the Wind spacecraft. The top panel plots contours of constant phase space density versus particle velocity with the horizontal axis being the velocity parallel to the background magnetic field and the vertical axis perpendicular to magnetic field in the plane defined by the magnetic field and solar wind velocity (black line projected onto contour). The bottom panel plots the parallel and perpendicular cuts of the distribution function. The wind ion measurements shown range from -2500 km/s to +2500 km/s. The color contours range from \( 10^{-12} \) to \( 10^{-10} \) s\(^{-3}\) km\(^{-3}\) cm\(^{-3}\).

Reflected ions streaming away from the shock along the magnetic field can interact with the incident solar wind through ion-ion instabilities producing large amplitude magnetic fluctuations both parallel and oblique to the magnetic field \cite{Akimoto1993, Gary1981} and low frequency oblique whistler waves \cite{Akimoto1989}. These hydromagnetic waves can pitch-angle scatter the reflected ion beams as they convect back towards the bow shock because their group and phase velocities are smaller than the solar wind bulk flow (typically). They are usually transverse at this stage with very small compressive components. As a consequence of the scattering, the ion beams broaden (in pitch-angle) into crescent-shaped ion distributions called intermediate ions, which are usually seen deeper (i.e. further from the sun) into the foreshock region. Eventually, it is thought, these wave-particle interactions scatter the ions to the point of suprathermal spherical shells in phase space called diffuse ion distributions \cite{Gary1981}.

### Intermediate and Gyrating Ions

Intermediate ion distributions represent a transition between reflected and diffuse ion distributions (see panel A in Figures 1.11 and 1.10) and appear as a crescent-shaped distribution with centers of curvature near the solar wind velocity. They have been observed in association with ULF foreshock waves and show a distinct boundary with
Figure 1.9: Here is an example of an ion beam seen upstream of the Earth’s bow shock by the Wind Pesa High instrument immediately after (i.e. closer to the sun) the intermediate distribution in Figure 1.10. The contour ranges are the same for both of these plots.
gyrating and diffuse ions [Meziane and D’Uston, 1998; Meziane et al., 2001]. Intermediate ions often show total energy densities that are nearly the same as the reflected ions observed adjacent to them (spatially) [Bonifazi and Moreno, 1981; Gary, 1985].

Figure 1.10 is an example of an intermediate ion distribution function. The format is similar to Figure 1.9. The intermediate ions can be seen as the crescent-shaped peak in the contour plot outlined by the red oval and in the parallel cut of the distribution by the red box. The main difference from a reflected ion distribution is that the intermediate ions are much broader in the direction perpendicular to the magnetic field in phase space. They are thought to result from the wave-particle scattering of reflected ion beams [Mazelle et al., 2003].

Gyrating ion distributions show signatures of gyromotion about the magnetic field and can be gyrophase-bunched (nongyrotropic) or nearly gyrotropic [Meziane et al., 2001]. The gyrotropic distributions are typically observed closer to the shock than the nongyrotropic ions. Gyrating ions are commonly observed in association with the magnetic foot and overshoot of quasi-perpendicular supercritical shocks due to specular reflection [Paschmann et al., 1980, 1982; Sckopke et al., 1983; Thomsen et al., 1985b]. Such distributions are thought to play an integral part in the processes which heat the ion distributions downstream of the shock [Sckopke et al., 1983]. The gyrophase-bunched ions can be created/formed by reflection at a collisionless shock or the disruption of an ion beam by waves generated from a beam-plasma instability [Akimoto et al., 1991; Meziane et al., 2001; Thomsen et al., 1985b]. The finite thermal velocity of the gyrophase-bunched ions should lead to gyrophase mixing of the distributions, which would result in a more gyrotropic distribution [Burgess and Schwartz, 1984]. Thus the observations of nongyrotropic distributions far upstream of the bow shock suggest local production mechanisms [Mazelle et al., 2003; Thomsen et al., 1985b].

Diffuse Ions

Diffuse ion distributions are a highly nonthermal, relatively isotropic spherical shell-like distribution in phase space extending up to ∼40 keV (see panel B in Figure 1.11). Diffuse ions often show anisotropies with pitch-angle distributions peaking at 90° [Gosling et al., 1984; Paschmann et al., 1981; Sentman et al., 1981] and have average thermal spreads ∼5-7 times greater than reflected ions [Bonifazi and Moreno, 1981].
Figure 1.10: Here is an example of intermediate ions seen upstream of the Earth’s bow shock by the Wind Pesa High instrument. The color contours range from $10^{-12}$ to $10^{-10}$ s$^3$ km$^{-3}$ cm$^{-3}$ and the UV contamination effects have been reduced.
They are likely the result of wave-particle interactions that pitch-angle scatter and heat an ion beam into a nearly isotropic shell, thus a shock generated source \cite{Gosling1984, Gosling1989, Winske1984}.

Figure 1.11 is two examples of an diffuse ion distributions upstream of an IP shock. The contour plots have the same format as those in Figures 1.9 and 1.10. The diffuse ions are seen as a nonthermal halo or shell-like distribution between 800-2500 km/s in both parallel and perpendicular directions. The theories behind their production and their importance will be discussed below.

Many theories, in conjunction with observations, suggest a spatial correlation with observations of diffuse ions and their production results from the interaction of low frequency hydromagnetic waves interacting with the foreshock ions \cite{Bame1980, Gary1981, Gosling1984, Hoppe1981, Mazelle2003}. Diffuse ions can dramatically alter the index of refraction of the local medium \cite{Hada1987, Omidi1990} because they can carry a significant fraction of the total energy density of the ions \cite{Bonifazi1981}. As a consequence of their influence on the medium, a class of highly steepened magnetic structures referred to as shocklets and/or SLAMS are always seen in association with diffuse ions (see Chapter 3 for further discussion) \cite{Hoppe1981, Hoppe1982, Hoppe1983, Thomsen1985, Thomsen1990, Wilson2009}.

In this thesis, I utilize ion distribution measurements to determine solar wind density and velocity, and examine evidence for ion reflection and acceleration.

1.5 Wave Particle Interactions

1.5.1 The Dispersion Relation

In order to discuss the interactions of waves and particles, one must begin by discussing a property called dispersion. Dispersion is a property of a medium where the frequency, $\omega$, depends on the wave number, $k$. A simple example of a dispersive medium would be a prism. If one sends multi-wavelength (e.g. white light) into the prism, one notices a rainbow when the light exits the opposite side. The incident light is undergoing a process called refraction where the light in the first medium with a specific wave number, $k_1$, enters a second medium that alters its wave number to $k_2$. The manner
Figure 1.11: An example of intermediate and diffuse ion distributions seen upstream of a IP shock on 04/06/2000. Panel A shows an example of an intermediate ion distribution mixed with a diffuse ion distribution while panel B is purely diffuse ions.
in which a wave propagates through any particular medium can be characterized by a property called the index of refraction, \( n^2 = k^2 c^2 / \omega^2 \), where \( c \) is the speed of light in vacuum. Thus one can see, \( n^2 = 1 \) for vacuum.

The simplest dispersion relation in a plasma is for a homogeneous, source free plasma with only small linear fluctuations/perturbations in any given quantity defined by their Fourier transform in the following manner:

\[
\delta A(k, \omega) = \sum_k A_k e^{i(k \cdot r - \omega t)} \tag{1.20}
\]

where \( \delta A(k, \omega) \) can represent any of the following relevant quantities: particle density (\( n_s \)), velocity (\( \mathbf{V}_s \)), electric or magnetic fields (\( \mathbf{E} \) or \( \mathbf{B} \)), etc. where the subscript \( s \) represents the particle species. If one assumes that the zeroth order functions have neither temporal nor spatial variations, then the operators \( \partial / \partial t \) and \( \nabla \) go to \(-i \omega\) and \(+i k\) respectively. This allows one to reduce Faraday’s and Ampere’s laws, to first order, to the following two equations:

\[
-i k \times \delta \mathbf{E} = (-i \omega) \delta \mathbf{B} \tag{1.21a}
\]
\[
i k \times \delta \mathbf{B} = -i \omega c^2 \tilde{\mathbf{K}} \cdot \delta \mathbf{E} \tag{1.21b}
\]

where \( \tilde{\mathbf{K}} \) is the dielectric tensor. Eliminate \( \delta \mathbf{B} \) between the two equations and remember that \( \mathbf{n} = c k / \omega \) to find:

\[
\mathbf{n} \times (\mathbf{n} \times \delta \mathbf{E}) + \tilde{\mathbf{K}} \cdot \delta \mathbf{E} = 0 \tag{1.22}
\]

which can be rewritten in a short form as \( \tilde{\mathbf{D}} \cdot \delta \mathbf{E} = 0 \), where setting the determinant of the tensor \( \tilde{\mathbf{D}} \) equal to zero yields a solution for \( \delta \mathbf{E} \). One should note that the determinant of \( \tilde{\mathbf{D}} \) is not, by definition, zero. When it is equal to zero, the solution is referred to as the dispersion relation, \( D(\omega, \mathbf{k}) \). Thus, the eigenmodes of the tensor, \( \tilde{\mathbf{D}} \), are the modes of the allowed waves in the given system. Or in other words, the solution(s) to the dispersion equation relate \( \omega \) to \( \mathbf{k} \) which describe the normal modes of the system [Gary, 1993].
1.5.2 Instability

In a dispersive plasma, waves can grow or damp depending on driving mechanisms and interactions with the medium itself. Plasmas are incredibly unstable and they are rarely, if ever, in thermodynamic equilibrium. Departure from thermodynamic equilibrium, \textit{i.e.} non-Maxwellian features in a velocity distribution, can be a source of free energy which needs to be dissipated \cite{Gary1993}. In order to dissipate the energy provided by a non-Maxwellian distribution, a plasma must transform the free energy into some other form which will allow the non-Maxwellian features to relax to a Maxwellian. This is done through normal modes of a system called instabilities.

An instability is a normal mode that can grow in space and time with well defined relationships between $\omega$ and $k$ \cite{Gary1993}. Each wave mode resulting from a source of free energy has its own dispersive properties much like any given medium has its own dispersive properties. If an instability grows in time or space it can couple to stable wave modes of the system \textit{i.e.} the instability uses radiation to dissipate the free energy. Thus, an instability is effectively the transfer mechanism by which the free energy can be dissipated. Given that instabilities depend on both $\omega$ and $k$, one can assume that there are going to be different ranges or scales for any given system.

For small wavelength(large $k$), we refer to these instabilities as \textit{microinstabilities} and for large wavelength(small $k$), we refer to these instabilities as \textit{macroinstabilities}. More quantitatively, a microinstability is any instability that satisfies $k \rho_i \gtrless 1$, where $\rho_i$ is the ion gyroradius. In this thesis, we will focus on microinstabilities but also look at some macroinstabilities.

In order for wave damping/growth to occur, the wave modes must have only an imaginary part to their frequency, since purely growing modes \textit{i.e.} real part of frequency is zero) can exist. Assuming a non-zero real part to the frequency, let us define the frequency of the wave as:

$$\omega(k) = \omega_r(k) + i\gamma(\omega_r, k)$$

(1.23)

where the real ($\omega_r$) and imaginary ($\gamma$) parts of the frequency are generally functions of $k$. However, it is often safe to assume that $\gamma = \gamma(k)$ or even just a constant. By this definition, if $\gamma > (-) 0$ the wave amplitude grows(decays) exponentially in time. One should note, however, that a solution to the dispersion relation of $\gamma > 0$ is not, in
itself, sufficient to produce an instability. Instabilities require a source of free energy, otherwise the growth implied in the result is an unphysical solution.

The intrinsically unstable nature of a plasma allows for a number of different instabilities, depending on the free energy source and the local plasma parameters of the medium. For instance, ion-acoustic waves can be excited by the following sources of free energy: cross-field currents [Lemons and Gary, 1978], ion velocity ring distributions [Akimoto et al., 1985; Akimoto and Winske, 1985], electron heat flux [Dum et al., 1980; Gurnett et al., 1979b], cold ion beam streaming through a warm electron background across the magnetic field [Moses et al., 1985, 1988], etc. Often times the instabilities are named after the source of free energy or the wave mode they excite which leads to a confusing misinterpretation of the physical differences between the free energy source, the instability, and the wave produced by the instability.

1.5.3 Landau Interactions

The collisionless nature of plasmas initially raised questions regarding the processes of energy dissipation. In 1946, Lev Davidovich Landau proposed a theory of collisionless dissipation where particles with velocities near the phase velocity of a plasma wave, $\omega/k$, interact resonantly with the wave fields [Landau, 1946]. The theory was slightly altered and corrected a few years later by [Jackson, 1960]. It is now known as Landau damping and was first shown to be real in experiments by [Malmberg and Wharton, 1964, 1966].

The physical interpretation of Landau damping was first explained by [Bohm and Gross, 1949a,b]. It can be thought of as an analog to elastic particle collisions, where one of the particles is the wave. One can think of particles moving in a frame of reference where the wave is at rest and consider the wave as a series of potential hills and troughs. In this frame, the particles with speeds greater(less) than the phase speed of the wave can transfer energy to(from) the wave, thus causing growth(damping) of the wave and the particles lose(gain) energy on average. The analogy often used is to consider a surfer riding water waves. If the surfer can speed up enough to just slightly less than the phase speed of the wave, they will be able to ride the wave, gain kinetic energy and thus surf. If they do not move, the wave will pass by them and they will rise and fall without any real energy transfer occurring between wave and surfer. If the surfer goes slightly too fast, they will transfer energy to the wave and never pass over
the wave only reaching just below the peak.

Landau interactions can also lead to phase trapping and mixing. In the wave frame, we consider a periodic wave to be a series of electric potential hills and troughs. In this frame, the particles with speeds greater than the phase speed of the wave can escape the troughs of the wave while those with speeds slightly less than the phase speed cannot. The particles with higher speeds than the wave will suffer small periodic fluctuations in their velocities, but nothing more. The particles which cannot escape the troughs will oscillate back and forth in the trough and are trapped. Note that the oscillation frequency, also called the bounce frequency, \( \omega_b \), depends on how nearly matched their velocity is to the resonant velocity. The difference in oscillation frequencies of the different particles leads to a phenomena called phase mixing. Phase mixing can also give rise to the dissipation implicitly implied in Landau damping.

It is important to note that the above discussion was only assuming linear Landau damping. During the linear phase, one can show that the total nonrelativistic energy of the particles is given by:

\[
W_{KE} = \frac{1}{2} m V_{\parallel_res}^2 + q \Phi \tag{1.24}
\]

where \( V_{\parallel_res} \) is the resonant velocity of the particles, \( q \) their charge, \( m \) their mass, and \( \Phi \) the electrostatic potential of the wave in the wave frame \[Gurnett and Bhattacharjee, 2005\]. During the linear phase, the total energy of the system is conserved. Therefore, if the kinetic energy of the particles increases, the potential energy of the wave must decrease. To determine whether the particle kinetic energy increases or decreases, one must consider the slope of zeroth order velocity distribution function, \( F_o(v) \), at the resonant velocity. If the slope is given by:

\[
\frac{\partial F_o(v)}{\partial v} \bigg|_{V_{res}} < 0 \tag{1.25}
\]

where \( V_{res} \) is the parallel resonant velocity, then the number of particles initially moving slower than \( V_{\parallel res} \) is higher than the number initially moving faster. Then the wave transfers energy to the particles in a similar fashion to how a surfer gains kinetic energy by riding a wave. The wave is initially faster than the particles and thus effectively runs into them, transferring momentum. The end result is particles gain kinetic energy and the wave loses amplitude due to a loss of its potential energy, thus the name Landau damping.
After some time, the interaction described above must change. The reason is that the bounce period of the trapped particles is often much greater than the period of the wave. Thus in the linear phase, the damping of the wave happens faster than the period of the trapped particles. During this time, one should also note that some energy was given to particles which weren’t trapped. This is important to mention because eventually the trapped particles begin to give energy back to the wave. Since some energy was given to untrapped particles, not all the energy can be transferred back to the wave. The trapped particles, by this time, have also begun to phase mix which reduces the effective transfer of energy back to the wave. Though the particles will not cause the wave amplitude to grow in this phase, they do slow the damping predicted by linear Landau damping. Eventually the damping ceases after many bounce periods and after the trapped particles become thoroughly phase mixed. Note, unless some dissipative/irreversible process occurs, the phase mixed particles are not truly randomly mixed. Their phase information can still be recovered.

In this thesis, I examine evidence for Landau damping of low frequency waves in Chapter 4.

1.5.4 Weak Growth Limit

Since nonlinear treatments of plasmas are often impossible to deal with analytically, linear approximations are considered. A specific example for linear waves is the weak growth approximation which assumes $|\gamma| \ll |\omega_r|$. The following derivations can be found in more detail in Gurnett and Bhattacharjee [2005].

One can then expand the dispersion relation in a Taylor series to obtain:

$$D(\omega, k) = D_r(\omega, k) + i\gamma \frac{\partial D(\omega, k)}{\partial \omega_r} = 0$$

(1.26)

where each part, $D(\omega, k)$ and $\partial D(\omega, k)/\partial \omega_r$, contains real and imaginary parts. Each term must be evaluated in the limit $\gamma \to 0$ to find solutions to the dispersion relation in the weak growth approximation. [Gurnett and Bhattacharjee, 2005]. Note that for the second term in this limit, $\gamma \partial D(\omega, k)/\partial \omega_r$, is negligible. Thus, only the real part is kept. The first term simplifies to:

$$D(\omega, k) = D_r(\omega, k) + iD_i(\omega, k)$$

(1.27)
where the subscripts, \( r \) and \( i \), refer to the real and imaginary parts respectively. Since we are interested in solutions of \( D(\omega, k) = 0 \), we solve the real and imaginary parts separately to give:

\[
D_r = 0 \\
\gamma \frac{\partial D_r(\omega_r, k)}{\partial \omega_r} + D_i = 0
\]

(1.28a)

(1.28b)

In the process of evaluating Equations 1.28a and 1.28b individually, one finds that \( D(\omega, k) \) has an integral form of:

\[
D(\omega, k) = 1 + \lim_{\gamma \to 0} \frac{\omega_{pe}^2}{k^2} \int_C \frac{\partial F_o}{\partial V} \frac{1}{||V|| - \omega_r/k - i\gamma/k} dV ||
\]

(1.29)

which requires the use of the Plemelj relation given by:

\[
\lim_{\epsilon \to 0} \int_{-\infty}^{\infty} \frac{f(x)}{x - (x_o \pm i\epsilon)} dx = P \int_{-\infty}^{\infty} \frac{f(x)}{x - x_o} dx \pm i\pi f(x_o)
\]

(1.30)

where \( \epsilon > 0 \) and the P refers to the principal value integral defined by:

\[
P \int_{-\infty}^{\infty} \ldots dx = \lim_{\delta \to 0} [\int_{-\infty}^{x_o - \delta} \ldots dx + \int_{x_o + \delta}^{\infty} \ldots dx]
\]

(1.31)

After applying these mathematical techniques, one finds the relationships for the real and imaginary parts of the dispersion relation reduce to:

\[
D_r = 1 - \frac{\omega_{pe}^2}{k^2} P \int_{-\infty}^{\infty} \frac{\partial F_o}{\partial V} \frac{1}{V || - \omega_r/k}
\]

(1.32a)

\[
D_i = -\pi \frac{k}{|k|} \frac{\omega_{pe}^2}{k^2} \frac{\partial F_o}{\partial V} ||_{V || = \omega/k}
\]

(1.32b)

The growth/damping rate, \( \gamma \), can be found from Equations 1.28b and 1.32b given by:

\[
\gamma \propto \left( \frac{\partial D_r}{\partial \omega_r} \right)^{-1} \frac{\partial F_o}{\partial V} ||_{V || = \omega/k}
\]

(1.33)

which is the classic result of Landau damping where \( \gamma \) depends on the slope of the distribution function at the phase velocity of the wave.
1.5.5 Cyclotron Resonance

Cyclotron interactions between waves and particles are similar to Landau interactions in some respects, but differ in one very important quality. The damping rate of cyclotron interactions, in the weak damping limit for an isotropic distribution, is directly proportional to the distribution function evaluated at a parallel resonance velocity, \(V_{\parallel \text{res}}\). Recall in the weak damping limit for Landau interactions, the damping rate was proportional to the slope at the resonance velocity. Thus, a greater number of particles can contribute to the damping in cyclotron interactions than in Landau damping because all the particles near \(V_{\parallel \text{res}}\) interact with the wave not just the particles with velocities just below \(V_{\parallel \text{res}}\). However, when the distributions become anisotropic, cyclotron damping suddenly depends on the slope of the distribution. In the weak growth limit, the growth rate goes as:

\[
\gamma = \frac{\pi}{\partial D_s/\partial\omega} \sum_s \left(\frac{\omega_{ps}}{\omega}\right)^2 \left[\frac{\omega}{|k_\parallel|} \int_0^\infty dv_\perp \ 2\pi v_\perp F_{so} \right] - \frac{\pi}{\partial D_s/\partial\omega} \sum_s \left(\frac{\omega_{ps}}{\omega}\right)^2 \left[\frac{k_\parallel}{|k_\parallel|} \int_0^\infty dv_\perp \pi v_\perp \left(\frac{\partial F_{so}}{\partial v_\perp} - \frac{\partial F_{so}}{\partial v_\parallel}\right)\right] |V_{\parallel \text{res}}| \tag{1.34}
\]

where \(F_{so} = \frac{f_{so}}{n_{so}}\) is the normalized zero-order distribution function of species \(s\) and \(\omega_{ps}\) is the plasma frequency of species \(s\). For an isotropic distribution, the second integrand in Equation 1.34 goes to zero because \(F_{so} = F_{so}(v)\), thus one can use \(v = (v_\parallel^2 + v_\perp^2)^{1/2}\) and the commutation property of partial derivatives to show:

\[
\frac{\partial F_{so}}{\partial v_\perp} = \frac{\partial v}{\partial v_\parallel} \frac{\partial F_{so}}{\partial v_\perp} = \frac{\partial v}{\partial v_\parallel} \frac{\partial F_{so}}{\partial v} = \frac{v}{v} \frac{\partial F_{so}}{\partial v_\perp} \tag{1.35a}
\]

If the distribution is anisotropic, then \(F_{so} \neq F_{so}(v)\) and the second integrand in Equation 1.34 does not go to zero (see Gurnett and Bhattacharjee [2003] for more detail). It was shown in detail by Kennel and Petscheck [1966] that if the anisotropy of the particles is great enough, the distribution can cause a cyclotron resonance that leads to growth.
The criteria for growth is given by:

$$A_s > \frac{1}{\Omega_{ce}/\omega - 1} \quad (1.36)$$

where $A_s$ is the anisotropy given by:

$$A_s = \frac{\partial F_{so}}{\partial v_\perp} + \frac{k_\parallel}{\omega} \left( v_\perp \frac{\partial F_{so}}{\partial v_\parallel} - v_\parallel \frac{\partial F_{so}}{\partial v_\perp} \right) \quad (1.37)$$

which reduces to $(T_{\perp,s}/T_{\parallel,s} - 1)$ for a bi-Maxwellian distribution. Anisotropic electron distributions are an important source of free energy for whistler waves [Gary et al. 1994, 1999].

In the solar wind, the plasma almost always satisfies the condition $\omega_{pe} \gg \Omega_{ce}$. In the high density limit and the cold plasma approximation, which assumes the particles are initially at rest with no thermal motions (i.e. $T_e = T_i = 0$), one can estimate the index of refraction for a RH-mode whistler wave. If the wave has frequency $\omega^2 \ll \omega^2_{pe}$ and propagates at an oblique angle to the magnetic field, $\theta_{kB}$, then the index of refraction for a RH-mode whistler wave is:

$$n^2 = \frac{\omega_{pe}^2}{\omega(\Omega_{ce}\cos\theta_{kB} - \omega)} \quad (1.38)$$

The angle at which $n^2$ goes to infinity, called resonance, is known as the resonance cone angle, $\theta_{kB} \rightarrow \theta_{res} \equiv \cos^{-1}(\omega/\Omega_{ce})$. An important aspect of a whistler mode wave is that their frequencies always satisfy $\omega < \Omega_{ce}$. Replacing $\omega_{pe}^2$ with $\zeta^2\Omega_{ce}^2$, where $\zeta^2 = (c^2\mu_0 m_e n_e)/B_o^2$ and after some manipulation, the index of refraction becomes:

$$n^2 = \frac{\zeta^2\Omega_{ce}}{\omega(\cos\theta_{kB} - \omega/\Omega_{ce})} \quad (1.39)$$

For particles to undergo resonance with a wave, their velocities parallel to the background magnetic field, $V_{||res}$, must satisfy the following condition:

$$V_{||res} = \frac{\omega + m\Omega_{ce}}{k_\parallel} \quad (1.40)$$

where the parallel wave number, $k_\parallel$, is given by $|k| \cos\theta_{kB}$ and $m = 0, -1, or +1$ (though higher integer values can exist, we only consider these for now). The different values of $m$ represent different resonances or different types of interactions. The first possible
value, 0, is for Landau resonance which is discussed in Section 1.5.3. The following two are two different types of cyclotron resonance called normal and anomalous, respectively.

Normal cyclotron resonance occurs between an electron and a RH-mode wave. In the case of the whistler mode, the electron cyclotron frequency always exceeds the whistler frequency. For normal cyclotron resonance to occur, the electron’s gyromotion must match the rotation frequency of the fields of the whistler mode. Thus, the electron’s guiding center velocity along the zero order magnetic field must Doppler shift the whistler’s frequency up to the electron cyclotron frequency (i.e. $\omega' = \omega - k_\parallel V_\parallel = \Omega_{ce}$). This implies that $k_\parallel V_\parallel < 0$ and whistler instabilities must be driven by positive anisotropies in the electron velocity distributions. In other words, the average kinetic energy perpendicular to the magnetic field should be larger than the average kinetic energy parallel to the magnetic field.

Anomalous cyclotron resonance applies a similar principle to the normal resonance, but with ions replacing electrons. Now the wave polarization in the particle’s guiding center rest frame must appear to be a LH-mode and the Doppler shifted frequency must match the ion cyclotron frequency. Thus, $k_\parallel V_\parallel > 0$ and $V_\parallel > \omega/k_\parallel$. One should note, however, that such resonance conditions are rarely met since $V_\parallel$ is often very high.

To estimate the resonant energies of particles with the whistler mode in the cold plasma approximation, we must take our arguments from above a little further. Recall that $n^2 = k^2 c^2/\omega^2$, which allows us to solve for $k^2$. After some algebra, we find:

\[
k^2 = \left( \frac{\mu_0 m_e n_e}{B_0^2} \right) \left( \frac{\omega \Omega_{ce}}{(\cos \theta_{kB} - \omega/\Omega_{ce})} \right). \tag{1.41}
\]

Now we replace $k_\parallel$ in Equation 1.40 by $|k| \cos \theta_{kB}$, square $V_{\parallel res}$, and replace $|k|^2$ with the result from Equation 1.41 to get:

\[
V_{\parallel res}^2 = \frac{(\omega + m \omega_{ce})^2}{k^2 \cos^2 \theta_{kB}} \tag{1.42a}
\]

\[
= \left( \frac{B_o^2}{\mu_0 m_e n_e \omega \cos^2 \theta_{kB}} \right) \left( \cos \theta_{kB} - \frac{\omega}{\Omega_{ce}} \right) \left[ m + \frac{\omega}{\Omega_{ce}} \right]^2. \tag{1.42b}
\]

In the nonrelativistic limit, the kinetic energy of a resonant particle is just the typical kinetic energy or $1/2 \ m_e \ V_{\parallel res}^2$, which gives:

\[
E_{\parallel res} = \left( \frac{B_o^2}{2 \mu_0 n_e} \right) \left( \frac{\Omega_{ce}}{\omega \cos^2 \theta_{kB}} \right) \left( \cos \theta_{kB} - \frac{\omega}{\Omega_{ce}} \right) \left[ m + \frac{\omega}{\Omega_{ce}} \right]^2. \tag{1.43}
\]
Figure 1.12: An example calculation of the parallel cyclotron resonance energy for normal \((m = -1)\) cyclotron resonance. The plot shows \(E_{||\text{res}}\) plotted versus \(\Omega_{ce}/\omega\) and \(\theta_{kB}\). The energy ranges from 10 eV to 10 keV and the values of \(\Omega_{ce}\) and \(V_{Ae}/c\) are shown at the top of the plot.

where the first factor, \((B_o^2/2\mu_0n_e)\) is \(1/2\) \(m_e\) \(V_{Ae}^2\) where \(V_{Ae}\) is the electron Alfvén speed. Though we should note that there is no electron Alfvén wave that propagates at this speed.

An example calculation of Equation \ref{eq:1.43} is shown in Figure \ref{fig:1.12}. \(E_{||\text{res}}\) is plotted versus \(\omega/\Omega_{ce}\) and \(\theta_{kB}\). The frequency dependence is weaker than the angle dependence at very oblique angles \((\gtrsim 60^\circ)\), but becomes dominant at lower angles. As one can see, at lower wave frequencies (high values of \(\omega/\Omega_{ce}\)) higher energy electrons are required to resonate with whistler waves through a normal cyclotron \((m = -1)\) resonance. The reason is that the electrons must move faster to Doppler shift the low frequencies up to the electron cyclotron frequency in the guiding center reference frame of the electrons.
Applications of Equation 1.43 have been used in research by Lengyel-Frey et al. [1994] and Wilson III et al. [2009] to estimate the energy range of electrons which would be resonant with the whistler waves they observed.

In this thesis, I examine evidence for cyclotron interactions with low frequency waves in Chapter 4 and higher frequency waves in Chapter 5.

1.6 Introduction to the Microphysics of IP shocks

Anomalous resistivity, through wave-particle interactions, has been thought to act as an energy dissipation mechanism in collisionless shock waves [Kellogg, 1965; Sagdeev, 1966]. It is thought to provide much of the dissipation at low Mach number shock waves [Gary, 1981], along with dispersive [Mellott and Greenstadt, 1984] and conductive [Kennel et al., 1985] effects. However, at higher Mach numbers other mechanisms become important, such as particle reflection [Edmiston and Kennel, 1984; Kennel, 1987]. Due to their generally lower Mach numbers at 1 AU and their larger radii of curvature, IP shocks provide an excellent opportunity to study the role of wave dissipation at low Mach number shocks.

The importance of wave-particle interactions in the total energy dissipation budget of collisionless shocks is not well known. To zeroth order and ignoring wave-particle effects, electrons conserve their first adiabatic invariants which states that the magnetic moment, $\mu$, is given by $\mu = B_o / (0.5 m_e V_{e,⊥}^2) = \text{constant}$. If the electrons conserve their magnetic moments, then one expects $V_{e,⊥}^2$ to increase (decrease) as $B_o$ increases (decreases). Wave-particle interactions can "de-magnetize" the electrons, or cause $\mu \neq \text{constant}$, which has a very important implication, namely irreversibility due to an increase in entropy. Recent observations of IP shocks [Wilson III et al., 2007, 2009, 2010] and simulation studies with a realistic mass ratio ($M_i / m_e \sim 1836$) [Petkaki et al., 2006; Petkaki and Freeman, 2008; Yoon and Lui, 2006, 2007] have found evidence to suggest that wave-particle interactions may be more important than previously thought. However, some previous studies of collisionless shocks have suggested that the electron heating can be adequately explained by an ES electric field in the shock ramp called the cross-shock potential [Scudder et al., 1986a]. For a more detailed discussion of macroscopic fields, see Section 3.3.
Quasi-linear theory estimates the effective resistivity produced by wave particle interactions:

$$\eta_{IA} = \frac{\nu_{IA}}{\varepsilon_0 \omega_{pe}^2}$$  \hspace{1cm} (1.44)

where $\nu_{IA}$ is the effective collision frequency between the wave electric fields and electrons given by:

$$\nu_{IA} = \omega_{pe} \varepsilon_0 \frac{|\delta E|^2}{2n_e k_B T_e}$$  \hspace{1cm} (1.45)

where $n_e$ is the electron number density, $T_e$ is the electron temperature, and $|\delta E|$ is the amplitude of the fluctuating electric field due to the wave. Vlasov simulations using realistic mass ratios have found that the results of Equation 1.44 can be up to 2-3 orders of magnitude smaller than the momentum transfers observed in the simulations [Petkaki et al., 2006; Petkaki and Freeman, 2008; Yoon and Lui, 2006, 2007]. Thus, the quasi-linear estimates of wave-particle collision rates can serve as a lower bound when estimating resistivities in observations.

The instabilities expected to occur in the ramps of collisionless shock waves depend on the shock geometry and Mach number. The main source of free energy in a low Mach number ($M_f \leq 3$) quasi-perpendicular ($\theta_{Bn} > 45^\circ$) collisionless shock wave is thought to be due to the relative drift between electrons and ions (either field-aligned or cross-field), a current [Lemons and Gary, 1978]. As a consequence of the free energy source, instabilities are thought to arise and contribute to resistive energy dissipation in the transition region of quasi-perpendicular collisionless shocks. Some commonly expected instabilities thought to be important in collisionless shocks are: ES IAWs, bipolar ES structures with Debye length scales parallel to the background magnetic field, called solitary waves or phase space holes, modified two stream instability (MTSI), lower-hybrid-drift instability (LHDI), and electron cyclotron drift instability (ECDI) [Wu et al., 1984].
1.7 Electromagnetic Waves: Background, Definitions, and Prior Work

1.7.1 Langmuir Waves

Langmuir waves have been studied extensively in the terrestrial foreshock and somewhat at IP shocks \cite{Bale1997, Fitzenreiter2003, Kellogg1999}. Langmuir waves are usually linearly polarized parallel to the ambient magnetic field with narrow frequency peaks near $f_{pe}$. They are capable of pitch-angle scattering electrons and perturb the background density levels \cite{Soucek2005}. Langmuir waves are also thought to be the progenitors of solar radio emissions, specifically Type II and Type III radio bursts \cite{Bale1999, Kellogg2003, Pulupa2008, Pulupa2010}. Though they are not directly related to collisionless shocks, the shock structure is related to the source of free energy for the Langmuir waves, electron beams \cite{Kellogg2003}. Langmuir waves are also thought to scatter off of density perturbations, which can be a useful tool in describing source regions of radio bursts or shock structure \cite{Krasnoselskikh2007}.

Figure 1.13 is an illustrative example of a linearly polarized Langmuir wave seen upstream of an IP shock on 04/06/2000 at 16:31:54.951 UT by the Wind spacecraft’s TDS detector \cite{Bougeret1995}. The figure is a $\sim 17$ ms waveform capture that has been rotated into field-aligned coordinates defined by the XY-GSE projection of the magnetic field (due to only two component observations). The top two panels are the parallel ($E_\parallel$ in red) and perpendicular ($E_\perp$ in blue) projections of the electric field while the bottom two are the Morlet wavelet transforms. As one can see, Langmuir waves are high frequency (typically $\sim 10$-50 kHz in the solar wind), linearly polarized, large amplitude ($> 10$ mV/m) waves. They are observed as both ES and electromagnetic in the solar wind \cite{Bale1999, Pulupa2008, Wilson2007}.

