Interstellar Shock Waves

Christopher F. McKee and Bruce T. Draine

An understanding of interstellar shock waves is crucial in determining the structure of the interstellar medium. By causing the gas to radiate, interstellar shocks provide astronomers with valuable diagnostics on both the physical conditions in the interstellar medium and the energy source that produced the shock. The complexity of the interstellar plasma—its degree of ionization, its molecular content, the presence of small dust grains and cosmic rays, and the magnetic field—leads to a rich variety of structures for interstellar shocks, which are being actively investigated both observationally and theoretically.

The interstellar medium (ISM) is a tenuous plasma filling the galactic disk that has a mean density of only about one particle per cubic centimeter. Its mass is divided almost evenly between atomic and molecular components, although the molecular gas occupies only a small fraction of the volume. The ionization of this plasma varies widely, from about 1 part in 10^7 in the molecular gas to virtually complete ionization in the hottest ionized gas. The molecular gas is typically at a temperature $T \approx 10$ K and a density $n \approx 10^3$ cm^{-3} and is concentrated in self-gravitating molecular clouds in which new stars form. The atomic gas is in three phases are in approximate pressure equilibrium: the cold ($T \approx 10^2$ K, $n \approx 40$ cm^{-3}), warm ($T \approx 10^4$ K, $n \approx 0.3$ cm^{-3}), and hot ($T \approx 10^6$ K, $n \approx 3 \times 10^{-3}$ cm^{-3}) phase. The composition of the ISM is complex: The gas is predominantly H with a smaller amount of He (~ 10% by number) and trace amounts of heavier elements (less than 0.1% by number). Some of the heavier elements, which are referred to as "metals," are in small dust grains that range in size from less than 1 nm to about 100 nm. A small fraction of the atoms in the ISM have been accelerated to relativistic energies; these are the cosmic rays. Although low in density; the cosmic rays have an energy density comparable to that of the rest of the gas, and they play an important role in governing the level of ionization in molecular clouds. Tying all these components together is a magnetic field that has a typical strength of a few microgauss, although it is stronger in denser regions (*1*).

The ISM is a violent place because of energy injection by stars (*2*). Nebulae, which are often depicted in books of astronomic photographs, are evidence of this violence. Some nebulae, such as the one in the constellation Orion, are produced by massive stars, which generate enormous fluxes of ionizing radiation that cause the nebula to glow. This ionizing radiation heats the surrounding gas and causes it to expand at velocities of tens of kilometers per second; such stars also have powerful stellar winds (with velocities of ~ 10^3 km s^{-1}) that abet the expansion. When these massive stars die, they explode as supernovae, releasing about 10^{51} erg of energy into the surrounding medium in ejecta moving at velocities of up to 10^4 km s^{-1}. About 10^4 years later, the remnant of a supernova may look like the Veil Nebula in Cygnus (part of the Cygnus Loop supernova remnant), which is expanding at several hundred kilometers per second. Radio and infrared observations have shown that there is another type of nebula that is normally obscured from our eyes by interstellar dust: high-velocity molecular flows from newly forming stars, which have typical injection velocities of order 200 km s^{-1}.

The effect of this energy injection depends on the velocity of the injected mass relative to the sound speed in the ambient medium. The ISM is generally transparent to the radiation that cools it, so that its temperature and thus its sound speed are relatively low. For example, molecular gas of cosmic abundances at a temperature of 10 K has an isothermal sound speed $(P/\rho)^{1/2} \approx 0.2$ km s^{-1} (P is pressure and ρ is density); warm ionized gas at a temperature of 10^4 K has a sound speed of about 10 km s^{-1}. Because the ISM is magnetized, signals can propagate at the Alfvén velocity $v_A = B/(4\pi\rho)^{1/2}$ as well (B is the magnetic field strength). In the atomic gas, v_A is generally comparable to the sound speed; in the molecular gas, it is typically a few kilometers per second. If the mass is injected into the ISM at a velocity exceeding the velocity at which signals can propagate, a shock wave forms.

Shock Waves

A shock is a hydrodynamic surprise. It travels at a velocity v_s that exceeds the velocity of propagation of signals in the medium; the Mach number of the shock M, which is the ratio of v_s to the signal velocity, is greater than 1. As a result, the gas ahead of the shock is unaware of its approach (although energetic particles or radiation from the shock may provide an early warning). When the shock hits, there is a sudden jump in the density, pressure, temperature, and velocity of the gas. The conditions behind the shock front (denoted by the subscript 1) are related to those ahead (denoted by the subscript 0) by a set of algebraic relations expressing the conservation of mass, momentum, and energy [the Rankine-Hugoniot, or jump, conditions (*1*)]. If the shock is weak ($M \simeq 1$), as in the case of a sonic boom, then the gas is almost unchanged across the shock front. On the other hand, if the shock is strong ($M \gg 1$), then the pressure behind the shock is about equal to the ram pressure, $P_1 \simeq \rho_0 v_s^2$ and is almost independent of the initial pressure. In a strong shock, the velocity of the shocked gas measured relative to the unshocked gas v_1 is comparable to but somewhat smaller than the shock velocity v_s. The density and temperature of the shocked gas depend on the nature of the shock: If all species of particles in the gas have a common velocity and kinetic temperature, as in the case of a highly ionized plasma, then the density in a strong shock increases by a factor of 4, and the temperature increases to

$$kT_{1,\mathrm{eq}} = \frac{3}{16}\mu v_s^2 \qquad (1)$$

so that the sound speed in the shocked gas is comparable to the shock velocity. Here k is Boltzmann's constant, μ is the mean mass per particle, and the subscript eq indicates that all the different

C. F. McKee is at the University of California at Berkeley, Berkeley, CA 94720. B. T. Draine is at Princeton University, Princeton, NJ 08540.

