The collisionless nature of high-temperature plasmas

T. M. O'Neil

Department of Physics, University of California at San Diego, La Jolla, California 92093

F. V. Coroniti

Department of Physics and Astronomy, University of California, Los Angeles, Los Angeles, California 90024

Two phenomena that illustrate the collisionless nature of high temperature plasmas are Landau damping (or, more generally, the resonant wave-particle interaction) and collisionless shock waves. The first half of this paper traces Landau's idea through the years as it is tested experimentally, extended nonlinearly, and applied. The second half traces the progress in understanding collisionless shocks in space and astrophysical plasmas. [S0034-6861(99)03002-0]

I. INTRODUCTION

An important property that distinguishes high temperature plasmas from normal fluids, even from conducting fluids such as liquid metals, is that the plasmas are to a first approximation collisionless. In a laboratory plasma, the mean-free-path between collisions can be much larger than the dimensions of the plasma. In space and astrophysical plasmas, the mean-free-path can easily exceed the dimensions of the structures of interest. The collisionless nature necessitates a kinetic treatment and introduces a variety of subtle new phenomena. For example, Landau damping (or growth) results from the resonant interaction of a wave with free streaming particles, a resonance that would be spoiled by collisions in a normal fluid. Also, the collisionless nature challenges us to find new descriptions for familiar phenomena. For example, what is the nature of a shock wave in a collisionless plasma?

This review provides a brief introduction to the collisionless nature of plasmas. Taking the resonant wave-particle interaction as characteristic, we follow Landau's idea through the years as it is tested experimentally, extended nonlinearly, and applied. We then describe the earth's bow shock, which is an important example of a collisionless shock wave. Finally, we touch on cosmic ray acceleration by supernova shocks.

II. WAVE-PARTICLE INTERACTIONS

A. Linear theory

The first proper treatment of modes in a collisionless plasma was provided by Landau (1946). Using the collisionless Boltzmann equation and Poisson's equation he obtained the dispersion relation for electron plasma oscillations (Langmuir oscillations). Tonks and Langmuir (1929) had described these simple electrostatic modes many years earlier using fluid equations. In fact, it was in this early paper that Tonks and Langmuir coined the name "plasma." Landau's main correction to the earlier fluid description was that the modes experience a collisionless damping (or growth). The electric potential for a mode that is characterized by wave number $\mathbf{k}=\hat{z}k$ damps (or grows) temporally at the rate

$$\gamma(k)=\frac{\pi}{2}\,\omega(k)\frac{\omega_p^2}{k^2}\,\frac{\partial f_0}{\partial v_z}\bigg|_{\omega(k)/k}, \tag{1}$$

where $\omega(k)$ is the mode frequency and $\omega_p=(4\pi ne^2/m)^{1/2}$, is the electron plasma frequency. Here, e and m are the electron charge and mass and n is the density. The function $f_0(v_z)=\int dv_x dv_y F_0(v_x,v_y,v_z)$ is the distribution of electron velocities parallel to the direction of propagation, where $F_0(v_x,v_y,v_z)$ is the distribution over all three velocity components. The subscript zero indicates that the distribution refers to the unperturbed equilibrium state, which is assumed to be spatially homogeneous. Langmuir waves are special in that the electrons dominate the dynamics; more generally, both electrons and ions contribute to $\gamma(k)$. Clearly, the damping (or growth) is associated with electrons that satisfy the relation $kv_z=\omega(k)$ [or equivalently, $\mathbf{k}\cdot\mathbf{v}=\omega(k)$]; these electrons maintain a constant phase relative to the wave and resonantly exchanging energy with the wave. For a Maxwellian velocity distribution, the derivative $\partial f_0/\partial v_z|_{\omega(k)/k}$ is negative so $\gamma(k)$ is negative and the mode damps. However, for a non-Maxwellian distribution, corresponding, say, to the case where a small warm beam drifts through the plasma, $\partial f_0/\partial v_z|_{\omega(k)/k}$ can be positive implying wave growth. Two caveats should be noted here. The first is that expression (1) is an approximate form for $\gamma(k)$ that is valid when the damping is weak (i.e., $|\gamma/\omega|\ll 1$). The second more important caveat is that Landau linearized the collisionless Boltzmann equation neglecting a term that is second order in the mode amplitude.