Figure 1.13 is an example of a waveform capture of an electric field measurement of a Langmuir wave observed upstream of an IP shock. The electric fields have been rotated into field-aligned coordinates where the top panel (red electric field) is the parallel component, $E_\parallel$, and the second panel (blue electric field) is the perpendicular, $E_\perp$. In the top two panels are vertical black arrows that define the relative amplitudes of
Figure 1.13: An example Langmuir wave observed upstream of the shock ramp of the 04/06/2000 event of [Wilson III et al. 2009]. The top two panels are the parallel ($E_\parallel$ in red) and perpendicular ($E_\perp$ in blue) projections of the electric field while the bottom two are the Morlet wavelet transforms. The wavelet transforms are labeled for each relative component. The wavelets are plotted from 100 Hz to 60 kHz. The hodogram to the right plots $E_\perp$ versus $E_\parallel$ in the region outlined by the magenta box with the solid green line representing XY-GSE projection of the shock normal vector in FACs.
each component of the electric field measured. Below the two waveforms are the corresponding frequency spectrum calculated using wavelet analysis (see Section 2.4.3 for more information). To the right is a hodogram plotting $E_\perp$ versus $E_\parallel$ for the time range outlined by the magenta box in the top two panels to the left. The solid green line in the hodogram represents the XY-GSE projection of the shock normal vector in FACs. As one can see, this particular Langmuir wave appears to be linearly polarized roughly parallel to the magnetic field.

### 1.7.2 Ion-Acoustic Waves

A number of authors [Gurnett et al., 1979b,a; Hess et al., 1998; Thomsen et al., 1985a] have concluded that IAWs are important in dissipating energy in lower Mach number shocks. Wave amplitudes in previous studies were found to be correlated with the electron to ion temperature ratio, $T_e/T_i$ [Gurnett et al., 1979a]. They tend to be broadband bursty waves with Doppler shifted frequencies between 1-10 kHz (typically $f_{\pi} < f < f_{pe}$) in the solar wind with a maximum intensity around 3 kHz [Gurnett et al., 1979b,a; Hess et al., 1998]. They are usually linearly polarized close to parallel or oblique to the ambient magnetic field [Akimoto et al., 1985; Akimoto and Winske, 1985; Fuselier and Gurnett, 1984]. In a shock, the instability is thought to be driven by a relative drift between electrons and ions [Mellott, 1985], with threshold drifts increasing for small $T_e/T_i$. A number of studies have concluded that IAWs are likely to be dominant in the terrestrial bow shock despite questions about high damping effects due to small $T_e/T_i$ [Akimoto et al., 1985; Akimoto and Winske, 1985; Fuselier and Gurnett, 1984]. Theoretical studies suggest temperature gradients and oblique propagation of the waves can reduce damping when $T_e \sim T_i$.

Figure 1.14 is an example of an IAW observed in the shock ramp of an IP shock on 1996-04-08. The format is the same as Figure 1.13. As one can see, the wave appears to be roughly linearly polarized parallel to the magnetic field in the plane of measurement, though their polarization can occasionally be quite oblique and more elliptical. They often exhibit a relatively narrow and well defined frequency peak, as illustrated in the wavelet transform plots below the waveforms.
Figure 1.14: An example IAW observed in the shock ramp of the 04/08/1996 event of Wilson III et al. [2009]. The format is similar to that of Figure 1.13, except the wavelets are plotted from 100 Hz to 10 kHz.
1.7.3 Electrostatic Solitary Waves

ES solitary waves (ESWs) are characterized as nonlinear ES Debye-scale bipolar electric field signatures parallel to the ambient magnetic field \cite{CattellEtAl2002a, CattellEtAl2002b, CattellEtAl2005, ErgunEtAl1998a, FranzEtAl2005, PickettEtAl2004}, often associated with electron beams \cite{CattellEtAl2005, ErgunEtAl1998a, FranzEtAl2005}. Phase space holes were first thought to be a nonlinear mode consistent with a BGK mode by \cite{MatsumotoEtAl1994}. Thus, the component of the electric field parallel to the magnetic field is seen as a bipolar pulse while the components perpendicular are monopolar, both are derivatives of a Gaussian. Solitary waves have been observed at the Earth’s bow shock \cite{BaleEtAl1998, BaleEtAl2002}, and at an IP shock near \(\sim 8.7\) AU \cite{WilliamsEtAl2005}, as well as within the magnetosphere at many boundaries \cite{CattellEtAl2002a, CattellEtAl2002b, CattellEtAl2005} possibly providing energy dissipation. Simulations have shown them to form in and around the ramp regions of high Mach number collisionless shock waves \cite{MatsukiyoAndScholer2006a, ShimadaAndHoshino2000}.

ESWs act like clumps of positive charge, if electron holes. In the frame of the electron hole, the ions can be incident on the hole at very large speeds relative to their thermal speed. The relative speed is large because electron holes travel at roughly the electron drift velocity or fractions of electron beam speeds which is much much larger than the ion thermal speed \cite{BehlkeEtAl2004, CattellEtAl2002a, CattellEtAl2003, CattellEtAl2005, ErgunEtAl1998a, ErgunEtAl1998b, FranzEtAl1998}. Since the structures are on the order of an electron Debye length, \(\lambda_{De}\), the transit times of incident ions will be relatively small compared to the local ion gyroperiod. Thus, the ions can become demagnetized if scattered. To visualize this, assume an ion is incident on an electron hole with an impact parameter of \(b\). Then the perpendicular impulse of the ion in response to the electron hole’s electric field can be shown as:

\[
M_i \Delta V_{i,\perp}(b) = e \int_{-\infty}^{\infty} dt \, E_{\perp}[b, z(t)]
\]

where \(M_i\) is the ion mass, \(\Delta V_{i,\perp}\) is the change in ion velocity, \(E_{\perp}\) is the electric field perpendicular to the magnetic field, and \(z(t)\) is the position along the direction parallel to the magnetic field at time \(t\). If we assume that the incident ion velocity in the electron hole frame is \(V_{ehole}\), and that the electron hole will not recoil upon impacting a single ion, then we can also estimate the parallel impulse of the ions as \(\Delta V_{i,\parallel} \simeq -\Delta V_{i,\perp}^2/(2 \, V_{ehole})\).
The net result is an exchange of momentum between the ions and the electrons both parallel and perpendicular to the magnetic field. If there is a relative drift between the two species, this momentum exchange can act to reduce the relative drift, thus reduce a current. In this way, electron holes can be effective waves for inducing anomalous resistivity. The end result is strong perpendicular ion heating (i.e. random perpendicular ion acceleration) and a significant amount of parallel momentum imparted upon the electrons from the ions. Observations have shown that $\Delta T_{\perp}$ across a train of electron holes can be a significant fraction of the ion thermal energy [Ergun et al., 1998b].

Since the holes act like clumps of positive charge, they have positive potentials and thus act to trap incident electrons [Dyrud and Oppenheim, 2006; Lu et al., 2008]. The particles that are trapped are the ones with energies below that of the max potential of the solitary structure. The accumulation of electrons trapped in the solitary structure acts to damp/saturate the instability driving the holes [Lu et al., 2008]. In addition to trapping electrons, the electron holes can create double-peaked electron distributions at low energies which are unstable to other wave modes [Berthomier et al., 2008; Matsukiyo and Scholer, 2006a].

The last, somewhat indirect, way in which solitary waves can heat/scatter particles is by coupling to other wave modes. Solitary waves can either couple to or directly cause IAWs [Dyrud and Oppenheim, 2006], whistler mode waves [Lu et al., 2008], and electron acoustic waves [Matsukiyo and Scholer, 2006a]. IAWs are known to heat electrons parallel to their fluctuating electric fields (typically along the magnetic field) [Dum et al., 1974]. Whistler waves are known to cause a perpendicular pitch-angle diffusion and heating of electrons [Brice, 1964; Kennel and Petscheck, 1966]. Electron acoustic waves are thought to produce strong parallel (with respect to the magnetic field) electron heating [Matsukiyo and Scholer, 2006a], but to the best of our knowledge, these modes have not been observed.

Figure 1.15 is an example of an ESW observed by the Wind spacecraft in the terrestrial magnetosphere. The format is the same as Figures 1.13 and 1.14. Though these wave modes are observed at collisionless shocks, they are typically much larger amplitude in the magnetosphere making the characteristic bipolar signature more obvious. Notice that the defining characteristic, parallel component (red) is bipolar while the perpendicular is monopolar (blue), is very obvious in this example.
Figure 1.15: An example solitary wave observed in the terrestrial magnetosphere. The format is similar to that of Figures 1.13 and 1.14. The parallel and perpendicular component wavelet transforms are labeled respectively.
1.7.4 Whistler Waves

Whistler waves were first discovered by Barkhausen [1919] while listening to signals from an antenna connected to a simple vacuum tube amplifier. The signals were heard to decrease in frequency with increasing time. Over thirty years later, Storey [1953] managed to explain these strange signals as being the result of lightning strikes. It is now known that whistler waves can exist as a RH electromagnetic mode [Kennel and Petscheck, 1966] or a slightly electrostatic mode and interact with both ions [Hoppe et al., 1982; Stasiewicz et al., 2003] or electrons [Brice, 1964; Kennel and Petscheck, 1966], and [Lyons et al., 1972]. Due to their capacity to resonantly and nonresonantly interact with particles, whistler waves are a topic of extreme interest in collisionless shock dissipation topics.

A specific class of whistler wave is often observed immediately upstream of a quasi-perpendicular collisionless shock wave in magnetic field data, called precursor whistler waves or just precursor waves. Their existence was theorized as a necessary part of the shock structure since collisionless shock waves were first predicted [Kellogg, 1962]. The necessity of their existence at low Mach number collisionless shock waves was more rigorously shown by Morton [1964] and Stringer [1963]. These waves are low enough frequency that they couple to the magnetosonic wave responsible for the shock ramp. Thus, they can provide energy dissipation in shock waves through dispersive effects and wave-particle interactions upstream of the shock ramp [Gary and Mellott, 1985]. Occasionally precursor whistlers phase stand with respect to the shock front [Fairfield and Feldman, 1975], which means their phase velocity matches the shock front phase velocity. Therefore, in the shock frame of reference, the waves appear as standing wave modes. This condition is satisfied at the bow shock when the phase speed of the precursor whistler matches the shock wave phase speed and the group velocity exceeds the incident solar wind speed [Greenstadt et al., 1975].

Figure 1.16 shows an example of a quasi-perpendicular, low Mach number IP shock with a precursor whistler wave (highlighted by purple box). The top panel shows the magnetic field magnitude and the bottom panel shows the three GSE components of the magnetic field. The fluctuations in the magnitude of the magnetic field illustrate the compressive and oblique nature of this class of whistler wave. This particular whistler was observed to propagate at an oblique angle to the magnetic field and has spatially
Figure 1.16: An example of an IP shock with an upstream precursor whistler wave. The top panel is the magnitude of the magnetic field and the bottom panel is the GSE components. The region with the precursor is outlined by the translucent blue box.

Fairfield and Feldman [1975] initially identified precursor whistler waves at the quasi-perpendicular bow shock using magnetometer data from Explorer 43 as phase standing. Later studies, using OGO 5 [Greenstadt et al., 1975] and ISEE 1 and 2 [Mellott and Greenstadt, 1984] magnetometer data, found many of the upstream waves to be inconsistent with phase standing whistlers. Mellott and Greenstadt [1984] found two different types of precursor whistler waves, a phase standing whistler wave propagating parallel to the shock normal and another whistler propagating parallel to the magnetic field. In the SC frame, the precursors propagating parallel to the magnetic field had higher frequencies (∼1 Hz) than the phase standing precursor whistlers (∼0.1 Hz). It is important to note that the ∼1 Hz waves studied by Hoppe et al. [1982] had relatively large $\theta_{kB}$ values, while the precursors of Mellott and Greenstadt [1984] were propagating parallel to the magnetic field, thus $\theta_{kB} \sim 0^\circ$. Mellott and Greenstadt [1984] proposed that the parallel propagating precursors were products of the phase standing precursors. The precursors propagating parallel to the shock normal (the phase standing
precursors) were found to have higher rest frame frequencies than the precursors propagating parallel to the magnetic field. The difference was due to their propagation with respect to the magnetic field. The Doppler effects on the parallel propagating precursors were negligible because the magnetic field was primarily directed in Y-GSE direction, roughly perpendicular to the solar wind velocity. Both the parallel propagating and phase standing precursors are characterized by a high degree of RH polarization and nearly monochromatic frequency spectrum. A more recent study by Farris et al. [1993] found that the comparison between observed and predicted wavelengths for phase standing precursor whistlers to be consistent with the results of Mellott and Greenstadt [1984]. However, the estimates by Farris et al. [1993] of the ratio of precursor whistler wavelength to shock thickness differed from those of Mellott and Greenstadt [1984]. Thus, Farris et al. [1993] concluded that the thickness of the shock ramp was not strictly dependent/related to the wavelength of the precursor whistler.

1.7.5 Waves at or Near the Electron Cyclotron Frequency

Electron cyclotron harmonic, electron Bernstein, \((n + 1/2)\), or "totem pole" waves have been observed throughout planetary magnetospheres by Barbosa et al. [1990] and Usui et al. [1999]. These emissions can be both broad or narrow in frequency range [Hubbard and Birmingham, 1978]. They are typically driven unstable by loss-cone or anisotropic electron distributions in the high energy hot halo in planetary magnetospheres. Usui et al. [1999], in a study near the terrestrial magnetopause, found the emissions to be associated with increases in the ratio of hot halo to cold core electron densities, \(n_h/n_c\). To the best of our knowledge, these emissions have not been observed previously in the solar wind.

Figure 1.17 shows two examples of electron cyclotron harmonic waves observed downstream of an IP shock by the Wind spacecraft. The top/bottom four panels correspond to the waveform observed at 16:32:25.358 UT (16:32:25.428 UT) on 04/06/2000. The top row of panels for each waveform with the red/blue lines correspond to the \(E_\parallel\)\((E_\perp)\) component of the wave electric field. The bottom two panels of each waveform show the power spectra (mV/m²/Hz) versus frequency (kHz) plots corresponding to the time range defined by the orange box in the top two panels. The vertical green/magenta lines overplotted on the power spectra correspond to integer/half-integer harmonics of
Figure 1.17: Two examples of electron cyclotron harmonic waves observed downstream of an IP shock. The top waveform was observed at 16:32:25.358 UT and the bottom at 16:32:25.428 UT. The top row contains $E_\parallel$ (red) and $E_\perp$ (blue). Below the waveforms are the power spectra (mV/m$^2$/Hz) versus frequency (kHz) plots corresponding to the time range defined by the orange box. The vertical green (magenta) lines overplotted on the power spectra correspond to integer(half-integer) harmonics of $f_{ce}$. The 16:32:25.428 UT has a similar format.
Note that the 16:32:25.358 UT waveform primarily shows enhanced power at integer harmonics of $f_{ce}$ while the 16:32:25.428 UT waveform shows mixtures of integer and half-integer harmonics of $f_{ce}$. The wave power enhancements shift dynamically in time throughout the waveform, thus why only small windows of time were used to calculate the power spectra. These waveforms are over two orders of magnitude above the background levels ($\sim 0.1 \text{ mV/m at 1 AU}$).

Simulations have found bipolar ES phase space holes form in and around the ramp regions of high Mach number collisionless shock waves [Dyrud and Oppenheim, 2006; Matsukiyo and Scholer, 2006a]. Due to their ability to efficiently exchange momentum between electrons and ions, the holes can heat and scatter particles. Simulations also show that the holes can also couple with other wave modes like IAWs and lower hybrid waves, providing resistive dissipation [Dyrud and Oppenheim, 2006; Matsukiyo and Scholer, 2006a]. Matsukiyo and Scholer [2006a] examined microinstabilities in the foot of supercritical collisionless shocks using a two dimensional PIC simulation with a realistic mass ratio ($M_i/m_e \sim 1836$). They observed six different types of instabilities excited in less than an ion gyroperiod with the dominant modes including ECDI, whistler instability, electron acoustic instability, and two different modified two-stream instabilities (MTSIs); MTSI-2 excited by relative drifts between incident electrons and reflected ions and MTSI-1 due to the relative drift between electrons and incident ions. Reflected ions cause the incident solar wind ions to decelerate in the shock foot, which locally decelerate the electrons to maintain current continuity in the shock normal direction. These instabilities give rise to waves which scatter and heat the plasma, thus dissipating energy.

In this thesis, we will discuss a type of cyclotron wave/instability in detail in Chapter 5 and their effect on local electron distributions.

### 1.7.6 Lower Hybrid Waves

Lower hybrid waves (LHWs) are typically an electrostatic (ES) mode propagating perpendicular to the magnetic field with a frequency given by:

$$f_{lh}^2 = \frac{f_{ce} f_{ci}}{1 + (f_{ce} f_{ci})/f_{pi}^2}$$

(1.47)
where $f_{cs}$ is the cyclotron frequency of species $s$ and $f_{pi}$ is the ion plasma frequency. In
the high density limit ($f_{pi}^2 \gg f_{ce} f_{ci}$), typical of the solar wind, $f_{th} \sim (f_{ce} f_{ci})^{1/2}$. LHWs
are capable of resonating with the bulk of the ion distribution, thus they can serve
as an effective ion heating mechanism [Davidson and Gladd, 1975; Lemons and Gary,
1978; Mellott, 1985]. In the ES limit and large propagation angles, the waves can
couple the parallel motion of the electrons to the perpendicular motion of the ions
[Marsch and Chang, 1983]. LHWs can be driven unstable by cross-field currents
[Lemons and Gary, 1978], electron heat flux in the solar wind [Marsch and Chang,
1982], and the lower-hybrid drift instability (LHDI) [Cairns and McMillan, 2005]. They
are an attractive candidate for resistive energy dissipation in collisionless shocks be-
cause their critical drift speeds, $V_{dcr} \approx V_{Ti}$, are much lower than that of an IAW,
$V_{dcr} \approx V_{Te}$. LHWs can couple to other wave modes like drift waves, modified two-
stream instabilities (MTSIs), etc., all of which cause significant wave-particle inter-
actions [Lemons and Gary, 1978]. LHWs can heat the ions transverse to the magnetic
field producing anisotropic ion distributions [Marsch and Chang, 1982]. Because LHWs
have $\omega/k_{\parallel} \gg \omega/k_{\perp}$, they tend to interact with the higher energy electrons producing
broadened high energy tails [Cairns and McMillan, 2005].

Electromagnetic lower hybrid waves (EMLHWs), or hybrid whistler waves, prop-
agate nearly perpendicular to the ambient magnetic field and can appear to have a
broadband frequency spectrum. One should note that in the limit of large $k_{\perp}$, LHWs
are on the same branch of the dispersion relation as whistler waves. As with the ES
LHWs, they are thought to be driven unstable by the solar wind electron heat flux
[Marsch and Chang, 1983]. They can, much like ES LHWs, heat the ions perpendicular
to the magnetic field. For these modes, $V_{i,\text{res},\perp} \approx \omega/k_{\perp} \ll \omega/k_{\parallel} \approx V_{e,\text{res},\parallel}$, which means
the Landau interactions are perpendicular for the ions and parallel for the electrons, with
respect to the magnetic field. They have frequencies of $f_{ci} \ll f \ll f_{ce}$ [Marsch and Chang,
1983]. These waves dissipate their wave energy through Landau interaction with the
ions producing perpendicular ion heating. They propagate very obliquely to the field
within a cone defined by $k_{i}/k_{\perp} \leq 1/5$ and $k_{i}/k_{\perp} \geq V_{Ti,\perp}/V_{Te,\parallel}$ [Marsch and Chang,
1983]. Zhang and Matsumoto [1998] observed EMLHWs at an IP shock using Geotail
and showed that the waves propagate nearly perpendicular to the magnetic field.

Another related wave mode is the LHDI, which in the presence of strong plasma
gradients, acts like a fluid instability excited through the coupling of a LHW and a drift wave \cite{Davidson and Gladd, 1975, Huba et al., 1978}. When the gradients are weak, the LHDI is a kinetic instability driven by a resonance between ions and a drift wave. When in the presence of a finite plasma $\beta$, the LHDI exists as an ES and electromagnetic mode \cite{Davidson and Gladd, 1975, Huba et al., 1978}. The growth rate of the LHDI peaks at $k \rho_e \approx 1$, for a broad range of frequencies near $f_{lh}$ \cite{Davidson and Gladd, 1975, Cairns and McMillan, 2005}. The mode is strongly unstable when the magnetic field gradient scale length, $L_B$, is comparable to $\rho_i$. The LHDI produces strong anomalous resistivity due to the wave’s electric fields, $\delta E_\perp$, perpendicular to the ambient magnetic field, $B_o$, which create $(\delta E_\perp \times B_o)$-drifts that transport particles across $B_o$. Thus, the LHDI causes cross-field diffusion which is an increase in entropy, thus irreversible and important for energy dissipation \cite{Coroniti, 1985}.

Figure 1.18 is a TDSS sample observed downstream of the shock ramp of an IP shock observed by the Wind spacecraft on 02/11/2000. The waveform is an example of an EMLHW. The left-hand side of the plot shows the three components of the magnetic field measured by the TDSS detector search coils on Wind in minimum variance coordinates (see Section 2.4.1 for more details). The right-hand side of the plot shows the corresponding wavelet transforms (see Section 2.4.3 for more details). To the right of the wavelets are the labels of the relevant frequencies for this wave, where $f_e$ is the electron cyclotron frequency and $f_{lh}$ is the lower hybrid resonance frequency. In the left-hand panel, one can see the angle of propagation with respect to the magnetic field is $\theta_{KB} \sim 85^\circ$, roughly perpendicular to the magnetic field. Note that at multiple points in the waveform, there appear to be low and high frequency signals intermixed. At roughly 400 ms, a higher frequency (at roughly 100 Hz), is superposed on the lower frequency ($\sim 10$ Hz or $\sim f_{lh}$) signal.

Figure 1.19 shows the result of filtering the signal for $7$ Hz $< f < 20$ Hz (see Section 2.4.2 for discussion of frequency filtering). Note that the mid-to-min eigenvalue ratio, $\lambda_2/\lambda_3$, is much higher now at $\sim 12$ compared to $\sim 3$ in Figure 1.18. Also, $\theta_{KB}$ has increased to $\sim 90^\circ$, consistent with previous observations \cite{Zhang and Matsumoto, 1998} and theory \cite{Marsch and Chang, 1983}. The amplitude of the low frequency component is $\sim 1$ nT.

Figure 1.20 shows the result of filtering the signal for $60$ Hz $< f < 200$ Hz. The
Figure 1.18: An example TDSS sample observed downstream of the shock ramp of an IP shock observed by the Wind spacecraft on 02/11/2000. The left-hand side of the plot shows the three components of the magnetic field measured by the TDSS detector search coils on Wind in minimum variance coordinates. The right-hand side of the plot shows the corresponding wavelet transforms. To the right of the wavelets are the labels of the relevant frequencies for this wave, where $f_{ce}$ is the electron cyclotron frequency and $f_{lh}$ is the lower hybrid resonance frequency. Other relevant information is given in the plot including the wave vector, $\mathbf{k}$, in GSE coordinates, the angle of propagation with respect to the magnetic field, $\theta_{kB}$, and the eigenvalue ratios, $\lambda_1/\lambda_2$ and $\lambda_2/\lambda_3$. 

$\mathbf{k}_{GSE} = \langle +0.714, +0.376, -0.590 \rangle \pm \langle 0.040, 0.091, 0.106 \rangle$

Figure 1.19: The low frequency filter (7 Hz $< f <$ 20 Hz) of the TDSS sample observed downstream of the 02/11/2000 IP shock ramp in Figure 1.18. The left-hand side of the plot shows the three components of the magnetic field measured by the TDSS detector search coils on Wind in minimum variance coordinates. The right-hand side of the plot shows the corresponding hodograms.
Figure 1.20: The high frequency filter ($60 \text{ Hz} < f < 200 \text{ Hz}$) of the TDSS sample observed downstream of the 02/11/2000 IP shock ramp in Figure 1.18. The left-hand side of the plot shows the three components of the magnetic field measured by the TDSS detector search coils on Wind in minimum variance coordinates. The right-hand side of the plot shows the corresponding hodograms.
format is the same as in Figure 1.19, but the amplitude of the wave is smaller and it is right-hand polarized with respect to the magnetic field, consistent with an electromagnetic whistler wave. The wave is propagating at a slightly oblique angle with respect to the magnetic field, which is roughly anti-parallel to the wave vector.

In this thesis, I utilize waveform captures of each type of wave reviewed in this section to look for evidence of wave-particle interactions observed as wave induced heating, pitch-angle scattering and/or diffusion, and particle acceleration.

1.8 Thesis Overview

Ever since the prediction of the existence of collisionless shock waves [Kellogg, 1962], intense interest has been focused on possible dissipation mechanisms [Sagdeev, 1966]. As discussed above, the energy dissipation must be irreversible for a true shock transition to occur. Possible energy dissipation mechanisms include wave dispersion [Mellott and Greenstadt, 1984], particle reflection [Edmiston and Kennel, 1984; Kennel, 1987], macroscopic field effects [Bale and Mozer, 2007; Hull et al., 2001; Walker et al., 2004; Wygant et al., 1987], and anomalous resistivity due to wave-particle interactions [Gary, 1981]. The relative significance of each type of energy dissipation can depend upon multiple factors, but they are not well understood. It is thought that particle reflection is more important at higher Mach number shocks, while lower Mach number shocks depend upon wave dispersion and anomalous resistivity [Mellott and Greenstadt, 1984]. Yet it should be noted that particle reflection alone cannot produce a complete shock transition [Kennel et al., 1985]. In space, the terrestrial bow shock has been extensively studied [Bale et al., 1997, 1998, 2000, 2002; Hull et al., 2006; Kellogg et al., 1999], but IP shocks are less well examined [Fitzenreiter et al., 2003; Gurnett et al., 1979; Hess et al., 1998; Thejappa and MacDowall, 2000; Wilson III et al., 2007]. Due to their lower Mach numbers, it is expected that IP shocks rely upon wave dispersion and anomalous resistivity for energy dissipation [Mellott and Greenstadt, 1984]. Thus, examining the waves in IP shocks is a necessary requirement for understanding the evolution of IP shocks as they propagate away from the sun.

This thesis examines energy dissipation mechanisms, shock structure, and particle heating in IP shocks using data from the Wind spacecraft. The thesis is organized as
1. Chapter 2 outlines and summarizes the Wind spacecraft instrumentation and the analysis techniques used herein.

2. Chapter 3 outlines the important low frequency waves seen in and around collisionless shock waves. This chapter also discusses some of the macroscopic fields observed in shock ramps and their effects on particle distributions.

3. Chapter 4 reports on the study done by Wilson III et al. [2009] on low frequency \((0.25 \, \text{Hz} < f < 10 \, \text{Hz})\) and summarizes their results and conclusions.

4. Chapter 5 reports on the studies done by Wilson III et al. [2007] and Wilson III et al. [2010] on high frequency \((\gtrsim 400 \, \text{Hz})\) waves in and around IP shocks.

5. Chapter 6 summarizes the conclusions of this thesis and outlines some future work.

This thesis includes material from the following papers: Wilson III et al. [2007], Wilson III et al. [2009], and Wilson III et al. [2010].
Chapter 2

Instrumentation, Measurements, and Analysis Techniques

In this chapter, the measurement and data analysis techniques needed for studies of IP shocks are described. Potential error sources are discussed, although some details are deferred to Appendix A. Illustrative examples from Wind measurements are presented.

2.1 Introduction to the Wind Spacecraft and instruments

The Wind spacecraft was launched on November 1st, 1994 by a Delta II rocket from Cape Canaveral Air Force Station in Merritt Island, FL. For the first two years of the mission WIND was in a highly elliptical orbit on the sunward side of the Earth with an apogee of 250 Earth radii ($R_E$) and a perigee of at least 5 $R_E$. Wind was the first of NASA’s Global Geospace Science (GGS) program, which was part of the International Solar-Terrestrial Physics (ISTP) Science Initiative, a collaboration between several countries in Europe, Asia, and North America. The aim of ISTP was to understand the behavior of the solar-terrestrial plasma environment in order to predict how the Earth’s magnetosphere responds to changes in solar wind conditions. WIND’s objective is to measure the properties of the solar wind before it reaches the Earth [Desch, 2005].

The Wind spacecraft has an array of instruments including: Konus [Antekar et al., 1995], the Wind Magnetic Field Investigation (MFI) [Lepping et al., 1995], the Solar
Wind and Suprathermal Ion Composition Experiment (SMS) [Gloeckler et al., 1995], the Solar Wind Experiment (SWE) [Ogilvie et al., 1995], a Three-Dimensional Plasma and Energetic Particle Investigation (3DP) [Lin et al., 1995], the Transient Gamma-Ray Spectrometer (TGRS) [Owens et al., 1995], and the Radio and Plasma Wave Investigation (WAVES) [Bougeret et al., 1995]. The Konus and TGRS instruments are primarily for gamma-ray and high energy photon observations of solar flares or gamma-ray bursts. The SMS experiment measures the mass and mass-to-charge ratios of heavy ions. The SWE and 3DP experiments were designed to measure the lower energy (below 10 MeV) solar wind protons and electrons. The WAVES and MFI experiments were designed to measure the electric and magnetic fields observed in the solar wind. All together, the Wind suite of instruments allows for a complete description of plasma phenomena in the solar wind in the ecliptic plane [Desch, 2005].

2.1.1 Wind 3DP Particle Detector

The Wind/3DP instrument was designed to make full three-dimensional measurements of the distributions of suprathermal electrons and ions in the solar wind. The instrument includes three arrays, each consisting of a pair of double-ended semi-conductor telescopes each with two or three closely sandwiched passivated ion implanted silicon detectors, which measure electrons and ions above ∼20 keV. The instrument also has top-hat symmetrical spherical section electrostatic analyzers with microchannel plate detectors (MCPs) are used to measure ions and electrons from ∼3 eV to 30 keV [Lin et al., 1995].

The two types of detectors have energy resolutions ranging from ΔE/E ≈ 0.3 for the solid state telescopes (SST) and ΔE/E ≈ 0.2 for the top-hat electrostatic (ES) analyzers. The angular resolutions are 22.5° × 36° for the SST and 5.6° (near the ecliptic) to 22.5° for the top-hat ES analyzers. The particle detectors can obtain a full 4π steradian coverage in one full(half) spin (∼3 s) for the SST(top-hat ES analyzers).

ES Analyzers

The arrays of detectors are mounted on two opposing booms, each 0.5 m in length. The top-hat ES analyzers are composed of four separate detectors, each with different geometric factors to cover different ranges of energies. The electron detectors, EESA,
and ion detectors, PESA, are each separated into low (L) and high (H) energy detectors. The H and L analyzers contain 24 and 16 discrete anodes, respectively. The anode layout provides a 5.6° angular resolution within ±22.5° of the ecliptic plane (increases to 22.5° at normal incidence to ecliptic plane). The analyzers are swept logarithmically in energy and counters sample at 1024 samples/spin (~3 ms sample period). Thus the analyzers can be set to sample 64 energy samples per sweep at 16 sweeps per spin or 32 energy samples per sweep at 32 sweeps per spin, etc. The detectors are defined as follows:

1. **EESA-L (EL):** covers electrons from ~3 eV to ~1 keV with a 11.25° spin phase resolution. EL has a total geometric factor of $1.3 \times 10^{-2}$ E cm$^{-2}$-sr (where E is energy in eV) with a nearly identical 180° field of view (FOV), radial to the spacecraft, to that of PESA-L.

2. **EESA-H (EH):** covers electrons from ~200 eV to ~30 keV (though typical values vary from a minimum of ~137 eV to a maximum of ~28 keV) in a 32 sample energy sweep each 11.25° of spacecraft spin. EH has a total geometric factor of $2.0 \times 10^{-1}$ E cm$^{-2}$-sr, MCP efficiency of about 70% and grid transmission of about 73%. EH has a 360° planar FOV tangent to the spacecraft surface which can be electrostatically deflected into a cone up to ±45° out of its normal plane.

3. **PESA-L (PL):** covers ions with a 14 sample energy sweep from ~100 eV to ~10 keV (often energies range from ~700 eV to ~6 keV) each 5.6° of spacecraft spin. PL has a total geometric factor of only $1.6 \times 10^{-4}$ E cm$^{-2}$-sr but an identical energy-angle response to that of PESA-H. While in the solar wind, PL reorients itself along the bulk flow direction to capture the solar wind flow which results in a narrow range of pitch-angle coverage.

4. **PESA-H (PH):** covers ions with a 15 sample energy sweep from as low as ~80 eV to as high as ~30 keV (typical energy range is ~500 eV to ~28 keV) each 11.25° of spacecraft spin. PH has a total geometric factor of $1.5 \times 10^{-2}$ E cm$^{-2}$-sr. These values vary from moment structure to moment structure depending on duration of data sampling, spacecraft potential, and whether in burst or survey mode. The typical range is ~5 eV to ~1.11 keV.

Note that in survey mode the data structures typically take 25 data points at 14 different energies while in burst mode they take 64 data points at 14 different energies.

Note that PH has multiple data modes where the number of data points per energy bin can be any of the following: 121, 97, 88, 65, or 56.
cm²-sr with a MCP efficiency of about 50% and grid entrance post transmission of about 75%.

Characteristics of the entire instrument suite are shown in Table 2.1 from Lin et al. [1995].

<table>
<thead>
<tr>
<th>Detector</th>
<th>Energy Range</th>
<th>Geometric Factor (cm²-sr)</th>
<th>FOV (°)</th>
<th>Dynamic Range</th>
</tr>
</thead>
<tbody>
<tr>
<td>EH/FPC¹</td>
<td>100eV - 30keV</td>
<td>0.1 E</td>
<td>360 × 90</td>
<td>~10⁹-10⁸</td>
</tr>
<tr>
<td>EL¹</td>
<td>3eV - 30keV</td>
<td>0.013 E</td>
<td>180 × 14</td>
<td>~10²-10⁹</td>
</tr>
<tr>
<td>PH²</td>
<td>3eV - 30keV</td>
<td>0.015 E</td>
<td>360 × 14</td>
<td>~10¹-10⁹</td>
</tr>
<tr>
<td>PL²</td>
<td>3eV - 30keV</td>
<td>0.00016 E</td>
<td>180 × 14</td>
<td>~10⁴-10¹¹</td>
</tr>
<tr>
<td>Foil¹</td>
<td>25-400keV</td>
<td>1.7 E</td>
<td>180 × 20</td>
<td>~10⁻¹-10⁶</td>
</tr>
<tr>
<td>Open²</td>
<td>20keV - 6MeV</td>
<td>1.7 E</td>
<td>180 × 20</td>
<td>~10⁻¹-10⁶</td>
</tr>
</tbody>
</table>

Solid-State Telescopes

The SST detectors consist of three arrays of double-ended telescopes, each of which is composed of either a pair or triplet of closely-sandwiched semi-conductor detectors. The center detector (Thick or T) of the triplet is 1.5 cm² in area, 500 µm thick, while the other detectors, foil (F) and open (O), are the same area but only 300 µm thick. One direction of the telescopes is covered in a thin lexan foil (SST-Foil) where the thickness was chosen to stop protons up to the energy of electrons (~400 keV). Electrons are essentially unaffected by the foil. On the opposite side (SST-Open), a common broom magnet is used to refuse electrons below ~400 keV from entering but leaves the ions essentially unaffected. Thus, if no higher energy particles penetrate the detector walls, the SST-Foil should only measure electrons and the SST-Open only ions. Each double-ended telescope has two 36° × 20° FWHM FOV, thus each end of the five telescopes can cover a 180° × 20° piece of space. Telescope 6 views the same angle to spin axis as telescope 2, but both ends of telescope 2 have a drilled tantalum cover to reduce the geometric factor by a factor of 10 to measure the most intense fluxes. The SST-Foil data

¹ The foil has ~1500 Å of aluminum evaporated on each side to completely eliminate sunlight.
structures typically have 7 energy bins each with 48 data points while the SST-Open has 9 energy bins each with 48 data points. Both detectors have energy resolutions of $\Delta E/E \approx 30\%$.