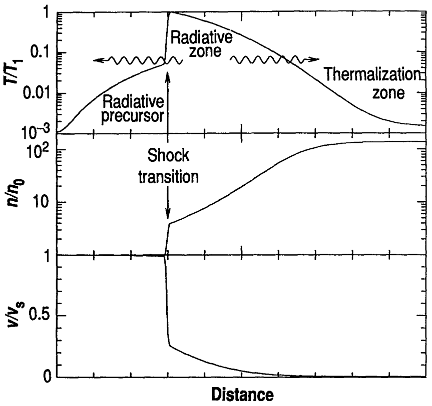

Fig. 1. Schematic structure of a strong (single-fluid) shock wave showing temperature T, density n, and velocity v (relative to the shock front).

species—electrons, ions, and neutrals—are in equipartition (that is, they have the same temperature); we have assumed that the adiabatic exponent is $\gamma = 5/3$. Numerically, this is $1.38 \times 10^5 v_{s7}^2$ K for a fully ionized plasma with the composition of interstellar gas (v_{s7} is the shock velocity measured in units of 10^7 cm s^{-1}, or 100 km s^{-1}). If, on the other hand, the shock occurs in a multifluid plasma, such as a weakly ionized plasma, then the temperature and velocity vary from one species to another, and radiative losses in the shock front may be important, causing the density jump across the shock front to exceed 4 and the temperature behind the shock front to be correspondingly lower than the value just cited.

The structure of a shock can be divided into four regions, which are not always distinct (*3–5*) (Fig. 1). The radiative precursor is the region upstream of the shock in which radiation emitted by the shock provides a premonition of the shock's arrival. The shock front, or shock transition, is the region in which the relative kinetic energy of the shocked and unshocked gas is dissipated. If the dissipation is due to collisions among the atoms or molecules of the gas, as in the case of a sonic boom, the shock is collisional. On the other hand, if the density is sufficiently low that collisions are unimportant and the dissipation is due to the collective interaction of the particles with turbulent electromagnetic fields, the shock is collisionless; Earth's bow shock, which is formed where the solar wind impacts with Earth's magnetosphere, is a good example of such a shock. Next comes the radiative zone, in which collisional processes cause the gas to radiate, possibly alter its chemical state, and increase in density as it cools. This region becomes important only if the age of the shock is comparable to or greater than the time for electrons to undergo radiative collisions. Finally, if the shock lasts long enough, there is the thermalization zone, in which radiation from the radiative zone is absorbed and then reradiated.

Astrophysical shocks are classified on the basis of the relative importance of collisional processes (Fig. 2). Consider a supernova explosion with an energy of 10^{51} E_{51} erg, where E_{51} is the energy expressed in units of 10^{51} erg, and for simplicity assume that the ambient medium is homogeneous with a hydrogen density n_0. Initially, the mass ejected by the supernova will drive a shock into the ambient medium at a velocity that is about the same as that of the ejecta, several thousand kilometers per second. As the shock expands, it will sweep up more and more mass until it begins to decelerate. So long as the energy of the supernova remnant remains nearly constant (the Sedov-Taylor stage of evolution), the shock is said to be nonradiative, even though the radiation emitted behind the shock provides astronomers with valuable diagnostics of the physical conditions there. In this stage, its velocity will decrease as $v_s \propto (E_{51}/n_0)^{1/5}t^{-3/5}$, and the average temperature will decrease as $T \propto v_s^2 \propto (E_{51}/n_0)^{2/5}t^{-6/5}$ (t is the age of the shock). Because the collision time scale t_{coll} for an ionized plasma is proportional to $T^{3/2}/n$, we see that the ratio of the age t to the collision time t_{coll} varies as $(n_0^2E_{51})^{1/3}v_s^{-14/3}$, so that there is a critical velocity below which collisions have enough time to be important behind the shock front. In particular, for $v_s < 180$ $(n_0^2E_{51})^{1/14}$ km s^{-1}, radiative losses due to collisions are sufficient to cool some of the gas in the remnant to the point that a dense shell forms (*6*). Subsequently, the shock is said to be radiative. As a rule of thumb, interstellar shocks that are much faster than 200 km s^{-1} are generally nonradiative, whereas those much slower are radiative. Radiative shocks are particularly important to astronomers because virtually all the energy flux crossing the shock front, which is expressed as $\rho_0 v_s^3/2$, is converted into radiation that is potentially observable.

As long as the shock velocity lies between 110 km s^{-1} and the upper limit depicted in Fig 2, there is enough ionizing radiation emitted from the cooling gas behind the shock front to create a photoionized precursor (*7, 8*). At lower velocities, the shock becomes increasingly neutral and more difficult to observe. If the ambient medium is atomic, the supernova remnant will slow down until it merges with the ISM.