Landau's work served as a model for the theoretical description of many kinds of plasma modes, both electrostatic and electromagnetic, and the Landau resonance for an unmagnetized plasma was generalized to the cyclotron resonance for a magnetized plasma (Stix, 1962, 1992). By the late 1950's, a large body of theory had been developed for the kinetic description of plasma modes. However, there was a concern that the theory had not been tested adequately, so several small scale laboratory experiments were developed to isolate and test the basic elements of the theory. Happily, scientific opportunity and availability of funds converged to make this a "golden era" for such small scale experiments.

Figure 1 shows the results of an experiment that in-

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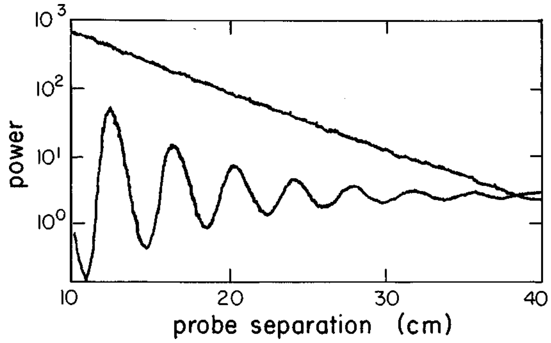

FIG. 1. Measurement of spatial Landau damping. The upper curve is the logarithm of the power measured by a receiver probe plotted as a function of the distance from the transmitter probe. The lower curve is the instantaneous wave form, obtained by operating the two-probe system as an interferometer. From Malmberg and Wharton, 1966.

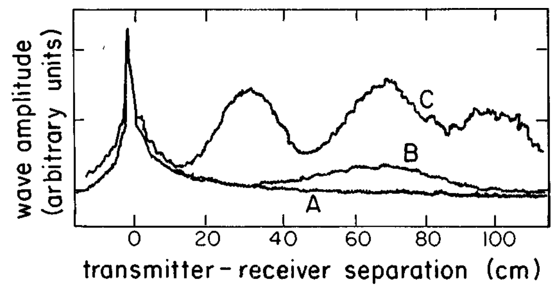

FIG. 2. Wave amplitude vs position. The transmitter voltage was 0.9, 2.85, and 9 V for curves A, B, and C, respectively. From Malmberg and Wharton, 1967.

vestigated Landau damping for the simple case of Langmuir waves (Malmberg and Wharton, 1966). The plasma was a steady-state (continuously produced and lost) 2-m-long column that was immersed in an axial magnetic field. A Langmuir wave was transmitted continuously by a probe (wire) that was inserted into the plasma and made to oscillate in potential. Waves propagated axially in both directions away from the transmitter damping spatially as they propagated. Landau also considered the case of spatial damping, predicting the spatial damping rate (imaginary wave number) $k_i(\omega) = \gamma[k(\omega)]/v_g$, where $\gamma(k)$ is the temporal damping rate and $v_g = d\omega/dk$ is the group velocity. The finite radial size of the column and the large axial magnetic field produce slight changes in the form of $\gamma(k)$ and of $\omega(k)$ relative to the results for an unmagnetized homogeneous plasma, but the changes are technical details, not matters of principle. The upper curve in Fig. 1 is the logarithm of the power measured by a receiver probe, plotted as a function of the distance from the transmitter probe. The oscillatory curve was obtained by operating the two-probe system as an interferometer. The nearly straight line dependence of the upper curve demonstrates that the damping was exponential, and the slope is twice the spatial damping decrement $k_i(\omega)$. The factor of 2 enters because power is proportional to the square of the wave amplitude. The measured decrement was far too large to be accounted for by collisional processes, but was in good agreement with the predictions of Landau's theory. A particularly convincing demonstration was that the damping ceased when the velocity distribution was manipulated to remove the resonant electrons. Also, Landau growth was observed when a warm beam was injected to make $\partial f_0/\partial v_z$ positive at the resonant velocity. In other early experiments, Landau's theory was tested using ion acoustic waves (Chen, 1984). By now, predictions of damping and growth due to the Landau and cyclotron resonances have been verified for many modes in a wide range of experimental settings.

