Hot Gas in Interstellar Space

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Abstract

Supernova explosions produce shock waves which heat the interstellar gas to temperatures exceeding 10^6K. This hot gas expands, interacts with clouds present between the stars, rises to appreciable distances from the galactic plane and generally affects the structure, dynamics and evolution of the interstellar medium. Recent theories of such processes are briefly reviewed.

I. *Introduction*

It is highly appropriate to include interstellar matter in a series of scientific papers in honour of Bengt Strömgren, whom I remember warmly as mentor, colleague and close friend. His two major theoretical contributions to this subject, made some 40 to 50 years ago, are still of fundamental importance. The differences which he pointed out between HII regions, in which hydrogen atoms are nearly all ionized, and HI zones, where they are nearly all neutral, pervade all our discussions of the interstellar gas. His interpretation of interstellar absorption lines has served as a model for subsequent investigations; his results on the chemical composition of the gas, especially on the overall ratio of calcium to sodium, are still valid qualitatively.

More recent work on interstellar problems has profited from new observational tools. In particular, the existence of a hot gas between the stars has been demonstrated in the last two decades by observations from instruments outside the Earth's atmosphere. To be sure, ground based detection of apparently normal absorbing clouds far from the galactic plane, in the galactic halo, had suggested much earlier the presence of a surrounding hot gas, whose pressure could keep these clouds from expanding. Definite identification of such a hot gas has been obtained not in the halo but in the galactic disc; this result was achieved with two types of observations from space: ultraviolet absorption lines of OVI and other highly ionized species, and soft X rays emitted by hot plasma in extended neighbouring regions of the Galaxy.

The extensive observational material concerning this hot interstellar gas, with a kinetic temperature in the range from 10^5 to 10^7K, has been clearly summarized in recent broad reviews by Cox and Reynolds (1987), Jenkins (1987) and Savage (1987). The present paper makes no attempt to duplicate these summaries, but treats instead some of the theoretical work that has been done, especially during the last few years, on the origin and development of this hot gas.

Understanding the processes which occur as the hot interstellar gas evolves is an ambitious goal which we are far from achieving. The dynamics of a compressible gas, subject to the photons and cosmic rays in interstellar space, is a complex topic. Some progress has been made through the development of idealized models, which are so simplified that one can hope to understand them and to compute their properties. In terms of such models one can distinguish three scenes in the unfolding drama of the hot interstellar gas. First the explosion of a supernova ejects a rapidly expanding envelope, whose interaction with the surrounding medium creates the hot gas in which we are interested. Next, as this heated gas expands it encounters regions whose internal density is well above the average. These regions, which we call clouds, are compressed by the hot gas, are heated by conduction and sometimes evaporate or are disrupted. In the final scene the remnant of heated gas surrounding one or more supernovae can rise to appreciable distances from the galactic plane, and may produce a hot galactic corona before it falls back down or escapes the Galaxy entirely.

In actuality these three scenes overlap so much that their mutual interactions are important. In most theoretical models these scenes have been kept somewhat separate to simplify the theory and to clarify what happens in at least a few highly simplified situations.

During the last few years theorists have constructed a number of such simplified models. To describe the details of all these models would require a substantial monograph. The present paper comments briefly on a few models, indicating the simplifications made and the general character of the results, together with some of the chief problems remaining. After discussions of the three evolutionary scenes listed above, a final section treats the vertical structure of the interstellar medium, again through discussion of simple models. The active dynamical processes in which the hot gas participates, often as the primary driving force, must play a major role in an overall account of the interstellar medium, particularly its structure and evolution.

II. *Expansion of Supernova Remnants*

The phenomena which follow a supernova explosion can in principle be followed in rather full theoretical detail if spherical symmetry is assumed. Such a spherical model is applicable if the stellar explosion itself produces this symmetry, if the initial properties of the surrounding interstellar gas are functions only of r, the distance from the supernova, and if the magnetic field **B** is ignored. In addition one must assume that no non-spherical instabilities will arise. Under these conditions, all quantities are functions of radius r and time t, and the relevant differential equations can be integrated. While this spherical symmetry is not likely to be realized in detail, it may provide an adequate first approximation, especially in those regions where the interstellar gas is resonably homogeneous; an initial particle density independent of r is usually assumed.