If the remnant is expanding into a molecular medium, however, a dramatic change can occur in the structure of the shock. For $v_s \gtrsim 40$ to 50 km s^{-1}, the shock is strong enough to dissociate the molecules (*9*), and the structure is qualitatively similar to that for an atomic shock. At lower velocities, however, the dissipation in the shock front is due to collisions between the neutrals and the rare ions; the shock front becomes thick, so that both radiative losses and chemical changes become important within the shock front itself (*10*). Such shocks are termed C shocks because the density increases continuously across the shock front. They are distinguished from the J shocks discussed above, in which the shock front is no thicker than a collisional mean free path, so that the density and so forth undergo a jump across the shock front.

In the past decade, observations and theory have combined to advance our understanding of many phenomena associated with interstellar shocks, including particle acceleration, the structure and chemistry of shocks in molecular gas, the effect of geometry on the spectra of shocks associated with H–H objects in outflows from young stellar objects, the heating and destruction of dust grains in shocks, the stability of shocks, maser emission, and so on (*11*). Here we focus on two particular cases of interest: shocks in ionized gas, such as occur in young supernova remnants (J shocks), and shocks in magnetized molecular gas.

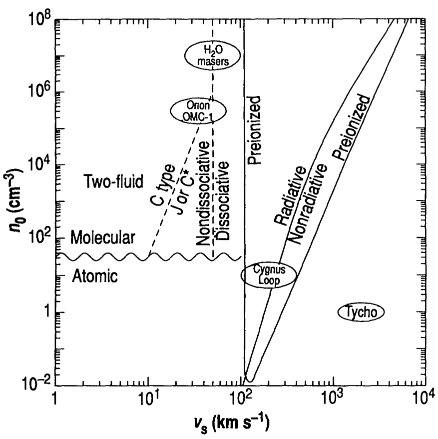

Fig. 2. Location of different shock types on the $v_s - n_0$ plane. The line labeled radiative/nonradiative separates radiative and nonradiative shocks (for shocks assumed to be driven by a 10^{51}-erg supernova explosion). The line labeled preionized encloses the conditions under which the radiation emitted by the shock wave is able to preionize initially neutral preshock gas. Strong two-fluid shocks are possible only if the gas is primarily molecular, which typically requires densities well in excess of 10 cm^{-3}; in atomic gas, strong shocks are J type. The approximate locations of the Cygnus Loop and Tycho supernova remnants, the H_2 line–emitting region in OMC-1, and the H_2O masers in star-forming regions (*48*) are shown.

Nonradiative Shocks in Ionized Gas

The shocks associated with supernova remnants produce large volumes of hot gas ($T \gtrsim 10^6$ K) in the ISM and are thus essential to the global structure of the ISM (*12*). By heating the gas crossing the shock to high temperatures, these shocks produce spectacular sources of both x-ray emission (as a result of Bremsstrahlung and line emission from collisionally excited ions) and infrared emission (as a result of collisional heating of the entrained dust). These shocks are also believed to be responsible for accelerating the cosmic rays, which are relativistic particles that pervade the galaxy (*13*). These effects are particularly important while the shocks are nonradiative, so that the remnant retains its initial energy, and in this stage the shocks generally propagate in ionized gas. Despite the importance of such shocks, much remains to be learned about their structure.

The basic issue, as yet only partly resolved, is the following question: what is the mechanism for dissipating the relative energy of the shocked and unshocked gas? For shocks in the atmosphere, this dissipation is due to collisions: The molecules in the shocked and unshocked gases bounce off each other, irreversibly converting the energy in relative motion into heat. At the low densities that characterize interstellar (and interplanetary) plasmas, however, the distance that a proton moving at a velocity $10^7 v_7$ cm s^{-1} must travel before it is deflected by 90° is almost $10^{15} v_7^4/n_e$ cm, where n_e is the electron density. In the solar wind, this distance exceeds the distance from Earth to the sun; for young supernova remnants in the ISM, which can expand at velocities exceeding 3000 km s^{-1} ($v_7 = 30$), this distance can exceed the thickness of the cold gas disk of the galaxy. Observations show that shocks form on much smaller length scales in both cases; the dissipation is effected by the interaction of the particles with turbulent electromagnetic fields arising from collective motions of the charged particles, and the shocks are said to be collisionless.

The best studied collisionless shock is Earth's bow shock, which is formed by the supersonic flow of the solar wind past Earth's magnetosphere. Collisionless shocks on a larger scale occur in the interplanetary medium, driven by variations in the solar wind. Observations of these shocks, together with the associated theory, have been summarized (*14*). The nature of the shock front is determined by the angle θ between the shock normal and the preshock magnetic field. If $\theta \gtrsim 45°$ to 50°, the shock is said to be quasi-perpendicular; otherwise, it is quasi-parallel. In a quasi-perpendicular shock, the magnetic field acts to insulate the upstream region from the downstream region; the motion of the charged particles is primarily gyration around the field rather than streaming along it. If the Alfvén Mach number $M_A = v_s/v_A$ exceeds a critical value between 1 and 3, depending on conditions, then the dissipation is associated with ions that are reflected by the compressed magnetic field in the shock before being transmitted downstream. Observations by the ISEE satellites (*15*) show that the magnetic field undergoes a large jump in a thin ramp; theory suggests that the structure of the ramp is determined by the electrons (*16*). Ahead of the ramp is the foot, which is associated with the reflected ions and in which the field begins to rise from its initial value; behind the ramp is the overshoot region, in which the magnetic field is larger than its final value by a factor that increases with M_A. The foot and the overshoot region each have a thickness on the order of the ion Larmor radius (*16*). The strongest shock observed to date is the bow shock around Uranus, with $M_A = 23$ (*17*); in principle, interstellar shocks can be far stronger, with M_A exceeding 100.