B. Nonlinear theory

We will consider three nonlinear extensions of Landau damping, discussing in each case the original work from the 1960s and then a modern incarnation (or application) of that work. The first two extensions illustrate complementary physical interpretations of Landau damping: damping as a result of energy exchange with electrons that "surf" on the wave field and damping of the wave field because of phase mixing in velocity space.

1. Trapped particle oscillations and the plasma wave accelerator

Consider a resonant electron that is trapped in the trough of a large amplitude Langmuir wave and is accelerated forward as it slides down one side of the wave trough. The electron gains energy by "surfing" on the wave. However, when the electron reaches the bottom of the trough and decelerates as it moves up the other side, it loses energy. Thus, we expect the wave damping decrement, $\gamma(t)$, to oscillate in time at the frequency of oscillation of an electron that is trapped in the trough of the wave, $\omega_{\rm osc} = (ek^2\delta\phi/m)^{1/2}$. Here, $\delta\phi$ is the amplitude of the wave potential. Formally, one can check that Landau's linearization procedure fails after a time $\omega_{\rm osc}^{-1}$. For a small amplitude wave [i.e., $\omega_{\rm osc} \ll |\gamma|$], the wave damps away long before the trapped electrons can complete an oscillation, so Landau's theory is valid. In the opposite limit of a very large amplitude wave [i.e., $\omega_{\rm osc} \gg |\gamma|$], the trapping oscillations stop the damping before the wave amplitude can change by a significant amount (Mazitov, 1965; O'Neil, 1965). In general trapped particle oscillations have been found to dominate the nonlinear wave-particle interaction in many situations. Most importantly, the Landau growth of a single wave (or narrow spectrum of waves) saturates nonlinearly when the amplitude is large enough that $\omega_{\rm osc} \sim |\gamma|$ (Drummond *et al.*, 1970; Onishchenko *et al.*, 1970).

Figure 2 shows measurements of spatial Landau damping when the transmitter power was turned up until trapped particle oscillations dominated the evolution (curve C) (Malmberg and Wharton, 1967). For spatial damping, the oscillations occur spatially and are characterized by the wave number $k_{\rm osc} = \omega_{\rm osc}/v_{\rm ph}$, since $v_{\rm ph} = \omega/k$ is the speed of the resonant electrons. The measured value of $k_{\rm osc}$ scaled with wave amplitude as $\sqrt{\delta\phi}$, as expected.

A modern version of this experiment is the plasma wave accelerator (Tajima and Dawson, 1979). Conventional accelerators are limited to acceleration rates of 100 MeV/m, the limit where radio frequency breakdown occurs. The longitudinal electric field of a plasma wave can be much larger than this, and the phase velocity can be relativistic ($v_{\mathrm{ph}} \simeq c$). In principle, trapped bunches of electrons can be accelerated resonantly to high energy in a relatively short distance, so there is the promise of a compact accelerator. In recent experiments (Everett *et al.*, 1995), a large-amplitude, relativistically propagating plasma wave was generated by beats between two co-propagating laser fields of slightly different frequency (i.e., $\Delta\omega \simeq \omega_p$). Trapping of externally injected electrons and an acceleration rate of 2.8 GeV/m were demonstrated. For these small scale experiments, the interaction length was only a cm, so the energy gain was modest (28 MeV).