Most theoretical models also make the restrictive "hydrodynamic" assumption that the mean free path of all particles is much less than r. This assumption has the great advantage that it yields the familiar equations of fluid dynamics, for which the techniques of numerical solution have been much studied. Physically the hydrodynamic assumption leads to a thin shock wave and to negligible conductive heat flow. In fact the mean free path of protons and electrons for 90° deflections in two-body encounters is many parsecs for a newly born supernova remnant, decreasing to about 1 pc when $T = 10^7$K, if the proton density is $0.1/cm^3$. A magnetic field restricts the travel of charged particles transverse to **B**, but does not yield the hydrodynamic approximation for motions parallel to **B**. Other processes have been suggested which can reduce the effective mean free path. Energetic particles moving through an ionized gas are sometimes slowed down by the plasma instabilities which they excite. This complex effect can perhaps provide full justification for the hydrodynamic assumption, which is also consistent (McKee and Hollenbach 1980) with the relatively sharp boundaries observed for the X-ray emission from some young supernova remnants, notably around the entire circumference of Cas A. However, definite confirmation is lacking.

We ignore initially here both the magnetic field, which is almost certainly present, and the uncertainties associated with the hydrodynamic assumption. The effects of thermal conductivity and of a **B** field are discussed briefly at the end of this section. Effects produced by the relativistic particles constituting cosmic rays have not been much considered in models of supernova remnants and are ignored here; in some circumstances such effects may be highly important. The processes which result when the initial ambient distribution is cloudy or has a vertical density gradient are treated in subsequent sections.

The spherical models based on these assumptions have yielded a substantial body of knowledge on supernova remnants. Most such models assume that there is no important energy source in the supernova core after the explosion; e.g., any radiation from a rapidly rotating neutron star, produced by the collapsing stellar core, is ignored. Three familiar evolutionary stages are then distinguished. First there is free expansion of the ejected material, a stage which lasts as long as the ejected mass is large compared to the mass of the swept-up interstellar gas. Subsequently, the interstellar mass swept up by the outwards moving shock, of radius r_s, much exceeds the ejected mass. As long as the kinetic temperature is high enough throughout the remnant that radiative cooling is slight, this is the well known "Sedov-Taylor stage" (see the monograph by Ostriker and McKee 1988); the density increases rather steeply outwards, with half the mass in the outer six percent of the radius, and the shock velocity V_s varies as $r_s^{-3/2}$, giving $r_s \propto t^{2/5}$.

With increasing r_s, the postshock temperature T_s decreases as V_s^2. When T_s falls below roughly 10^6K the increased rate of radiation then cools the postshock gas to a temperature below 10^4K, and the density increases by a large factor behind the shock, forming a cold shell much thinner than before; this is the third or "snowplow stage." As long as the internal pressure deep within the remnant remains high, the outwards momentum of the cold shell gradually increases, and V_s now varies as $r_s^{-5/2}$ (Cox 1972), giving $r_s \propto t^{2/7}$.

For an actual remnant these stages are approximations. An exact numerical solution of the fluid dynamical equations (Cioffi et al. 1988), including radiative emission in the appropriate energy equation, shows spherical disturbances moving inward and outward, resulting in large part from transitions between successive stages. As a result of these disturbances and the limited duration of each stage, there is only rough agreement with the predictions for the various stages in isolation. The combination of numerical calculations with approximate analytic results gives a reasonably complete understanding of the idealized spherical model, based on the hydrodynamic assumption and an initially uniform interstellar gas density.

In subsequent discussions we shall take as typical parameters for supernova remnants the results obtained in this numerical model, with an ambient particle density of 0.1 atoms/cm^3 and an initial kinetic energy of 0.93×10^{51} ergs in an ejected envelope of mass $3M_\odot$. In this model the swept-up and ejected masses are equal at roughly 10^3 years. The cold shell forms during the interval from 1.2 to 1.7×10^5 years as the shock velocity drops below about 120 km/s; the shock radius r_s is about 55 pc during this interval.

We turn now to the effects which thermal conduction and magnetic fields can produce in supernova remnants. A model which takes into account thermal conduction by electrons as well as energy exchange between ions and electrons shows (Cowie 1977) that after the initial free expansion stage the electron temperature T_e becomes nearly constant with radius inside the remnant, a marked change from the Sedov-Taylor solution. When the electrons are nearly isothermal the positive ion temperature, T_i, assumed to equal T_e immediately behind the shock, is found to increase inward. For the particular case treated the rate of expansion is not much altered by such effects, with r_s at each time increased by not more than 8 percent. These detailed results depend on the assumption that no heat conduction occurs through the moving shock front. More detailed computations (Cox and Edgar 1983 and 1984) show the effects of thermal conduction and of electron-ion energy exchange on various properties of an evolving supernova remnant.

If a magnetic field **B** is present, as seems very likely, thermal conduction transverse to **B** will be almost completely suppressed. Thus T_e will tend to be constant along **B**, but to show a variation of Sedov-Taylor type in planes transverse to **B**. The positive-ion temperature T_i will exceed T_e deep within the remnant, since approach to equipartition through electron-ion encounters is relatively slow (see refs. in preceding paragraph); conduction of heat by positive ions contributes significantly to holding down dT_i/dr along **B**.