Quasi-parallel shocks are much thicker than quasi-perpendicular shocks. Quasi-parallel shocks have an extended foreshock containing electromagnetic waves together with energetic ions that have leaked upstream, a localized jump in density with a thickness comparable to the ion gyroradius, and a downstream region with large-amplitude magnetic turbulence (in which the ions become isotropic). Parker (*18*) proposed that shocks of this type would have a structure determined by the magnetic analog of the firehose instability: The shocked and unshocked ions constitute two counterstreaming beams with a large stress along the field; just as the rapid flow of water in a hose causes it to flap about, so this flow of ions along the field causes it to become unstable when the difference between the parallel and perpendicular pressures exceeds $B^2/8\pi$. Quest (*19*) has generalized this model to include the instability that occurs when the waves can resonate with the gyration of the particles in the foreshock. He has shown that the resulting model is consistent with both numerical simulation and the available observations, which are of shocks of relatively modest Mach number.

Because the energetic particles that leak upstream determine the amplitude of the waves that eventually scatter the ions in the localized jump, quasi-parallel shocks are natural sites of particle acceleration. Energetic particles in the shock can bounce back and forth from the downstream region to the upstream region, gaining energy at each bounce. The result is that the energetic particles acquire a power-law distribution of particle energies that depends only on the compression ratio $r = n_1/n_0$ of the shock; for relativistic particles, one finds that $dN/dE \propto E^{-\alpha}$, where N is the number of energetic particles with energy greater than E, and $\alpha = (r + 2)/(r - 1)$ (*13, 20*). This process is termed first-order Fermi acceleration or diffusive shock acceleration. The energy distribution of the shocked particles differs greatly from an equilibrium (maxwellian) distribution because the number of degrees of freedom excited in a collisionless shock is many orders of magnitude smaller than in a collisional one (*20*). In the solar system, this process produces many of the energetic particles that are observed [see, for instance, (*21, 22*)]; it is conjectured that in the ISM it produces the cosmic rays. The energy required to accelerate the particles is extracted from the converging flow of the shocked and unshocked plasmas. In principle, so much energy can go into energetic particles that the structure of the shock is modified (*23*); the fate of the thermal particles entering the shock is governed, at least in part, by relativistic particles.

Although this theory has been quite successful, there are several outstanding problems. What is the structure of a quasi-parallel shock with a substantial amount of energy in accelerated particles? Which interstellar shocks are the ones responsible for accelerating the cosmic rays? How is it possible to account for the observed distribution of cosmic rays, which extends smoothly up to about 10^{15} eV before showing a break? How is the electron component of the cosmic rays accelerated? How does the efficiency of the acceleration depend on the angle θ between the shock normal and the magnetic field (for example, *24*)?

One of the central unsolved problems in the theory of collisionless shocks is the determination of the relative temperatures of the electrons and the ions. The shock jump conditions determine the mean temperature behind the shock front $T_{1,eq}$ (see Eq. 1) but not the temperatures of the individual species; note that the mean temperature includes the energy in energetic particles, which is also unknown. Because the electrons are generally responsible for the cooling of the plasma, the value of the electron temperature T_e is crucial to an understanding of both the dynamics of the shock and its emitted spectrum. Most of the energy in the preshock flow is contained in the ions; the kinetic energy in the streaming of the electrons is less than 1/2000 of the total. To decelerate the ions, large electric fields are developed in the shock front with potentials that are a significant fraction of $\frac{1}{2}m_i v_s^2$ (m_i is the ion mass), and it has been suggested that turbulent fields of this magnitude could result in substantial electron heating, with $T_e \sim T_{1,eq}$ (*25*). For quasi-

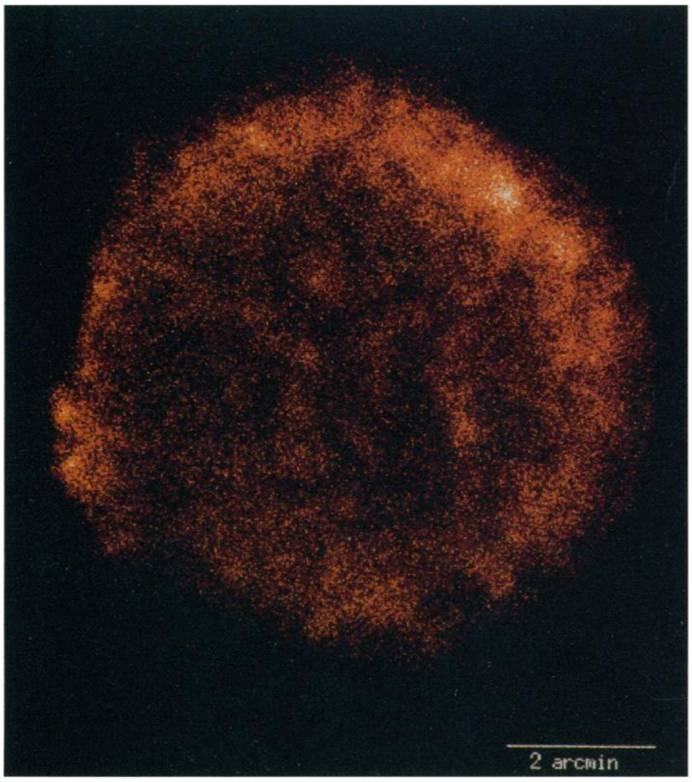

Fig. 3. An x-ray image of Tycho's supernova remnant taken with the ROSAT satellite. The remnant is bounded by a shock wave moving at several thousand kilometers per second that is heating the gas to temperatures of order 10^8 K. The sharp edge of the remnant shows that the shock is collisionless. [Photo courtesy of the Max-Planck-Institut für Extraterrestrische Physik, Garching, Germany]