2. Plasma wave echoes and beam echoes

The plasma wave echo (Gould, O'Neil, and Malmberg, 1967) is another nonlinear extension of Landau's theory. The echo explicitly demonstrates that the free energy associated with the wave is not dissipated in collisionless damping, but is stored in the distribution function and can reappear later as a wave electric field. The plasma wave echo is closely related to other echo phenomena such as the spin echo. Landau's analysis shows that macroscopic quantities such as the electric field or charge density damp away, but that the perturbation in the distribution function, $\delta f(z,v_z,t)$, oscillates indefinitely. Since the perturbed electron density is given by $\delta n(z,t) = \int dv_z \delta f(z,v_z,t)$, one may think of Landau damping as a phase mixing of different parts of the distribution function. When an electric field of spatial dependence $\exp(-ik_1 z)$ is excited and then Landau damps away, it modulates the distribution function leaving a perturbation of the form $\delta f = f_1(v_z)\exp[-ik_1 z + ik_1 v_z t]$. This perturbation propagates at the local streaming velocity in phase space, v_z. For large t, there is no electric field associated with the perturbation since a velocity integral over the perturbation phase mixes to zero. If after a time Δt an electric field of spatial dependence $\exp[ik_2 z]$ is excited and then damps away, it moderates the unperturbed part of the distribution leaving a first-order term $f_2(v_z)\exp[ik_2 z - ik_2 v_z(t-\Delta t)]$. However, it also modulates the perturbation due to the first field leaving a second-order perturbation of the form

$$\delta f^{(2)} = f_1(v_z) f_2(v_z) \exp[i(k_2-k_1)z + ik_2 v_z \Delta t - i(k_2-k_1)v_z t]. \qquad (2)$$

The coefficient of velocity in this exponential vanishes when $t = \Delta t k_2/(k_2-k_1)$, so at this time a velocity integral over this second-order perturbation does not phase mix to zero and an electric field (the echo) reappears in the plasma. This is a temporal echo, but there are also spatial echoes, where the wave fields damp spatially, and the echo is separated spatially.

Soon after they were predicted, spatial echoes were observed experimentally using both Langmuir waves and ion-acoustic waves (Chen, 1984). Also, echoes were used to make very sensitive measurements of small angle scattering due to Coulomb collisions (and to stochastic fields). The echo depends on very fine scale structure in the phase space distribution, which is easily smoothed out by small angle scattering. For example, the plasma column that was used to make the measurements shown in Figs. 1 and 2 was only 2 m long, but Langmuir wave echoes in this plasma were used to measure an effective mean free path of 2 km.

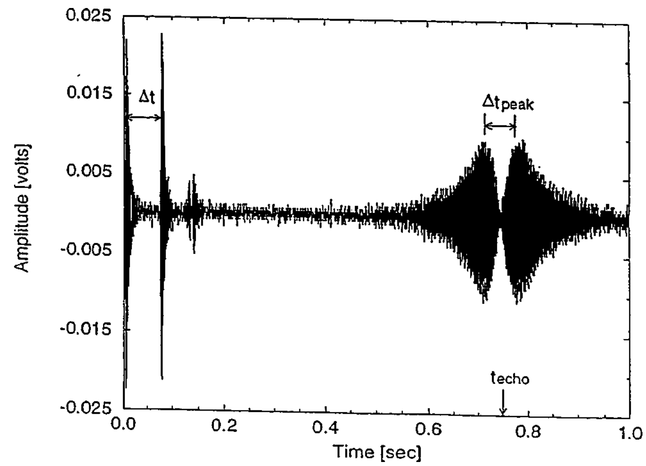

FIG. 3. Temporal echo observed on a stored, coasting antiproton beam in the Fermilab Antiproton Accumulator. The two excitations and the echo are shown. From Spentzouris *et al.*, 1996.