2.1.2 WAVES

The electric field detectors of the Wind WAVES instrument [Bougeret et al., 1995] are composed of three orthogonal electric field dipole antenna, two in the spin plane (roughly the plane of the ecliptic) of the spacecraft and one along the spin axis. The complete WAVES suite of instruments includes five total receivers including: Low Frequency FFT receiver called FFT (0.3 Hz to 11 kHz), Thermal Noise Receiver called TNR (4-256 kHz), Radio receiver band 1 called RAD1 (20-1040 kHz), Radio receiver band 2 called RAD2 (1.075-13.825 MHz), and the Time Domain Sampler called TDS ($\leq 7.5$ kHz in slow mode and $\leq 120$ kHz in fast mode). The longer of the two spin plane antenna, defined as $E_x$, is 100 m tip-to-tip while the shorter, defined as $E_y$, is 15 m tip-to-tip. The spin axis dipole, defined as $E_z$, is roughly 12 m tip-to-tip. When accounting for spacecraft potential, these antenna lengths are adjusted to $\sim 41.1$ m, $\sim 3.79$ m, and $\sim 2.17$ m, respectively (P.J. Kellogg, Personal Communication, 2007). The magnetic field detectors of the Wind WAVES instrument are composed of three orthogonal search coil magnetometers (designed and built by the University of Iowa). The XY search coils are oriented to be parallel to the XY dipole antenna. Thus, in Figure 2.1, the X-component search coil is at an angle $\theta$ from the X-GSE direction. The search coils allow for high frequency magnetic field measurements (defined as $B_x$, $B_y$, and $B_z$). The WAVES Z-Axis is anti-parallel to Z-GSE direction. Thus any rotations can be done about the Z-Axis in the normal Eulerian sense followed by a change of sign in the Z-component of any GSE vector rotated into WAVES coordinates.

Figure 2.1 is a schematic used to illustrate the relationship between the various field instruments on the Wind spacecraft and the GSE-coordinate system. The point of view is from the negative Z-GSE direction (i.e. below the plane of the ecliptic) looking toward the positive Z-GSE direction (i.e. above the plane of the ecliptic). The spacecraft rotates in a counterclockwise direction in this coordinate system. The house keeping information which informs a user where each boom is with respect to the sun direction is returned as an eight bit integer, providing $\sim 1.5^\circ$ accuracy. The angle $\theta$ is relevant to
Figure 2.1: This is a schematic used to illustrate the relationship between the various field instruments on the Wind spacecraft and the GSE-coordinate system. The relative angles are determined by an eight bit integer. As seen in the image, when this integer equals zero, the magnetic field boom (the x-framed boom along the X-GSE axis) is pointed towards the sun (X-GSE). When the integer reads 224, the positive X-antenna (electric field, $E_x$) is pointed towards the sun. The magnetic field boom is roughly 45° from either X (red line) or Y-antenna (blue line). As indicated by the figure, the satellite rotates in a counter-clockwise direction in this coordinate system.
the software which retrieves TDS samples from the VAX/ALPHA systems (see Section 2.1.2).

**Time Domain Sampler Receiver**

![Diagram](image)

Figure 2.2: This is a schematic cartoon used to illustrate the rotation from a GSE basis to the WAVES antenna basis. The Z-component of the magnetic field must change a sign after rotation about the Z-GSE axis. The figure shows a rotation about the Z-Axis by an angle $\phi$ (blue) which results in the X-antenna aligning with the new X’-Axis. The X’Y’-projection of the magnetic field in this basis is the rotated field one uses when plotting hodograms of the electric field from WAVES.

Electric (and magnetic) field waveform captures can be obtained from the Time Domain Sampler (TDS) receiver *Bougeret et al. 1995*. In the highest sampling rates, the TDS samples are $\sim 17$ ms waveform capture of 2048 points (120 kHz) for the Fast (TDSF) sampler and 7.5 kHz for the Slow (TDSS) sampler. The TDS receiver has four possible modes (see Table 2.2). In the solar wind, TDSF is usually set to sample in the 120 kHz mode but was changed to the 7.5 or 1.9 kHz modes during the petal orbits.
through the Earth’s magnetosphere. TDSF samples (see Figure 2.3) are two components of the electric field (typically in the XY-GSE plane), defined as $E_x$ and $E_y$ (occasionally chosen to be $E_z$ prior to 1996). For TDSF samples, we define $|E_{xy}| = \sqrt{E_x^2 + E_y^2}$ as the peak-to-peak (pk-pk) amplitude of the waveform. When operating at 120 kHz, the TDSF receiver has little to no gain below $\sim 120$ Hz, thus the data is cutoff below 150 Hz when performing gain corrections on ground (P.J. Kellogg, Personal Communication, 2007). Nearly all of the data presented in this thesis is from the TDSF receiver.

Figure 2.2 is a schematic cartoon used to illustrate an example rotation from the GSE basis to the WAVES antenna basis. The blue lines represent the WAVES basis at rotated about the negative $Z$-GSE axis by an angle $\phi$. The angle $\zeta$ represents the angle between the projection of the magnetic field in GSE coordinates ($B_{GSE}$) onto the XY-GSE plane and the X-GSE axis. The rotation matrix given in the figure will rotate $B_{GSE}$ into WAVES coordinates resulting in, $B_{WAVES}$. This is calculation necessary for analysis of wave polarizations. Also, for lower sampling rates (e.g. $\leq 7.5$ kHz), this rotation should be done on each data point since the angle $\phi$ can change by up to $\sim 120^\circ$ during on TDS sample when sampling at 1875 Hz.

The TDSS samples return 4 field vectors, either three electric and one magnetic

<table>
<thead>
<tr>
<th>Table 2.2: Wind WAVES TDS Specs</th>
</tr>
</thead>
<tbody>
<tr>
<td>Speed</td>
</tr>
<tr>
<td>-------</td>
</tr>
<tr>
<td>A</td>
</tr>
<tr>
<td>B</td>
</tr>
<tr>
<td>C</td>
</tr>
<tr>
<td>D</td>
</tr>
<tr>
<td>E</td>
</tr>
<tr>
<td>F</td>
</tr>
</tbody>
</table>

or one magnetic and three electric field measurements. The TDSS receiver also has four possible sample rates, as shown in Table 2.2 from Bougeret et al. [1995]. The gain for the TDSS search coils rolls off below 3.3 Hz. When returning three electric fields, the TDSS samples are often contaminated by spin effects due to different levels of photoelectron currents on the antenna in sunlight versus shadow (i.e. only affects the X
and Y components) and a glitch associated with the Z-antenna (P.J. Kellogg, Personal Communication, 2007). However, if the electric fields are large enough and Wind is in the shadow of the Earth, these effects are negligible. When returning three magnetic and one electric field, the TDSS receiver can return well defined waveforms with only small noise and glitch effects. The TDSS receiver is rarely used because the triggering mechanism did not work correctly. Later in its operation, the TDSS receiver was set to trigger off of the TDSF receiver (P.J. Kellogg, Personal Communication, 2010). This can result in a failure to observe large amplitude wave modes. On occasion, the TDSS receiver does return large amplitude wave modes which are of interest because they are often below the low frequency cutoff of TDSF (see Figures 1.18 through 1.20).

Radio Receivers

The TNR measures $\sim 4$-256 kHz electric fields in up to 5 logarithmically-spaced frequency bands, though typically only set at 3 bands (K. Goetz, Personal Communication, 2007), from 32 or 16 channels per band, with a $7 \text{nV}/\sqrt{\text{Hz}}$ sensitivity, 400 Hz to 6.4 kHz bandwidth, and total dynamic range in excess of 100 dB [Bougeret et al., 1995]. The data are taken by two multi-channel receivers which nominally sample for 20 ms at a 1 MHz sampling rate (see Table 2.3 for more information). The TNR is often used to determine the local plasma density by observing the plasma line, an emission at the local plasma frequency due to a thermal noise response of the wire dipole antenna. One should note that observation of the plasma line requires the dipole antenna to be longer than the local $\lambda_D$ [Meyer-Vernet and Perche, 1989]. For typical conditions in the solar wind, the wire dipole antenna on Wind easily satisfy this condition.

<table>
<thead>
<tr>
<th>Band</th>
<th>Range (kHz)</th>
<th>Sampling Rate (kHz)</th>
<th>Measurement Time (ms)</th>
</tr>
</thead>
<tbody>
<tr>
<td>A</td>
<td>4-16</td>
<td>64.1</td>
<td>320</td>
</tr>
<tr>
<td>B</td>
<td>8-32</td>
<td>126.5</td>
<td>160</td>
</tr>
<tr>
<td>C</td>
<td>16-64</td>
<td>255.7</td>
<td>80</td>
</tr>
<tr>
<td>D</td>
<td>32-128</td>
<td>528.5</td>
<td>40</td>
</tr>
<tr>
<td>E</td>
<td>64-256</td>
<td>1000.0</td>
<td>20</td>
</tr>
</tbody>
</table>
2.1.3 Magnetic Field Instrument

The magnetic field instrument (MFI) on board Wind [Lepping et al., 1995] is composed of dual triaxial fluxgate magnetometers. The MFI has a dynamic range of ±4 nT to ±65,536 nT, digital resolution ranging from ±0.001 nT to ±16 nT, sensor noise level of < 0.006 nT (r.m.s.) for 0-10 Hz signals, and sample rates varying from 44 samples per second (sps) in snapshot memory to 10.87 sps in standard mode. The data are also available in averages at 3 seconds, 1 minute, and 1 hour. The data sampled at higher rates (i.e. >10 sps) will be referred to as High Time Resolution (HTR) data from here on.

2.2 Wind Data Analysis and Calibrations

![Figure 2.3](image.png)

Figure 2.3: This is an example of a TDS sample from the Wind/WAVES instrument. The two panels on the left are the X and Y-antenna measurements of the electric field in the WAVES coordinate system. The hodogram on the right is a plot of $E_x$ vs. $E_y$ with the associated magnetic field rotated into the proper coordinate system. Notice the angle between the X-antenna and the sun-direction ($\theta$ from Figures 2.1 and 2.2) is roughly $-43.96^\circ \leq \theta \leq -45.94^\circ$ over the duration of the TDS sample. The angle of the XY-GSE projection of the magnetic field from the X-GSE direction is roughly $-93.4^\circ$. 
2.2.1 Magnetic Field Rotations

The magnetic field data is typically obtained in GSE coordinates and thus we need to rotate the data into WAVES coordinates, as explained in Section 2.1.2 for polarization analysis. Figure 2.3 shows an example of TDSF sample in 120 kHz sample rate. The event, shown in WAVES coordinates, is an example of a large amplitude solitary wave observed just downstream of an IP shock observed on 2000-04-06. On the left, the top(bottom) panel is a plot of $E_x (E_y)$. To the right is the hodogram, plot of $E_x$ vs. $E_y$, for the region outlined by the red box in the left two panels. Overplotted on the hodogram, with a red line, is the XY-projection of the background DC magnetic field. The angle of the X-antenna from the sun direction varies from $-43.96^\circ$ to $-45.94^\circ$ over the course of the 17 ms TDS sample. In GSE coordinates, the XY-GSE projection of the DC magnetic field is roughly $-93.4^\circ$ away from the sun direction, thus varies away from the X-antenna by roughly $44.8^\circ$ (shown in red in the hodogram). Notice that the bipolar signature of the solitary wave is roughly aligned with the magnetic field, consistent with magnetospheric observations [Ergun et al., 1998a; Franz et al., 2000].

2.3 Particle Data Analysis

2.3.1 Electron Distributions

Due to their much lighter mass than hydrogen and helium, electrons carry the bulk of thermal energy flow away from the sun. Non-Maxwellian aspects of electron distribution functions (eDFs) have been theorized to produce a number of plasma instabilities [Gary et al., 1975; Gary, 1981] and supported by observations [Gurnett et al., 1979b; Thomsen et al., 1983a, 1987; Wilson III et al., 2009]. Thus, it is important to examine the multi-component characteristics of the solar wind eDFs.

Due to the existence of multiple components of electrons in the solar wind, there have been many approaches to modeling the observed eDFs. [Feldman et al., 1983a] modeled the parallel and perpendicular (with respect to the magnetic field) cuts of the electron distribution functions as one dimensional functions (see Section 1.4.2 for more details). Fitting of each component individually tends to be much easier than attempting to simultaneously fit both. The problem with one dimensional fits of the eDFs is
that the amplitudes are inseparable functions of temperature (if anisotropic) as shown by Thomsen et al. [1983a]. Therefore, estimation of temperature components (parallel or perpendicular to the magnetic field) become subject to scrutiny. However, one can still estimate an anisotropy for the electrons by using the ratio of the thermal speeds:

$$\left\| \frac{v_{T_{s,\perp}}}{v_{T_{s,\parallel}}} \right\|^2 \sim \frac{T_{s,\perp}}{T_{s,\parallel}}$$

(2.1)

where \(v_{T_{s,\parallel}}(v_{T_{s,\perp}})\) is the corresponding parallel(perpendicular) thermal speed of species \(s\) (e.g. core) and \(T_{s,\parallel}(T_{s,\perp})\) is the effective parallel(perpendicular) temperatures. The reason that one would like an estimate for the temperature anisotropy of the halo (or core) electrons is that it is a potential free energy source for instability [Gary et al., 1994, 1999].

Figure 2.4 shows an example of an electron distribution seen by the Wind spacecraft in the solar wind with fit lines superposed. The actual data is shown as the asterisks while the fits for the core (blue) and halo (red) are solid lines. The different components of the distribution are outlined by color-coded boxes. The cold core (outlined by blue box in the bottom panel) is modeled as a bi-Maxwellian (Equations 2.3a through 2.4). The hot halo (outlined by the green and gray boxes in the bottom panel) is modeled as a modified Lorentzian (Equations 2.5 and 2.6). In addition to the power-law tail of the halo electrons, Figure 2.4 shows an additional higher energy component of the solar wind electrons known as the super halo [Lin et al., 1996]. Note that the perpendicular cut appears far more symmetric about zero velocity than the parallel cut. Though this observation can be a real phenomena, the symmetry in these cuts is due to an assumed gyrotropy in the creation/calculation of the distribution functions. Regardless, it is generally true that the perpendicular components are more symmetric than the parallel component in the solar wind due to the strahl component (outlined by the green box in the top panel). As previously mentioned in Section 1.4.2 each component of the distributions are thought to provide information regarding their origin/progenitors [Feldman et al., 1975; Lin et al., 1996].

In the analysis done in this thesis, the energies used to estimate the halo and core electron temperatures for all the EL distributions were determined by fitting the core to a Maxwellian velocity distribution and the higher energy halo to a modified Lorentzian Thomsen et al. [1983a]. The point where the Lorentzian begins to dominate the overall
Figure 2.4: An example of the different components of a solar wind electron distribution. The top plot is the parallel cut of the distribution and the bottom plot is the perpendicular cut.
distribution is defined as the break energy, used as the upper bound on the core electrons and the lower bound on the halo electrons. The moments were then calculated directly from the full 3D electron distributions using my modified version of the Wind/3DP SSL IDL Software package (original version provided by R.P. Lin). The strahl component introduces a highly anisotropic peak in the parallel cuts of distribution functions which can increase the difficulty fitting a function to the halo electron distribution. Thus, the strahl electrons were removed in the halo electron fits. The relevant parameters are then calculated from the original electron distributions using the energy bins below the break energy for the core and the energy bins above the break energy for the halo electrons. One should note that the use of energy bin cutoffs instead of the fit functions can lead to increased uncertainty in the estimates of core and halo parameters. The core and halo components overlap in energy, thus one may have core(halo) electrons in their halo(core) moment calculations. These details will be discussed later in Section 2.3.2.

This approach, often used in solar wind studies, is used to find temperatures for each respective component [Feldman et al., 1983a; Thomsen et al., 1983a]. The measured electron moments are used to estimate temperatures, velocities, fluxes, etc. The core and halo components were determined by fitting the data to a Maxwellian (for the core) and a modified Lorentzian or kappa distribution (halo). The distributions are estimated by separately fitting the core and the halo components then adding the resultant distributions together as:

\[ f_e(v) = f_L(v) + f_M(v) \] (2.2)

where \( f_L(f_M) \) represent the modified Lorentzian(Maxwellian) components of the distribution. For the Maxwellian core, we model the parallel and perpendicular components separately [Feldman et al., 1983a; Thomsen et al., 1983a] as individual one dimensional functions in the following manner:

\[ f_{\parallel M}(v) = A_{\parallel M} e^{-\left(\frac{v_{\parallel} - v_o}{v_T\parallel}\right)^2} \] (2.3a)

\[ f_{\perp M}(v) = A_{\perp M} e^{-\left(\frac{v_{\perp}}{v_T\perp}\right)^2} \] (2.3b)

where \( A_{\parallel M}(A_{\perp M}) \) is the parallel(perpendicular) amplitude of the one dimensional Maxwellian, \( v_{\parallel}(v_{\perp}) \) is the parallel velocity, \( v_o \) is the parallel drift velocity (typically
zero because one transforms into the solar wind rest frame), and \( v_{T\parallel}(v_{T\perp}) \) is the parallel(perpendicular) thermal speed. It would be more appropriate to model both the parallel and perpendicular components as one bi-Maxwellian with:

\[
f_M(v) = N_M \left( \frac{m_e}{2\pi T_{M}} \right)^{m_e/2} e^{-\frac{m_e}{2k_B T_{M}} \left( (v_{\parallel}-v_o)^2/T_{\parallel M} + v_{\perp}^2/T_{\perp M} \right)}
\]

where \( N_M \) is the number density, \( m_e \) is the electron mass, \( T_{\parallel M}(T_{\perp M}) \) is the parallel(perpendicular) electron temperature, \( v_{\parallel}(v_{\perp}) \) is the parallel(perpendicular) velocity as in Equations 2.3a and 2.3b and \( v_o \) is the parallel drift velocity is the same as in Equation 2.3a. This is not done in our analysis because we are interested in determining the break energies for the parallel and perpendicular components. Attempting to fit the multi-component anisotropic solar wind electron distributions to a bi-Maxwellian poses a number of problems due to beam-like features and the solar wind strahl component. Thus, we used one dimensional functions to fit the parallel and perpendicular components to find the break energies after fitting the halo components, discussed below.

The hot halo electrons were modeled as a one dimensional modified Lorentzian function [Feldman et al., 1983a]:

\[
f_L(v) = A_L \left[ 1 + \left( \frac{v_{\perp}}{v_{\perp,L}} \right)^2 \right]^{\frac{1}{m}} \left( \frac{m+1}{m} \right) \left( \frac{v_{\parallel}}{v_{\parallel,L}} \right)^{2m} \]

where \( v_{\perp,L} \) are speeds which characterize the break points (i.e. points at which a Maxwellian no longer dominates the distribution) in the distribution perpendicular and parallel to the ambient magnetic field. The exponent, \( m \), characterizes the sharpness of the break and the hardness of the high energy power law tail. Of course, one should model these in two dimensions, as [Thomsen et al., 1983a] did, to examine flat-topped distributions and/or just the halo component of the electron distributions which is shown as:

\[
f_L(v) = A_L \left[ 1 + \left( \frac{v_{\perp}}{v_{\perp,L}} \right)^2 + \left( \frac{v_{\parallel}}{v_{\parallel,L}} \right)^{2m} \right]^{\frac{1}{m}} \left( \frac{m+1}{m} \right) \]

where again \( v_{\perp,L} \) are speeds which characterize the break points in the distribution perpendicular and parallel to the ambient magnetic field and \( A_L \) is the amplitude defined by:

\[
A_L = \frac{N_L m \sin (\pi/(2m))}{\pi^2 v_{\perp,L}^2 v_{\parallel,L}^2}.
\]
One can also use a kappa distribution in two dimensions, as seen here:

$$f(v^2, v_∥) = \left(\frac{\pi}{\kappa}\right)^{-3/2} \frac{v^2 \Gamma(\kappa + j + 1)}{\theta_\perp^{2(j+1)} \theta_\parallel \Gamma(\kappa - 1/2)} \left(1 + \frac{v^2}{\kappa \theta_\perp^2} + \frac{v_\parallel^2}{\theta_\parallel^2}\right)^{-(\kappa + j + 1)}$$  \tag{2.8}

where $\kappa$ is some index, $j$ is an index which indicates the strength of a loss-cone ($j = 0$ indicates no loss-cone), $\Gamma()$ are gamma functions. One should note that as $\kappa \to \infty$, $f(v^2, v_∥) \to$ Dory-Guest-Harris loss-cone bi-Maxwellian distribution $\text{Mace, 1998}$. The parameters $\theta_\perp, \theta_\parallel$ are effective thermal speeds given by:

$$\theta_\perp = \left(\frac{2\kappa - 3}{\kappa}\right)^{1/2} \sqrt{\frac{k_B T_\perp}{(j + 1) m_e}} \tag{2.9a}$$

$$\theta_\parallel = \left(\frac{2\kappa - 3}{\kappa}\right)^{1/2} \sqrt{\frac{k_B T_\parallel}{m_e}} \tag{2.9b}$$

where $T_\perp, T_\parallel$ are the respective temperatures perpendicular and parallel to the ambient magnetic field $\text{Mace, 1998}$. An electron (kappa dependent) plasma beta is defined when using this distribution by:

$$\beta_e = \frac{\theta_\parallel^2 \omega_{pe}^2}{c^2 \Omega_{ce}^2} \tag{2.10a}$$

$$= \left(\frac{\kappa_e - 3/2}{\kappa_e}\right) \frac{n_e T_\parallel}{B_o^2/(2\mu_o)} \tag{2.10b}$$

where $B_o$ is the average magnetic field amplitude. This reduces to the usual $\beta_e$ in the limit as $\kappa_e \to \infty$.

As mentioned earlier, there is a certain type of electron distribution often seen downstream of strong quasi-perpendicular shocks called a flat-topped (or flattop) distribution $\text{Feldman et al., 1983a; Thomsen et al., 1983a}$. These distributions can be modeled with the modified Lorentzian (shown in Equation 2.6) or with a modified exponential given by:

$$f(v) = A_f e^{-\left(\frac{|v_\perp/v_{Th,\perp}|^{5} + |v_\parallel/v_{Th,\parallel}|^{5}}{\kappa_e}\right)} \tag{2.11}$$

This specific distribution is thought to be produced by interactions with strong current-driven ion acoustic waves $\text{Dum et al., 1974; Dum, 1975; Feldman et al., 1983a}$ or due to the cross-shock potential $\text{Scudder et al., 1986a}$. The distribution is referred to as a self-similar distribution and the exponent in its general form should be proportional to
\[(v_{\perp}/v_{Th,\perp})^x\] where \(x \to 5\) in the quasi-linear limit \[\text{Dum et al., 1974}\].

In most cases presented in this thesis, the electron distributions upstream of IP shocks are modeled with one dimensional functions of a bi-Maxwellian core and bi-Lorenztian (typically \(m = 2\) to \(3\)) halo. Most of the flattop distributions were modeled with a self-similar distribution \((x = 4)\) for the core and bi-Lorenztian halo. An automated fit routine was used to gain a zeroth order estimate of the relevant fit parameters (\(e.g. T_{\perp,||} etc.\)) and then corrections were made by hand if necessary. Typically the core and halo fits were adjusted to correct errors caused by the strahl component or beam-like features. Once the functional forms were determined for the respective components, the break energy was determined. The break energy defined the separation between the core and halo electrons in 3DP data structures. Thus, we defined the relevant energy bins for each component, computed the moments of the distribution function, and analyzed the results. Since we were often interested in the respective temperatures, or average kinetic energy per particle, the moment calculations done by the Wind/3DP software was an appropriate approximation to actual temperatures because it directly calculates the average kinetic energy from the second moment. The software does not ignore the strahl nor does it assume any functional form for the data. Calculation of the moments directly from the data yielded more accurate results than fitting cuts of the distribution function to model functions like a bi-Maxwellian.

### 2.3.2 Temperature Anisotropies

Temperature anisotropies are known to be sources of free energy for instabilities in magnetized plasmas \[\text{Gary et al., 1975}\]. They have been examined between 0.3 and 0.8 AU using Helios 1 and 2 data \[\text{Pilipp et al., 1987b}\], in the near Earth solar wind using ISEE data \[\text{Phillips et al., 1989a,b}\], and at IP shocks \[\text{Wilson III et al., 2009}\]. Therefore, estimating an effective temperature anisotropy can allow one to examine the potential to excite waves.

Figure 2.5 presents an example of Wind data used to show the comparison between the modeled distribution function cuts (solid blue lines) and the actual data (* for parallel cut and ♦ for perpendicular cut) for the distribution shown in Figure 1.8. To argue for the use of my modified Wind/3DP SSL IDL Software package in calculating the moments of the distribution function, I compared the temperature anisotropies calculated
(seen in bottom panel) by the software with the values calculated (seen in top panel) from the modeled functions. In every case examined, the modeled distribution functions show a stronger anisotropy (i.e. $T_{\perp}/T_{\parallel}$ is larger) in both the core and halo components than the program estimates. As previously discussed, the larger anisotropies in the fit function estimates result from the neglect of the strahl component when fitting the data. The fit estimates for the distribution in Figure 2.5 are $T_{c,\perp}/T_{c,\parallel} \sim 1.43$ and $T_{h,\perp}/T_{h,\parallel} \sim 0.99$. The computed estimates of the temperature anisotropies from the actual data are $T_{c,\perp}/T_{c,\parallel} \sim 0.99$ and $T_{h,\perp}/T_{h,\parallel} \sim 0.84$.

The consistent difference between modeled fit function and program estimates in temperature anisotropies due to the solar wind strahl (see Figure 2.4) component, skews the parallel temperature in such a way as to reduce the halo (and often core) temperature anisotropy, $T_{h,\perp}/T_{h,\parallel}$. As a consequence, the strahl component is either removed or not included in the modeled fits used in this thesis due to its highly non-Maxwellian behavior and intermittent occurrence. Since instability estimates often ignore the strahl component when examining the solar wind heat flux [Gary et al., 1994, 1999], the fit estimates in Figure 2.5 are relevant to at least the threshold estimates of [Gary et al., 1994] and [Gary et al., 1999]. One should also note that the average temperature estimates for the core electrons tend to be higher for the 3DP software estimates than the model fits because of the strahl contamination.

As previously discussed, the strahl electrons can often skew the anisotropies away from distributions which may be unstable to the whistler heat flux instability [Gary et al., 1994, 1999]. Thus, when calculating the temperatures of both the core and halo components using the 3DP software, one finds that $T_{h,\perp}/T_{h,\parallel}$ is lower than from fit estimates. As a consequence, the 3DP software estimates were used as a lower bound on the temperature anisotropy estimates by Wilson III et al. [2009] when comparing to the instability thresholds of [Gary et al., 1994].

2.3.3 Ion Distributions

The Pesa Low detector is a spherical top-hat ES particle detector which covers energy ranges typically between $\sim 700$ eV to $\sim 6$ keV when in the solar wind. The detector is designed to detect the solar wind beam and thus has a narrow field of view. It is also
Figure 2.5: An example of the parallel and perpendicular cuts of the distribution function. The asterisks in the top panel and diamonds in the bottom panel are the real parallel and perpendicular cuts from the bottom panel of Figure 1.8. The solid lines are the composite fit examples. The top panel shows the solar wind strahl as the asymmetry anti-parallel to the magnetic field.
prone to saturation and innaccurate measurements if the solar wind parameters change abruptly, like in IP shocks. If the ions get too hot, the detector can saturate as well. However, for most of the duration of Wind’s experience in the solar wind the detector measures nearly all of the solar wind beam.

The Pesa High detector is a spherical top-hat ES particle detector which covers energy ranges typically between $\sim 500 \text{ eV}$ to $\sim 28 \text{ keV}$ when in the solar wind. The detector has a relatively high threshold at low energies making it difficult to detect background ions below $\sim 500 \text{ km/s}$ unless there is a well defined peak. In other words, below $\sim 500 \text{ km/s}$ the one-count levels of the detector, which scales as $E^{-2}$, are too high to observe the particles often found in this energy range. At the lowest energies, below $\sim 200 \text{ km/s}$, one can see the thermal core (for an example, see the bottom two panels of Figure [A.1]). The thermal core often saturates the detector and is often contaminated with UV light. The detector is not, however, by any means useless. It offers a full $4\pi$ steradian coverage every $\sim 3$ seconds in burst mode and has been used to detect gyrating ion distributions upstream of the terrestrial bow shock and their association with low frequency waves [Meziane et al., 2001]. However, with any instrument there are idiosyncracies which must be regarded before interpreting data. As was previously mentioned, the low energies of the Pesa High detector have relatively high one-count levels. Thus, if the particles in that low energy population are shifted to higher energies by say, a convectional electric field, they can appear as an asymmetric particle distribution function even though the distribution may really be symmetric.

2.4 Waveform Analysis

2.4.1 Minimum Variance Analysis

Minimum variance (MV) analysis utilizes a property of plane polarized linear electromagnetic waves which allows one to assume that fluctuations in the electric ($\delta E$) and magnetic ($\delta B$) fields are are in a plane orthogonal to the direction of propagation ($\hat{k}$) [Khrabrov and Sonnerup, 1998]. Using a plane wave assumption for $\nabla \cdot B = 0$ results in the linearized equation $\hat{k} \cdot \delta B = 0$. The analysis is performed by minimizing the
variance matrix of the magnetic field given by:

\[ S_{pq} = \left\langle \left( B_p - \langle B_p \rangle \right) \left( B_q - \langle B_q \rangle \right) \right\rangle \]  

(2.12)

where \( \langle B_p \rangle \) is the average of the \( p \)th component of the magnetic field. We assume \( S_{pq} \) to be a non-degenerate matrix with three distinct eigenvalues, \( \lambda_3 < \lambda_2 < \lambda_1 \), and three corresponding eigenvectors, \( e_3, e_2, e_1 \). Thus the minimum variance eigenvalue and eigenvector are \( \lambda_3 \) and \( e_3 \). The propagation direction is assumed to be along \( \hat{e}_3 \) if one assumes small isotropic noise and the condition \( \lambda_2/\lambda_3 \geq 10 \) is satisfied. Then the uncertainty in this direction is given by [Kawano and Higuchi 1995]:

\[ \delta \hat{k} = \pm \left( \hat{e}_1 \sqrt{\frac{\delta \lambda_3}{\lambda_1 - \lambda_3}} + \hat{e}_2 \sqrt{\frac{\delta \lambda_3}{\lambda_2 - \lambda_3}} \right) \]  

(2.13)

where \( K \) is the number of vectors used and \( \delta \lambda_3 \), the uncertainty in the \( \lambda_3 \) eigenvalue, is given by:

\[ \delta \lambda_3 = \pm \lambda_3 \sqrt{\frac{2}{K-1}} . \]  

(2.14)

In general, the uncertainty of \( \delta \lambda_i \) is given by:

\[ \delta \lambda_i = \pm \sqrt{\frac{2\lambda_3(2\lambda_i - \lambda_3)}{(K-1)}} . \]  

(2.15)

Another useful quantity to know is the angle between the local ambient magnetic field and the propagation direction, \( \theta_{kB} \). This can be calculated in the typical manner, \( \theta_{kB} \equiv \cos^{-1} \left( \hat{k} \cdot \hat{b} \right) \), with associated uncertainties of:

\[ \delta \theta_{kB} = \pm \sqrt{\frac{\lambda_3 \lambda_2}{(K-1)(\lambda_2 - \lambda_3)^2}} \]  

(2.16)

MV analysis can be done on specific time ranges to determine the wave vector, \( \mathbf{k} \), and the polarization. The use of bandpass filters can increase the accuracy of MV analysis by removing low frequency compressional components (typical of magnetosonic whistler waves in the solar wind). Using the wave vector from MV analysis, the angle of propagation for each wave with respect to the shock normal vector, \( \theta_{kn} \), upstream averaged solar wind velocity, \( \theta_{kV} \), and the magnetic field, \( \theta_{kB} \), were examined. These angles are determined by the following equation:

\[ \theta_{kj} \equiv \cos^{-1} \left( \frac{\mathbf{k} \cdot \hat{j}}{||j||} \right) \]  

(2.17)
where \( j \) can be \( b, n \), or \( V_{sw} \).

In our use of MV analysis, we define the eigenvalues of the spectral matrix, from minimum to maximum, as \( \lambda_3, \lambda_2, \) and \( \lambda_1 \). As a general rule for determining whether the MV analysis has yielded a well determined plane circularly polarized wave, we require \( \lambda_2/\lambda_3 \geq 10.0 \) and \( \lambda_1/\lambda_2 \sim 1.0 \) if less than 50 field vectors were used in the analysis. For the case where \( \lambda_2/\lambda_3 \geq 10.0 \) but \( 1.0 < \lambda_1/\lambda_2 \ll \lambda_2/\lambda_3 \), the wave is elliptically polarized. If \( \lambda_1/\lambda_2 \gg \lambda_2/\lambda_3 \sim 1.0 \) then the wave is linearly polarized and the \( \mathbf{k} \)-vector cannot be trusted. These assumptions hold for data with small isotropic Gaussian noise [Khrabrov and Sonnerup, 1998]. Single satellite measurements introduce another complication. Though the plane orthogonal to the \( \mathbf{k} \)-vector may be well determined, the sign of the vector cannot be known without at least one component of the electric field or another satellite measurement [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983].