perpendicular shocks of relatively low Mach number, such as Earth's bow shock, detailed modeling has shown that the electrons should acquire only a small fraction of the ion energy (*26*); simulations have shown that the fraction of the energy transferred to the electrons should increase with the strength of the shock (*27*). In a simulation of a strong shock ($M_A = 50$) in which the electrons were treated as a resistive fluid, Cargill and Papadopoulos (*28*) found substantial electron heating, with $T_e/T_{1,eq} \simeq 0.36$.

There are several observational paths available for determining the amount of collisionless heating of the electrons in shocks. Observations of Earth's bow shock are generally consistent with the heating expected from adiabatic compression and do not indicate any significant collisionless dissipation of energy into the electrons (*14*). The bow shocks at the outer planets are stronger ($M_A \sim 20$) and have lower upstream electron temperatures, however, and they do indicate significant collisionless heating of the electrons, with $T_e/T_{1,eq} \sim 0.2$ (*17, 29*). In a survey of 82 shocks, primarily terrestrial bow shocks, but including interplanetary shocks and the jovian and uranian bow shocks, Schwartz *et al.* (*30*) found a good correlation between the electron heating ΔT_e and the ion heating ΔT_i, with $\Delta T_e \sim 0.25\Delta T_i$. For the highest Alfvén Mach numbers, the heating was found to be somewhat less, consistent with the result $\Delta T_e \sim 0.1\Delta T_i \sim 0.2\ T_{1,eq}$ for the bow shocks of the outer planets. None of these solar system shocks was observed to be supersonic relative to the electrons (that is, v_s was not greater than the electron thermal velocity), although such shocks are expected in supernova remnants.

The most direct observations of the electron temperature in strong interstellar shocks come from x-ray observations of young supernova remnants. Inferring the postshock value of T_e is made difficult by the complex structure of many of these remnants, which often have plasmas at several different temperatures. The remnant of the supernova observed by Tycho in 1572 (Fig. 3) appears to have one component of x-ray emission at a temperature of 7×10^7 to 8×10^7 K, which is generally attributed to the shocked ISM (*31–33*); a lower temperature component associated with the shocked ejecta from the supernova is also indicated (*33*). The velocity of the shock has been variously estimated as 2200 to 3500 km s^{-1} from observations of the optical lines (*34*) to 5200 km s^{-1} from modeling of the x-ray spectrum (*35*). Even if we adopt the highest estimate for the shock velocity, the lower bound on $T_e/T_{1,eq}$ in Tycho is 0.2. A substantial portion of this electron heating must be collisionless. One can show that collisions with the hot shocked ions can raise the electrons to a temperature of only about $1.8 \times 10^7\ (n_0 t_3 v_{s8}^2)^{2/5}$ K (where t_3 is the age of the remnant in units of 10^3 years), and for Tycho the data (*32, 35*) imply that this is less than about 3×10^7 K. There is some evidence for electrons that are substantially hotter than 7×10^7 K: Pravdo and Smith (*31*) reported evidence for a high-temperature component, with $T_e \gtrsim 2 \times 10^8$ K. The interpretation of such a hot component is complicated by the fact that collisionlessly heated electrons need not have a maxwellian distribution [compare this with the discussion of Asvarov *et al.* (*36*)].

Supernova remnants such as Tycho have unusual optical spectra that are dominated by the Balmer lines of hydrogen, and the analysis of these spectra provides the only method for inferring the temperature of the shocked protons (*37*). The shock waves in young remnants such as Tycho are moving so fast that they do not preionize the gas ahead of the shock (see Fig. 2). If a fast shock propagates through a partially ionized medium, the ionized component is promptly heated by collisionless processes as described above. The neutral atoms can interact with the shocked ionized component of the plasma only through collisions. Some of these collisions will ionize the atoms; some will excite the atoms, producing Balmer lines with a narrow line profile characteristic of the ambient gas; and some will lead to charge transfer between the stationary neutral atoms and the fast ions, producing a population of fast neutral atoms. These fast atoms can themselves be excited, and their emission has a line profile that is directly related to the velocity distribution of the collisionlessly heated protons. The Balmer lines from nonradiative shocks in partially ionized gas thus have two components: a narrow line from the excited neutrals at rest, and a broad line from the fast neutrals produced by charge exchange with shock heated protons. The ratio of the intensities of these two components depends on the ratio of the electron and ion temperatures behind the shock. The most recent analysis of the spectra from Tycho and other similar remnants suggests that the electrons undergo substantial collisionless heating in the shock but that they do not reach the full equipartition temperature (*34*). The Balmer-dominated supernova remnants such as Tycho provide a unique window on the structure of strong astrophysical shocks, and further research in this area should be fruitful.