A modern version of this experiment was used recently to measure the energy diffusion rate (due to intrabeam Coulomb collisions) of a coasting antiproton beam in the Fermilab Antiproton Accumulator (Spentzouris, Ostiguy, and Colestock, 1996). Figure 3 shows the signal for a temporal echo on the beam. These echoes are the same as plasma echoes except that the beam particles are relativistic and that collective fields can be ignored in the dynamics. By measuring the decay of the echo amplitude as a function of time to the echo, an intra-beam collision frequency of $(3.0 \pm 0.8) \times 10^{-4}$ Hz was obtained. More recently, an effective collision frequency of 10^{-13} Hz was measured for a higher-energy coasting proton beam at CERN (Brüning *et al.*, 1997). Clearly, the echo provides an exquisitely sensitive measure of small angle scattering.

3. Quasilinear theory and current drive

The most widely used of the nonlinear extensions is quasilinear theory (Drummond and Pines, 1962; Vedenov *et al.*, 1962). This physically appealing theory provides a simplified description of the nonlinear wave-particle interaction for the case of a broad spectrum of randomly phased waves. The auto-correlation time for the field as seen by a resonant particle is assumed to be short compared to the time for a trapped particle oscillation. The sign of the field experienced by a resonant particle then undergoes rapid random changes, and the particle experiences a kind of Brownian diffusion in velocity space. For the simple case where all of the waves propagate in the z direction (or a strong magnetic field constrains the particle motion to the z direction), quasilinear theory predicts that the equilibrium velocity distribution evolves according to the diffusion equation

$$\frac{\partial f_0}{\partial t}(v_z,t)=\frac{\partial}{\partial v_z}D(v_z,t)\frac{\partial f_0}{\partial v_z}(v_z,t). \tag{3}$$

In the resonant region, the diffusion coefficient $D(v_z,t)$ is proportional to the energy in the waves that are characterized by phase velocity $\omega/k=v_z$; these waves satisfy a Landau resonance with particles at velocity v_z. Equations (1) and (3) govern the evolution of the wave energy [or, equivalently, of $D(v_z,t)$] and of $f_0(v_z,t)$. This truncated description conserves particle number, momentum, and energy, and is in reasonably good agreement with experiment (Roberson and Gentle, 1971; Hartmann *et al.*, 1995). Quasilinear theory has generated a vast literature, including many applications and many attempts to place the theory on a stronger theoretical foundation. Efforts to understand quasilinear theory from a first principles dynamical perspective helped to motivate early work on the dynamical origins of chaos.

A modern application of the wave-particle interaction, where the quasilinear diffusion equation is used to describe the theory, is rf current drive in tokamaks (Fisch, 1984). As described by Fowler (see article in this volume), a tokamak is a toroidal magnetic confinement device for high temperature plasmas, and is the leading contender to be a fusion reactor. Confinement in a tokamak requires that the plasma carry a toroidal electric current, and in a conventional tokamak this current is driven by an inductive electric field that is directed toroidally. The plasma is the secondary in a transformer circuit where the primary passes through the hole in the torus. Since the magnetic flux in the primary is finite, a substantial inductive electric field can be maintained only for a finite time (about an hour for a reactor scale tokamak). However, there would be technological advantage in the steady-state operation of a tokamak reactor.

In recent years, experiments on tokamaks have demonstrated that the required steady state current can be driven with the wave-particle interaction. Mega-amps of current have been driven by this method in large tokamaks. A phased array of wave guides is used to launch lower hybrid plasma waves so that they propagate in a particular direction around the torus. The waves Landau damp on the tail of the electron distribution, transferring wave momentum to the resonant electrons. In this way, electrons are pulled out from the Maxwellian to produce a high velocity tail (or plateau) in the direction of wave propagation. From a quasilinear perspective, the spectrum of waves diffusively sweeps particles down hill on the Maxwellian distribution forming the high velocity plateau. The current resides in this high velocity plateau. Of course, the steady state shape of the plateau is determined by a balance between momentum deposition by the waves and collisional drag on the ions and slow electrons. Incidentally, it is advantageous to deposit the momentum in fast electrons since the collisional drag on these electrons is less than that on thermal electrons. When research on current drive was beginning, critics worried that the high rf power levels and the high velocity electrons would produce anomalous (collective) processes and that these would confuse the theory and spoil the efficiency of current drive. However, this has not been the case; traditional theory (e.g. quasilinear theory and the classical collision operator) provides a good description of experimental results over a wide parameter range.