The effects which a magnetic field produces on a supernova remnant are more conspicuous if the energy density $B^2/8\pi$ is at least comparable with the material pressure nkT, where n is the total number of particles per cm^3. An equivalent condition is that the Alfvén speed $V_A = B/(4\pi\varrho)^{1/2}$ be at least comparable with the isothermal sound speed $C_s = (kT/m)^{1/2}$, where $m = \varrho/n$ is the mean mass per particle. This situation can arise even if the magnetic field in the gas surrounding the supernova is initially very weak.

One such case for which detailed calculations have been made (Kulsrud et al. 1965) is the hydromagnetic flow around a conducting spherical shell, or "piston," assumed to be expanding at a constant rate into a conducting gas permeated by an initially weak and uniform magnetic field. The lines of force which have been pushed outward by the piston tend to accumulate in a thin boundary layer, where the magnetic field grows steadily until its energy density becomes comparable with that of the streaming gas. In an actual remnant the ionized ejected gases from the supernova may take the place of the expanding piston. The high magnetic pressure in the surrounding boundary layer would then tend to decelerate the inner ejected gases and the Rayleigh-Taylor instability should occur. While the analysis is evidently idealized and other effects will certainly be present, such a process may play a part in producing the filaments of high magnetic field observed in the inner region of the Crab supernova. The shock which moves outward from the piston into the surrounding gas is not much affected by the magnetic field, whose energy density just behind the shock is relatively small.

A uniform interstellar **B** field can produce important effects during the snowplow stage of supernova expansion. With representative parameters for the warm interstellar medium ($n_H = 0.15$ cm^{-3}, T = 6300K, B = 3μG) the Alfvén speed V_A is 14 km/s, substantially greater than the isothermal sound speed C_s of 6.1 km/s. As pointed out above, when the cold radiative shell starts to form, the shock velocity V_s is about 120 km/s; we adopt 40 km/s as a representative value of V_s during the snowplow phase. For these parameters the relative increase of density across an isothermal shock (Spitzer, 1978) is a factor $(V_s/C_s)^2 = 43$ if **B** vanishes or is parallel to $\mathbf{V}_s$, but is only 3.4 (approximately $2^{1/2}\,V_s/V_A$) for propagation transverse to the assumed magnetic field. This latter compression is not only weak but nearly reversible, since the energy stored in compressing the magnetic field can drive a reexpansion when the pressure falls. If the high compression "parallel shocks" are assumed to be relatively infrequent and the low-compression transverse ones are regarded as dominant, one may conclude (Cox 1986 and 1988) that the late expansion stages of a supernova remnant produce only a very modest compression of the interstellar medium.

To evaluate the compression expected in an actual interstellar situation one must consider oblique shocks, with a wave normal at some arbitrary angle to the magnetic field. The physical principles governing such shocks are summarized in the accompanying Appendix. It turns out that for the conditions specified above, a single shock, even if parallel to the magnetic field, produces a relatively low compression (a factor 8.2). However, two successive parallel shocks (with an inclined magnetic field between them) can produce the same high compression found above for a single shock with $B = 0$, and constituting effectively the "parallel shock" discussed above. Similar results are presumably possible for two successive shocks within some range of directions relative to **B.** Until such possibilities have been explored, the average compression in the late stage of a supernova shock is essentially unknown, except that it presumably lies between $(V_s/C_s)^2$ and about $2^{1/2}V_s/V_A$.

The amount of this compression has important effects on the structure and dynamics of the interstellar gas. A familiar picture of the interstellar medium, often used as a standard of reference, is based on the sweeping synthesis by McKee and Ostriker (1977), which brings into one theoretical framework many different aspects of this medium. If the compression in an isothermal shock during the snowplow phase were much reduced by magnetic forces, with the compression ratio decreasing from 4 to 1 as V_s decreases from 40 down to about 14 km/s, some aspects of this picture would require modification (Cox 1986). In particular, the warm neutral medium (WNM) would not be swept up into dense shells but would occupy an appreciable part of the remnant's volume, reducing the fraction of this volume occupied by the hot gas. The overall filling factor f_h of the hot gas (the fraction of the volume of the galactic disc occupied by this gas) would be correspondingly reduced. In addition, the evolution of the remnant after expansion ceased would be greatly altered in detail.