### 2.4.2 Bandpass Filters and Spectral Analysis

High time resolution (HTR) magnetic field data were obtained from the Wind spacecraft within ±1 hour of the IP shocks analyzed. The HTR MFI data used in this thesis were sampled at two different rates: \( \sim 22 \) samples/s (for 04/03/1996, 04/08/1996, 12/02/1996, 01/10/1997, and 10/24/1997) and \( \sim 11 \) samples/s (02/27/1997, 12/10/1997, 04/23/1998, 04/30/1998, 05/15/1998, 08/26/1998, 02/11/2000, 02/20/2000, and 04/06/2000). HTR magnetic field data, was used to define the ramp region, or transition region, of the IP shock (see Figure 1.3 for visual reference). The wave vector and other wave properties were determined using Minimum Variance (MV) analysis [Khrabrov and Sonnerup, 1998], as described in detail in Section 2.4.1. The magnetic field fluctuations were identified and analyzed using a bandpass filter, as shown in Figure 2.7.

The frequency ranges for each bandpass filter was determined from spectral analysis. Spectral analysis is simply an examination of the data in Fourier space as seen in Figure 2.6 which can illustrate where the greatest power is located with respect to frequency. Finding a frequency peak, as seen in the black box of Figure 2.6, can help one identify the particular wave mode being examined or other processes occurring in the plasma. Spectral analysis was used to determine an approximate range of frequencies to use for the initial guess of the bandpass filter frequencies then refined by maximizing the ratio.
Figure 2.6: A plot of the magnetic field power spectrum as a function of frequency for 16:27:12-16:28:45 UT on 04/06/2000. The three colors represent the X (red), Y (green), and Z (blue) GSE components of the magnetic field. Black box illustrates the range of frequencies used for the bandpass filter in panel A of Figure 4.4.

of intermediate to maximum eigenvalues, $\lambda_2/\lambda_3$. The data is plotted and adjustments are made to the frequency range of the filter until the wave envelope resembles something similar to the three higher frequency filters seen in Figure 2.7.

Figure 2.7 is an example of four different bandpass filters being applied to the magnetic field data shown in the top panels measured by the Wind spacecraft upstream of an IP shock. The example shown is a plot of a shocklet (see Section 3.2.1) observed roughly five minutes upstream of the 2000-04-06 event discussed in detail in Chapter 4. All six panels plot the same time range and the upper left is a plot of the magnetic field magnitude and upper right are the XYZ-GSE components. For this event, the MFI data was sampled at $\sim 10$ Hz. The bottom two rows are plots of the resulting XYZ-GSE components after filtering the data. The lower frequency components (upper left panel
Figure 2.7: An example plot of bandpass filtered HTR MFI data on 04/06/2000. The top two panels have not been filtered and represent the magnitude and vector components near a shocklet upstream of the 04/06/2000 IP shock event. In every vector plot, red is the X-GSE direction, green the Y-GSE, and blue the Z-GSE directions. The frequencies at which each of the bottom four panels were filtered is labeled in each plot. The relative magnitude of the fluctuations are also labeled with the vertical black arrows in each filtered plot.

of filtered signals) often have larger amplitudes and in the case of dispersive waves like shocklets, have different spatial locations than their high frequency counterparts. The spatial dispersion can be seen by comparing the upper left filtered panel to the three other filtered data panels. Notice that the high frequency wave envelope occurs roughly 30 seconds ahead of the lower frequency amplitude peak in the upper left filtered data plot.

Figure 2.8 is an example of MV analysis done on low frequency magnetosonic whistler waves upstream of an IP shock. The left hand set of wave events plot the GSE (gray scale) components, the middle panels plot the MV (color scale) components, and the right hand set of panels show the hodograms, $B_y$ vs. $B_x$, $B_z$ vs. $B_x$, and $B_z$ vs. $B_y$. The time for each wave event in Figure 2.8 is defined by the color-coded boxes outlined in
Figure 2.8: An example of MV analysis on the leading whistler waves of a low frequency magnetosonic whistler wave, called a shocklet (discussed later). The frequency ranges and angles of propagation are: $f > 0.6$ Hz and $\theta_{kB} = 41^\circ (139^\circ)$ for A, $0.25$ Hz $< f < 0.45$ Hz and $\theta_{kB} = 38^\circ (142^\circ)$ B, $f > 1.0$ Hz and $\theta_{kB} = 41^\circ (139^\circ)$ for C, and $0.6$ Hz $< f < 3.0$ Hz and $\theta_{kB} = 27^\circ (153^\circ)$ for D. The eigenvalue ratios from the MV analysis are also shown with the MV estimate of the $k$-vector direction in GSE coordinates above each hodogram. The purple arrows indicate the direction of rotation for each respective plot.
the center plot of the magnetic field magnitude and GSE components. The difference in polarization, seen in the hodograms, between wave events A and C in Figure 2.8 can be explained by projection effects due to single satellite measurements using only magnetic field measurements [Hoppe et al., 1981]. In the spacecraft frame, wave events A and C in Figure 2.8 are LH polarized with respect to the propagation direction, but all wave events in Figure 2.8 show a RH sense with respect to the magnetic field, characteristic of whistler modes. Wave events B and D in Figure 4.3 are RH polarized both with respect to the wave vector and the magnetic field.

2.4.3 Wavelet Analysis

Wavelet transforms are square integrable basis functions used as a mathematical technique to analyze waveform data in both space and scale [Farge, 1992; Hudgins et al., 1993; Lau and Weng, 1995; Torrence and Compo, 1998a]. Wavelets are basis functions based on group theory and are used for multiple purposes including signal analysis, image coding, and renormalization theory in quantum electrodynamics [Farge, 1992]. Traditionally Fast Fourier Transforms or FFTs have been used to analyze electric and magnetic field data in space physics. Wavelets have only recently become a popular analysis tool, and a great deal of confusion and apprehension surround this mathematical tool.

Let us first consider the traditional waveform analysis technique, the Fourier transform. Fourier transforms are a mathematical technique where one attempts to reconstruct an input signal using the superposition of a semi-infinite or infinite set of trigonometric functions. There are advantages and disadvantages to using trigonometric functions in signal decomposition. In particular, a disadvantage is that the basis functions oscillate forever. Thus any localized signal, which we’ll define as a data glitch of non-continuous behavior, in space will be spread out over all Fourier coefficients. One should also note that any signal which has an abrupt change in amplitude with respect to time, such as a solitary wave, would show a similar result on a Fourier transform. This results in an inability to filter out the glitch [Farge, 1992]. For reference, we define local as the subdomain of time about which the transform is being applied. For instance, let us assume we have some input signal, f(t), which depends on time and has an equal time spacing of \( \delta t \). The generalized Fourier series of f(t) is given by:
\[ f(t) = \frac{1}{2} a_0 + \sum_{n=1}^{\infty} a_n \cos nt + \sum_{n=1}^{\infty} b_n \sin nt \]  

(2.18)

where \( a_n \) and \( b_n \) are the Fourier coefficients of \( f(t) \). The Fourier transform of \( f(t) \) is given by:

\[ F(\omega) = \sqrt{\frac{1}{2}} \int_{-\infty}^{\infty} f(t) e^{-i\omega t} dt \]  

(2.20)

We return to the example of a glitch in a finite time series. One can see that there is no local information stored in the Fourier coefficients defined in Equations 2.19a through 2.19c. The main consequence is that the glitch is found in the phase of every Fourier coefficient. In contrast, wavelet transforms intrinsically retain localized information allowing one to isolate and remove data glitches or more clearly analyze real signals that are localized.

In 1980 a French researcher, named Jean Morlet, working on seismic data collaborated with a theoretical physicist, named Grossmann, to develop a geometrical translation and dilation [Farge, 1992]. They were interested in decomposing seismic signals into both space and scale, an analysis technique of extremely useful potential. The wavelet transform removes the behavior of a signal at infinity, which an FFT cannot do. The wavelet coefficients retain the locality of the input signal allowing for a reconstruction of the original input signal locally. Meaning, if there is a glitch at one point in an otherwise smooth/continuous input signal, the amplitude of the wavelet coefficients will increase near the glitch in time and frequency, but be small elsewhere. This allows one to reconstruct the wavelet locally without requiring the entire transform.

I will focus my efforts on the most commonly used wavelet transform in this thesis,
the un-normalized Morlet wavelet transform \cite{Morlet et al. 1982, Morlet 1982} defined by:

\[
\psi_o(\eta) = \pi^{-1/4} e^{i\omega_o \eta} e^{-\eta^2/2}
\] (2.21)

where \(\omega_o\) is a non-dimensional frequency (taken to be 6 to satisfy the admissibility condition set by the uncertainty principle) and \(\eta\) is a non-dimensional time parameter. When we scale and translate \(\psi_o(\eta)\), the wavelet becomes a function of scale, \(s\), and time, \(n\delta t\). The Fourier transform of the scaled and translated wavelet is given by:

\[
\hat{\psi}(s\omega_k) = \left(\frac{2\pi s}{\delta t}\right)^{1/2} \hat{\psi}_o(s\omega_k)
\] (2.22)

where \(s\) is the wavelet scale, \(\delta t\) is the sample period, and \(\omega_k\) is defined by:

\[
\omega_k = \begin{cases} 
\frac{2\pi k}{N\delta t} & \text{for } k \leq N/2, \\
-\frac{2\pi k}{N\delta t} & \text{for } k > N/2.
\end{cases}
\]

Every unscaled \(\psi_o\) must have unit energy to satisfy the admissibility condition seen as:

\[
\int_{-\infty}^{\infty} |\hat{\psi}_o(\omega')|^2 d\omega' = 1
\] (2.23)

which results in there being a factor \(N\) in the normalization shown by:

\[
\sum_{k=0}^{N-1} |\hat{\psi}(s\omega_k)|^2 = N
\] (2.24)

where \(N\) is the number of points in the input array. The Fourier transform of \(\psi_o\) is defined as:

\[
\hat{\psi}_o(s\omega) = \pi^{-1/4} H(\omega) e^{-(s\omega - \omega_o)^2/2}
\] (2.25)

where \(H(\omega)\) is the Heaviside step function, \(H(\omega) = 1\) for \(\omega > 0\), \(H(\omega) = 0\) otherwise. Thus, if one has an input time series \(x_n\), with a uniform sample rate of \(1/\delta t\), and time index \(n = 0, 1, ..., N-1\). We define the Fourier transform of \(x_n\), \(\hat{x}_k\), as:

\[
\hat{x}_k = \frac{1}{N} \sum_{n=0}^{N-1} x_n e^{-2\pi i kn/N}
\] (2.26)

where \(k = 0, 1, ..., N-1\) is the frequency index. From this, we can construct the wavelet transform as a function of wavelet scale, \(s\), to be:

\[
W_n(s) = \sum_{k=0}^{N-1} \hat{x}_k \hat{\psi}^*(s\omega_k) e^{i\omega_k n\delta t}
\] (2.27)
where $\omega_k$ is defined as in Equation 2.4.3 \cite{Torrence_and_Compo_1998a}.

Figure 2.9 is an example image of a Morlet wavelet transform. The figure is produced from a real TDS event, but it is used here to illustrate the format of the wavelet transforms presented later. The color scale is a relative intensity of the wave mode as a function of frequency versus time. The plot was produced using software produced by \cite{Torrence_and_Compo_1998b}. There are two significance tests shown in Figure 2.9: the cone of influence and the 95% confidence level. The cone of influence is the bowl-shaped line that outlines the region of the plot where edge effects become important due to the finite sample time. The 95% confidence level is a statistical test of significance comparing the relative amplitudes of the wave mode at different frequencies and times. The 95% confidence level is seen as the contour outlining the most intense part of the wavelet transform. It represents the region of the data composed of the top 5% of all the amplitudes in the wave. Meaning, inside the contour for a particular time, all the amplitudes are larger than 95% of the other amplitudes at any other frequency outside of the contour. Finally, the two horizontal lines at the bottom of the wavelet transform represent the electron cyclotron frequency and its first harmonic.

Figure 2.10 is an illustrative comparison between a Morlet wavelet transform and
three different size windowed FFTs. The windowed FFT was originally implemented to force a locality into the time/frequency decomposition of time series analysis prior to wavelet transforms [Hudgins et al., 1993]. All four spectral power density plots are plotted on the same scale. The three different sized windows are shown as the magenta colored boxes on the left-hand side of the image. The sizes are labeled at the top of each windowed FFT panel as a fraction of the total time range of the input signal. One should note that there is no common standard for defining a useful window width for a windowed FFT. Given the arbitrary choice in window width, one can see that the interpretation of an input signal can vary drastically depending on the choice in size of FFT window. For instance, the 1/64th window is so narrow, it fails to resolve the actual frequency of the input sawtooth signal. As with the wavelet transform, there is a balance between the number of frequency bins and time steps for a windowed FFT. The ultimate lower bound is set by the uncertainty principle [Torrence and Compo, 1998a].

Figure 2.10 illustrates how an arbitrary choice in window size could lead to a dramatically different interpretation of the waveform. The Morlet transform, however, very distinctly resolves the fundamental frequency of the sawtooth waveform. Due to the advantages listed above and the clear difference in the resolving power of the Morlet wavelet over a windowed FFT (seen in Figure 2.10), I will primarily rely on Morlet wavelets throughout this thesis for dynamic waveform analysis.
Figure 2.10: An example of a sawtooth wave analyzed with a Morlet wavelet transform and three different windowed FFTs. The relative floating FFT window sizes are shown on the left-hand side by the magenta colored boxes. The actual frequency of the sawtooth wave is shown as the white line in the top wavelet transform panel (and subsequent FFT panels).
Chapter 3

Previous Studies of Low Frequency Waves and Macroscopic Fields

This chapter presents an overview of previous work done on low frequency (<20 Hz) magnetic fluctuations and macroscopic (scales > \( \rho_g \)) electromagnetic fields in the solar wind and planetary foreshocks. The purpose of this chapter is to provide some background for arguments and analyses made in Chapters 4 and 5. The discussion of low frequency waves presents a history of observations with detailed figures to serve as illustrative examples. The discussion regarding macroscopic fields in shock waves highlights the work of previous studies and presents arguments for considering higher frequency fluctuations when examining the shocked particle distribution functions.

3.1 Introduction to Low Frequency Waves at Collisionless Shocks

Irregular turbulence upstream of planetary bow shocks has been observed for over 40 years [Fairfield, 1969]. These waves are thought to play an intrinsic role in particle acceleration, heating, and energy dissipation in collisionless shock waves by allowing the shock to communicate with the upstream plasma [Hada et al., 1987; Stasiewicz et al.].
The frequencies of this turbulence were often observed to fall in the range of the ion cyclotron frequency. Thus, the magnetic turbulence upstream of the Earth’s bow shock was initially examined in association with ion particle data \cite{Paschmann1981, Sentman1981}. Waves in the frequency range around the ion cyclotron frequency are important dynamically for shock formation and stability. By radiating energy away from the shock itself, both transverse Alfvénic and compressional magneto-sonic modes allow a collisionless shock to communicate with the upstream plasma by pre-heating or decelerating the incoming plasma, thus altering the Rankine-Hugoniot conditions \cite{Scholer1971}. The waves in this frequency regime upstream of collisionless shocks are thought to play an intrinsic role in particle acceleration, heating, and energy dissipation \cite{Hada1987, Stasiewicz2003}.

3.1.1 ULF Foreshock Waves

\cite{Fairfield1974} separated the magnetic turbulence observed upstream of the terrestrial bow shock into two categories using Imp 6 magnetometer data, due to possible generation mechanisms and polarization: low frequency (0.01-0.05 Hz) and high frequency (0.5-4.0 Hz), which will be discussed in the next section. We will briefly discuss the 0.01-0.05 Hz class of waves in this section but this thesis does not focus on them. However, they are an important aspect of collisionless shock dynamics so we cannot ignore their possible impacts on the energy dissipation occurring in the shock.

The 0.01-0.05 Hz class of waves discussed by \cite{Fairfield1974} has been extensively studied. They were observed to be mostly transverse, with $\delta B \sim$ several nT, and a left-hand polarization in the spacecraft frame. These waves were associated with the bow shock since they are only observed on magnetic field lines connected to the bow shock. \cite{Hoppe1983} studied ultra-low frequency (ULF) waves using the ISEE 1 and 2 spacecraft. The waves were observed to be associated with intermediate and diffuse ions, and \cite{Hoppe1983} classified them as mixtures of transverse Alfvénic and compressional magneto-sonic modes with rest frame frequencies $\sim 0.1 \Omega_{\text{cp}}$ ($\sim 0.005$-0.01 Hz in solar wind) and wavelengths $\sim 6000$ km. More recently, \cite{Eastwood2005a, Eastwood2005b} examined a specific class of foreshock ULF waves they called the 30 second quasi-monochromatic wave using multi-spacecraft observations from Cluster. They
found, through the use of cross-correlation and MVA that the waves had wavelengths \( \sim 600 \text{ km} \), right(left)-hand polarized in plasma(spacecraft) frame, and \( \theta_{kB} = 21^\circ \pm 14^\circ \).

They also observed that when the angle between the interplanetary magnetic field (IMF) and the X-GSE direction were large, the wave vectors were confined to the plane defined by the IMF and X-GSE vectors. They could not conclude why the waves were observed to propagate obliquely to the magnetic field, however the waves were typically observed simultaneously with intermediate or gyrophase bunched ion distributions. The 30 second ULF waves are thought to be driven unstable by the interaction of backstreaming (with respect to the solar wind flow) ions and incident solar wind ions through the right-hand ion-ion cyclotron beam instability [Gary, 1993]. From all of these observations, Eastwood et al. [2005a] concluded that the 30 second quasi-monochromatic foreshock ULF wave is actually a kinetic fast magnetosonic ULF wave. More recent studies with higher resolution particle instruments have found that the ULF wave boundary of the foreshock coincides with an inner (i.e. closer to Earth) boundary of field-aligned ion beams [Meziane et al., 2004a], consistent with theory [Gary, 1993].

Such low frequency waves have been observed upstream of the terrestrial bow shock by Archer et al. [2005], Balikhin et al. [1997a, b, 2003], Bavassano-Cattaneo et al. [1983, 2003, 2004], Behlke et al. [2003, 2004], Eastwood et al. [2005a, b], Hoppe et al. [1981, 1982], Hoppe and Russell [1980, 1983], Russell et al. [1971], and Schwartz et al. [1992], the Mercurian bow shock by Orlowski et al. [1990], the Venusian bow shock by Orlowski and Russell [1991] and Orlowski et al. [1994], the Martian bow shock by Brain et al. [2002], the Jovian bow shock by Bavassano-Cattaneo et al. [1987], Bertucci et al. [2005], and Sonnerup et al. [1981], the Saturnian bow shock by Bertucci et al. [2007] and Orlowski et al. [1992], the Uranian bow shock by Smith et al. [1989, 1991], and upstream of comets by Brinca and Tsurutani [1989], Galeev et al. [1989, 1991], Kennel et al. [1986], Klimov et al. [1986], Le et al. [1989], Tsurutani et al. [1987, 1989a, b]. Their ubiquitous nature, potential to alter shock parameters, and contribution to energy dissipation make these waves of extreme interest in collisionless shock studies. We will not focus on these low frequencies in this thesis, but we acknowledge their importance in shock wave dynamics.
3.2 Shocklets, SLAMS, and Whistler Foreshock Waves

In the previous section we highlighted the studies done on some of the lowest frequency waves of interest connected to energy dissipation and wave-particle interactions in the foreshock region of collisionless shocks. This section will examine some of the higher frequency (still $< 10$ Hz) foreshock waves often observed in the foreshocks of bow shocks and IP shocks. Initial observations found these waves to be composed of multiple wave modes and properties making them more difficult to analyze than their lower frequency counterparts [Fairfield, 1974]. Here we will highlight some previous work and examine the importance of each wave mode in the 0.5-4.0 Hz frequency range of foreshock waves.

Fairfield [1974] initially found that the 0.5-4.0 Hz waves had a left-hand polarization below 2.5 Hz and right-hand above 2.5 Hz. This led them to conclude that the waves were all right-hand polarized waves in the plasma frame, but the waves with frequencies below 2.5 Hz were being Doppler shifted causing a reversal of polarization. The Doppler effect is observed as a reversal of polarization in the spacecraft frame because the waves with phase velocities slower than the solar wind are convected back over the spacecraft. The wave frequencies in the plasma frame, $\omega_o$, are observed as $\omega' = \omega_o + k \cdot V_{sw}$ in the spacecraft frame. A reversal of polarization occurs when $\omega' < 0$, assuming $\omega_o > 0$, which implies that $k \cdot V_{sw} < 0$ and $|k \cdot V_{sw}| > \omega_o$. Since whistler waves are known to have higher phase velocities at higher frequencies in this frequency range, the corresponding right-hand polarization of the higher frequency waves observed by Fairfield [1974] led him to conclude that the observed modes were in fact Doppler shifted whistler waves. Fairfield [1974] determined, using cold plasma dispersion, that the waves were propagating at angles between $20^\circ \leq \theta_{kB} \leq 40^\circ$ and their wavelengths were $\sim 100$ km.

Later studies began to find that the 0.5-4.0 Hz waves were actually composed of two different waves with slightly different characteristics. In the 0.5-4.0 Hz frequency range, the lower frequencies were found to have a larger amplitude and more compressive component than the higher frequency end of this frequency range. Both modes were observed to have a left-hand (LH) polarization in the spacecraft (SC) frame, but a RH polarization in the plasma rest frame [Hoppe et al., 1981, 1982]. The higher frequency modes were also shown to be a whistler wave, but were observed farther upstream of the
bow shock in association with ion beams [Hoppe et al., 1981, 1982] (see Section 1.7.4).

There has been some confusion regarding the terminology surrounding the lower fre-

Figure 3.1: The irregular low-frequency fluctuations in this image are examples of the multi-frequency shocklets observed by [Russell et al., 1971] using the OGO 5 spacecraft. This is adapted from Figure 3 in [Russell et al., 1971].

quency larger amplitude mode in this frequency range. The waves were first described by [Russell et al., 1971] and called a discrete wave packet. Figure 3.1 is an adapted image from [Russell et al., 1971] showing a few examples of what they referred to as discrete wave packets observed by OGO 5 upstream of the terrestrial bow shock. The top line corresponds to the magnitude of the magnetic field and the next three lines are the spacecraft coordinate components of the magnetic field. The image is roughly three minutes of data sampled at 56 samples per second. The waves are compressive and propagating at large $\theta_{k_B}$'s (e.g. $>30^\circ$).

Later, [Hoppe et al., 1981; Stasiewicz et al., 2003; Thomsen et al., 1990] all referred to these same magnetic fluctuations as shocklets. Figure 3.2 is an adapted image from [Hoppe et al., 1981] showing four examples of shocklets. All four panels show the three spacecraft components of the magnetic field followed by the magnitude of the magnetic
Figure 26: Individual examples of the low-frequency structures of November 3, 1977, seen plotted on our standard scale in Figure 21, are presented on expanded time scales. The plots present the three vector components (in spacecraft coordinates) and magnitude of the magnetic field measured at ISEE 1.

Figure 3.2: Examples of shocklets with (D) and without (A-C) leading high frequency wave packets. This is adapted from Figure 26 in [Hoppe et al. 1981].

...field (nT). The different panels show examples of shocklets with (D) and without (A-C) leading whistler wave trains. As one can see, the wave modes can exhibit multiple characteristics which led to confusion in the literature (e.g. [Lucek and Balogh 1997] referred to similar fluctuations as discrete wave packets while [Schwartz et al. 1992] called them ULF Waves).

Further examination of foreshock waves led [Schwartz and Burgess 1991] and [Schwartz et al. 1992] to identify another type of compressive wave they defined as short large-amplitude magnetic structures (SLAMS). They readily admit that the only identifiable difference (in their data set) between shocklets (what they call foreshock ULF waves) and SLAMS is the amplitude of SLAMS tends to be much greater than that
of shocklets. Regardless, a cursory view of the various examples used in the literature suggests strong similarities between these phenomena, though they have been given multiple names. We will discuss, in detail, all of these foreshock modes in detail in the following sections.

3.2.1 Shocklets

Hoppe et al. [1981], using the ISEE 1 and 2 satellites, identified the structures they called shocklets as ULF magnetosonic waves associated with a leading magnetosonic whistler wave train, though they have been observed without a whistler wave train [Hoppe et al., 1981; Le et al., 1989]. The whistler wave train is thought to result from wave spreading due to the difference in phase velocities between lower and higher frequency whistler waves. Thus, the higher frequency waves can out-run the lower frequency waves leading to spatially dependent frequency spectrum within the shocklets [Omidi and Winske, 1990]. Observations found that the high frequency whistler wave train on the leading edge of the shocklets have rest frame frequencies of $0.1 < \omega/\Omega_{cp} < 40$ ($\sim 0.001-0.4$ Hz in solar wind), wavelengths of $30 \text{ km} < \lambda < 2100 \text{ km}$ and propagation angles of $\theta_{kB} \sim 20^\circ-30^\circ$ [Russell et al., 1971; Hoppe et al., 1981].

Shocklets appear to be a commonly observed magnetic fluctuation upstream of quasi-parallel bow shocks and very rarely observed upstream of IP shocks. They have been observed upstream of the quasi-parallel terrestrial bow shock [Hoppe et al., 1981], the Saturnian and Jovian bow shock [Bertucci et al., 2005], cometary foreshocks [Le et al., 1989; Tsurutani et al., 1987, 1990a,b], one observation at a quasi-parallel IP shock [Lucek and Balogh, 1997], and 12 shocklets were observed upstream of one supercritical quasi-perpendicular IP shock [Wilson III et al., 2009] (see Chapter 4). Planetary foreshock studies [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983; Thomsen et al., 1985b, 1990] and one IP shock study [Wilson III et al., 2009] have shown shocklets to occur in association with diffuse ion distributions, suggesting such ion distributions may be a necessary component for shocklet formation.

Figure 3.3 shows four examples of shocklets observed upstream of a quasi-perpendicular IP shock on 2000-04-06. In each panel, the top plot is the magnitude of the magnetic field and the bottom plot shows the three GSE components of the magnetic field (X-Red, Y-Green, and Z-Blue). All plots are shown in units of nanotesla.
Figure 3.3: Four examples of shocklets upstream of an IP shock on 04/06/2000. The IP shock was observed by Wind at 16:32:09.237 UT. The shocklets are labeled with the various substructures observed simultaneously. The first example in panel A is the only example here without the leading high frequency magnetosonic whistler.

These plots highlight the various forms shocklets may exhibit. For instance, shocklet A is an example without a leading high frequency ($\gtrsim 1$ Hz) whistler. It is observed further upstream from the three other events, but under similar plasma densities and magnetic field strengths. The increase in magnetic field strength in all four shocklets is observed concurrently with an increase in plasma density (see Figure 4.1), consistent with magnetosonic waves. Shocklets C and D are examined in greater detail in Chapter 4.

Shocklets have been examined in association with high frequency electric fields
finding the downstream region (earthward side) of the shocklet to contain strong ion-acoustic-like emissions [Stasiewicz et al., 2003]. In the foot (sunward side) of the shocklets, high frequency plasma emissions were observed. Both wave types are capable of heating the plasma, thus possible dissipation mechanisms for the entire shocklet structure. In Chapter 4 I will discuss the detailed analysis of 12 shocklets observed upstream of a quasi-perpendicular shock and their effects on particle distributions.

3.2.2 SLAMS

In 1991, [Schwartz and Burgess, 1991] proposed a new framework for explaining the complex and diverse phenomena observed upstream of the quasi-parallel bow shock, in which they coined the acronym SLAMS. They proposed that ULF waves grow out of background fluctuations and steepen into the structures referred to as SLAMS. SLAMS have a short duration, \~10 seconds, are generally embedded in the other foreshock waves though they can be isolated, have wavelengths \~0.5 R_E (or half that of the upstream waves observed with them) along the bulk flow direction, generally share the polarization characteristics of the ULF waves surrounding them (RH-polarized in SC frame but LH in rest frame), have phase speeds which increase(decrease) with amplitude in the rest(SC) frame, and amplitudes which are generally larger than what is expected from adiabatic compression (i.e. ratio of magnetic field strengths exceeds four) [Schwartz and Burgess, 1991; Schwartz et al., 1992].

Figure 3.4 shows two examples of SLAMS observed upstream of the terrestrial bow shock observed by the Wind spacecraft. The top panel plots the magnitude of the magnetic field and the bottom panel plots the three GSE components. Red boxes are used to outline the two SLAMS in both panels. Notice that the difference between the maximum value of the magnetic field in the SLAMS and the average value upstream of the SLAMS exceeds 15 nT in both cases. To contrast, the same difference for the shocklets observed in Figure 3.3 rarely exceeds 5 nT.

Since the introduction of the term SLAMS, a great deal of effort has been put forth to distinguish between characteristics of structures given different names for a phenomena which appear to result from the same source. [Lucek et al., 2002], using data from the four Cluster spacecraft, concluded that SLAMS have shorter scale lengths than
Figure 3.4: The figure above is an example of a bow shock crossing by Wind on 12/02/1996 at 16:56:20.015 UT. Two examples of SLAMS are seen upstream of the bow shock. Note the magnitude of the change of the magnetic field from the average upstream values to the maximum value in SLAMS A, B, and the shock itself are $\sim$15 nT, $\sim$19 nT, and $>$40 nT respectively. The ratio of the maximum field magnitude to the upstream averages are $\sim$2.3, $\sim$2.7, and $>$4 for the same three events.

shocklets and that the two structures had many similarities. Others have simply analyzed the data to determine relevant and critical properties useful for identification and/or differentiation. Behlke et al. 2003, using the four Cluster spacecraft, examined the convective electric fields associated with SLAMS at the terrestrial bow shock. The convective electric showed that the SLAMS were moving with the solar wind plasma and they also observed a downstream wake of decreased density. Behlke et al. 2004 observed electric field signatures inside of SLAMS consistent with BGK phase space holes moving above the ion thermal speed with negative potential structures but were unable to determine their impact on the structure of the SLAMS.
No experimental study, however, has shown any truly distinct characteristics that distinguish SLAMS from shocklets (or discrete wave packets). Only their scale lengths have been reported to be different \cite{Lucek et al., 2002}, and this is not a truly distinguishing characteristic in a dispersive medium. One characteristic, which seems to be prevalent in the observational studies, is the change in magnitude of the magnetic field across each structure. As one can see from Figures 3.1 through 3.3, the maximum amplitude achieved by any particular shocklet rarely exceeds twice that of the average upstream field magnitude. In almost every example figure presented by \cite{Schwartz et al., 1992}, the average upstream field magnitude is \(\sim 5\text{-}10\) nT and the peak amplitude in the SLAMS vary from as low as 25 nT to as high as \(>60\) nT. The ratio of amplitudes, peak in SLAMS to average upstream, is often in excess of a factor of five, well beyond what is expected for adiabatic compression. The two examples in Figure 3.4 have magnetic compression ratios greater than two.

### 3.2.3 1 Hz Whistlers

The higher frequency (\(\geq 1\) Hz) waves of Fairfield’s high frequency category (0.5-4.0 Hz) have been identified as whistler waves \cite{Fairfield, 1974, Hoppe et al., 1982}. A nearly monochromatic whistler wave was discovered by \cite{Hoppe et al., 1982} far upstream (\(>1\) \(R_E\)) of the quasi-parallel bow shock in association with reflected ion beam distributions. \cite{Hoppe et al., 1982} showed that the rest frame frequencies were \(20 < \omega/\Omega_{cp} < 100\) (\(\sim 0.2\text{-}10\) Hz in the solar wind), wavelengths were \(\sim 100\) km, and propagation angles with respect to the magnetic field \(\theta_{kB} \sim 20^\circ\text{-}45^\circ\). Oblique whistler waves with \(f \sim 1\) Hz observed by \cite{Farris et al., 1993} upstream of low \(\beta\) quasi-perpendicular subcritical and marginally supercritical bow shocks were consistent with the \(\sim 1\) Hz whistlers reported by \cite{Russell et al., 1971} and \cite{Hoppe et al., 1981, 1982}. Whistler waves observed upstream of the Mercurian, Venusian, and Saturnian bow shocks were found to have characteristics similar to the "one-Hz" whistler waves observed by \cite{Hoppe et al., 1981, 1982}.
3.2.4 ULF Foreshock Waves at IP shocks

There are fewer observations of shocklets, SLAMS, or foreshock whistlers observed upstream of IP shocks. This may be due to the fact that IP shocks are unlikely to produce conditions conducive to the production of these waves due to their massive scales, larger radius of curvature at 1 AU, tendency to have a quasi-perpendicular geometry, and typically lower Mach numbers. Not only are quasi-perpendicular shocks less likely to produce the diffuse ion distributions thought to be necessary for shocklet generation and growth, they are not intrinsically unstable to reformation when sub-critical \cite{Farris:1993}. Both \cite{Hada:1987} and \cite{Omidi:1990} suggested shocklet generation mechanisms which require conditions more likely to occur in the foreshocks of planetary or cometary bow shocks connected to the quasi-parallel section.

This is not to say that ULF foreshock waves are not observed upstream of IP shocks, rather that the phenomena are far more rarely observed. Using multi-satellite measurements upstream from quasi-perpendicular and quasi-parallel IP shocks, \cite{Russell:1983} observed two distinct ULF foreshock wave types, a whistler precursor near the ramp and a 30 second wave they called irregular turbulence farther upstream which had a nearly featureless frequency spectrum. \cite{Tsurutani:1983}, in a study of \sim100 quasi-parallel (defined by the authors as $\theta_{Bn} < 65^\circ$) IP shocks, found low frequency waves (\sim0.05 Hz) to propagate within 15° of the ambient magnetic field.

3.2.5 Simulations and Summary of ULF Foreshock Waves

Simulations have attempted to clarify the distinguishing characteristics between foreshock ULF waves, shocklets, and SLAMS. \cite{Omidi:1990} used an electromagnetic hybrid code to study shocklets and they showed that the leading wave train of the shocklets were a consequence of wave spreading due to dispersive effects. They initialized an elliptically polarized ULF wave in a nonperiodic box and allowed it to interact with a proton beam with a spatially dependent density. The interaction between their artificial solar wind and the gradient in beam density caused the ULF waves to steepen and form kinetic magnetosonic waves. As the magnetosonic waves steepened, magnetosonic whistler waves grew just downstream of the steepened portions and began to propagate upstream away from the steepened edge of the magnetosonic waves.
Omidi and Winske [1990] found that the beam provides the free energy for wave growth and steepening and that the propagation of the wave results in the wave train due to wave spreading. Thus, in their simulation the beam is a necessary source of free energy for the formation of the steepened magnetosonic waves with leading whistler wave train. Akimoto et al. [1991] performed a 1-D electromagnetic hybrid simulation to investigate the steepening of parallel propagating hydromagnetic waves (ULF foreshock waves) into magnetic pulsations driven by ion beams. The waves were found to strongly interact with an ion beam, trapping it, and then gyrophase-bunching it which locally generates more waves. The waves could also resonantly interact with the shorter wavelength waves, adding to the total wave energy present [Akimoto et al., 1991].