III. COLLISIONLESS SHOCKS IN SPACE PLASMAS

In the late 1950's, the question of whether shocks exist in collisionless plasmas posed a great challenge to the developing discipline of plasma physics. Gas dynamic shocks and magnetohydrodynamic fast and slow shocks form as the steepened limit of a nonlinear compression wave in which the thickness of the shock layer (shock front) is determined by the characteristic dissipation length associated with the particle collision mean free path. Since early laboratory experiments with plasma shocks were partially collisional, the collisionless shock challenge was first met by the satellite study of space plasmas.

The Earth's dipole magnetic field is immersed in the supersonic, super-Alfvenic solar wind whose low density ($5\ \mathrm{cm}^{-3}$) and moderate temperature (2×10^5 K) result in a 1 AU mean free path. Nevertheless, plasma physicists speculated that a bow shock would stand in the solar wind upstream of the Earth's magnetosphere. In late 1964, using the magnetometer measurements from NASA's IMP 1 satellite, Ness *et al.* (1964) unambiguously identified the magnetic compression signature of a thin magnetohydrodynamic fast mode collisionless bow shock; the IMP 1 plasma measurements subsequently confirmed the shock heating and the slowing of the solar wind. Today high Mach number bow shocks have been detected at all the planets that have been visited by spacecraft and around three comets. Although tantalizing evidence of slow magnetohydrodynamic shocks was found in solar wind magnetic structures, the unambiguous detection of a magnetohydrodynamic collisionless slow shock did not come until Feldman *et al.* (1984) used plasma and magnetic field measurements from NASA's ISEE 3 spacecraft to verify the slow shock Rankine-Hugoniot relations at the plasma sheet boundary in the distant geomagnetic tail.

A. Early collisionless shock models

For collisionless shocks, the critical question is, What dissipation mechanisms replace particle-particle collisions? Viewed broadly, the answer is the wave-particle interaction discovered by Landau, although a wide variety of different wave modes with different dispersive properties are involved even for a single type of shock. In the late 1950's, Adlam and Allen (1958) and R. Z. Sagdeev (reported in Sagdeev, 1966) showed that the thickness of a steepening nonlinear fast magnetohydrodynamic wave that propagates perpendicular to the magnetic field would be limited by the finite inertia of the plasma electrons at the dispersive scale length c/ω_p, the so-called collisionless skin depth. The balance between nonlinear steepening and dispersion results in a steady compressive soliton that propagates in the ideal fluid plasma without dissipation. With the addition of resistivity, however, the soliton converts into a sharp leading front, in which the magnetic field strength rises to above the downstream Rankine-Hugoniot value, followed by a train of trailing wave oscillations that resistively damp to the downstream Rankine-Hugoniot state. Recall that the Rankine-Hugoniot relations, which are robust consequences of conservation theorems applied across the shock, determine the downstream state in terms of the upstream state independent of the details of the dissipation mechanism in the shock interior. Sagdeev argued that, for a collisionless shock, the resistivity would be provided by ion-acoustic wave turbulence which is self-consistently excited by the strong cross-field current drift of the electrons in the magnetic field ramp at the leading edge of the shock; in quasilinear theory, ion acoustic waves elastically scatter the current-carrying electrons and transfer electron momentum to ions. For oblique fast shocks, ion inertial dispersion speeds up the fast wave, so that the soliton is a rarefaction pulse; with the addition of resistivity, the soliton becomes a leading whistler wave train that propagates into the upstream region ahead of the magnetic ramp, and is spatially damped with increasing upstream distance.