Shocks in Molecular Gas

The neutral gas in the ISM is largely clumped into clouds. In the low-density clouds, with densities of H nuclei $n_H \sim 10$ to 10^2 cm^{-3}, the gas is for the most part atomic, and the fractional ionization is $\sim 10^{-3}$. In the more dense clouds, with $n_H \approx 10^2$ to 10^6 cm^{-3}, most of the mass is in the form of H_2, and fractional ionizations range from 10^{-4} to 10^{-8}. Interstellar magnetic fields are difficult to measure, but measured Zeeman splitting of the 21-cm line of atomic H and of the 18-cm transitions of OH suggests that interstellar magnetic fields in clouds vary approximately as $B \propto (n_H)^{1/2}$, with a strength such that the Alfvén speed $v_A = B/(4\pi\rho)^{1/2} \approx 1$ to 3 km s^{-1} (*38*).

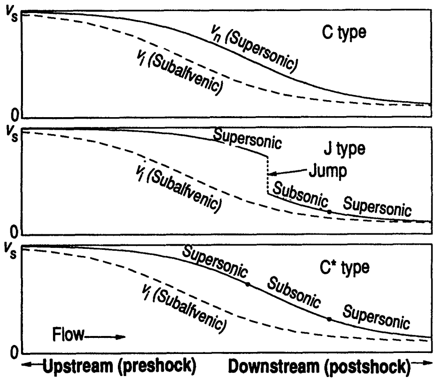

Fig. 4. Schematic structure of two-fluid shock waves showing velocities of the neutral and ionized fluids (relative to the shock front). Three types of solutions are possible: C type, J type, and C* type.

Although the ions and electrons in these clouds contribute only a small fraction of the mass and pressure, the plasma is dynamically important because only the charged particles couple to the magnetic field, and the magnetic field is generally strong enough to affect the dynamics. When the fractional ionization x_e is low, the neutral-ion collision rate is small, and the coupling between ions and neutrals becomes weak enough that it is important to think of the ions and the neutrals as distinct but interpenetrating fluids (*10, 39*).

For a shock to be present, the shock speed must exceed the speed for long-wavelength compressive waves in the medium (in this case, approximately the Alfvén speed). Because the fractional ionization x_e is low, the ion fluid itself has little inertia (its density ρ_i is much less than the total density ρ), and short-wavelength magnetic disturbances can propagate with a large signal speed $c_i = B/(4\pi\rho_i)^{1/2} \approx 70(x_e/10^{-4})^{-1/2}$ km s^{-1}. The effect of ion-neutral collisions is to damp these magnetohydrodynamic (MHD) waves, but because of this large effective sound speed c_i for the ions they are able to communicate ahead of the shock, provided that the shock speed $v_s < c_i$. From the standpoint of the ions, there is then no hydrodynamic surprise, and the ion density and velocity vary smoothly through the shock.

Because the ions and the neutrals are differentially accelerated, the ions stream through the neutrals in the shock, a phenomenon that is referred to as ion-neutral slip. Ion-neutral collisions, then, both heat and accelerate the neutral fluid. The resulting shock structure will take one of three generic forms (Fig. 4). If the neutral fluid remains cold (either because the shock is weak or because radiative cooling is able to remove the heat dissipated by ion-neutral collisions), then the neutral fluid itself may remain everywhere supersonic (in the frame of reference where the shock front is stationary), and in this case the neutral fluid velocity and density also vary continuously through the shock transition; such shocks are termed C type (*10*) because all the flow variables are continuous. If, however, the heating by ion-neutral collisions raises the neutral gas temperature, and therefore the sound speed, enough that the neutral flow (which was initially supersonic) becomes subsonic, then two possibilities exist. The neutral fluid may make the supersonic→subsonic transition by means of a collisional subshock (with a thickness on the order of the neutral mean free path and with energy dissipation and entropy generation within the subshock being due to molecular viscosity); such a shock is termed J type because it contains a jump in the neutral hydrodynamic variables. Under some circumstances, however, it is possible for the neutral fluid to make the supersonic→subsonic transition smoothly; such shocks are termed C* type (*40, 41*). The domains of C-type, J-type, and C*-type shocks have been delineated only approximately (*41–44*); the boundary within which C-type shocks occur is indicated roughly in Fig. 2.

It should be stressed that both far upstream and far downstream from the shock the ion and neutral fluids are always moving together and close to equipartition, so that in the absence of radiative cooling (or effective cooling due, for example, to molecular dissociation) these shocks would obey the same overall jump conditions as single-fluid MHD shocks. What makes these multifluid shocks different is that ion-neutral slip provides a frictional mechanism that can convert the directed kinetic energy of the upstream gas into heat at a rate that (because of the rarity of ions) is low enough that radiative cooling may be able to maintain the fluids at temperatures far lower than the temperature $T_{1,eq}$ that would have been attained in the absence of cooling. In effect this provides a situation where the shock transition and the radiative cooling zone (Fig. 1) coincide. As a result, molecules can be accelerated to high velocities without being dissociated, and the power radiated by such shocks will largely appear in atomic and ionic fine structure lines and in rotation-vibration transitions of molecules. Because these long-wavelength photons are able to penetrate the dusty clouds in which these molecular shock waves are generally located, they render these shocks accessible to detailed astronomic study.