Scholer [1993], also using an electromagnetic hybrid code, examined ULF waves, excited by ion beams, steepening upstream of a quasi-parallel shock as they convect into regions of increasing diffuse ion density. The diffuse ions were shown to have a significant impact on the index of refraction for the whistler and magnetosonic waves. As the ULF waves steepen, they dispersively radiate a whistler wave. Scholer [1993] also found that the leading whistler train was phase standing with respect to the shock front, which is different from the results of Omidi and Winske [1990]. Scholer [1993] found that the existence of diffuse ions is a necessary component for the formation of the steepened magnetosonic waves with leading whistler wave train.

Two more recent 1-D PIC simulation studies focused on the evolution of SLAMS excited by diffuse ion distributions [Scholer et al., 2003; Tsubouchi and Lembège, 2004]. Both studies found SLAMS to result from steepened ULF foreshock waves and the leading whistler train to result from the radiation of the dispersive waves by the steepened edge of the SLAMS. Some of the first studies referred to the phenomena as shocklets [Omidi and Winske, 1990; Scholer, 1993] while later studies called them SLAMS [Scholer et al., 2003; Tsubouchi and Lembège, 2004]. Regardless of the terminology they used, Scholer [1993]; Scholer et al. [2003]; Tsubouchi and Lembège [2004] all concluded that ULF foreshock waves were the source of the shocklets/SLAMS and that they were all the same phenomena, just at different stages in their evolution, consistent with the predictions of Schwartz and Burgess [1991].

To summarize, the evolution of the low frequency foreshock waves occurs in the following manner:
1. Initially low frequency ULF waves, both magnetosonic and Alfvénic in nature, are excited by field-aligned proton beams in cyclotron resonance with the waves far upstream of the bow shock (i.e., greater than an ion gyroradius) near the ion foreshock boundary [Hoppe et al., 1982; Meziane et al., 1999, 2004b; Mazelle et al., 2003].

2. The foreshock ULF waves have a phase velocity less than that of the solar wind, which causes the waves to be convected back towards the bow shock where they interact with changing ion distributions [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983; Meziane and D’Uston, 1998]. As the ULF waves convect and interact with the ion beams, they pitch-angle scatter the beam producing intermediate and/or gyrating ion distributions [Mazelle et al., 2003].

3. These anisotropic and nongyrotropic ion distributions alter the index of refraction of the waves leading to wave growth and/or steepening [Akimoto et al., 1991; Schwartz and Burgess, 1991; Scholer, 1993].

4. The steepened waves convect still further into the foreshock and the ion distributions become broader to the point of becoming diffuse in phase space.

5. When the steepened waves convect into a region of diffuse ions (see Section 1.4.3), the index of refraction becomes even more strongly dependent on these suprathermal diffuse ions [Omidi and Winske, 1990; Scholer, 1993] causing the waves to steepen more still to the point of breaking. The breaking waves are the result of insufficient energy dissipation needed to limit the steepening.

6. As the waves break, they dispersively radiate whistler waves upstream of the steepened edge because the whistlers in this frequency range have larger phase velocities than the lower frequency magnetosonic waves. The frequency dependence of the phase velocity leads to the leading whistler wave train, with the highest frequencies observed at the greatest spatial distance from the steepened edge of the magnetosonic wave. This process leads to the shocklets of the type seen in Figure 3.3.

7. As the shocklets continue to convect towards the bow shock, the wave amplitude increases still more due to pressure gradients [Giacalone et al., 1993] and the
structures eventually grow to become SLAMS, extracting energy from suprathermal ions \cite{deWit1999}, and diffuse ions \cite{Scholer1993, Scholer2003, Tsubouchi2004}. These structures are then thought to play an important role in quasi-parallel shock reformation and energy dissipation \cite{Scholer2003, Tsubouchi2004}.

### 3.3 Introduction to Cross-Shock Potential

The phenomena responsible for changing the electron distribution functions (eDFs) across collisionless shock waves has been a topic of considerable debate since the prediction of the existence of collisionless shock waves. Some theories invoke microscopic wave-particle interactions to explain the observed heating \cite{Dum1974, Forslund1970, Fredricks1965, Gary1970, Lampe1972, Scarf1965, Wong1970}, while others claim that only macroscopic effects are necessary to explain the heating \cite{Hull1998, Hull2000, Hull2001, Kuncic2002, Scudder1986}. Early observations of weak IP shocks found that electrons were preferentially heated perpendicular to the magnetic field, consistent with adiabatic compression and conservation of the first electron adiabatic invariant \cite{Feldman1983}. Observations of strong shocks, however, showed a preference for parallel heating with downstream distributions resembling an inflated Maxwellian with a flattened characteristic at low energies in the parallel cut of the eDFs, called a flattop distribution \cite{Feldman1983}. Similar eDF features were observed near the terrestrial bow shock \cite{Feldman1983}, which is often a much higher Mach number shock than IP shocks.

To produce a flattop eDF, some suggested that a macroscopic \((\gtrsim\text{scale size of shock thickness})\) quasi-static electric field could produce the observed distributions \cite{Feldman1982, 1983b, Scudder1986a, Thomsen1987}. \cite{Scudder1986a} showed, using ISEE observations, that this electric field was frame dependent and that the most important reference frame for the electrons was the de Hoffmann-Teller frame of reference. The de Hoffmann-Teller frame is the frame of reference where the incident flow is parallel to the magnetic field, thus no upstream convective electric field exists in this frame. The electric potential due to the quasi-static electric
field in the shock ramp is called the de Hoffmann-Teller potential or $\Delta \Phi_{dHT}$. Numerous studies have shown evidence to suggest that $\Delta \Phi_{dHT}$ plays an important role in shaping the eDFs downstream of collisionless shock waves [Feldman et al., 1982, 1983a,b; Scudder et al., 1986b,c,a; Thomsen et al., 1987].

To zeroth order, $\Delta \Phi_{dHT}$ can accelerate the incident electrons through the shock resulting in a peak (e.g. drifting Maxwellian) in the eDFs offset parallel to the magnetic field. In the absence of any wave-particle interactions, this peak in the eDF would remain constant through the shock and into the downstream. However, the combination of conservation of magnetic moment, $\mu$, shocked electrons attempting to move back upstream, and wave-particle scattering/heating are thought to give rise to the observed downstream eDFs [Scudder et al., 1986a; Thomsen et al., 1987]. Some have shown evidence to suggest that wave-particle interactions are not necessary to fill the voids in phase space resulting from acceleration of an entire distribution by a quasi-static electric potential [Hull et al., 1998, 2000; Hull and Scudder, 2000; Hull et al., 2001] assuming the maximal trapping approximation developed by Morse [1965]. If one assumes a constant phase space density along a phase space trajectory, then one requires that the downstream phase space density cannot exceed the maximum upstream phase space density along that same trajectory. Maximal trapping is an assumption whereby one assumes that changes in background field parameters are slow enough to allow a sufficient number of electrons to fill the voids in phase space produced by the quasi-static electric potential in the shock ramp resulting in a phase space density that approaches the maximum upstream phase space density [Morse, 1965]. This assumption is thought to be satisfied because shocked electrons in the downstream region have thermal speeds much greater than the shock speed and solar wind speed, thus a fraction of the thermal electrons can attempt to move back upstream. The electric field produced by $\Delta \Phi_{dHT}$ acts to accelerate upstream electrons into the downstream, so the shocked electrons attempting to return upstream are decelerated by this same potential and in some cases trapped in the downstream. Also, time-dependent magnetic fields and electric potentials in the shock ramp can act to increase the number of trapped electrons [Morse, 1965]. In summary, the process by which $\Delta \Phi_{dHT}$ can lead to heating and flattop eDFs is:
1. Incident electrons are accelerated parallel to the magnetic field by $\Delta \Phi^{dHT}$ producing a peak in the eDF offset along the magnetic field.

2. This peak in phase space density decreases in amplitude as the spacecraft transits the shock \cite{Feldman1983}, suggesting the accelerated particles are relaxing toward an equilibrium state which is thought to occur through either wave-particle interactions \cite{Thomsen1983a} and/or maximal trapping effects \cite{Morse1965}.

3. Shocked electrons in the downstream region attempt to return upstream and are subject to wave-particle effects \cite{Thomsen1983a} and/or trapping in the ES fields due to $\Delta \Phi^{dHT}$ \cite{Scudder1986a}, which has been observed near the terrestrial bow shock \cite{Larson1996}.

4. These combined effects can result in the parallel flattop eDF. The perpendicular heating observed in collisionless shocks often can be explained by a conservation of the first adiabatic invariant (or magnetic moment) for thermal electrons, thus $T_{e,\perp} \propto B_0$. However, conservation of magnetic moment can lead to mirroring which should result in phase space voids inaccessible in the downstream eDFs. To explain this, \cite{Hull2000} and \cite{Scudder1986a} suggest that wave-particle interactions only act to fill the voids in phase space and smooth the downstream eDFs. Other studies have gone further and shown evidence to suggest that macroscopic fields due to $\Delta \Phi^{dHT}$ and gradients in the magnetic field magnitude can explain the downstream eDFs without any wave-particle interactions \cite{Hull1998, HullScudder2000, Hull2001, HullScudder2001}. In the following sections we will explain why many of the assumptions made when calculating $\Delta \Phi^{dHT}$ are not valid for most collisionless shock waves. However, we cannot definitively show that wave-particle interactions are absolutely necessary to explain the downstream eDFs. In other words, the invalidation of the assumptions made by \cite{Hull1998, Hull2001, HullScudder2000, Hull2001, Scudder1986a} may not invalidate their theory. A corollary is observing whistler waves in finite temperature plasmas yet their characteristics conform to the cold plasma dispersion relation \cite{Fairfield1974}. 
3.3.1 Expected Observations

Flattop parallel eDFs observed downstream of collisionless shocks present an argument in favor of quasi-static electric field driven electron heating [Scudder et al., 1986a]. However, these same eDFs have been shown to be produce by IAWs in the quasi-linear limit [Dum et al., 1974]. The question we address in this section is, which mechanism dominates the downstream eDF structure, the quasi-static macroscopic fields or the high frequency microscopic fields?

Figure 3.5 is an example of a typical electron heating signature observed at a supercritical quasi-perpendicular IP shock. The IP shock, observed by the Wind spacecraft on 09/24/1998 at 23:20:37 UT, had $M_f = 2.87 \pm 0.08$, $\theta_{Bn} = 78.6^\circ \pm 2.4^\circ$, $N_2/N_1 = 2.17 \pm 0.38$, and $\beta_1 = 0.215 \pm 0.214$ [Kasper, 2007]. Each eDF has a similar format to that of Figure 1.8 except that the parallel cuts are red, perpendicular are blue and the one-count levels are included. The top panels are two-dimensional contours of constant phase space density in the plane defined by the solar wind velocity and magnetic field, showing the projections of the shock normal vector (blue dotted line), $\hat{n}_{sh}$, the solar wind velocity (solid black line), $V_{sw}$, and the electron heat flux vector for these specific distributions (solid red line), $Q_f$. The data ranges are $\pm 20,000$ km/s for the parallel (horizontal), $V_\parallel$, and perpendicular (vertical), $V_\perp$, axes. The bottom panels contain the parallel (red line) and perpendicular (blue line) cuts of the eDFs along with the one-count (green line) levels for uncertainty estimates. In between the eDFs is a vertical line used as an illustration to distinguish between the upstream eDF on the left and downstream eDF on the right. The upstream magnetic field for this IP shock was mostly in the positive Y-GSE and negative X-GSE directions. Thus, the deviation in the parallel cut of the eDF on the left from isotropy is likely due to a combination of the solar wind strahl and mirrored electrons.

As illustrated in the parallel cut (red line) of the eDF on the right in Figure 3.5, the downstream electron distribution can have a flattened characteristic at low energies, known as a flattop distribution. The specific example shown also has a beam-like feature on top of the flattop distribution, a feature often seen downstream of the terrestrial bow shock [Thomsen et al., 1983a]. The large strahl component observed in the parallel cut in the left eDF (23:20:31 UT) skews the estimate of the electron temperature anisotropy ($T_\perp/T_\parallel$) for the core, halo, and entire distribution. In the eDF on the right (23:22:11
Figure 3.5: The plot shows two EL eDFs upstream and downstream of a supercritical quasi-perpendicular IP shock observed by the Wind spacecraft on 09/24/1998 at 23:20:37 UT. Each eDF has a similar format to that of Figure 1.8 except that the parallel cuts are red, perpendicular are blue and the one-count levels are included. In between the eDFs is a vertical line used as an illustration to distinguish between the upstream eDF on the left and downstream eDF on the right. The upstream magnetic field for this IP shock was mostly in the positive Y-GSE and negative X-GSE directions.
UT), one can see that at very low energies (\(\lesssim 45\) eV or \(\sim 4000\) km/s), the electrons are more anisotropic parallel to the magnetic field but at slightly higher energies (\(\sim 70\) eV or \(\sim 5000\) km/s) the eDF starts to shift towards \(T_\perp > T_\parallel\). The halo electrons (\(\gtrsim 100\) eV or \(\sim 6000\) km/s) become more and more perpendicularly anisotropic with increasing energy. The anisotropy values for the core (subscript c) and halo (subscript h) for both distributions are \(T_{\perp,c}/T_{\parallel,c} \sim 0.82, 0.93\) and \(T_{\perp,h}/T_{\parallel,h} \sim 0.81, 1.77\). Note that for the eDF on the right (23:22:11 UT), the halo is almost twice as anisotropic as the core. Though the low energy parallel flattop is consistent with previous bow shock observations, many supercritical IP shocks show very anisotropic halo electrons and nearly isotropic core electrons near the downstream edge of the shock and the halo electrons relax to more isotropic distributions further downstream. For instance, the eDF at 1998-09-25/00:08:53 UT (not shown) had \(T_{\perp,c}/T_{\parallel,c} \sim 0.94\) and \(T_{\perp,h}/T_{\parallel,h} \sim 0.97\). We will show an example of an IP shock where the halo and core electrons do not isotropize within an hour of crossing the shock ramp in Chapters 4 and 5, suggesting local processes pitch-angle scattering the electrons in the downstream.

### 3.4 Electric and Magnetic Field Measurements

#### 3.4.1 Electric Fields in Shock Ramps

This section will highlight some of the direct measurements of electric fields which can exist in the shock transition region \cite{BaleMozer2007}. The illustrations will show that the electric fields within a shock layer consist of a spectrum of frequencies which accumulate to make up the cross-shock potential. Thus, we will argue that one cannot use only the quasi-static fields in the shock ramp to describe the fundamental physics of energy dissipation. For future reference, we will define the electrons as being demagnetized if \(\vec{E} + \vec{v}_e \times \vec{B} \neq 0\).

\cite{Wygantetal1987} measured low frequency (\(\lesssim 32\) Hz) electric fields within the shock ramp using the ISEE-1 instrument. They presented the first measurements of high time resolution \((32\) samples/s) 1-D electric field data at a collisionless shock supplemented with low time resolution \((3\) second averages) 2-D electric field data. The low frequency measurements were the first measurements of the cross-shock quasi-static
electric field, measured in both the normal-incidence and de-Hoffmann-Teller frame of reference. The quasi-static electric fields were on the order of 4-10 mV/m resulting in potentials >300 Volts while the higher frequency data showed intense spikes with amplitudes ranging from 10-100 mV/m. The lower limit on the potential of the largest spike was estimated to be $\sim$8 Volts (assuming 80 m wavelength). When they assumed the waveform was driven by the lower hybrid instability, for which the wavelength should be roughly an electron gyroradius, $\sim r_{pe}$, or roughly 1 km, the resulting potential drop was roughly 50 Volts. Though this value is smaller than the quasi-static potentials, it is not insignificant and should be considered when estimating the electric fields in the ramp of collisionless shocks.

Walker et al. [2004] measured the quasi-static electric field perpendicular to the magnetic field, $E_\perp$, finding the amplitudes to exceed 70 mV/m and the scale sizes to be as small as an electron inertial length, $c/\omega_{pe}$. They also found that as the shock normal angle increased to 90°, the scale length of these electric field structures decreased. The two dimensional measurements were made at 25 samples per second with the four Cluster spacecraft, slightly lower than the highest frequency 1-D measurements made by Wygant et al. [1987]. The sharp gradients observed by Walker et al. [2004] could potentially demagnetize the electrons, changing the shock dynamics.

Hull et al. [2006] measured 3-D electric fields in the bow shock using the Polar spacecraft, finding ion-acoustic-like waves in the ramp with amplitudes exceeding 80 mV/m. The data was sampled at up to 8000 Hz (occasionally at 1600 Hz), but the frequencies of the ion-acoustic-like waves extended beyond 100 Hz, above the sampling frequencies of Wygant et al. [1987] and Walker et al. [2004]. It may be that the measurements made by Wygant et al. [1987] and Walker et al. [2004] were attenuated by aliasing implying their electric field measurements underestimated the electric fields in the shock ramp.

Bale and Mozer [2007] measured parallel and perpendicular electric fields, $E_\parallel$ and $E_\perp$, at 1600 Hz using the 3-D electric field instrument on the Polar spacecraft. They reported on the first direct measurements of $E_\parallel$ at a collisionless shock. At nearly perpendicular shocks, $E_\perp$ will act to slow the incident solar wind while the existence of $E_\parallel$ in a shock ramp is significant since it can directly demagnetize electrons. Bale and Mozer [2007] observe $E_\parallel \leq 100$ mV/m and $E_\perp \leq 600$ mV/m corresponding to localized potential wells of 10’s of volts and 100-1000’s of volts, respectively. The scale lengths of
Figure 3.6: The plot shows six different TDS samples from the Wind/WAVES instrument observed during a bow shock crossing on 12/02/1996 at roughly 16:56:20 UT (see Figure 3.4). The left hand column (red waveforms) correspond to the X-antenna and the right hand column (blue waveforms) correspond to the Y-antenna. Each waveform is labeled with a respective letter and amplitude illustrated by the vertical black line on the left hand of each panel. Notice that at least one component of each waveform exceeds 100 mV/m peak-to-peak. On the right hand side are lists of the cyclotron frequency, $f_{ce}$, and the range of frequencies for the peak power of the wave normalized by $f_{ce}$. All the waveforms are oblique linearly polarized waves (hodograms not shown). The sample rate was $\sim 120$ kHz.
the observed $E_\parallel$ were as small as fractions of $c/\omega_{pe}$ while $E_\perp$ scale lengths were on the same order. Sharp gradients in the electric field could demagnetize electrons even in the absence of large $E_\parallel$.

An illustrative example of higher frequency waveforms can be seen in Figure 3.6. The waves have frequency peaks between 1-10 kHz, are oblique linearly polarized waves consistent with Doppler shifted ion-acoustic waves \cite{Gurnett et al., 1979b}, \cite{Wilson III et al., 2007}. \cite{Wilson III et al., 2007} calculated quasi-linear estimates of the anomalous resistivities produced by large amplitude ion-acoustic waves in the ramps of IP shocks finding they ranged from $\sim 1$–856 $\Omega \cdot m$ ($\sim 10^7$ times greater than classical Spitzer particle collision estimates). However, recent Vlasov simulations have compared quasi-linear estimates of anomalous resistivities produced by ion-acoustic waves with the momentum exchange calculated from their simulation results. They found the simulation results can be 2-3 orders of magnitude greater than the quasi-linear estimates \cite{Dyrud and Oppenheim, 2006}, \cite{Petkaki et al., 2003, 2006}, \cite{Petkaki and Freeman, 2008}, \cite{Watt et al., 2002}, suggesting microscopic wave-particle interactions may be more important than previously thought.

In this section we have presented multiple examples from previous studies and our own observations of electric fields within the ramp region of collisionless shocks. Exactly how the cross-shock potential arises is still poorly understood and how it scales with upstream parameters is even less so. Regardless, the examples presented herein show that the electric fields can produce both large potential drops and localized potential wells which can affect the local particle distributions. Thus, observed changes in the eDFs downstream of collisionless shocks may not be entirely due to macroscopic quasi-static fields.

### 3.4.2 Magnetic Fields in Shock Ramps

This section highlights some of the measurements of magnetic fields within shock ramps. These illustrative observations will show that the assumptions of laminar, steady-state changes used in calculation of $\Delta \Phi^{dHT}$ may not be valid for many IP shocks. An illustrative example will be used to explain why the use of DC measurements in shock analysis can be misleading. The quantitative implications for theories relying upon monotonic shock transitions, however, will not be addressed here.
A cursory examination of most IP or bow shock crossings with magnetic field data

Figure 3.7: The plot shows the magnitude and the GSE coordinates of the magnetic field for the IP shock on 10/24/1997 in 3 second (top two panels) and 22 samples/second (bottom two panels) data. The upstream flow speed in the shock frame is roughly 117 km/s which means that roughly each sample for the two different sample rates corresponds to length scales of roughly 351 km and 5.32 km, respectively. For reference, the average upstream electron inertial length is roughly 10 km.

Sampling at rates higher than 1 Hz will show that fluctuations on the same order as the amplitude of the shock ramp can occur, even at weak IP shocks (see example in Wilson III et al. [2009]). Figure 3.7 is a perfect example of how a shock crossing may look very different when examining at different sample rates. It is also clear from the figure that the amplitude of the high frequency (~1-3 Hz) magnetic fluctuations can be similar to that of the shock itself, which immediately throws into question any of
the small amplitude fluctuation arguments (i.e. monotonic magnetic field profile, time-stationary, and planar shock arguments). The 10/24/1997 event had a Mach number of \( \sim 1.7 \), shock normal angle \( \sim 68^\circ \), compression ratio \( \sim 2.5 \), average upstream electron inertial length \( \sim 10 \) km, and normal thickness of roughly 35 km (\( \sim 3.5 <c/\omega_{pe,1}> \) or average upstream electron inertial lengths). Also, the 10/24/1997 event had an upstream normal flow speed of roughly 117 km/s. This implies that when sampling at 22 samples/second, each sample corresponds to roughly 5.32 km (\( \sim 1/2 <c/\omega_{pe,1}> \)) of solar wind plasma flowing past the spacecraft. For example, in one cycle of a 3 Hz wave, \( \sim 4 <c/\omega_{pe,1}> \) of solar wind plasma has traversed the satellite (a similar number is found when estimating the whistler wavelength in the Alfvén Mach number limit [Mellott, 1984]). Often the assumption is made that all fluctuations and/or relevant scale lengths in a shock are much greater than the electron inertial length which allows one to drop electron inertial terms in certain calculations. However, in this case the relevant scale lengths are on the same order as \( <c/\omega_{pe,1}> \).

In addition, not all electrons will interact in the same manner with a given wave, as seen in Sections 1.5.3 through 1.5.5 [Saito and Gary, 2007] found that higher energy electrons were more efficient at interacting with whistler waves, reinforcing the treatment of the solar wind as a multi-component kinetic plasma, not a single electron fluid. The magnetic fluctuations shown in Figure 3.7 combined with the electric fields observed in the multiple previous studies highlighted in Section 3.4.1 raise doubts about whether one can assume electrons stay magnetized everywhere within a shock. The magnetic field measurements in Figure 3.7 bring into question whether the shock is time-stationary. Regardless, the purpose of these illustrations are to exhibit why the detailed microphysics of collisionless shock waves is important.

### 3.5 Assumptions For de-Hoffmann-Teller Potential Calculations

If one transforms into the de-Hoffmann-Teller frame of reference, theory shows that a DC electric field should exist across the shock, producing the de-Hoffmann-Teller potential, \( \Delta \Phi^{dHT} \). In this frame of reference, the convection electric field (\( \mathbf{E} = -\mathbf{V} \times \mathbf{B} \)) of the
solar wind disappears in the upstream region. The electrons are assumed to be magnetized everywhere in this frame of reference. Numerous studies have calculated $\Delta \Phi_{dHT}$, given specific assumptions, many of which may not be physically realistic for many shocks [Hull et al., 1998, 2000; Hull and Scudder, 2000; Hull et al., 2001; Kuncic et al., 2002; Scudder et al., 1986a]. However, it is difficult to determine the relative importance and validity of each assumption without doing thorough simulations and comparisons with data. The following is a list of assumptions that have been made in order to estimate $\Delta \Phi_{dHT}$:

1. electrons are magnetized everywhere (i.e. changes in magnetic field should be larger than the electron inertial length, $c/\omega_{pe}$)
2. neglect electron heat flux, assume isotropic pressure, and adiabatic gas law
3. field-aligned flow for electrons
4. shock is time-stationary and planar
5. shock has a monotonic jump in the magnetic field magnitude
6. flat-topped electron distributions downstream of the shock (at the same time as assuming the weak shock limit)
7. $\Delta \Phi_{dHT} \propto \Delta B$
8. the total energy of the electrons in the de-Hoffmann-Teller frame is a constant
9. steady-state Vlasov fluid equations

In Section 3.4.1, previous observations which contradict Assumption 1 (thus also Assumption 3) were presented, including the existence of sharp electric field gradients and $E_B$. Assumption 2 can be problematic when dealing with IP shocks since there are rarely events without a finite electron heat flux near 1 AU in the solar wind. Furthermore, the isotropic pressure assumption can be significantly inaccurate as observations of strong anisotropies in the electron distributions have been presented [Pilipp et al., 1987b; Thomsen et al., 1985a; Wilson III et al., 2009]. A cursory examination of most IP or bow shock crossings with magnetic field data sampling at rates higher than 1...
Hz will show that most events show Assumptions 3 through 5 to be inaccurate, where Assumptions 4 and 5 can break down even for some weak shocks (see examples in Wilson III et al. [2009]). For most IP shocks, Assumption 6 is not accurate, yet studies often make this assumption even when invoking the weak shock limit. Demagnetization of electrons due to stochastic wave-particle interactions can affect the validity of Assumptions 8 and 9.

In summary, all of the assumptions used to calculate $\Delta \Phi^{dHT}$ are often inaccurate for IP shocks and planetary bow shocks. Whether the inaccuracy of these assumptions is great enough to invalidate the calculations is not known. Scudder et al. [1986a] presented one example where the calculations matched their observations within uncertainties. Thus, we believe that the downstream eDFs result from some combination of quasi-static macroscopic and high frequency microscopic fields affecting the particles in the shock ramp.

3.5.1 Calculation of the de Hoffmann-Teller Potential

The de Hoffmann-Teller potential, $\Delta \Phi^{dHT}$, was first shown to be the relevant cross-shock potential when considering eDFs by Scudder et al. [1986a]. They estimated $\Delta \Phi^{dHT}$ by integrating the measured electron pressure across the shock ramp, given by:

$$
\Delta \Phi^{dHT} = \int_{x_1}^{x_2} dx \frac{\nabla \cdot P}{n_e e}
$$

(3.1)

where $x$ is a measure of distance across the shock in the shock normal direction and $\Delta$ represents the total difference between upstream and downstream. The discussion in Appendices A.3.2 and A.3.3 is a specific example of the moment calculations done using particle distribution function measurements. The purpose of referencing those sections is to highlight the fact that moment software often uses the assumption $<T_e> = \text{Tr}[P_{i,j}]/n_e$ to calculate the components of the "temperature" directly from the particle distribution function measurements. If one makes this assumption, then all calculations using the electron pressure should be considered nearly the same as calculations done using the electron temperature. However, some studies calculate $\Delta \Phi^{dHT}$ using the electron pressure measurements and then correlate $\Delta \Phi^{dHT}$ to $T_e$, which we do not feel as being physically significant. This is not to say that $\Delta \Phi^{dHT}$ is insignificant nor do
we dispute the theory. We simply do not agree with the arguments that suggest the observed downstream eDFs can be entirely explained by the macroscopic DC fields in the shock ramp \cite{Hull1998, HullScudder2000, Hull2001} nor do we agree with the interpretation of the correlation observed between $\Delta \Phi^{dHT}$ and $T_e$ reported by \cite{Hull2000}. In summary, we do not believe that changes in the eDFs across collisionless shock ramps are adequately explained using only macroscopic DC fields. In this thesis, we will examine in detail the eDFs and higher frequency waveforms and argue that our observations provide evidence to support the theory that wave-particle interactions are important in collisionless shock waves.
Chapter 4

Shocklets and Low Frequency Whistlers at an Interplanetary Shock

4.1 An Atypical IP Shock


In the examination of low frequency waves using high time resolution magnetometer data from the Wind spacecraft at ten IP shocks, five quasi-perpendicular events had foreshock or precursor waves which were examined in detail by Wilson III et al. [2009]. They observed unusual waveforms upstream of one IP shock on 04/06/2000, finding the waves to be consistent with shocklets. This event was the strongest shock examined in this study with $M_f \sim 4$, $\theta_{Bn} \sim 68^\circ$, and $N_{i2}/N_{i1} \sim 4$ [Kasper, 2007]. The shock also showed strong electron and ion heating that far exceeded that of the other four more typical IP shocks. Within $\pm 5$ seconds of the shock crossing, four large amplitude ($> 15 \text{ mV/m peak-to-peak}$) waveforms were observed, two of which are consistent with solitary waves. The unusual upstream shocklets and strong electron and ion heating led to another study [Wilson III et al., 2010] which examined the microphysics of the shock ramp region.

Figure 4.1 shows the shock on 04/06/2000 with the 12 observed shocklets labeled
Figure 4.1: A plot of the electric field intensity as a function of frequency and time, the magnetic field magnitude (3 second), and the ion density from PL on 04/06/2000. The vertical blue lines indicate the location of the 12 shocklets observed upstream of this event. One can see that the magnetic field magnitude and density/thermal line are in phase, consistent with magnetosonic waves.

with blue lines, showing the relationship between magnetic field magnitude and ion density. The top panel shows the electric field power from the WAVES TNR detector. Upstream of the shock one can easily see the plasma line which is proportional to the square root of the plasma density. Note that the TNR data shown is one minute averages and on a log scale, thus the relative changes in phase with the magnetic field (panel directly below) are not always obvious. The bottom panel of Figure 4.1 plots the ion density from the Pesa Low detector on the Wind 3DP instrument [Lin et al., 1995]. One can clearly see that the magnetic field magnitude and ion density are in phase at the main shock front and at each shocklet, consistent with magnetosonic waves.

Figure 4.2 shows the same event as Figure 4.1 on a shorter time scale, with the magnetic field magnitude scaled to emphasize the shocklets seen upstream (indicated by the blue arrows). Wind crossed the main shock ramp at 16:32:09.237 UT (i.e. the far right-hand side of the top panel of Figure 4.2 or roughly 3 seconds after shocklet B).
Figure 4.2: A plot of the magnetic fields for the shock on 04/06/2000. The shock arrival time was 16:32:09.237 UT. The top panel is the magnetic field strength, |B| (nT), followed by the field components in GSE coordinates from 15:39:00 - 16:36:00 UT. In the top panel, the location of each shocklet observed for this event is labeled with a blue number and an arrow. The two shocklets which will be analyzed in detail in this paper are labeled A and B (Figures 4.3 and 4.4 show a more detailed picture of each shocklet). The top two panels have been scaled down to show the shocklets more clearly because the IP shock itself jumps to over 30 nT downstream (only a few seconds after the shocklet B).
The bottom two panels are examples of two shocklets with waves on the leading edge consistent with a RH-whistler mode. The typical structure of the shocklets is labeled in panel \(B\). Shocklets \(1, 3-5, 7, \) and \(9-12\) all had clearly formed waves on their leading edges. These waves had RH polarizations with respect to the magnetic field and an increase in ion density coincident with the increase in \(|B|\), consistent with magnetosonic whistlers and magnetosonic waves.

Figures 4.3 and 4.4 show examples of the analysis done on the leading whistler wave of each shocklet. The left hand set of wave events plot the GSE (gray scale) components, the middle panels plot the MV (color scale) components, and the right hand set of panels show the hodograms, \(B_y\) vs. \(B_x\), \(B_z\) vs. \(B_x\), and \(B_z\) vs. \(B_y\). The time ranges for the selected subintervals seen in wave events \(A-C\), were chosen to maximize the intermediate to minimum eigenvalue ratio, seen in red in each wave event. In every wave analysis presented and every wave examined (125 different analyses for all five IP shocks), \(\lambda_2/\lambda_3 \geq 10.0\) and 56 cases had \(\lambda_2/\lambda_3 \geq 50.0\).

The use of multiple band pass filters on the shocklets revealed that their wave vectors remained relatively unchanged (within \(\sim 10^\circ-15^\circ\)) between the different frequency bands chosen for our filters for similar time intervals. This is illustrated clearly in events \(A\) and \(C\) of Figure 4.3 and events \(A, C,\) and \(D\) of Figure 4.4. The time interval length for each event are similar, but each event was filtered over a different frequency bin. The wave vector is the same for each event within uncertainties (see Equations 2.13 through 2.17 for uncertainty calculation). We also observed a clear dependence of the frequency at peak power on the distance from the steepened edge of the shocklets, with the higher frequency waves being seen first, followed by the lower frequency waves as the shocklets convected over the satellite. This is consistent with the frequency dependence of magnetosonic whistlers, whose group velocities increase with increasing frequency. This result is also seen in simulations \(\text{Omidi and Winske}, 1990; \text{Scholer}, 1993; \text{Scholer et al.}, 2003; \text{Tsubouchi and Lembège}, 2004\).

The difference in polarization between wave events \(A, C,\) and \(D\) in Figure 4.4 can be explained by projection effects due to single satellite measurements using only magnetic field measurements. In the spacecraft frame, wave events \(A, B,\) and \(D\) in Figure 4.3 are LH polarized with respect to the propagation direction, but all wave events in Figures 4.3 and 4.4 show a RH sense with respect to the magnetic field, characteristic of whistler
Figure 4.3: An example of MV analysis on the leading whistler waves of a shocklet shown in wave event A, Figure 4.2. The figure has the same format as Figure 2.8. The frequency ranges and angles of propagation are: $0.5 \, \text{Hz} < f < 1.0 \, \text{Hz}$ and $\theta_{kB} = 35^\circ (145^\circ)$ for A, $0.5 \, \text{Hz} < f < 1.0 \, \text{Hz}$ and $\theta_{kB} = 14^\circ (166^\circ)$ for B, $0.6 \, \text{Hz} < f < 3.0 \, \text{Hz}$ and $\theta_{kB} = 41^\circ (139^\circ)$ for C, and $0.6 \, \text{Hz} < f < 3.0 \, \text{Hz}$ and $\theta_{kB} = 18^\circ (162^\circ)$ for D. The eigenvalue ratios from the MV analysis are also shown with the MV estimate of the k-vector direction in GSE coordinates above each hodogram. The purple arrows indicate the direction of rotation for each respective plot.
Figure 4.4: Another example of MV analysis on the leading whistler waves of a shocklet (See panel B in Figure 4.2). The format of this figure matches that of Figure 4.3 but with different frequency ranges. The frequency ranges and angles of propagation are: $0.6 \text{ Hz} < f < 3.0 \text{ Hz}$ and $\theta_{kB} = 27^\circ (153^\circ)$ for A, $0.45 \text{ Hz} < f < 1.0 \text{ Hz}$ and $\theta_{kB} = 49^\circ (131^\circ)$ B, $f > 0.6 \text{ Hz}$ and $\theta_{kB} = 25^\circ (155^\circ)$ for C, and $f > 1.0 \text{ Hz}$ and $\theta_{kB} = 26^\circ (154^\circ)$ for D.
modes. Wave event C in Figure 4.3 and wave events A and B in Figure 4.4 are RH polarized both with respect to the wave vector and the magnetic field.