A quite different approach to the collisionless shock dissipation mechanism was proposed by E. Parker and H. E. Petschek. Parker (1961) was interested in high-β plasmas (β is the ratio of the plasma to magnetic pressure). He argued that the shock would consist of interpenetrating upstream and downstream ion beams that would excite the beam-firehose instability, an instability driven by velocity space anisotropy. A similar quasiparallel shock theory was proposed by Kennel and Sagdeev (1967) in which the shock compression creates a plasma distribution with a higher parallel than perpendicular temperature; this anisotropy excites magnetosonic Alfvén turbulence via the temperature anisotropy version of the firehose instability. They developed this model into a full quasilinear theory of weak, high-β quasiparallel shocks in which the scattering of the upstream ions by the fluctuating wave magnetic fields provides the shock dissipation.

Petschek (1965) was interested in high-Mach-number (supercritical) quasiperpendicular shocks for which resistivity alone (as in the Sagdeev model) cannot provide sufficient dissipation to satisfy the Rankine-Hugoniot relations. In Petschek's shock model, the increased entropy associated with the whistler turbulence replaces the thermal heating of the plasma; plasma heating occurs as the Alfvén waves are gradually absorbed by Landau damping on spatial scales that are much greater than the shock thickness.

B. Fast shocks

After the discovery of the Earth's bow shock, subsequent NASA spacecraft, especially OGO 5, with its high time resolution magnetometer and the first electrostatic wave measurements, compiled an observational database on collisionless fast shock structure for a wide range of Mach numbers, propagation angles, and plasma betas. In addition, a powerful new diagnostic—the plasma numerical simulation of collisionless shocks—was pioneered at Los Alamos by D. Forslund, and developed into a technique that permitted a detailed comparison of theory and spacecraft observations. Today, although many details remain, a broad outline of collisionless shock dissipation mechanisms has been established.

The observed structure of low-Alfvén-Mach-number ($M_A<2$) oblique fast shocks consists of the predicted leading whistler wave train. However, the source of the anomalous resistivity that converts the rarefaction soliton into the shock wave train is still uncertain. In the solar wind the electron to ion temperature ratio is too low for the electron drifts associated with the whistler's magnetic field oscillations to destabilize the ion acoustic wave. Other modes, such as the lower hybrid drift wave, could be excited for lower electron drift speeds, or some other process may be responsible for introducing irreversibility into whistler wave trains.

At higher Alfvén Mach numbers, the ISEE 1 and 2 spacecraft typically observed the following: a large magnetic overshoot of a factor of 2 or so above the downstream field strength, very little evidence of an upstream whistler wave train, and a double-peaked downstream ion distribution. These observations were beautifully explained by the plasma simulations of supercritical shocks (Leroy *et al.*, 1982). In addition to the magnetic ramp, supercritical shocks have a net electrostatic potential jump across the shock front. Roughly, due to their large inertia, the ions plough through the sharp magnetic ramp whereas the electrons remain attached to the field lines; thus a charge separation electric field develops to slow the incoming ions. The combined $\mathbf{v}\times\mathbf{B}$ and normal electric field force reflects a fraction of the incoming ions back upstream; the reflected ions then gain energy by drifting parallel to the tangential shock electric field and penetrate through the potential barrier into the downstream region on their next gyro-encounter with the shock front. Downstream, the reflected ion population and the still-unshocked upstream ions form a plasma distribution with a velocity space gyrating beam, which appear to spacecraft as a double-peaked distribution. Subsequent 2D and 3D simulation studies have shown that the beam ions ballistically mix and eventually thermalize with the directly transmitted upstream ions via turbulent interactions with electromagnetic ion cyclotron waves; the reflected ions create a distribution that is more perpendicular than parallel and a temperature that destabilizes the ion cyclotron waves. In the 2D and 3D simulations, ion reflection produces a shock that is steady only in an average sense; the number of reflected ions, the size of the magnetic overshoot, and the thickness of the magnetic ramp actually vary in time and in position along the shock front. Observations at the bow shocks of the outer planets and some simulation studies suggest that very high Mach number ion reflection shocks may be intrinsically and violently unsteady.