A number of star-forming molecular clouds are known to contain high-velocity molecular flows (that is, 10 to 10^2 km s^{-1}) sometimes with associated emission from vibrationally excited H_2. The best studied example is OMC-1, a dense star-forming clump of dust and gas in the Orion molecular cloud. The powerful H_2 line emission from OMC-1, as well as emission from rotationally excited OH and high rotational states of CO, has been attributed to a C-type shock wave propagating outward into cool, dense molecular gas (*9, 45*). In these models the preshock gas has $n_0 \approx 4 \times 10^5$ cm^{-3}, $B_0 \approx 10^{-3}$ G, and $v_s \approx 37$ km s^{-1}. A conventional single-fluid shock of 37 km s^{-1} in a gas with a mean molecular weight of 2.3 amu and an adiabatic exponent $\gamma = 5/3$ would have $T_{1,eq} = 72{,}000$ K; in the C-type shock models for OMC-1, the neutral gas temperature remains below 3000 K. These C-type shock models predict that there should be powerful emission from rotationally excited H_2O in the shocks (this is unobservable from the ground because of strong telluric H_2O absorption); this prediction will soon be tested by far-infrared spectroscopy from the Infrared Space Observatory (ISO) satellite, to be launched by the European Space Agency in 1993. Maser emission from rotationally excited H_2O has been observed in both the radio and the submillimeter regions of the spectrum (*46, 47*), and has been attributed to shocks (*48*).

Although there are fundamental reasons for advocating multifluid shock waves to explain the H_2 emission from OMC-1, there are also several disturbing problems with these models. In the C-type shock models, the gas temperature rises to a peak value that depends on the fractional ionization, the shock speed, the magnetic field strength, and the angle θ between the shock normal and the direction of the preshock magnetic field. Because these parameters (particularly θ) would naturally be expected to vary from point to point on the extended shock front, it was expected that various temperature-sensitive line ratios would show spatial variability. Careful observational studies (*49*) have instead found remarkable uniformity in the ratios of intensities from the 0–0S(13) and 1–0O(7) lines. In addition, the C-type shock models for OMC-1 do not predict sufficient emission from either high (for example, $v = 4$, $J = 5$) or low (for example, $v = 0$, $J = 4$) energy levels of H_2 (*50*). Finally, the observed broad H_2 emission line profile in OMC-1 [full width, $\gtrsim 140$ km s^{-1} (*51*)] is not explained by the simplest C-type shock models. To account for these observations, Brand *et al.* (*50*) have proposed nonmagnetic single-fluid shock models, but these shock models have their own difficulties (*52*).

Some of the failures of the C-type shock models for OMC-1 may be attributable to trying to model a complex outflow with a single, spherically symmetric shock front. Models in which the H_2 emission originates in bow shocks ahead of fast-moving gas clumps have been considered for OMC-1 as well as for Herbig-Haro objects (*53*). Nevertheless, one suspects that the shock modeling may still lack some important physical element. Has some important molecular physics or chemistry been overlooked? In low–fractional ionization gas (such as in OMC-1) the dust grains play an important (but complicated) role in the dynamics (*43, 54*); are we treating the dust grains satisfactorily? Perhaps the plane-parallel, steady-state conditions favored by the shock modelers are far from the truth.

Wardle (*55*) has pursued this last possibility by studying the dynamic stability of plane-parallel steady C-type shock waves. He discovered that these shocks are generally stable for $M_A \lesssim 5$ but are unstable to a dynamic instability for $M_A \gtrsim 5$. The mechanism driving the Wardle instability is as follows: Consider a plane-parallel shock with $\mathbf{B}_0 \cdot \mathbf{v}_s = 0$ and with low–fractional ionization, so that the ion fluid has negligible inertia [that is, $B_0/(4\pi\rho_i)^{1/2} \gg v_s$], in which case the net force per unit volume on the ionized fluid at each point must be essentially 0. The dominant forces on the ion-electron fluid are $\mathbf{J} \times \mathbf{B}/c = -\nabla(B^2/8\pi)$ and the drag force due to ion-neutral slip, which is proportional to $\rho_n\rho_i(\mathbf{v}_n - \mathbf{v}_i)$.

Now consider what would happen if the straight magnetic field lines of the plane-parallel steady solution were to be perturbed as in Fig. 5. The drag force will now have a component parallel to the local magnetic field that cannot be balanced by $\mathbf{J} \times \mathbf{B}$ forces, and ions will therefore be accelerated along the field lines to collect in the magnetic valleys. As a consequence ρ_i will increase in the valleys, the drag force proportional to ρ_i will increase, and the distortion of the field lines may be further increased. Wardle has demonstrated the existence of this instability through a linear stability analysis, but the nonlinear development remains to be investigated, so that it is not yet known to what degree these unstable shocks will differ from the idealized steady-flow solutions that have been studied numerically. The perturbation grows most rapidly for wavelengths approximately equal to the thickness of the shock transition. It seems entirely possible that, after this instability has grown to saturation, the emission spectrum from the shock front (even after time averaging or spatial averaging over the shock front) may be quite different from the spectra computed for the unperturbed steady-flow solution. Because the flow in OMC-1 is likely to be highly unstable, it is perhaps not surprising that steady C-type models are unable to reproduce the observed spectrum.

In the J-type shock structure of Fig. 4, the region ahead of the jump is referred to as a magnetic precursor. Even relatively high-velocity shock waves in molecular gas should have such a precursor, in which the low-ionization preshock gas may be heated slowly enough that the molecular material can survive and radiate until arriving at the jump, at which point the gas will be suddenly heated to temperatures at which molecules may be rapidly destroyed.