### 4.2 Comparison of the 04/06/2000 event to the four typical events

As mentioned above, four of the IP shocks had waves with characteristics consistent with previous shock studies. Figure 4.5 shows the magnetic field magnitude and GSE components for the four IP shocks with typical characteristics. The four events with precursor waves are much lower Mach number shocks \( (M_f < 2.3) \) than the unusual event \( (M_f \sim 4) \), consistent with theory [Morton, 1964; Stringer, 1963] and previous observation [Mellott and Greenstadt, 1984].

The events shown in the upper panels of Figure 4.5 are examples of lower Mach

![Figure 4.5: Four IP shocks with upstream precursor whistler waves, typical of subcritical quasi-perpendicular shocks. The top two panels (04/03/1996 and 04/08/1996) are examples of laminar quasi-perpendicular shocks with leading precursor whistler waves. The bottom two panels (10/24/1997 and 12/10/1997) show a far more turbulent transition from upstream to downstream.](image-url)
number quasi-perpendicular shocks with a leading wave train and a relatively stable transition from up to downstream. The events in the lower panels are examples of higher Mach number shocks with a much more turbulent transition. The relevant shock parameters are given in the green box. The waves in Figure 4.5 are similar in frequency and their propagation angle with respect to the shock normal, $\theta_{kn}$ to the precursor whistler waves observed by Russell et al. [1983]. Since $\theta_{kn}$ is not small it is not likely that these waves are phase standing with respect to the shock. The precursors had nearly circular RH polarization with respect to the magnetic field, and propagate obliquely to the field with $> 95\%$ having propagation angles $\theta_{kB} > 20^\circ$. However, Russell et al. [1983] found that $75\%$ of the precursors had propagation angles $\theta_{kB} < 20^\circ$.

The whistler precursor waves and shocklets shared some characteristics. The range of $\theta_{kB}$ values can be seen in Figure 4.6. There is an obvious difference between the whistlers upstream of the four typical events (bottom panel) and the 12 shocklets (top panel) observed upstream of the unusual event. The shocklets have a much broader range of $\theta_{kB}$ and tend to be more oblique than the precursor whistler waves. If the shocklets are in fact magnetosonic whistlers, the higher values of $\theta_{kB}$ would be consistent with their more compressive nature than that of the precursor whistlers. Almost 80\% of the shocklets observed for the 04/06/2000 event had SC frame frequencies $f > 0.45$ Hz had $\theta_{kB} < 45^\circ$, consistent with bow shock observations [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983; Russell et al., 1971] and cometary foreshock observations [Le et al., 1989]. Over 90\% of the whistler observed for the 4 IP shocks without shocklets had $\theta_{kB} < 45^\circ$, consistent with theory [Gary et al., 1994, 1999] and observations of whistler precursor waves at IP shocks [Russell et al., 1983]. There were no distinguishing characteristics in $\theta_{kn}$ or $\theta_{kV}$ between the shocklets and precursor whistlers.

4.3 Analysis of Particle Data for this Study

Estimates of the electron temperature anisotropies in both the cold dense core (subscript c) and the hotter more tenuous halo (subscript h) can be obtained from full 3D electron distributions (see Chapter 2 for more details). For both EL and EH distributions, average temperatures, parallel (subscript ||), and perpendicular (subscript \perp) to the magnetic field are computed. Temperature anisotropies, $T_{\perp j}/T_{|| j}$ ($j = c$ or h), were
Figure 4.6: Histograms comparing the angle of propagation with respect to the magnetic field for the IP shock precursor whistler waves and the shocklet whistlers. The top panel shows the range of angles for all bandpass frequency bins greater than 0.45 Hz but only for the 12 shocklets (each shocklet has multiple waves as seen in Figures 4.3 and 4.4) observed on 04/06/2000. The bottom panel shows the range of angles for all bandpass frequency bins greater than 0.6 Hz for the 4 IP shocks without shocklets. The horizontal and vertical axes are on the same scales for both panels. Multiple frequency ranges were chosen for each IP shock.
computed for each PAD and compared to threshold conditions for whistler heat flux and anisotropy instabilities of Gary et al. [1994] and Gary et al. [1999].

The wave events of Figures 4.3 and 4.4 show a clear relationship with the low to mid energy (∼few eV to 9 keV) electron distributions. Previous studies of whistler waves at shocks suggested a relationship between whistler mode generation and electrons [Tokar et al., 1984; Tsurutani et al., 1983; Gary et al., 1994] determined the threshold conditions for whistler heat flux and whistler anisotropy instabilities for typical solar wind conditions. They found that the instabilities were strongly dependent on the core parallel plasma beta, $\beta_{\parallel c}$, the ratio of parallel halo temperature to parallel core temperature, $T_{\parallel h}/T_{\parallel c}$, and the temperature anisotropy of the halo, $T_{\perp h}/T_{\parallel h}$. Using linear Vlasov theory, Gary et al. [1999] showed that the halo temperature anisotropy has a larger effect on the heat flux instability than the core temperature anisotropy. They also found that the heat flux-driven whistler mode was always unstable for $T_{\perp h}/T_{\parallel h} > 1.01$ and always stable for $\beta_{\parallel c} \leq 0.25$. In the cases where $T_{\perp h}/T_{\parallel h} > 1.01$ but $T_{\parallel h}/T_{\parallel c}$ is small, Gary et al. [1994] suggested that a whistler anisotropy instability may be excited even in the absence of a relative drift between the core and halo electrons. Thus if the halo electrons initially meet this criteria, the whistler anisotropy instability would dominate over the whistler heat flux instability. The whistler anisotropy instability acts to reduce the relative drift between the halo and core electrons (if present) and isotropize the halo temperatures. $T_{\perp h}/T_{\parallel h}$ would reduce faster than the halo/core drift (and $T_{\parallel h}/T_{\parallel c}$) could reduce causing the electron distributions to become unstable to a whistler heat flux instability. In the case of large $T_{\parallel h}/T_{\parallel c}$ and small $T_{\perp h}/T_{\parallel h}$, the whistler heat flux instability would initially dominate over the whistler anisotropy instability. This instability would increase $T_{\perp h}/T_{\parallel h}$ slower than it could reduce $T_{\parallel h}/T_{\parallel c}$ [Gary et al., 1994].

The electron pitch-angle distributions (PADs) for three wave events are plotted in Figures 4.7, 4.8, and 4.9. The primary influence of whistler heat flux instability is to pitch-angle scatter the halo electrons through a cyclotron resonance. There is a clear increase in the halo electron temperature anisotropy, $T_{\perp h}/T_{\parallel h}$ (see Table 4.1), as one crosses each wave, consistent with normal cyclotron resonance increasing the transverse energy of the electrons. $T_{\perp c}/T_{\parallel c}$ follows the same pattern, but the increase is not
Figure 4.7: Comparison of wave polarization and electron distributions (wave event D from Figure 4.3) with six pitch-angle distributions (PADs) from the Eesa Low (EL) and High (EH) instruments. Each PAD is plotted in number ($# \text{s}^{-1} \text{sr}^{-1} \text{cm}^{-2} \text{eV}^{-1}$) and energy (eV $\text{s}^{-1} \text{sr}^{-1} \text{cm}^{-2} \text{eV}^{-1}$) flux. The energies plotted range from 27 eV - 1113 eV for EL and 137 eV - 8875 eV for EH. The frequency range is $0.6 \text{ Hz} < f < 3 \text{ Hz}$ for the bandpass filter used on the MFI data and $\theta_{KB} = 40^\circ (140^\circ)$. The vertical lines on the PADs represent an average estimate of the propagation angle, $\theta_{KB}$, for the wave shown. The electron temperature anisotropies and other parameters can be found in Table 4.1.
as dramatic. The threshold conditions for a whistler heat flux or anisotropy instability (Figures 7 and 8 of [Gary et al. 1994]) are met by most of the EL PADs up and downstream of the waves in these three figures. There are, however, differences in our estimates of $n_{he}/n_e$, distributions used to model the halo electrons, and definition of heat flux from those of [Gary et al. 1994]. Our estimates of $n_{he}/n_e$ were often a factor of 10 or more smaller than the estimates used by [Gary et al. 1994] (~0.05) for the PADs shown herein (see Table 4.1). [Gary et al. 1994] used bi-Maxwellian distributions to model both core and halo electrons whereas we fit the halo distributions to modified Lorentzian distributions. [Gary et al. 1994] used a simplified version of the heat flux from [Feldman et al. 1975] whereas we calculated the full heat flux tensor, assuming it symmetric, and derived a vector from that tensor. The main difference is a loss of the vector direction associated with the heat flux, since [Gary et al. 1994] assumed it parallel to the ambient magnetic field. The magnitude estimates should not be affected significantly.

One can see that $T_{\perp h}/T_{\parallel h}$ increases across the waves in Figures 4.8 and 4.9 (from 0.55 for PAD 16:31:32-16:31:35 UT to 1.02 for PAD 16:31:44-16:31:47 UT shown in Table 4.1, an increase of ~85%). [Gary et al. 1994] found that the whistler heat flux instability reduced $T_{\parallel h}/T_{\perp c}$ but increased $T_{\perp h}/T_{\parallel h}$ at a faster relative rate. We observed $T_{\parallel h}/T_{\perp c}$ to decrease across the waves (from 13.4 for PAD 16:31:32-16:31:35 UT to 11.4 for PAD 16:31:44-16:31:47 UT, a decrease of ~15%). The same change occurred across the wave in Figure 4.7, $T_{\perp h}/T_{\parallel h}$ increases across the wave (0.65 for PAD 16:27:18-16:27:21 to 0.91 for PAD 16:30:36-16:30:39, an increase of ~40%) and $T_{\parallel h}/T_{\perp c}$ decreases (13.7 for PAD 16:27:18-16:27:21 to 10.9 for PAD 16:30:36-16:30:39, a decrease of ~20%). Notice in both cases $T_{\perp h}/T_{\parallel h}$ increases at a faster relative rate than $T_{\parallel h}/T_{\perp c}$ decreases, consistent with the simulation results by [Gary et al. 1994].

If we assume that the rest frame frequencies of the observed waves are consistent with [Hoppe et al. 1982] and use our measured $\theta_{hB}$, then the resonant energies for normal cyclotron resonance can be calculated from Equation 1.43. The resonant energies are $250 \text{ eV} \leq E_{res} \leq 4 \text{ keV}$ for the event in Figure 4.7 and $200 \text{ eV} \leq E_{res} \leq 3 \text{ keV}$ for the wave in Figure 4.8. In Figure 4.7, most of the energy bins (65 eV - 689 eV for EL and 136 eV - 3 keV for EH) of the PADs which undergo the greatest change across the wave are within the estimated resonant energy range. The PADs in Figures 4.8 and 4.9 are
Figure 4.8: This figure has the same format as Figure 4.7. The frequency range is $f > 1.0$ Hz and $\theta_{LB} = 25^\circ (155^\circ)$. The electron temperature anisotropies and other parameters can be found in Table 4.1.
Figure 4.9: This figure has the same format as Figure 4.8 except that the wave occurs roughly 2.5 seconds later. The frequency range is $0.45 \, \text{Hz} < f < 1.0 \, \text{Hz}$ and $\theta_{kB} = 49^\circ (131^\circ)$. The electron temperature anisotropies and other parameters can be found in Table 4.1.
also consistent with the resonant energies showing the greatest changes from the PAD at 16:31:32 UT in Figure 4.8 to the PAD at 16:31:44 UT in Figure 4.9. The average increase in pitch-angle of the electrons in this energy range across the waves would be consistent with pitch-angle diffusion were we observing the same distribution in time. It is difficult to say whether the strong anisotropies in the electron PADs downstream of the shocklets are a consequence of their traversal of the shocklets or if their downstream location with respect to the shocklet results in the isotropization through leakage. Leakage is the process whereby the downstream electrons with large pitch-angles are restricted to the downstream by the shocklet’s magnetic fields, while the lower pitch-angle electrons can move freely upstream [Larson et al., 1996; Thomsen et al., 1983b].

The electron heat flux, or more appropriately, the kinetic energy flux in the plasma rest frame, was calculated to more thoroughly examine the instability thresholds discussed by Gary et al. [1994, 1999]. Each 3DP electron distribution was first transformed into the solar wind frame, including the effects of the spacecraft potential, and then the first four moments of the distribution function were calculated. The heat flux, calculated from Equation 1.19, changed both in angle with respect to magnetic field and wave vector and in magnitude across each shocklet. The angle between the wave vectors and heat flux vectors calculated for each PAD change from $\sim 20^\circ$ to $\sim 27^\circ$ for 16:31:32-16:31:44 UT. Also, the magnitude of the heat flux changes from $\sim 49$ to $\sim 85$ keV cm$^{-3}$ km/s for 16:31:32-16:31:44 UT. Thus, the heat flux magnitude, angle between wave vector and heat flux, and angle between heat flux vector and magnetic field all peak in the 16:31:44-16:31:47 UT PAD which is just downstream of the peak amplitude of the shocklet in panel B of Figure 4.2.

Most studies of shocklets at the terrestrial bow shock focused on ion distributions. Shocklets were observed to have a location dependence in the terrestrial ion foreshock with respect to the sunward edge (boundary between red and yellow regions in Figure 1.5). The location dependence is due to the evolution of the observed ion distributions in the foreshock. Shocklets are observed in association with diffuse ion distributions, a characteristic distribution seen in deeper regions (i.e. farther from the sun) of the foreshock [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983]. Figure 4.10 shows examples of diffuse ions seen simultaneously with the shocklet in panel B of Figure 4.2. The plots are Pesa High Burst (PHB) distribution functions plotted with the horizontal axis
Figure 4.10: This figure is a series of contour plots of the ion distribution function from PHB from 16:31:29 UT to 16:31:47 UT. The high energy diffuse ions are seen as the semi-isotropic ring between 1000 and 2000 km/s. These distributions are observed simultaneously with the shocklet in panel B of Figure 4.2.
parallel to the ambient magnetic field in the plane created by the magnetic field and solar wind velocity. In each plot, the solar wind direction (black line) and shock normal direction (red line) are projected onto the distributions for reference.

### 4.3.1 Comparison of Electron Heating

The electron distributions at the unusual event show strong heating in the downstream region. The downstream region of the 04/06/2000 event had broad flattop distributions downstream, thought to result from strong current-driven ion-acoustic waves [Thomsen et al., 1983a]. The 12/10/1997 event showed weak flattop distributions downstream for a few seconds, followed by a Maxwellian that was hotter than the upstream distributions. The flattop distributions lasted for over an hour downstream of the 04/06/2000 IP shock.

Figure 4.11 shows the temperature anisotropies for both the halo and core electrons within $\pm 1$ hour of each of the five IP shocks examined. Each panel is labeled with a capital letter differentiate each event, where **A** is for 1996-04-03, **B** is for 1996-04-08, **C** is for 1997-10-24, **D** is for 1997-12-10, and **E** is for 2000-04-06. The level of unity (i.e. isotropy) is marked by a horizontal blue line. The red vertical bars in the right hand side of each panel denotes the magnitude of $\sim 0.2$ for quick comparison between each panel. The important thing to note is that the red bar is the smallest in both panels of the 2000-04-06 event, which shows that the changes in temperature anisotropies are the largest for that event. The 2000-04-06 event is also the only event to show a global increase in $T_{\perp h}/T_{\parallel h}$.

The halo electrons have strong anisotropies in the downstream region of the 04/06/2000 event. The halo electrons show strong heating perpendicular to the magnetic field with a remarkably low change in the parallel halo temperature. The preferential perpendicular heating of the halo electrons may be explained by the efficiency of the pitch-angle scattering discussed by Saito and Gary [2007] which showed a preferential efficiency with higher kinetic energy electrons. The reason for the low heating in the parallel halo electrons is not known at this time.

The fact that shocklets were observed upstream at only one of the five IP shocks raises the question of what characteristics of the 04/06/2000 event might lead to their generation. It is highly likely that the 04/06/2000 event is a supercritical shock (i.e.
Figure 4.11: A summary plot of temperature anisotropies for both the halo and core electrons for the five events of interest in Wilson III et al. [2009]. Each event is labeled with a capital letter. The left column is for the core electrons and the right column is the halo electrons. The blue horizontal line marks the 1.0 level and the red bar is used to show the relative amplitude (equal to ~0.2) of any change between each plot.
requires particle reflection for energy dissipation), which may explain why shocklets are observed upstream of this event and none of the others. Particle reflection could explain the difference in heating between the 04/06/2000 event and the other four. Recent observations [L.B. Wilson III et. al., Large amplitude electrostatic waves observed at a supercritical interplanetary shock, 2009] have shown evidence of the modified two stream instabilities discussed by Matsukiyo and Scholer [2006a] which were shown to strongly heat the electrons and ions. The core electrons show significant heating \( \frac{T_{c2}}{T_{c1}} \geq 3.5 \) which one would expect from an interaction with electrostatic waves \[ \text{Thomsen et al., 1985a} \] and/or electromagnetic modes like the modified two stream instability \[ \text{Matsukiyo and Scholer, 2006a} \]. Thus, ion reflection is likely playing a more significant role in energy dissipation than at the other lower Mach number events \[ \text{Thomsen et al., 1985a} \].

Figure 4.12 shows the perpendicular temperatures for both the halo and core electrons within ±1 hour of each of the five IP shocks examined. The format is same as in Figure 4.11 except the vertical red bars denoting the relative magnitudes are different between halo (\( \sim 30.0 \)) and core (\( \sim 3.0 \)). Again, the relative magnitude is the largest for the 2000-04-06 event compared to the other four events.

Figure 4.13 shows the parallel temperatures for both the halo and core electrons within ±1 hour of each of the five IP shocks examined. The format is same as in Figure 4.12. Here the vertical red bars denoting the relative magnitudes for the halo and core are \( \sim 30.0 \) and \( \sim 3.0 \), respectively. Again, the relative magnitude is the largest for the 2000-04-06 event compared to the other four events. Also notice that the 2000-04-06 event is the only event to show a global decrease in \( T_{\parallel h} \) as seen in the right hand panel E.

One can see from examining Figures 4.11 through 4.13 that there is a preference for perpendicular heating in the 2000-04-06 event in both the core and halo electrons. The change in core temperatures for the 2000-04-06 event is also well beyond the adiabatic limit, the only one of the five events shown in these figures (see Section 5.2 for more details). The differences in anisotropic heating is well illustrated in Figure 4.11.
Figure 4.12: A summary plot of perpendicular temperatures for both the halo and core electrons for the five events of interest in Wilson III et al. [2009]. The format is the same as in Figure 4.11 except the vertical red bars are different in magnitude between the halo and core.
Figure 4.13: A summary plot of parallel temperatures for both the halo and core electrons for the five events of interest in Wilson III et al. [2009]. The format is the same as in Figure 4.12.
4.4 Conclusions

The strongest event in our low frequency upstream waves study [Wilson III et al., 2009] was the most unusual of the five because it was the only event with shocklets. We observed 12 shocklets upstream ($\leq$ 1 hour of shock ramp) of the shock. Almost 80% of the shocklets had $\theta_{kB} \leq 45^\circ$, consistent with the cometary bow shock study by [Le et al., 1989] and terrestrial bow shock studies [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983; Russell et al., 1971]. It is likely that shocklets only occurred at the 04/06/2000 event because of its unusually high Mach number ($M_f \sim 4$). The high Mach number and quasi-perpendicular nature of the shock suggest that it is a supercritical shock, thus requiring ion reflection to dissipate energy. Ion reflection has been shown to be an important aspect of ULF wave generation in observations [Hoppe and Russell, 1983; Thomsen et al., 1983b; Meziane and D’Uston, 1998] and simulations [Omidi and Winske, 1990; Scholer et al., 2003; Tsubouchi and Lembege, 2004]. Shocklets are often seen in association with diffuse ion distributions [Hoppe et al., 1981; Hoppe and Russell, 1983] and the ULF waves observed by [Hoppe and Russell, 1983] were seen to steepen into shocklets with associated leading magnetosonic whistler wave packets when in regions of diffuse ions. Simulation studies have supported these observations and suggest that diffuse ions may be a necessary factor for ULF waves steepening into shocklets [Omidi and Winske, 1990; Scholer et al., 2003; Tsubouchi and Lembege, 2004]. Thus, the simultaneous observation of the diffuse ion distributions in Figure 4.10 and the wave in panel B of Figure 4.2 supports our hypothesis that these magnetic structures are in fact shocklets.

The major differences in electron moments between the 04/06/2000 event and the other four occurred primarily in the halo electrons. The only IP shock to show a strong global increase ($\sim$ a factor of 2) in $T_{\perp h}/T_{\parallel h}$ from upstream to downstream was the 04/06/2000 event. The relative heating between the electron halo and core ($T_{\parallel e}/T_{\parallel c}$) components is more dramatically affected in the 04/06/2000 event than any other studied decreasing by a factor of $\gtrsim 4$ across the IP shock (other events increase by $\leq 2$). The decrease is due to the large increase in $T_{\parallel c}$ and slight decrease in $T_{\perp h}$ across the shock. The $\beta_{\parallel c}$ is almost always $\leq 1.0$ upstream (within an hour of the ramp) of the 04/06/2000 event, while of the four other IP shocks examined, only the 10/24/1997 event has $\beta_{\parallel c} \leq$
1.0 anywhere upstream (for ~24 minutes immediately upstream of the ramp). The low
$\beta_{ec}$ may be a necessary condition for the excitation of shocklets. Recent observations
[Wilson III et al., 2010] have found evidence to suggest that the stronger heating of the
core electrons at the 04/06/2000 event may have resulted from the microinstabilities of
[Matsukiyo and Scholer, 2006a]. The instabilities are excited by reflected ions, known
to exist upstream of supercritical quasi-perpendicular shocks, interacting with the incident
solar wind. Thus the likely supercritical nature of the 04/06/2000 event makes it
a prime candidate for these instabilities. To summarize, the the 04/06/2000 event was
the most unusual in the study for the following reasons: 1) it had the lowest upstream
$\beta_e$, 2) the largest change in $T_i$ and $T_e$ across the shock, 3) the highest $M_f$, $n_2/n_1$, and
$U_n$, 4) only event to strong anisotropic preferential heating of the electrons, and 5) large
amplitude atypical electric field waveforms observed in the ramp regions.

The electron distributions showed a clear relationship between the core/halo temper-
ature anisotropy and shocklets. Regardless of how the halo electrons became anisotropic
$T_{\perp h}/T_{\parallel h}$ downstream of the shocklets, it exceeded the threshold estimated by [Gary et al.,
1994, 1999] for excitation of the whistler heat flux instability. One should also note
$T_{\perp h}/T_{\parallel h}$ increases more dramatically than $T_{\perp c}/T_{\parallel c}$ across almost every shocklet ob-
served with a leading magnetosonic whistler, consistent with the simulation results of
[Gary et al., 1994, 1999; Saito and Gary, 2007]. Thus it appears that the higher en-
ergy halo electrons may have experienced a more efficient pitch-angle diffusion than the
lower energy core. This suggests that the anisotropies often observed in the the ambient
solar wind electron distributions may be more unstable than previous estimates which
used the entire distribution to estimate $T_{\perp}/T_{\parallel}$. Another interesting observation is that
nearly every distribution within 30 seconds of each IP shock ramp exceeded the $T_{\parallel h}/T_{\parallel c}$
threshold for whistler heat flux instability estimated by [Gary et al., 1994].

In addition to particle distributions unstable to the whistler heat flux instability, an
electron heat flux was observed. The heat flux itself is the source of the free energy for
the instability while the halo electron anisotropies change the threshold and growth rate
of the instability. The the heat flux magnitude, angle between wave vector and heat flux,
and angle between heat flux vector and magnetic field all peak in the 16:31:44-16:31:47
UT PAD which is just downstream of the peak amplitude of the shocklet in panel B of
Figure 4.2. Also, the angle between the wave vectors and heat flux vectors calculated
for each PAD change from \(\sim 20^\circ\) to \(\sim 27^\circ\) for 16:31:32-16:31:44 UT. The absolute angle between the magnetic field and heat flux vector changes from \(\sim 9^\circ\) to \(\sim 24^\circ\) for the same PADs, consistent with the magnetosonic whistler pitch-angle scattering the heat flux carrying electrons [Gary et al., 1994].

The PADs downstream of the shocklets are suggestive of perpendicular heating and pitch-angle diffusion. However, it is unclear whether the unstable electron distributions seen in the downstream region play any role in the shocklet formation. If the unstable distributions do excite a whistler heat flux instability, the resultant waves could potentially propagate upstream of the steepened edge of the shocklets producing the observed wave train. The waves could prolong or increase the perpendicular electron heating, thus producing more unstable distributions. This would lead to a cyclical behavior of wave formation, propagation, and damping that would be self-reinforcing. However, it is also possible that electrons with smaller pitch-angles were able to return upstream of the shocklets while the higher pitch-angle electrons could not [Larson et al., 1996]. This could also explain the strong parallel anisotropies in the electron PADs observed upstream of the shocklets.

The other four IP shocks discussed herein showed characteristics typical of sub-critical to marginally-critical quasi-perpendicular shocks [Parris et al., 1993]. They each had precursor whistler waves upstream of the ramp and in the ramp. Almost all the precursor whistlers observed had \(20^\circ \leq \theta_{kB} \leq 45^\circ\), consistent with theory [Gary et al., 1994, 1999]. The precursor waves observed in this study actually show more similarities to the \(\sim 1\) Hz whistlers of [Hoppe et al., 1982; Sentman et al., 1983]. The waves are far more oblique than the previous observations of precursors at IP shocks observed by Russell et al. [1983] or the upstream whistlers at IP shocks observed by Tsurutani et al. [1983]. More than 90% of the waves had \(29^\circ \leq \theta_{kn} \leq 75^\circ\), thus it is not likely that they are phase standing. The waves are also seen in association with anisotropic electron distributions, though less oblique than previous observations of whistlers associated with anisotropic electrons [Sentman et al., 1983]. We found that the anisotropic electron distributions exceed the thresholds for the whistler heat flux instability estimated by Gary et al. [1994, 1999]. The threshold conditions for a whistler heat flux instability are met by almost all electron distributions within 30 seconds of every IP shock in this study.
In this chapter we presented observations on two classes of low frequency (0.25 Hz < f < 10 Hz) waves at five quasi-perpendicular IP shocks [Wilson III et al., 2009]. The first class of low frequency waves is a non-phase standing precursor whistler observed just upstream of four of the IP shocks examined in this study. The second class is a steepened magnetosonic wave, with a leading magnetosonic whistler wave train, called a shocklet. The shocklets and precursor whistlers are observed in association with electron distributions unstable to whistler heat flux and/or whistler anisotropy instabilities.
Table 4.1: Wind 3DP Electron Stats on 04/06/2000

<table>
<thead>
<tr>
<th>Start-End Time</th>
<th>$T_{ee}$ (eV)</th>
<th>$T_{he}$ (eV)</th>
<th>$T_{ee}/T_{he}$</th>
<th>$T_{\perp h}/T_{\parallel h}$</th>
<th>$n_{ee}$ (cm$^{-3}$)</th>
<th>$n_{he}$ (cm$^{-3}$)</th>
<th>$T_{\parallel h}/T_{\perp h}$</th>
<th>$\beta_{he}$</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Eesa Low</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>16:27:18-16:27:21</td>
<td>10.80</td>
<td>125.34</td>
<td>0.86</td>
<td>0.66</td>
<td>5.95</td>
<td>0.030</td>
<td>13.69</td>
<td>0.86</td>
</tr>
<tr>
<td>16:28:57-16:29:00</td>
<td>10.56</td>
<td>125.52</td>
<td>0.99</td>
<td>0.84</td>
<td>9.95</td>
<td>0.039</td>
<td>13.19</td>
<td>0.71</td>
</tr>
<tr>
<td>16:30:36-16:30:39</td>
<td>11.14</td>
<td>117.85</td>
<td>0.95</td>
<td>0.91</td>
<td>9.03</td>
<td>0.051</td>
<td>10.86</td>
<td>0.70</td>
</tr>
<tr>
<td><strong>Eesa Low Burst</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>16:31:32-16:31:35</td>
<td>10.02</td>
<td>117.29</td>
<td>0.70</td>
<td>0.55</td>
<td>5.02</td>
<td>0.026</td>
<td>13.42</td>
<td>1.16</td>
</tr>
<tr>
<td>16:31:35-16:31:38</td>
<td>10.40</td>
<td>120.33</td>
<td>0.73</td>
<td>0.59</td>
<td>4.74</td>
<td>0.026</td>
<td>13.07</td>
<td>1.02</td>
</tr>
<tr>
<td>16:31:38-16:31:41</td>
<td>10.28</td>
<td>127.57</td>
<td>0.73</td>
<td>0.55</td>
<td>5.16</td>
<td>0.023</td>
<td>14.53</td>
<td>1.14</td>
</tr>
<tr>
<td>16:31:41-16:31:44</td>
<td>11.54</td>
<td>129.47</td>
<td>0.94</td>
<td>0.84</td>
<td>6.08</td>
<td>0.039</td>
<td>12.03</td>
<td>0.79</td>
</tr>
<tr>
<td>16:31:44-16:31:47</td>
<td>11.16</td>
<td>126.61</td>
<td>1.02</td>
<td>1.02</td>
<td>10.42</td>
<td>0.050</td>
<td>11.37</td>
<td>0.75</td>
</tr>
</tbody>
</table>
Chapter 5

High Frequency Waves at Interplanetary Shocks

5.1 Statistical Study of High Frequency Waves at Collisionless Shocks

In a study of high frequency \((f \gtrsim 600 \text{ Hz})\) waves at 67 IP shocks with \(1 < M_f < 6\), Wilson III et al. [2007] observed three electrostatic (ES) wave modes in the ramp region of the IP shocks: Langmuir waves, IAWs, and solitary waves. Thirty-three of the 67 IP shocks had at least one waveform capture in the ramp region. The average number of TDS samples in the 33 IP shock ramps analyzed was \(\sim 3\). The IAWs were observed to have Doppler shifted frequencies ranging from \(f_{pi} \lesssim f \lesssim 10 \text{ kHz}\), where \(f_{pi}\) is the ion plasma frequency (\(\sim 500 \text{ Hz}\) in solar wind at 1 AU). The Langmuir waves are always above \(f_{pe}\), but typically the shift away from \(f_{pe}\) due to thermal effects is small.

Table 5.1 summarizes the statistics on the waves in the upstream, ramp, and downstream regions for the 33 IP shocks. The last three rows were normalized by the average duration for their respective regions: up/downstream \(\bar{\tau}_u = \bar{\tau}_d \simeq 3596 \text{ s}\), and ramp \(\bar{\tau}_r \simeq 8 \text{ s}\). Though it is known that the shock ramp is much shorter than 8 seconds, the use of 3s magnetic field data restricted our estimation of the ramp durations. To normalize our data, we take the number of events observed in each region, \(n_{\text{region}}\), and divide by the duration of each region to get an expected observation frequency, \(\eta_{\text{region}}\). Each \(\eta_{\text{region}}\)
is summed to give a total, $\eta_T$, which is then used to get a percentage by: $\eta_{\text{region}}/\eta_T \times 100\%$. Thus, we have a normalized percentage of occurrence per region. One can easily see the ramp region has the highest probability of wave occurrence. Roughly 87\% of all the waves seen in the 33 IP shock ramps were IAWs and 90\% of these were large ($\geq 5$ mV/m). The dominant wave mode upstream is large amplitude Langmuir waves ($\sim 53\%$ of all the large waves upstream for the 33 shocks with ramp waves), consistent with other observations at IP shocks \cite{Fitzenreiter2003,Thejappa2000} and the terrestrial bow shock \cite{Bale1997}. The downstream region is dominated by IAWs ($\sim 82\%$ of all the downstream waves for the 33 IP shocks); however, $\sim 68\%$ of those were small IAWs. The dominant wave mode (for $f_{\pi} < f < f_{pe}$) in the ramp region is large IAWs. The most striking observation is that essentially 100\% of the large amplitude IAWs occur in the ramp when normalized by time, supporting theories on dissipation in low Mach number shocks \cite{Matsukiyo2006b,Shimada2000}. Wilson III et al. \cite{Wilson2007} also showed that the amplitude of large amplitude ($\geq 5$ mV/m peak-to-peak) IAWs increased with increasing Mach number and the shock strength, consistent with larger shock strengths causing larger cross-field currents which may provide free energy for wave generation, consistent with theory \cite{Gary1981}.

Because of the high probability of occurrence for large amplitude IAWs in the ramp regions, the dependence of wave amplitude ($|E_{xy}|$) on various shock parameters including: $\theta_{Bn}$, $n_2/n_1$, $M_A$, $M_f$, $M_{es}$, and $T_e/T_i$ were tested for only the largest IAWs in each

<table>
<thead>
<tr>
<th>All Waves</th>
<th>Large ($\geq 5$ mV/m)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Lang. IAW Solitary</td>
<td>Lang. IAW Solitary</td>
</tr>
<tr>
<td>$\ddagger$ Up</td>
<td>52</td>
</tr>
<tr>
<td>$\ddagger$ in Ramps</td>
<td>5</td>
</tr>
<tr>
<td>$\ddagger$ Down</td>
<td>8</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Normalized Probability of Occurrence</th>
</tr>
</thead>
<tbody>
<tr>
<td>Up</td>
</tr>
<tr>
<td>Ramp</td>
</tr>
<tr>
<td>Down</td>
</tr>
</tbody>
</table>

Table 5.1: Wave Occurrence by Regions for the 33 IP Shocks with Ramp TDS Samples
Figure 5.1: Correlation plot for the largest IAWs in the ramp regions of 33 IP shocks relating wave amplitude to the shock compression ratio ($n_2/n_1$).

shock ramp. The largest correlation is between $|E_{xy}|$ and $n_2/n_1$ (seen in Figures 5.1 and 5.2), consistent with larger shock strengths resulting in larger cross-field currents which may provide free energy for wave generation. A slightly weaker correlation is seen between $|E_{xy}|$ (mV/m) and the fast mode Mach number ($M_f$). Similar correlations were seen for the other Mach numbers. No correlation between $T_e/T_i$ and wave amplitude was observed. Polarization vectors for IAWs are consistent with a cross–field current source. Wilson III et al. [2007] found the waves to produce resistivities of $\sim$1–856 $\Omega\cdot$m, using Equation 1.44 which is roughly $10^7$ times greater than the classical Spitzer electron-ion collision estimates. One should also note that numerous simulation studies have found Equations 1.44 and 1.45 to under estimate the actual resistivity due to IAWs by up to 3 orders of magnitude [Dyrud and Oppenheim, 2006; Hellingar et al., 2004; Petkaki et al., 2003, 2006; Petkaki and Freeman, 2008; Watt et al., 2002; Yadav et al., 2009; Yoon and Lui, 2007]. These results are consistent with the theory that wave-particle interactions are important for dissipation at subcritical shocks [Matsukiyo and Scholer, 2006b, a].
5.2 Analysis of Microphysics

[Work presented in Wilson III et al. 2010]

This study is motivated by the examination of an unusual supercritical IP shock observed by the Wind spacecraft on 04/06/2000 [Wilson III et al., 2009]. The previous study found unusual waveforms, called shocklets, upstream of the shock and strong electron heating across the shock ramp. This section will focus on the unusually high electron heating observed in the ramp region. Strong perpendicular electron heating could be evidence of resistive heating due to wave-particle interactions [Thomsen et al., 1985a]. Within \( \pm 5 \) seconds, corresponding to \( \pm 14 \frac{c}{\omega_{pi}} \) (upstream ion inertial lengths where \( c \) is the speed of light and \( \omega_{pi} \) is the average upstream ion plasma frequency) of the shock crossing, four waveform captures were obtained by the Wind spacecraft, two of which have frequency spectra consistent with the ECDI and two with solitary wave
signatures. The results of this section were presented by Wilson III et al. [2010].