For quasiparallel shocks, since the downstream flow velocity is less than the speed of sound, the shocked ions can readily travel back upstream along the magnetic field. The interpenetrating inflowing and backstreaming ions create a beam-firehose distribution (as envisioned by Parker), but the most unstable waves are actually fast magnetosonic modes (Krauss-Varban and Omidi, 1993). For quasiparallel terrestrial bow shocks, the amplitudes of the observed magnetic field fluctuations can be comparable to the DC field, an example of order-one magnetic turbulence.

C. Cometary bow shocks

Near perihelion cometary nuclei emit large fluxes of neutral hydrogen and water group molecules that are then ionized by the solar UV, forming a large halo around the comet. Unless the solar wind velocity is exactly parallel to the interplanetary magnetic field, a newborn ion is first accelerated by the solar wind $\mathbf{v}\times\mathbf{B}$ electric field, thereby forming a cold gyrating beam in velocity space. The energy and momentum of this ion pick-up process must come at the expense of the solar wind flow. Thus the pick-up ions inertially load and slow the solar wind, thereby forcing a bow shock to form around the comet. Since the interaction is collisionless, the coupling between the pick-up ions and the solar wind is mitigated by the unstable excitation of low-frequency hydromagnetic waves that scatter the pick-up ions in pitch angle and energy and extract momentum from the solar wind. The spacecraft that encountered comets Giacobini-Zinner, Halley, and Grigg-Skjellerup observed large-amplitude magnetic turbulence both upstream and downstream of the cometary bow shocks.

D. Cosmic ray acceleration by supernova shocks

The galactic component of cosmic rays extends from about 1 GeV to at least 10^6 GeV with a momentum distribution that decreases as a simple power law. The maintenance of the cosmic rays against escape losses from the galaxy requires an energy input of roughly 10^{41} ergs/sec whose only known source is blast waves from supernova explosions. The problem of how shock energy can be efficiently converted into very high energy particles was solved independently by G. F. Krimskii (1977), I. W. Axford *et al.* (1977), and R. D. Blandford and J. Ostriker (1978). They recognized that a quasiparallel collisionless shock could establish the physical conditions for a particle to undergo type-I Fermi acceleration.

In a quasiparallel shock, suprathermal particles (seed cosmic rays) can rather freely travel upstream along the magnetic field. As a streaming or beamlike distribution, these particles excite parallel propagating magnetosonic waves that scatter the streaming particles in pitch angle. In the shock frame, the unstable waves, which travel at the Alfvén speed relative to the plasma at rest, are convected into the shock by the upstream flow. Once the streaming particles are scattered through 90° in pitch angle, they also travel back to the shock and into the downstream flow region where they can again be scattered in pitch angle. As a net result of these wave-particle scattering interactions, the particles diffuse back and forth across the shock. Since the downstream flow speed is less (typically 1/4 for a strong shock) than the upstream flow speed, to the particles, the wave scattering centers appear to converge; thus the particles are effectively trapped between converging mirrors and experience type-I Fermi acceleration. Rather amazingly, the steady-state momentum distribution that the particles acquire from the shock acceleration process is a power law whose spectral index is independent of the pitch angle or spatial diffusion coefficient (i.e., the intensity of the unstable waves), and only depends on the shock compression ratio. The predicted power spectral index closely agrees with the observed cosmic-ray spectrum from 1 GeV to about 10^6 GeV.

Since the original model was proposed, the theory of cosmic-ray accelerating shocks has been greatly elaborated and refined. In particular, the energy density of the accelerated cosmic rays can evolve to become comparable to the flow energy associated with the shock; thus the cosmic rays become part of the overall shock structure, so that supernova shocks actually extend over vast distances. The acceleration theory was extended to solar flare blast waves traveling in the interplanetary medium by M. A. Lee (1983), and was thoroughly and successfully tested by Kennel *et al.* (1984) in the spacecraft study of a particle-accelerating interplanetary shock. Finally, shock acceleration may also explain the production of relativistic electrons in synchrotron extragalactic radio jets. Today, the shock acceleration of cosmic rays stands as, perhaps, the one major successful application of collisionless plasma physics to astrophysics.

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