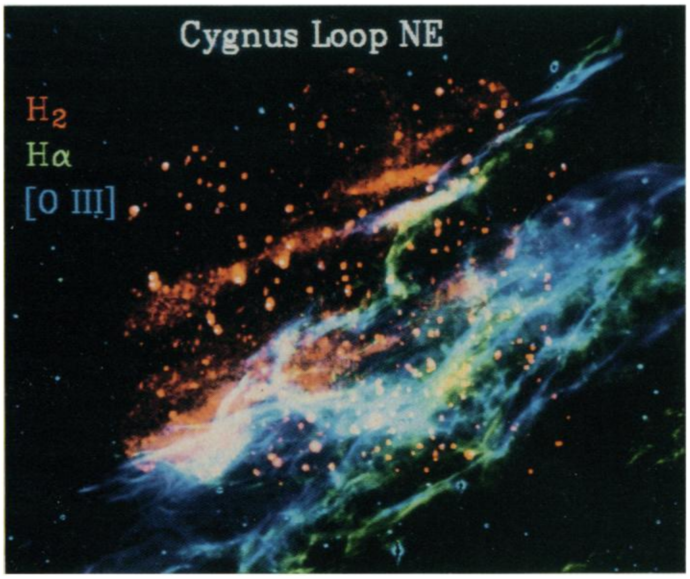

Fig. 6. A small portion of the bright filaments to the northeast of the Cygnus Loop supernova remnant imaged in three spectral lines: $H_2$1-0*S*(1) (red), Hα (green), and [OIII]5007 (blue). The center of the remnant is to the lower right. The H_2 emission may be produced in a magnetic precursor to a fast shock wave; the Hα and [OIII] emission originate in postshock gas. [Reprinted from (*56*) with permission]

Recent observations of the Cygnus Loop supernova remnant by Graham *et al.* (*56*), which are shown in Fig. 6, have revealed extended emission from vibrationally excited H_2. Graham *et al.* have interpreted this emission as emanating from magnetic precursors to shock waves of about 200 km s^{-1} propagating into diffuse molecular clouds at the edge of the supernova remnant. This interpretation predicts that the H_2 emission should lie just outside regions of strong Hα emission (in agreement with Fig. 6) and that the H_2 itself should be at nearly the same velocity as the unshocked material; this has recently been observationally confirmed by high-resolution spectroscopy (*57*). From a theoretical standpoint it is somewhat surprising that the fractional ionization ahead of the shock is sufficiently low for the magnetic precursor to exist in these clouds because ionizing radiation from the shocked gas would be expected significantly to ionize the preshock gas. Further studies of this region will be of considerable interest.

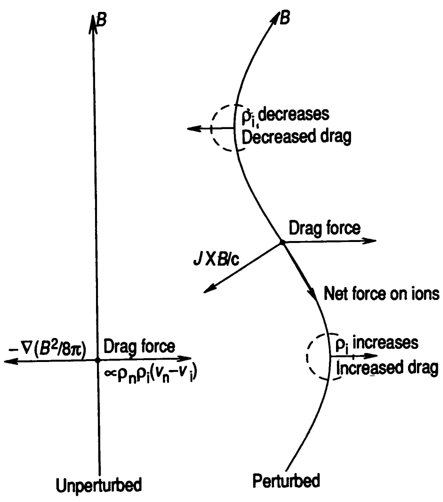

Fig. 5. Mechanism for the Wardle instability in two-fluid MHD shock waves.

Future Directions

Shock waves are a central element in interstellar physics. Existing theory allows us to understand quantitatively many interstellar shock phenomena (such as the sizes and rates of expansion of supernova remnants and the spectra of radiative shocks in atomic gas), but a great deal of the physics of interstellar shock waves remains to be elucidated.

One of the subtleties in the structure of interstellar shocks that has become apparent in the past decade is the importance of trace constituents in the interstellar plasma. In collisionless shocks, the acceleration of a small fraction of the particles to high energies can both alter the structure of the shock and produce the cosmic rays; in shocks in molecular gas, the coupling of the small concentration of ions to the dominant neutral gas through the magnetic field completely alters the structure and the emitted spectrum of the shock. Theoretical work, primarily through increasingly powerful computer simulations, should clarify the outstanding issues associated with the structure of these shocks, including collisionless heating of the electrons, diffusive shock acceleration, and the stability of two-fluid MHD shocks. Spectral lines already provide powerful diagnostics of interstellar shocks; further progress in measurement or calculation of various poorly determined inelastic cross-sections (for rotational excitation of H_2 by H atoms, for instance) is essential if we are to be able to draw more precise inferences from the observations.

Finally, we can expect rapid observational progress. The incorporation of large-format charged-coupled devices and infrared detector arrays into cameras with narrow-band filters or tunable Fabry-Perot interferometers is allowing astronomers to obtain spectral line images of faint shocked regions with excellent spatial resolution, thereby shedding light on details of shock structure. Continued progress in detectors is allowing sensitive spectroscopy in the infrared from the ground and from airplanes. Finally, observations from orbiting observatories such as Hubble Space Telescope, ISO, Space Infrared Telescope Facility, Far Ultraviolet Spectroscopy Explorer, and Advanced X-Ray Astrophysics Facility will permit sensitive spectroscopy at wavelengths blocked by Earth's atmosphere, thereby broadening the confrontation between theory and observation.

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