5.2.1 Data Sets and Analysis

Waveform captures were obtained from the Wind/WAVES instrument [Bougeret et al., 1995], through the time domain sampler (TDS) receiver, which provides a \( \sim 17 \) ms waveform capture of 2048 points (from here on, a waveform capture is called a TDS sample). TDS samples utilized herein provide two components of the electric field in the XY-GSE plane, called \( E_x \) and \( E_y \). The spin axis (roughly the Z-GSE component) electric field was not sampled for these TDS samples. The TDS buffer stores and evaluates waveforms based upon their amplitude. Thus if small amplitude waves are observed, they will not be stored and transmitted to the ground if larger amplitude waves fill the buffer. We define \( |E_{xy}| = \sqrt{E_x^2 + E_y^2} \) as the maximum peak-to-peak (pk-pk) amplitude of the TDS samples. We only measure two components of the electric field vector, roughly in the XY-GSE plane. This limits our ability to measure electric fields parallel(perpendicular) to the magnetic field direction. Thus, when rotating the electric field components into magnetic field-aligned coordinates (FACs), we rotate the fields by the angle between the positive X-antenna and the XY-GSE projection of the magnetic field vector. In FACs, we define the subscript \( \parallel (\perp) \) as the direction parallel(perpendicular) to the XY-GSE projection of the magnetic field direction. However, if a significant fraction of the magnetic field is in the Z-GSE direction, the measured \( E_\parallel \) may be significantly different from the total electric field along the magnetic field. The converse applies to the measured values of \( E_\perp \). The lack of full 3D electric field measurements can influence the determination of polarization as well.

To analyze the waveforms dynamically in time and frequency, we computed the Morlet wavelet transform [Torrence and Compo, 1998b] for the four waveforms examined in this study. The wavelet transform has a number of distinct advantages over windowed FFTs, but the two most important for our analysis are the resolved dynamic power spectra at low frequencies (\( \sim 200 \) to 1000 Hz) and the analysis of non-stationary power intensifications at many different frequencies simultaneously [Torrence and Compo, 1998a]. Windowed FFTs have imposed periodicity and have a single time scale while wavelets do not. However, two sanity checks are performed to expose noise or numerical artifacts
due to the interpolation routines. The cone of influence is a calculation done to determine the time-frequency region of the wavelet transform which is subject to edge effects. Thus, values above (with respect to frequency and time) this line can be trusted to not have artificial effects due to the finite time range of the data. The second calculation determined the 95% confidence level (also called "significant at the 5% level"). The 95% confidence level outlines regions of the wavelet transform which enclose intensities above 95% of the data. The 95% confidence level can be calculated from a theoretical red or white noise spectrum, but here we used the actual waves since they are clearly well above the background noise level of the solar wind ($\lesssim 0.1$ mV/m at 1 AU). Both significance tests are used to increase our ability to make quantitative analysis using the wavelet transforms.

The magnetic field instrument on board Wind [Lepping et al., 1995] is composed of dual triaxial fluxgate magnetometers. High time resolution (HTR) magnetic field data, sampled at $\sim 11$ samples/s, were used to define the ramp region, or transition region, of the IP shock, as well as to identify low frequency magnetic fluctuations upstream of the ramp [Wilson III et al., 2009]. The ramp is defined as the interval from the point of lowest magnetic field immediately preceding the discontinuity (in magnetic field amplitude) to the point of highest magnetic field immediately following the discontinuity [Farris et al., 1993].

Low energy ($<30$ keV) electron and ion distributions were obtained from the Wind 3DP EESA and PESA particle detectors [Lin et al., 1995]. The EESA Low (EL and ELB in burst mode) instrument can measure electrons at 15 different energies from a few eV to a little more than a keV. The PESA Low (PL and PLB in burst mode) instrument measures ions at 14 different energies from as low as 100 eV to as high as 10 keV (typical range in the solar wind is 700 eV to 6 keV). The instrument is used primarily for bulk solar wind properties like ion velocity, density, and temperature. The PESA High (PH and PHB in burst mode) instrument measures ions at 15 different energies from as low as 80 eV to as high as 30 keV (typical range in the solar wind is 500 eV to 28 keV). The time resolution of each instrument depends on whether the instruments are in burst mode or not. In burst mode, both EESA and PESA instruments return full three dimensional particle distributions every three seconds (\~1 spin period). Electron and ion distributions were examined for possible wave free energy sources and evidence
of heating. Estimates of the electron temperature anisotropies in both the cold dense core (subscript c) and the hotter more tenuous halo (subscript h) can be obtained from full 3D electron distributions. The method for determining the break energy between halo and core electrons is outlined by Wilson III et al. [2009].

High energy (>30 keV) electron and proton measurements were obtained from three pairs of double-ended solid state telescopes (SSTs), each with a pair or triplet of closely stacked silicon semiconductor detectors [Lin et al., 1995]. The SSTs provide a full 4π steradian coverage with a 22.5° × 36° angular resolution and ΔE/E ≈ 0.2 energy resolution. One side of each detector is covered in a thin lexan foil (SST Foil) to stop protons up to ∼400 keV while leaving electrons relatively unaffected. The opposite end of the detector is left open (SST Open) using a common broom magnet to sweep away electrons below ∼400 keV while leaving the protons relatively unaffected. Thus, in the absence of any higher energy (penetrating) particles, the SST Foil counts only electrons and the SST Open counts only ions.

The relevant shock parameters, determined by Kasper [2007], are the shock normal angle, θ_Bn, fast mode Mach number, M_f, shock normal vector, n̂, upstream normal flow velocity in the shock frame, U_n, upstream solar wind velocity, V_sw, and shock strength, N_i2/N_i1. The values for the 04/06/2000 event are: M_f ∼ 4, θ_Bn ∼ 68°, U_n ∼ 278 km/s, V_sw ∼ <-370,-18,-22> km/s (GSE coordinates), n̂ ∼ <-0.98,-0.08,-0.16> (GSE coordinates), and N_i2/N_i1 ∼ 4.

5.2.2 Observations

Figure 5.3 shows a plot of the magnitude (first panel) and the GSE components of the magnetic field (labeled with color coded component letters, second panel) for the 04/06/2000 event between 16:32:03 and 16:32:15 UT. The shock arrival time, or middle of the magnetic ramp, is 16:32:09.2 UT. The shaded regions correspond to the time ranges of each particle distribution found in Figure 5.4. The vertical color-coded lines labeled with capital letters indicate the locations of the TDS samples shown in Figure 5.5.

Figure 5.4 shows four ion (top row) and electron (bottom row) distribution functions plotted with the horizontal axes corresponding to the direction parallel to the magnetic field. The plots are projected into the plane produced by the solar wind velocity and the
Figure 5.3: The plot shows the magnitude and the GSE coordinates of the magnetic field for the IP shock on 04/06/2000. The shaded regions correspond to the time ranges of each particle distribution found in Figure 5.4. The vertical color-coded lines indicate the locations of the TDS samples shown in Figure 5.3, labeled with the respective capital letters. The time range of the plot is 16:32:02-16:32:16 UT and the shock arrival time is 16:32:09.2 UT.

local magnetic field. The electron and ion velocity ranges in the plots are ±20,000 km/s and ±2,500 km/s, respectively. The phase space density color scales for each instrument are the same for the four different distributions in each row.

A summary of the relative changes in moments for the four electron distributions shown in Figure 5.4 can be found in Table 5.2. The relative change in any given moment is defined as $\Delta Q_s = ((Q_f - Q_i)/Q_i)_s \times 100\%$, where the subscripts $f$ and $i$ represent the final and initial state, respectively and the subscript $s$ represents the particle species (e.g. core). The final and initial state for the calculations in Table 5.2 are specific only to the four distributions shown in Figure 5.4. Thus, the final state refers to the distribution starting at 16:32:12 UT and the initial state refers to the distribution starting at 16:32:03 UT. The top half of the table shows the values for $\Delta T_{es}$, $\Delta T_{ls}$,
Figure 5.4: Four particle distributions from both the PHB (top row) and ELB (bottom row) detectors on the Wind 3DP instrument. The horizontal axis of each plot is parallel to the magnetic field and the vertical axis is perpendicular in the plane created by the solar wind velocity and magnetic field direction. Each pair of samples, ELB and PHB, are outlined by a colored box which corresponds to the shaded regions in Figure 5.3. The black solid line represents the projection of the solar wind velocity and the red dotted line represents the shock normal vector projection. Note the sample times are the same for each instrument.

\[ \Delta T_{\parallel, s}, \Delta (T_{\perp}/T_{\parallel})_s \] for the core, halo, and entire distribution (i.e. core and halo combined). The bottom of Table 5.2 shows the global changes (downstream, subscript 2, over upstream, subscript 1) across the 04/06/2000 event for the core, halo, and entire distribution. Here, the subscripts 1 and 2 refer to asymptotic values for the upstream and downstream estimated by averaging the quantities over 10 minute intervals in each region (i.e. outside of the time period shown in Figure 5.3 and calculated for more than just the four distributions shown herein). The quantities for global changes shown in Table 5.2 are the average, perpendicular, and parallel temperatures.

The core electrons show the strongest heating in bulk, parallel, and perpendicular, components while the halo dominates in the change in temperature anisotropy. Note also that there appears to be a preference toward perpendicular heating as indicated by
Table 5.2: Wind 3DP ELB stats across the 04/06/2000 IP shock

<table>
<thead>
<tr>
<th>Species</th>
<th>$\Delta T_{es}$</th>
<th>$\Delta T_{\perp}$</th>
<th>$\Delta T_{\parallel}$</th>
<th>$\Delta (T_{\perp}/T_{\parallel})_s$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Eesa Low Burst (Core/Halo)</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Core</td>
<td>150%</td>
<td>174%</td>
<td>110%</td>
<td>30%</td>
</tr>
<tr>
<td>Halo</td>
<td>42%</td>
<td>67%</td>
<td>5%</td>
<td>58%</td>
</tr>
<tr>
<td>Eesa Low Burst (Entire Distribution)</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Entire</td>
<td>137%</td>
<td>164%</td>
<td>96%</td>
<td>35%</td>
</tr>
</tbody>
</table>

Temperature Anisotropies ($T_{\perp}/T_{\parallel}$)  

<table>
<thead>
<tr>
<th>Time</th>
<th>$T_{\perp}/T_{\parallel}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>16:32:03</td>
<td>0.81</td>
</tr>
<tr>
<td>16:32:06</td>
<td>0.75</td>
</tr>
<tr>
<td>16:32:09</td>
<td>0.73</td>
</tr>
<tr>
<td>16:32:12</td>
<td>1.06</td>
</tr>
</tbody>
</table>

Global Changes Across Shock  

<table>
<thead>
<tr>
<th>Time</th>
<th>$T_{es}/T_{es}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>16:32:03</td>
<td>3.3</td>
</tr>
<tr>
<td>16:32:06</td>
<td>3.8</td>
</tr>
<tr>
<td>16:32:09</td>
<td>2.5</td>
</tr>
</tbody>
</table>

$\Delta T_{es}$ for all three electron components. The middle part of Table 5.2 shows the specific values of $(T_{\perp}/T_{\parallel})_s$ for the core, halo, and entire energy range for each distribution shown in Figure 5.4. Notice that all three temperature anisotropies increase across the shock ramp, with the core increasing by $\gtrsim 30\%$ and the halo by $\gtrsim 58\%$. The global changes across the shock show a preference for perpendicular heating as well, increasing by almost a factor of four (see bottom of Table 5.2).

To be more quantitative, let us define the following parameter, $\alpha$, used by Omidi and Winske [1990] as a test of whether or not the shock is heating the particles adiabatically:

$$\alpha \equiv \left(\frac{T_{es,2}}{T_{es,1}}\right) - \left(\frac{N_2}{N_1}\right)^{\gamma-1}$$

(5.1a)

$$\alpha_\perp \equiv \left(\frac{T_{es,2\perp}}{T_{es,1\perp}}\right) - \left(\frac{N_2}{N_1}\right)^{\gamma-1}$$

(5.1b)

where $T_{es,j}$ is the asymptotic estimate of the electron temperature in for the regions defined by the subscript, $j$, which represents the upstream ($j = 1$) and downstream ($j = 2$), $T_{es,j,\perp}$ are the perpendicular asymptotic estimates of the electron temperature, $N_j$.
are the asymptotic estimates of the particle density, and \( \gamma \) is the ratio of specific heats (here we used 5/3). If the factor \( \alpha \) is equal to zero, then the electrons were heated adiabatically. If the \( \alpha > (\leq) 0 \), then the electrons were over(under) heated across the shock.

All the asymptotic values of interest are given in Table 5.3, for comparison with the same definition for the \( \Delta Q \) quantities as in Table 5.2. As one can see, the only shock with positive values of \( \alpha \) for the entire electron distribution is the 2000-04-06 event. Oddly enough, the two weakest events, 1996-04-03 and 1996-04-08, have \( \alpha > 0 \) for both the core and halo electrons but \( \alpha < 0 \) when the entire distribution is considered. Although the values of \( \Delta B \) and \( \Delta T_{\perp,e} \) for the 2000-04-06 event are roughly the same, this result is unique to this event.

The differences in electron heating between the 04/06/2000 event and the four more-

<table>
<thead>
<tr>
<th>Species</th>
<th>( \Delta T_{cs} )</th>
<th>( \Delta T_{\perp,s} )</th>
<th>( \Delta T_{\parallel,s} )</th>
<th>( \Delta B )</th>
<th>( \alpha )</th>
<th>( \alpha_{\perp} )</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>1996-04-03 Event</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Core</td>
<td>45%</td>
<td>46%</td>
<td>45%</td>
<td>0.15</td>
<td>0.15</td>
<td></td>
</tr>
<tr>
<td>Halo</td>
<td>134%</td>
<td>125%</td>
<td>153%</td>
<td>1.03</td>
<td>0.95</td>
<td></td>
</tr>
<tr>
<td>Entire</td>
<td>-14%</td>
<td>-15%</td>
<td>-13%</td>
<td>57%</td>
<td>-0.45</td>
<td>-0.45</td>
</tr>
<tr>
<td><strong>1996-04-08 Event</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Core</td>
<td>54%</td>
<td>48%</td>
<td>65%</td>
<td>0.12</td>
<td>0.07</td>
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<tr>
<td>Halo</td>
<td>119%</td>
<td>108%</td>
<td>144%</td>
<td>0.78</td>
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<tr>
<td>Entire</td>
<td>-15%</td>
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<td><strong>1997-10-24 Event</strong></td>
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<tr>
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<td>57%</td>
<td>49%</td>
<td>-0.28</td>
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<tr>
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<tr>
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<td>92%</td>
<td>-0.79</td>
<td>-0.72</td>
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<td><strong>1997-12-10 Event</strong></td>
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<tr>
<td>Core</td>
<td>12%</td>
<td>13%</td>
<td>9%</td>
<td>-0.72</td>
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<tr>
<td>Halo</td>
<td>44%</td>
<td>38%</td>
<td>55%</td>
<td>-0.40</td>
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<tr>
<td>Entire</td>
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<td>6%</td>
<td>2%</td>
<td>121%</td>
<td>-0.79</td>
<td>-0.78</td>
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<td><strong>2000-04-06 Event</strong></td>
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<tr>
<td>Core</td>
<td>228%</td>
<td>281%</td>
<td>153%</td>
<td>0.83</td>
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<tr>
<td>Halo</td>
<td>37%</td>
<td>84%</td>
<td>-23%</td>
<td>-1.08</td>
<td>-0.61</td>
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<tr>
<td>Entire</td>
<td>209%</td>
<td>267%</td>
<td>128%</td>
<td>266%</td>
<td>0.63</td>
<td>1.22</td>
</tr>
</tbody>
</table>

typical events are: 1) a global decrease in \( T_{h,\parallel} \) across the shock, 2) an increase by over a
factor of three in $T_e$ and $T_i$ across the shock, 3) a global increase in both $T_{c,\perp}/T_{c,\parallel}$ and $T_{h,\perp}/T_{h,\parallel}$ across the shock, 4) a global decrease in both $T_{h,\parallel}/T_{c,\parallel}$ and $T_{h,\perp}/T_{c,\perp}$ across the shock, and 5) sustained flattop electron distributions downstream for over an hour. We suggest that these differences are the result of wave-particle heating.

Figure 5.5 shows four TDS samples plotted in FACs at the times shown by the vertical lines in Figure 5.3. For each TDS sample, the magnetic field estimates were determined by averaging the HTR MFI data over the time range of the TDS sample. The angle of the magnetic field unit vector out of the XY-GSE plane, $\theta_{Bxy}$, for each TDS sample is: $+10^\circ$ for A, $+10^\circ$ for B, $-10^\circ$ for C, and $-31^\circ$ for D. The relative scales for both $E_\parallel$ (shown in red) and $E_\perp$ (shown in blue) are shown with the vertical arrows in each plot (e.g. wave D is $\sim 160$ mV/m pk-pk for $E_\perp$). The peak values of $|E_{xy}|$ for the four TDS samples are $\sim 39$, $\sim 35$, $\sim 20$, and $\sim 166$ mV/m. Note that these values of $|E_{xy}|$ are a lower bound to the actual maximum amplitudes of the waves since, as discussed in Section 5.2.1, we only measure two components of the electric field. Below the four waveforms are their associated hodograms. The time ranges of the hodograms are defined by the magenta boxes overlaying the waveforms. The solid green line corresponds to the XY-GSE projection of the shock normal vector, $n_{xy}/|n|$. The projection of the shock normal vector was scaled to the maximum value of $|E_{xy}|$ for each TDS sample for ease of comparison to the polarization of the electric fields.

Each panel in Figure 5.6 represents a Morlet wavelet transform power spectral density plot of a single component of a waveform from Figure 5.5, $E_\parallel$ on the left and $E_\perp$ on the right. In each panel, two horizontal lines, a bowl shaped line, and multiple contours are plotted. The two horizontal lines correspond to the fundamental and first harmonic of the electron cyclotron frequency. The cyclotron frequency and first harmonic were calculated by interpolating the magnitude of the magnetic field over the duration of each TDS sample. The bowl-like line corresponds to the cone of influence and the contours correspond to the 95% confidence level. The top three rows share the same power range of 0.01 to 1000 (mV/m)$^2$/Hz for the wavelets, while waveform D ranged from 0.05 to 45000 (mV/m)$^2$/Hz because of its much larger amplitude.

Further examination of the wavelet transforms of waveforms A through D show enhanced power near the fundamental and first harmonic of the electron cyclotron frequency ($f_{ce}$) for $E_\perp$ and $E_\parallel$. Half-integer harmonic power intensifications are most easily
Figure 5.5: Four waveform captures in and around the ramp region of the 04/06/2000 IP shock. The waves in Panels A (16:32:09.380 UT) and B (16:32:09.447 UT) are trains of steepened waves while the waves in Panels C (16:32:09.886 UT) and D (16:32:12.498 UT) are solitary-like waves. The top panels of each waveform (red) corresponds to $E_\parallel$ and the bottom (blue) corresponds to $E_\perp$. Their respective peak-to-peak amplitudes are marked by the vertical black arrows. Below the four waveforms are their respective hodograms labeled with the same capital letters corresponding to the time range defined by the magenta boxes in each TDS sample. The solid green line in each hodograms represents the XY-projection of the shock normal vector in FACs.
Figure 5.6: The Morlet wavelet power spectrum for the four waveforms in Figure 5.5, labeled A through D accordingly. The left column is for $E_\parallel$ and the right for $E_\perp$. The top three rows share the same spectral range of 0.01 to 1000 for the wavelets while waveform D ranged from 0.05 to 45000, while all four have been normalized to conserve energy in the wavelet transformations. In each panel, two horizontal lines mark the fundamental and first harmonic of the electron cyclotron frequency. The bowl-like line plotted in each spectra marks the cone of influence while the contours mark the 95% confidence levels [Torrence and Compo, 1998a].
seen in the $E_\parallel$ component of waveform $B$, outlined by the black box in Figure 5.6. However, the large amplitude of the higher frequency components makes a detailed analysis difficult for these events. Figures 5.7 through 5.10 present examples of snapshots of windowed FFTs for the four waveforms from Figure 5.5 to provide a complementary way to examine the power spectra. The left-hand column of panels shows $E_\parallel$, the power spectra (mV/m$^2$/Hz) versus frequency (kHz) for the orange box, and the power spectra versus frequency for the black box, respectively. The right-hand column shows the same for $E_\perp$. The vertical lines in the power spectra correspond to integer (green) and half-integer (purple) multiples of the electron cyclotron frequency. All power spectra plots in Figures 5.7 and 5.8 range from $10^{-5}$ to $4 \times 10^{-1}$ while in Figures 5.9 and 5.10 they range from $10^{-6}$ to $2 \times 10^0$.

Figures 5.7 through 5.10 show that waveforms $A$ through $D$ have mixtures of integer half-integer multiples of the electron cyclotron frequency, $f_{ce}$, that change throughout the TDS samples. Note that the maximum Doppler shifted IAW frequency for waveform $A$ ($B$) was estimated to be roughly 7.5 kHz (8.5 kHz), assuming wavelengths consistent with the measurements of Fuselier and Gurnett [1984], and the background noise level for this event is roughly $10^{-6}$ mV/m$^2$/Hz. The images clearly illustrates how the waveforms shift power between integer and half-integer harmonics of $f_{ce}$ during the duration of the TDS samples. If we calculate the power spectrum using the entire time range for each TDS sample, the shifting peaks smear together into one peak. Thus, the combination of wavelet and windowed FFT analysis shows that waveforms $A$ through $D$ exhibit mixtures of integer and half-integer multiples of $f_{ce}$ throughout the samples. Waveforms $C$ and $D$ differ from waveforms $A$ and $B$ in that they do not show broad power enhancements above 4 kHz. The difference is likely due to the fact that waveforms $A$ and $B$ are composed of IAWs and cyclotron harmonic waves whereas the power spectra for waveforms $C$ and $D$ are dominated primarily by the solitary waves. Also, waveforms $A$ and $B$ often have $E_\perp > E_\parallel$, consistent with cyclotron harmonic or Bernstein-like waves [Usui et al., 1999]. Note that $1/f_{ce} \sim 1.95$ ms (1.62 ms) for waveform $C$ ($D$). Also, the solitary wave in waveform $C$ near 4 ms (8 ms) has a period of $\sim 1.33$ ms (1.44 ms), slightly smaller than the cyclotron period. The solitary wave in waveform $D$ near 5 ms (8 ms) has a period of $\sim 1.56$ ms (1.67 ms), almost identical to the cyclotron period. Thus, the Bernstein-like emissions are still simultaneously observed with the solitary waves.
Figure 5.7: Example snapshots of windowed FFTs of waveform A from Figure 5.5. The left-hand column of panels shows $E_{\parallel}$, the power spectra (mV/m$^2$/Hz) versus frequency (kHz) for the orange box, and the power spectra versus frequency for the black box, respectively. The right-hand column shows the same for $E_{\perp}$. The vertical lines in the power spectra correspond to integer (green) and half-integer (purple) multiples of the electron cyclotron frequency. All power spectra plots in this figure range from $10^{-5}$ to $4 \times 10^{-1}$. 
up to \( \sim 14 \ c/\omega_{pi} \) away from the shock ramp. These integer and/or half-integer harmonic intensifications are consistent with previous observations of cyclotron harmonic or Bernstein-like waves [Usui et al., 1999].

Figure 5.11 shows the 2D Hammer-Aitoff projection of 27 and 40 keV electrons from

![Figure 5.8: Example snapshots of windowed FFTs of waveform B from Figure 5.5. The format is the same as Figure 5.7. All power spectra plots in this figure range from \( 10^{-5} \) to \( 4 \times 10^{-1} \).](image)

the SST Foil instrument on Wind between 16:31:01 UT and 16:32:15 UT. The Hammer-Aitoff projections display a full \( 4\pi \) steradian of the measured particles. The data are plotted in units of phase-space density \( (s^3 \text{ cm}^{-3} \text{ km}^{-3}) \) with ranges of \( 5 \times 10^{-22} \) (purple) to \( 2.2 \times 10^{-21} \) (red) for the 27 keV electrons and \( 5 \times 10^{-22} \) (purple) to \( 1.1 \times 10^{-21} \) (red) for the 40 keV electrons. Different scales were used for the two energies to highlight the intensification near the center of each plot. The SST Foil distributions are sampled every 12-13 seconds (\( \sim 4 \) spin periods) for this event. Four symbols are projected onto the 3D
maps corresponding to relevant vectors including: +, the magnetic field direction, ◊, the anti-parallel magnetic field direction, ●, the solar wind velocity direction (roughly negative X-GSE for this event), and △, the shock normal vector. The dotted lines denote 30° increments in both latitude and longitude. The beam-like feature, observed in every SST distribution of Figure 5.11 parallel to the magnetic field is also intermittently observed for nearly 20 minutes upstream of the shock ramp (not shown). It is not clear whether the beam-like feature is a contributing source of free energy for the solitary waves, but electron beams have been associated with solitary wave observations in many studies (e.g. Ergun et al. [1998a]). Also, a beam-like feature roughly 30° to the right of the shock normal and solar wind direction is observed in both energies with strong enhancements in the last two panels. Correspondingly, there is another beam-like enhancement roughly 180° to the right of the beam-like feature near the shock normal and solar wind directions, suggesting they are contaminated data bins.

5.2.3 Discussion

Because the full three dimensional electric field is not measured, we can infer some properties of the waves but cannot definitively determine the polarization. Recall that θ_{Bxy} is not zero for any of the four waveforms. If we assume the solitary waves propagate along the magnetic field, the relative magnitude of \( E_\perp \) to \( E_\parallel \) should be smaller for waveform D than C. The reason is that we should be measuring a unipolar \( E_\perp \) and bipolar \( E_\parallel \), but the observation of a bipolar \( E_\perp \) we believe is the result of not measuring \( E_z \). Thus, the smaller bipolar \( E_\perp \) to \( E_\parallel \) in waveform D compared to C is likely due to the larger \( \theta_{Bxy} \).

Typical TDS samples in most IP shock ramps are consistent with Doppler shifted IAWs [Wilson III et al., 2007]. For comparison to waveforms A and B, Figure 1.14 shows an example of a typical IAW observed in the 04/08/1996 IP shock ramp, one of the four lower Mach number shocks examined by Wilson III et al. [2009]. The parallel/perpendicular components are plotted in red/blue with the associated hodogram to the right. The hodogram has the XY-GSE projection of the shock normal vector plotted as a solid green line, as in the hodograms of Figure 5.5. The wavelet transforms for the parallel and perpendicular components are plotted below the waveform captures. The IAW wavelet shows a strong isolated band of power between \( \sim 2-5 \) kHz for both
Figure 5.9: Example snapshots of windowed FFTs of waveform C from Figure 5.5. The format is the same as Figure 5.7. All power spectra plots in this figure range from $10^{-6}$ to $2 \times 10^0$. 

Electric Field vs. Time on: 04-06-2000 at: 16:32:09.886 UT

Frame: 238

$\sim 15 \text{ mV/m}$

Frame: 665

$\sim 15 \text{ mV/m}$

Frame: 238

Frame: 665
Figure 5.10: Example snapshots of windowed FFTs of waveform D from Figure 5.5. The format is the same as Figure 5.7. All power spectra plots in this figure range from $10^{-6}$ to $2 \times 10^0$. 
components. This isolated, relatively narrow, band of emission is typical of the IAWs observed at the four lower Mach number IP shocks. The wavelets for waveforms A through D are clearly different, supporting our argument that these waveforms are not simple IAWs, rather mixtures of multiple modes. We also examined snapshot FFTs of the IAWs in the four more-typical IP shock ramps and found that the waves show slight enhancements below \( \sim 1 \text{ kHz} \), but no noticeable enhancements near integer and half-integer harmonics were observed. Waveforms A and B are quite different than the typical IAWs observed at IP shocks [Wilson III et al., 2007] but similar to the waves observed by Hull et al. [2006]. However, we argue that these waves are not simple IAWs for the following reasons: 1) neither waveform electric field component oscillates symmetrically about zero; 2) both waveforms have mixtures of frequencies consistent with IAWs and electron cyclotron harmonics in their power spectra; and 3) the polarizations are correlated with the shock normal vector, not the magnetic field.

The 04/06/2000 event was the only shock of the five studied in detail to show strong sustained core ion heating across the shock. Although the PL detector has a narrow field of view \( \sim 180^\circ \times 14^\circ \) compared to PH \( \sim 360^\circ \times 14^\circ \), which could limit the accurate measurement of the solar wind in the immediate downstream region of strong shocks. We determined the relative accuracy of the PL measurements by comparing the downstream density measured by the PL detector to the downstream density estimated from the plasma line (indicative of local density) seen by the WAVES thermal noise receiver [Bougeret et al., 1995]. We only use the temperature increase determined from the PL detector as a qualitative proxy for the bulk ion temperature increase. The sustained ion temperature increase, measured with PL, across the 04/06/2000 shock was roughly a factor of 7 for \( T_{i,\perp} \) and 5 for \( T_{i,\parallel} \) with spikes in downstream temperatures exceeding factors of 8 and 6 respectively, consistent with previous observations [Thomsen et al., 1985a]. However, the the four more-typical events showed ion temperature increases that were less than a factor of three. The PH distributions showed qualitatively similar changes.

A distinct difference in both electron and ion properties is observed across the 04/06/2000 shock compared to the other four events of Wilson III et al. [2009]. The ions show strong heating in both the bulk of the ion distribution and high energy tail (observed as diffuse ions for the 04/06/2000 event), while all four more-typical events only
Figure 5.11: Hammer-Aitoff projections of SST Foil observations from 16:31:01 UT to 16:32:15 UT of field-aligned enhancements of 27-40 keV electrons. The + symbol represents the magnetic field-aligned direction, diamond the anti-parallel field direction, asterisks is the solar wind direction, and the triangle represents the shock normal direction. The left column plots the 27 keV energy bin while the right plots the 40 keV energy bin. All images in each column have the same color scales, shown at the bottom of each column.
showed evidence for slight heating in the high energy ion tails. The 04/06/2000 event is also the only event to show sustained (over an hour) electron heating downstream of the shock ramp observed as flattop distributions, while the four more-typical events showed weak heating consistent with adiabatic compression. The core and halo electron heating is strongly anisotropic ($T_\perp > T_\parallel$) in the 04/06/2000 event, while the electron heating in the four more-typical events showed no particular preference towards parallel or perpendicular, consistent with previous studies of marginally critical shocks \cite{Thomsen1985a}. Typical supercritical quasi-perpendicular shocks exhibit perpendicular heating due to adiabatic compression and parallel heating due to a two-step process whereby the cross-shock potential accelerates the electrons parallel to the magnetic field and then the free energy associated with this accelerated beam excites microinstabilities, which redistribute the electrons in phase space to form flattop distributions \cite{Thomsen1987}. This two-step process leads to a roughly isotropic electron distribution in the downstream region of typical supercritical quasi-perpendicular shocks. The isotropy increases as one moves further into the downstream region, consistent with relaxation of the distributions due redistribution in phase space \cite{Thomsen1985a, Thomsen1987}. The core electrons in the 04/06/2000 event, however, are observed to become more anisotropic ($T_{e,\perp}/T_{e,\parallel}$ increases) as one progresses further downstream for up to an hour after the shock encounter (not shown). Adiabatic compression due to the conservation of the first adiabatic invariant cannot explain this observation since the magnetic field magnitude does not correlate with $T_{e,\perp}$ in the downstream, as shown quantitatively in Table 5.3. These observations suggest that wave heating is important, supported by the observation of electron cyclotron harmonic waves (see Figure 1.17) and solitary waves (not shown) similar to those of waveforms C and D, which would act to increase this anisotropy.

5.3 Summary and Conclusions

The first part of this chapter focused on results from a statistical study of high frequency ($\gtrsim 1$ kHz) waves at IP shocks. The results of the statistical study showed that the amplitude of large amplitude ($\geq 5$ mV/m peak-to-peak) IAWs increased with increasing Mach number and the shock strength, consistent with larger shock strengths
causing larger cross-field currents which may provide free energy for wave generation. The study also found that large amplitude ($\geq 5$ mV/m peak-to-peak) IAWs had the highest probability of occurrence in the ramp region, consistent with theory [Gary, 1981].

The second part of this chapter focused on the conclusions from the study of a specific IP shock event, namely the atypical IP shock of 04/06/2000 [Wilson III et al., 2010]. The waveforms were the first observations of large amplitude (>100 mV/m pk-pk) solitary waves and large amplitude (~30 mV/m pk-pk) waves exhibiting characteristics consistent with electron Bernstein waves at an interplanetary shock. Waveforms A through D in Figures 5.7 through 5.10 all clearly show enhanced power near integer and half-integer harmonics of the cyclotron frequency. Waveforms A and B show significant power along the shock normal and are obliquely polarized with respect to the magnetic field, consistent with the ECDI. If we use Equations 1.44 and 1.45 to estimate the contribution of anomalous resistivity due to waveforms A and B we find $\eta_{IA} \sim 920$ and 740 $\Omega$m, respectively. These values are over $10^8$ times the classical Spitzer collisional estimates, consistent with our arguments for their relative importance to the observed heating. Strong particle heating in both the halo and core of the electrons and ions is observed near these waves, consistent with the simulation results of Matsukiyo and Scholer [2006a].

Waves with power spectra exhibiting characteristics of both IAWs and Bernstein-like emissions are consistent with the predicted spectrum of the ECDI. The IAWs are Doppler shifted and resonantly interact with the Bernstein-like emissions, coupling to form a time-dependent diffuse frequency and wave vector spectrum [Matsukiyo and Scholer, 2006a]. The ECDI is an attractive candidate for the event herein because it can explain both the particle heating and the atypical waveforms. The current produced by the relative drift between incident electrons and reflected ions is unstable to the ECDI. Our observations suggest that this current is the source of free energy for the waveforms observed in the 2000-04-06 event.

Further evidence to suggest that our observations are consistent with the ECDI is shown using a simple test of linear instability. Using the ECDI instability criterion, at the Debye length cutoff or $(k \lambda_D) > 1$, determined from Equation 15 in Forslund et al.
under the observed conditions, we found that the ECDI instability criterion estimates are a factor of 20 or more above the threshold. Therefore, we argue that the instability criterion determined by Forslund et al. [1972] is easily satisfied.

To further examine the consistence of the observed waves with the ECDI, we looked at ion measurements from 10’s of eV to a few MeV, using the PH and SST Open detectors, for reflected ion signatures, the source of free energy for the ECDI. No beam-like or gyrophase-restricted features could be definitively discerned in high energy ion SST Open measurements near the shock ramp. The SST measurements did show enhancements in the energy flux of 1-6.7 MeV ions and 100-500 keV electrons (not shown herein) upstream of the 04/06/2000 event concurrent with the 12 shocklets observed by Wilson III et al. [2009]. However, simultaneous increases in high energy particle fluxes in and around low frequency waves is not unusual [Sanderson et al., 1985]. We also examined distributions from the PH instrument (shown in Figure 5.3). We did not detect reflected ion beams using the PH detector due to the combination of two factors: 1) the 2000-04-06 event had very atypical ion distributions upstream called diffuse ions, thought to be remnants of reflected ion beams scattered by wave-particle interactions, [Paschmann et al., 1981] and 2) UV-light contamination. Note 12 shocklets of the type observed upstream of this event have been shown to have a one-to-one correlation with reflected ions observed as diffuse ion distributions [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983], as observed in this case as well. Thus, the simultaneous observation of shocklets with diffuse ions is evidence that reflected ions exist upstream of the 04/06/2000 IP shock. A cursory comparison of the shock parameters for this event with the critical Mach number estimates of Edmiston and Kennel [1984] suggest that it is almost certainly supercritical, which would also entail ion reflection [Greenstadt and Mellott, 1987]. Since the percentage of reflected ions relative to incident increases with increasing Mach number [Kennel et al., 1985], the 2000-04-06 event likely reflects more ions than the four more-typical events. It is possible, therefore, that the 2000-04-06 event surpassed a threshold for the minimum percentage of reflected ions relative to incident ions necessary for the ECDI to become unstable. We also observe a beam-like feature seen in the high energy electrons; although it likely has a very low density compared to the background density, it may contribute to the free energy needed to drive the observed waveforms. Note that we observe this beam-like feature for over
20 minutes upstream modulated in intensity concurrently with shocklets.

Some of the ion heating may be due to the large amplitude solitary waves (electron phase space holes), which act like clumps of positive charge scattering and heating ions. Observations have shown that the change in perpendicular ion temperature across a train of electron holes can be comparable to the initial ion thermal energy [Ergun et al., 1998b]. The solitary waves may also be contributing to the observed changes in the electron distributions across the shock since their positive potentials can trap incident electrons [Dyrud and Oppenheim, 2006; Lu et al., 2008]. Solitary waves can either couple to or directly cause the growth of IAWs [Dyrud and Oppenheim, 2006], whistler mode waves [Lu et al., 2008], and/or electron acoustic waves [Matsukiyo and Scholer, 2006a]. Thus, solitary waves can directly heat/scatter particles or indirectly heat/scatter particles through the generation of or coupling to secondary waves.

This is the first study to report on the simultaneous observation of electron beams with large amplitude Bernstein-like waves and solitary waves at an IP shock. The preference for perpendicular ion heating is consistent with previous bow shock observations [Thomsen et al., 1985a], but the perpendicular electron heating is not. The polarizations and frequencies of the wave modes observed are inconsistent with previous observations of waves at IP shocks [Wilson III et al., 2007, 2009]. These results suggest a need for further investigation into the detailed microphysics of collisionless shock dissipation, particularly for higher Mach number events.
Chapter 6

Conclusion and Discussion

Collisionless shock waves have been studied for over 40 years, and still to this day scientists debate about the possible mechanisms which limit wave steepening. Shock wave formation requires a nonlinearly steepened wave to reach a point where the wave steepening is balanced by some form of irreversible energy dissipation. We observe shock waves in the IP medium, but we still do not have a complete understanding of the microphysics of these shocks. The ultimate goal of IP collisionless shock research is to fully understand the mechanisms behind energy dissipation (e.g. particle reflection/acceleration) that affect shock wave evolution/propagation for predictive capabilities applied to less accessible shocks. This thesis has increased the understanding of energy dissipation in low Mach number (Mf $\sim$1-4) IP shocks using electric and magnetic field and particle distribution data from the Wind spacecraft. The increased understanding was gained through the study of low frequency ($\sim$0.2-10 Hz) magnetosonic whistler waves upstream of the ramp regions, high frequency ($\gtrsim$1 kHz) electric field measurements in the ramp regions, and high frequency instabilities driven by reflected ions in the foot, ramp, and immediate downstream regions of IP shocks.

Large scale ($\lambda \sim R_E$) low frequency ($\sim$0.2-10 Hz) upstream waves are important in the total energy dissipation budget of collisionless shocks. They are produced by particles reflected by the shock, thus the particles carry energy away from the shock front and radiate the waves. The waves can then scatter and diffuse particles resulting irreversible energy dissipation upstream or downstream of the shock. These effects can also alter the Rankine-Hugoniot conservation relations. Part of this thesis focused
on Wind spacecraft observations of five IP shocks with low frequency upstream waves [Wilson III et al., 2009]. Two types of low frequency waves were observed, both of which lie on the magnetosonic whistler branch: precursor whistlers and shocklets. The precursor whistler waves propagated at angles with respect to the magnetic field of 20° to 50° and large propagation angles with respect to the shock normal, thus they did not appear to be phase standing, consistent with previous studies [Russell et al., 1983]. In this study, we presented the first observation of shocklets upstream of a quasi-perpendicular IP shock. Almost 80% of the shocklets had $\theta_{kB} \leq 45^\circ$, consistent with the cometary bow shock study by Le et al. [1989] and terrestrial bow shock studies [Hoppe et al., 1981, 1982; Hoppe and Russell, 1983; Russell et al., 1971]. The shocklets appeared to have higher frequencies leading the lower frequency waves, consistent with the theory that ULF waves driven by reflected ions steepen into the shocklets [Omidi and Winske, 1990; Scholer et al., 2003; Tsubouchi and Lembège, 2004]. Also, shocklets are only supposed to exist upstream of quasi-parallel shocks with small radii of curvatures. Since the spacecraft was over 50 $R_E$ in the X-GSE direction and over 35 $R_E$ in the Y-GSE direction away from Earth, the shocklet source must be the IP shock and not the terrestrial bow shock. Thus, our observation of 12 shocklets observed over an hour upstream of a quasi-perpendicular IP shock has implications for shock structure and dynamics. A possible explanation is that the shock front may not be of planar geometry. Meaning, the IP shock could have quasi-parallel geometries adjacent to the quasi-perpendicular crossing location. It is also possible, due to the supercritical nature of this IP shock, that the shock itself was not stable and reforming. If the shock was undergoing reformation, then the local geometry could oscillate between quasi-parallel and quasi-perpendicular. Regardless of how the shocklets were generated, the observation was the first of its kind.

In addition to the observation of shocklets simultaneous with diffuse ion distributions upstream of a quasi-perpendicular IP shock, we also examined low energy ($\leq 1$ keV) electron PADs for evidence of wave-particle interactions. The electrons showed evidence of cyclotron interactions through pitch-angle diffusion and/or scattering. The study was the first to report on electron distributions unstable to the whistler anisotropy and/or heat flux criterion of Gary et al. [1994, 1999] observed simultaneously with whistler waves at an IP shock. The associated electron heat flux calculated for each distribution changed both in angle with respect to magnetic field and wave vector and in magnitude.
across each shocklet. The maximum heat flux magnitude was observed just downstream of the leading whistler wave packets on the upstream edge of the shocklets, consistent with an electron heat flux free energy source. The angle between the heat flux vector and magnetic field direction also peaks just downstream of the steepened edge of the shocklets, consistent with the shocklets pitch-angle scattering the heat flux vectors. This was the first study to observe shocklets and precursor whistlers in association with electron distributions unstable to whistler heat flux and/or whistler anisotropy instabilities simultaneous with finite electron heat fluxes at quasi-perpendicular IP shocks.

Closer to the shock, we examined high frequency ($\gtrsim 200$ Hz) electric field measurements in the ramp regions of IP shocks \cite{WilsonIII2007, WilsonIII2010}. The first of two studies was a statistical study of high frequency ($\gtrsim 1$ kHz) electric field measurements in the ramp regions of IP shocks \cite{WilsonIII2007}. We established the first statistical evidence to suggest a dependence of large amplitude ($>5$ mV/m peak-to-peak) IAWs in the ramp regions on the shock compression ratio and fast mode Mach number. The dependence on shock compression ratio is consistent with larger shock strengths causing larger cross-field currents, which may provide free energy for wave generation. The study also showed the first statistical evidence that large amplitude ($\gtrsim 5$ mV/m peak-to-peak) IAWs had the highest probability of occurrence in the ramp region, consistent with theory \cite{Gary1981}. In addition to their relatively high frequency of occurrence, the waves were found to produce large anomalous resistivities. Quasi-linear estimates of the anomalous resistivities produced by these large amplitude IAWs ranged from $\sim 1$–856 $\Omega \cdot m$ ($\sim 10^7$ times greater than classical estimates). Recent Vlasov simulations using realistic mass ratios have found resistivities that are 2-3 orders of magnitude larger than their equivalent quasi-linear estimates \cite{Petkaki2006, Petkaki2008, Yoon2006, Yoon2007}. The argument is then that the effects due to these high frequency waves in the ramp regions of collisionless shocks may be more important than previously thought.

The second study of high frequency waves at IP shocks focused on the microphysics of an atypical supercritical IP shock \cite{WilsonIII2010}. It provided the first observations of large amplitude ($>100$ mV/m pk-pk) solitary waves and large amplitude ($\sim 30$ mV/m pk-pk) waves exhibiting characteristics consistent with electron Bernstein waves at an IP shock. The observed waveforms all clearly show enhanced power near
integer and half-integer harmonics of the cyclotron frequency. Two of the waveforms showed significant power along the shock normal and are obliquely polarized with respect to the magnetic field, consistent with the ECDI. Strong particle heating in both the halo and core of the electrons and ions is observed near these waves, consistent with the simulation results of Matsukiyo and Scholer [2006a]. The discussions in Section 5.2 indicate that the waves are consistent with the ECDI because: 1) the polarizations are primarily aligned with the shock normal direction, not the magnetic field direction; 2) the frequency spectrum shows integer and half-integer cyclotron harmonics; 3) the broad power intensifications at frequencies above 1 kHz of waveforms A and B are inconsistent with typical IAWs; and 4) we observed strong and preferentially perpendicular electron heating that is consistent with cyclotron heating as described in Forslund et al. [1972] and Matsukiyo and Scholer [2006a]. The study showed the first observational evidence of the ECDI at IP shocks and we presented the first observations of electron cyclotron harmonic waves downstream of an IP shock in this thesis (Figure 1.17).

An examination of electron distributions at other supercritical shocks (not shown) showed particle heating resulting in parallel flattop and perpendicular heated Maxwellians. The heating was predominantly in the parallel direction immediately downstream of the shock but isotropized as one progressed further into the downstream, in agreement with previous observations [Feldman et al., 1983a; Thomsen et al., 1987]. Strong quasi-static electric fields have been observed at collisionless shocks [Wygant et al., 1987; Walker et al., 2004; Bale and Mozer, 2007] and used to help explain particle heating [Scudder et al., 1986a; Thomsen et al., 1987]. Particle heating due to the cross-shock potential through a two step process of acceleration and relaxation is expected to create stronger heating of the core electrons parallel to the magnetic field than perpendicular [Thomsen et al., 1987]. In contrast, the 04/06/2000 event showed perpendicular heating dominating in the downstream and $T_{e,\perp}/T_{e,\parallel}$ increasing as one progressed farther downstream, suggesting wave heating is important. This study provided the first observation evidence to definitively suggest that wave heating may dominate over the heating due to the cross-shock potential.
6.1 Summary

The work presented in this thesis has helped increase the understanding of the microphysics of IP shocks in addition to raising new questions regarding the energy dissipation mechanisms dominating in the ramp regions. This was performed by providing the first detailed analysis of high time resolution waveform captures and particle distributions observed simultaneously in IP shocks. The initial work focused on a statistical study of high frequency waveforms in IP shock ramps. The study results suggested a re-evaluation of the importance of anomalous resistivity due to wave-particle interactions relative to dispersion, particle reflection, and macroscopic DC field effects. Further evidence to support the need for a re-evaluation of wave-particle significance was reported through the observational evidence of pitch-angle scattering and diffusion of electrons by low frequency magnetosonic whistler waves upstream of IP shocks. The importance of anomalous resistivity due to wave-particle interactions was examined in further detail in the study of the atypical electron heating observed at a supercritical IP shock, which we claimed clearly showed a dependence on the observed waveforms. The nearly ubiquitous observations of large amplitude IAWs in the ramp regions of IP shocks raise doubts about ignoring these high frequency fluctuations when examining the heating of particle distributions and particle energization in collisionless shocks.

Thus, we conclude that in the analysis of IP shocks the microphysics can no longer be disregarded. The results of the work presented in this thesis suggest a need for further investigation into the detailed microphysics of collisionless shock dissipation, particularly for higher Mach number events.


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Appendix A

Wind Measurements and Data Analysis Issues

A.1 The Pesa High Detector: The Glitch

Table A.1: Pesa High Bad Data Bins

<table>
<thead>
<tr>
<th>Mapcode (Hex)</th>
<th>Mapcode (Long)</th>
<th># of Bins</th>
<th>Bin Elements</th>
</tr>
</thead>
<tbody>
<tr>
<td>D4A4</td>
<td>54436</td>
<td>121</td>
<td>[0,1,8,9,16,17,24,25,35,36,43,44,51,52,59,60]</td>
</tr>
<tr>
<td>D4FE</td>
<td>54526</td>
<td>97</td>
<td>[0,1,5,6,10,11,15,16,23,24,28,29,33,34,38,39]</td>
</tr>
<tr>
<td>D5EC</td>
<td>54764</td>
<td>56</td>
<td>[0,1,2,8,9,10,19,20,21,27,28,29]</td>
</tr>
<tr>
<td>D6BB</td>
<td>54971</td>
<td>65</td>
<td>Unknown</td>
</tr>
<tr>
<td>D65D</td>
<td>54877</td>
<td>88</td>
<td>Unknown</td>
</tr>
</tbody>
</table>

On occasion, the PH detector has a glitch in the data which always occurs in the same data bins, regardless of date/time. The data bins which are affected are specific to the mode that PH is in, as seen in Table A.1. As a consequence, the set of data bins are essentially useless because they contain incorrect energy levels, among other things. The glitch appears to connected to a saturation due to the solar wind and solar UV contamination.

1 This connection was subjectively made upon plotting the PH distribution functions. I found
The glitches occur consistently in the same data bins, regardless of day or year. The

glitches, however, depend upon the mode that the Pesa High detector is in. I have found
the following bins to be an issue for the sample modes seen in Table A.1. The only
two modes I have yet to find a glitch in are the two where the number of bins equals
65 or 88. This is largely due to the fact that I almost never find Pesa High data with
this number of data bins. Also, the glitch appears always at the same GSE azimuthal
angles, φ (degrees), in the Pesa High data structures. In each case a repeating pattern
is seen, as with the data bins, consistent with the spacecraft rotation rate remaining
that the excess observed in the contour plots along the projection of the solar wind velocity decreased
dramatically when these bins were removed from the data. The UV contamination seen in the thermal
core also changed as a result of these alterations.

Figure A.1: Here is an example of an ion distribution seen upstream of an IP shock by
the Wind Pesa High instrument. The figure illustrates the differences between before
and after the glitch and solar wind noise were removed.
approximately the same, within the angular resolution of the Pesa High detector.

Figure A.1 shows an example of two distributions from PHB where I have attempted
to correct/remove the glitch and reduce the solar wind noise. The top two distributions,
panels A and B, have not been corrected and show the characteristics of the UV and
solar wind contamination while the bottom two distributions, panels C and D, show
reduced levels of both. Though there still exists contamination, the reduced effects help
one to recognize more important and real phenomena in the distributions. The contour
plots illustrate typical diffuse ion distributions, as seen in Figure 1.11. The large ge-
ometry factor and response of the PH instrument causes the thermal core measurement
to saturate the detectors, thus the isolated peak in the cuts of Figures 1.9 and 1.10.

One can also see remnants of UV contamination that have not been removed. Notice
that there appears to be a beam-like feature at an oblique angle to the magnetic field in
both distributions, which I have outlined in panels A and B. This is not a real feature,
rather an artifact of the glitch. It can be misleading, but the beam-like feature follows
the sun direction in every distribution with the glitch. The feature is seen far enough
upstream that one could argue that the particles are not magnetically connected with
the shock, thus it is not a signature of any sort of particle reflection/energization. The
intensity of this effect varies from event to event depending on location and orientation
of the satellite.

Another aspect of particle distributions which can cause misinterpretations is con-
tour closing algorithms used in IDL. Though the contour levels are determined through
a systematic and automated algorithm, they can alter interpretations of the data. As
one examines panels A and C, one can see that after correcting for the noise a beam-like
feature appears parallel to the magnetic field at roughly 450 km/s, and a weaker and
more diffuse signal peaked at roughly 850 km/s. The feature is likely a real signal, but
it is not long lived. Regardless, if one examines panel A, the feature is not noticeable
and would otherwise be missed.

A.2 SCET Corrections

The TDS samples are waveform captures of electric and magnetic field data. The
data is triggered by the largest amplitude waves which exceed a specific threshold and
are then stored in a memory buffer. The TDS datation time is sampled after the event is acquired which requires a delay buffer. The datation time requires two corrections. The first correction arises from tagging the TDS datation with an associated spacecraft major frame in house keeping (HK) data. The second correction removes the delay buffer duration. Both inaccuracies are essentially artifacts of on ground derived values in the archives created by the WINDlib software (K. Goetz, Personal Communication, 2008).

The WAVES instrument’s HK mode sends relevant low rate science back to ground once every spacecraft major frame. If multiple TDS events occur in the same major frame, it is possible for the WINDlib software to assign them the same SCETs. One can correct these issues to within +10 ms. This is often not necessary since the highest sampling rate of the HTR MFI data is only 22 Hz, however there are occasions where the 0.33 Hz data differs from the HTR data by > 20°. As a consequence, the polarization analysis (see hodogram in Figure 2.3) can be misleading.

The TDS receiver onboard calculations force the peak amplitude of the data to occur roughly in the center of the waveform capture time window. Thus the data is sampled from both before and after the peak. The time stamps associated with the data sampled after the peak is easily obtained but the times before the peak require the delay buffer memory. The delay buffer duration depends on sampling rate and waveform duration. WINDlib attempts to fix the delay buffer uncertainty but it cannot correct the built-in timing error of ∼300 ms when sampling at 120 kHz. One should note that for lower sample rates, the timing error is larger (K. Goetz, Personal Communication, 2008).

The TDSF receiver (sampled at 120 kHz) time stamps retrieved from WINDlib, before corrections, are accurate to +300 ms. The 300 ms uncertainty, due to the HK corrections mentioned above, results from WINDlib trying to recreate the time stamp after it has been telemetered back to ground. If an event stays in the TDS buffer for extended periods of time (i.e. >2 days), the interpolation done by WINDlib can make mistakes in the 11th significant digit. The positive definite nature of this uncertainty is due to rounding errors associated with the onboard DPU clock rollover. The DPU clock is a 24 bit integer clock sampling at ∼50,018.8 Hz. The clock rolls over at
The sample rate is a temperature sensitive issue and thus subject to change over time. From a sample of 384 different points on 14 different days, a statistical estimate of the rollover time is $5366.691124061162 \pm 0.000478370049$ seconds.

The method by which to correct the SCETs is as follows:

1. Retrieve the DPU clock times, SCETs, UR8 times, and DPU Major Frame Numbers from the Windlib libraries on the VAX/ALPHA systems for the TDSS(F) data of interest.

2. Retrieve the same quantities from the HK data.

3. Match the HK event with the same DPU Major Frame Number as the TDSS(F) event of interest.

4. Find the difference in DPU clock times between the TDSS(F) event of interest and the HK event with matching major frame number (Note: The TDSS(F) DPU clock time will always be greater than the HK DPU clock if they are the same DPU Major Frame Number and the DPU clock has not rolled over).

5. Convert the difference to a UR8 time and add this to the HK UR8 time. The new UR8 time is the corrected UR8 time to within +10 ms.

6. Find the difference between the new UR8 time and the UR8 time WINDlib associates with the TDSS(F) event. Add the difference to the DPU clock time assigned by WINDlib to get the corrected DPU clock time (Note: watch for the DPU clock rollover).

7. Convert the new UR8 time to a SCET using either the IDL Windlib libraries or TMLib libraries of available functions. This new SCET is accurate to within +10 ms.

---

2 The calculation is done by $(16*2^{24})/(50,018.8 \text{ Hz})$
A.3 Explicit Explanation of Actual Data Calculations

A.3.1 Wind 3DP Software Calculations

Two tests were performed to determine the stability of our original temperature anisotropy estimates and to observe the effects on our comparisons to the instability estimates of [Gary et al. 1994, 1999]. The first test was to lower and raise our cutoff energy bin for the minimum energy of the halo electrons to raise and lower our estimate of $n_{he}/n_e$. Occasionally our estimates of $n_{he}/n_e$ were a factor of 5-10 less (for the 12/10/1997 and 04/06/2000 events) than the estimate used by [Gary et al. 1994], $n_{he}/n_e = 0.05$. The second test was a rigorous examination of our fit estimates in comparison to the anisotropy estimates calculated directly from the 3D P particle data. Note, since [Gary et al. 1999] explicitly removed the strahl electrons from their calculations, we ignored the component of the strahl electrons and only fit to the more isotropic halo components when calculating temperature anisotropies for comparison to the threshold estimates estimated by [Gary et al. 1994, 1999].

The results of the first test showed the largest relative change in the relevant parameters of [Gary et al. 1994, 1999] occurred in $\beta_{\parallel c}$ and $T_{\parallel h}/T_{\parallel c}$. When we lowered(raised) the cutoff energy, both $\beta_{\parallel c}$ and $T_{\parallel h}/T_{\parallel c}$ decreased(increased). However, for most distributions $T_{\perp h}/T_{\parallel h}$ increased(decreased). Since $T_{\perp h}/T_{\parallel h}$ changed inversely with respect to $T_{\parallel h}/T_{\parallel c}$, we concluded that changing the energy bins did not dramatically affect the instability estimates. Almost all of the PADs which were unstable in our original estimates were still unstable after lowering the lowest energy of the halo electrons. In fact, many of the PADs for the four IP shocks without shocklets were more unstable when we lowered the energy cutoff. For both estimates of break energies, almost all the distributions within 30 seconds of all five IP shock ramps were unstable to the whistler heat flux instability estimated by [Gary et al. 1994].

The results of the second test almost always showed a stronger anisotropy in the core and halo electrons than the direct calculations from the data. We found that the strahl electrons, found almost entirely in the parallel component of the electron distribution, lowered estimates of $T_{\perp}/T_{\parallel}$ for both the core and halo. The effect is seen more strongly in the halo electrons, and [Gary et al. 1994, 1999] suggested the halo is more important for the whistler heat flux and anisotropy instabilities than the core. This implies that
the electron distributions we present may be more unstable than our estimates. Therefore, we used our moment calculations as a lower bound on the halo/core temperature anisotropies.

### A.3.2 Specific Software Implementations

This section will illustrate how typical particle software calculates the particle moments in detail, specifically for the Wind/3DP software (consequently, this same software is used for every Berkeley particle instrument, e.g. Cluster, STEREO, THEMIS etc.).

To calculate the distribution function, one must first define a weighting factor, \( w_t \) (unitless). \( w_t \) depends upon the energy, \( E \), spacecraft potential, \( \phi_{sc} \), particle charge, \( q \) (in units of fundamental charge), and differential energy, \( \delta E \). When calculating the weighting factor for real data, for mathematical and practical purposes, \( w_t \) is constrained by defining it as:

\[
wt = 0.0 < \left( \frac{E + \phi_{sc}}{\delta E} + 0.5 \right) < 1.0
\]  
(A.1)

where Equation A.1 defines the mathematical result of \( w_t \). The next step is to shift the energy of the particles by \( \phi_{sc} \) to avoid using photo-electrons from the satellite in the calculations. This is done by defining the differential velocity, \( \delta v \left( \sqrt{eV} \right) \), as:

\[
\delta v \equiv \sqrt{\left( E + \frac{\phi_{sc}}{q} \right)}
\]  
(A.2)

where \( \delta v \) is constrained to be \( > 0 \) always. Before going any further, one must convert the data units to energy flux (Energy cm\(^{-2}\) sr\(^{-1}\) s\(^{-1}\) eV\(^{-1}\)) so that our distribution function is a function of energy and solid angle. At this point, one can calculate the differential distribution function in the following manner:

\[
\delta f \equiv \frac{\delta v}{E} \ast \frac{d(E,\Omega)}{10^5} \ast \frac{\delta E}{E} \ast d\Omega \ast wt
\]  
(A.3)

where \( d\Omega \) is the differential volume and the factor \( 10^5 \) is a conversion factor to multiply by cm/km to get the resulting units eV\(^{-1/2}\) km/s cm\(^{-3}\) sr\(^{-1}\). One should note, that at this point, \( \delta f \) is a 2-dimensional \( N_E \times N_e \)-array, where \( N_E \) corresponds to the number of energy bins and \( N_e \) corresponds to the number of data bins. Integrating over \( d\Omega \)
(2nd dimension, or $N_b$) gives an estimate of the number density per root mass (units of $eV/c^2$ and $c$ is in km/s) per energy bin. Summing over the energy bins and dividing by a normalization factor given by:

$$n_o = \sqrt{\frac{2 \times q}{M_s}} = 593.09544 \text{ (for electrons)} \quad (A.4)$$

yields the number per centimeter cubed, or the number density. Generally one would refer to this as the zeroth order moment calculation. The mass is given by:

$$
\text{mass} = \frac{510.990.6eV/c^2}{(2.99792458 \times 10^8 \text{km/s})^2} \\
= 5.6967578 \times 10^{-6} \text{ eV/(km/s)}^2. \quad (A.5b)
$$

and a particle charge normalization is given by:

$$0.010438871 \frac{eV}{(\text{km/s})^2} = \frac{938.27231 \times 10^6 eV/c^2}{(2.99792458 \times 10^5 \text{km/s})^2} \quad (A.6)$$

which is used for determining the fractional charge (units of proton charge as $eV/(\text{km/s})^2$) of the input particle type. For electrons, the rounded value of the particle mass over this charge conversion constant is zero, thus the program assigns a value of -1 to the variable $\text{charge}$.

The first moment comes from integrating (summing in IDL) $\delta f \delta v \hat{r}$, where $\hat{r}$ is defined as:

$$\hat{r} = (\cos \theta \cos \phi) \hat{x} + (\cos \theta \sin \phi) \hat{y} + (\sin \theta) \hat{z} \quad (A.7)$$

where $\theta$ and $\phi$ are the latitude and azimuthal angles in which ever coordinate system you happen to be in. After some calculations, the first moment gives a flux density ($\text{km/s cm}^{-3}$) as a vector in the IDL routine.

The second moment is calculated in a similar manner but now we must define a new variable, $\delta f_1$, as $\delta f \delta v$. So now we see that the second moment can be calculated as:

$$\vec{V}_{l,m} = \int d^3v (\delta f_1 * \delta v * (\vec{r})(l,m)) \quad (A.8)$$

where $\vec{V}_{l,m}$ is the $(l,m)$-component of the velocity flux ($\sqrt{eVkm/s \text{ cm}^{-3}}$).

The third moment is calculated in a similar manner but now we must define a new
variable, $\delta f_2$, as $\delta f \delta v^2$. Thus one can show the third moment goes as:

$$\overline{Q}_{l,m,n} = \int d^3v (\delta f_2 \ast \mathbf{f} \ast (\mathbf{r} \mathbf{r} \mathbf{r})_{l,m,n})$$

(A.9)

where $\overline{Q}_{l,m,n}$ is the general form of the heat flux. However, one assumes that the order of the indices does not matter (due to symmetry assumption) which reduces the possible combinations from 27 down to 10. The 10 variations of $(l,m,n)$ are: $Q_{x,x,x}$, $Q_{x,y,y}$, $Q_{x,z,z}$, $Q_{x,x,y}$, $Q_{x,x,z}$, $Q_{x,y,z}$, $Q_{y,y,z}$, $Q_{y,z,z}$, $Q_{y,y,y}$, and $Q_{z,z,z}$. Generally in any particle code, one reduces this even further to a simple rank-2 tensor or $3 \times 3$ matrix where the sum of the $i^{th}$ row results in the $i^{th}$ component of the resultant heat flux vector. The above method is the same as the typical definition of the heat flux given by Equation (A.10):

$$\mathbf{q} = \frac{m_e}{2} \int d^3v \mathbf{v} \mathbf{v} f(\mathbf{x}, \mathbf{v}, t)$$

(A.10)

where $m_e$ is the electron mass (in the program this is done by multiplying the result of Equation (A.5b) by the square of the result of Equation (A.4)), $\mathbf{v}$ the velocities, and $f(\mathbf{v}, \mathbf{x}, t)$ represents a general form of the distribution function. Note, the heat flux is only relevant in the plasma rest frame so that $f(\mathbf{v}, \mathbf{x}, t)$ should be appropriately transformed before calculating Equation (A.10).

### A.3.3 Pressure and Temperature Calculation

This section will illustrate how typical particle software calculates the particle pressure and temperature in detail. In theory, the particle pressure need not explicitly depend upon a parameter called the temperature. In practice, however, the assumption is made that the temperatures come directly from the pressure tensor. It is important to note that the particle temperature, or average kinetic energy in the bulk flow frame, is derived from the second moment of the distribution function. It should be derived from carefully calculated numerical fits of the $\parallel$ or $\perp$ cut of the reduced distribution function for every data structure (see Equations 2.2 through 2.11). However, this is rarely done and often the temperature is calculated from the moments derived in the particle software. There is a fundamental difference in the two calculations which should be carefully regarded because it has some important implications on the interpretation of the data.
The pressure tensor is calculated from the second moment of the distribution function, $V_{l,m}$ (Equation A.8), where $l$ and $m$ correspond to the rows and columns of the second moment tensor. To calculate the pressure tensor, one multiplies $V_{l,m}$ by two normalization constants (corresponding to a charge and mass normalization) in the software. In the software, the only indices $(l,m)$ used for $V_{l,m}$ are the following: $V_{x,x}$, $V_{y,y}$, $V_{z,z}$, $V_{x,y}$, $V_{x,z}$, and $V_{y,z}$ (thus a symmetric tensor is assumed). In typical software calculations, the total pressure is calculated by taking one third of the trace of the pressure tensor. Similarly, the average temperature is calculated by that same quantity divided by the number density. Thus, $\langle T_e \rangle = \text{Tr}[P_{i,j}]/n_e$. 
Appendix B

Glossary and Acronyms

- **Coronal Mass Ejection (CME)** – A large eruption of plasma from the surface of the sun releasing upwards of 50 billion tons of mass and propagating through the interplanetary medium at up to 3000 km/s.

- **Cyclotron(Gyro) Frequency** ($\Omega_{cs}$ for angular in radians per second or $f_{cs}$ for circular in Hz) – The characteristic frequency associated with a particle in circular motion in a uniform static magnetic field.

- **Cyclotron(Gyro) Radius** ($\rho_{gs}$) – The characteristic radius of a charged particle orbit with a velocity perpendicular to the magnetic field.

- **Electron Cyclotron Drift Instability (ECDI)** – Instability driven by a relative drift between reflected ions and incident electrons at a quasi-perpendicular collisionless shock. This results in a resonant interaction between Doppler-shifted ion-acoustic waves and electron Bernstein (or cyclotron) waves.

- **Electron Cyclotron Harmonic Waves (ECHWs)** – A group of wave modes that exhibit power enhancements at integer and half-integer harmonics of $f_{cs}$ and can include $(n + 1/2)$, electron Bernstein-like, or "totem pole" emissions which are typically banded electrostatic waves propagating nearly perpendicular to the magnetic field.

- **Electrostatic[Electromagnetic] (ES[EM])** – A wave that has its wave vector
parallel (perpendicular) to the electric field fluctuations without (with) corresponding magnetic field fluctuations, which is also a longitudinal (transverse) fluctuation.

- **Electrostatic Solitary Waves (ESWs)** – A nonlinear wave with a bipolar electric field signature parallel to the background magnetic field that is consistent with a BGK phase space hole.

- **Interplanetary Coronal Mass Ejection (ICME)** – (See CME) A CME that leaves the sun’s surface.

- **Interplanetary Shock (IP Shock)** – A shock wave typically propagating away from the sun due to an ICME or a corotating interaction region acting as the driver.

- **Ion-Acoustic Wave (IAW)** – A plasma oscillation which is composed of the inertia carrying ions fluctuating as the wave propagates with the electrons providing the charge balance. In the electrostatic limit, the waves are linearly polarized parallel to the magnetic (or at slightly oblique angles) field and propagate parallel to the electric field oscillations.

- **Lower Hybrid Wave (LHW)** – A plasma oscillation that is typically ES with phase fronts propagating nearly perpendicular to the ambient magnetic field carried primarily by demagnetized ion motions with a frequency near the lower hybrid resonance, \( f_{lh}^2 = f_{ce} f_{ci} \).

- **Macro(Micro)instability** – A fluid(kinetic) instability which generally has the largest growth rates when \( k \rho_i \ll 1 (k \rho_i \gg 1) \).

- **Modified Two-Stream Instability (MTSI)** – Instability driven by relative drift between electrons and ions across a magnetic field.

- **Plasma Frequency** (\( \omega_{ps} \) for angular in radians per second or \( f_{ps} \) for circular in Hz) – The frequency associated with a small net displacement of a particle species, \( s \), from equilibrium due to an electric field restoring force.

- **Right[Left]-Hand Polarized Wave (RH[LH])** – A transverse electromagnetic wave which has \( \hat{k} \parallel B_o \) but \( \delta B \) and \( \delta E \) rotate in a clock(antclock)wise sense
when looking along $B_z$.

- **Shocklet** – A nonlinearly steepened magnetosonic wave which often exhibits a high frequency wave packet on its leading edge and has a magnetic compression ratio $\leq 2$. They are thought to result from ultra low frequency foreshock waves that interact with intermediate (or gyrating or gyrophase bunched) and diffuse ion distributions which alter the index of refraction, thus steepening the waves.

- **Short Large Amplitude Magnetic Structures (SLAMS)** – Similar to shocklets except they have steepened to much larger amplitudes with magnetic compression ratios often exceeding a factor of 4.

- **Soliton** – A nonlinear wave whose different Fourier components propagate at different speeds for small amplitudes and have a constant shape as they propagate because the nonlinear steepening is balanced by dispersive effects.