In reality strong compression is likely to occur along some lines of the magnetic field. Coupling between motions parallel and transverse to **B** may conceivably convert the enhanced magnetic energy of the compressed WNM into radiation from dense clumps of cooling gas. In any case, magnetic tensions and pressures in the complex interstellar medium are likely to produce unexpected consequences. Computers are reaching the power necessary to follow such hydromagnetic processes approximately. Until more knowledge is available either from analysis or from simulations, theory cannot indicate the hot gas filling factor. The topology of this gas – isolated hot bubbles at one extreme and isolated cooler clouds (either warm or cold) at the other – is equally unclear. Observational information on these questions is also not definite.

III. *Interaction of Remnants with Clouds*

When a star explodes, the expanding remnant may sweep through a gas quite different from the uniform interstellar medium discussed above. It is well known (Spitzer 1985) that the gas between the stars has inhomogeneities with a wide variety of sizes, ranging from filaments less than a parsec across to giant molecular clouds and cloud complexes a hundred parsecs in size. If the exploding star was particularly luminous or its companions had exploded a short time before, the local gas may have been greatly modified and perhaps somewhat homogenized by ultraviolet photons and by expanding hot gases. We do not treat here the details of these complex environments, but consider some of the physical processes which can occur when a supernova remnant engulfs a cloud. Since a full review of such processes has recently appeared (McKee 1988), the present discussion is brief and highly selective.

Strong compression of clouds is believed to be a major effect produced by supernova remnants. A combination of analytic theory and numerical simulation gives an approximate indication of how this "cloud crushing" might proceed. A passing supernova shock generates a slower shock in the denser cloud material. Behind this cloud shock the velocity field leads both to compression and shear, and can produce instabilities which may disrupt the cloud at least in part. An analysis of cloud compression when a magnetic field is present indicates (Oettl et al. 1985), as expected, that if the magnetic pressure is dominant motions transverse to **B** are suppressed. The effect of instabilities may also be less when the compression is one-dimensional, though further study would be needed to establish this conclusion.

Another major effect produced by remnants is cloud evaporation as a result of thermal conduction from the hot gas. The shock itself is not directly involved in this process, which is usually modelled with the hot gas in pressure equilibrium with the cloud and with no systematic velocity of the cloud with respect to the gas. The nature of this process depends on a global saturation parameter σ_0 (Cowie and McKee 1977), essentially the ratio of the electron mean free path to the cloud radius a, and equal to 0.4 $T_{f7}^2/n_{ef}a_{pc}$; T_{f7} is the asymptotic temperature of the hot gas, far from the cloud, in units of 10^7K, n_{ef} is the asymptotic electron density and a_{pc} is the cloud radius in pc. When σ_0 is small compared to unity, the thermal conductivity is given by its classical value, the temperature distribution (in this three-dimensional situation) has a quasi-steady state, and the velocity of the gas in the expanding envelope is subsonic everywhere, However, if σ_0 is too small, less than about 0.03, (corresponding to a_{pc} exceeding about 10 pc in a typical situation) the heat flow is inadequate to offset radiative cooling, and condensation of the hot gas replaces evaporation of the cloud (McKee and Cowie 1977). On the other hand, when σ_0 exceeds unity the physical situation becomes more complicated; the heat flow in this "saturated" condition can be estimated from observations of the solar wind (Cowie and McKee 1977) and the resultant mass loss computed.

A recent investigation (Draine and Giuliani 1984) indicates that for large σ_0 the effect of viscosity, produced by atomic ions, must also be considered. Calculations based on the simplifying assumption that $T_i = T_e$ show that for $\sigma_0 \geqslant 100$ the pressure in the cold cloud can substantially exceed that in the hot gas, and that the viscosity can increase the mass loss rate by about an order of magnitude above the inviscid case; the viscosity requires a pressure increase within the cloud to drive the flow, and at the resultant higher density the heat flow and the resultant mass-loss rate are increased. If the positive ions are heated only by two-body encounters with electrons, T_i will be much less than T_e and these viscous effects will be much reduced.

A further modification in the theory is needed if a magnetic field is present; as pointed out above, there is virtually no conductive flow of heat across an interstellar magnetic field. Limitation of the conductive flux to the direction of **B** has been taken into account (Balbus 1986) for the evaporative flow from an infinite plane surface, initially separating a cold gas on one side from a hot gas extending infinitely far on the other. In this model all variables depend only on t and on z, the distance from the plane; both $B^2/8\pi$ and ϱv^2 are assumed small compared to p. No steady state is possible in this one-dimensional situation (unless a surface at some fixed temperature is located a finite distance away). Instead a self-similar solution is obtained, with a similarity variable proportional to $z/t^{1/2}$; thus as t increases the mass loss rate decreases and the conductive front thickens.

An interesting result of this analysis is that the initial rate of evaporation varies roughly as $\cos^2\Theta_\infty$, where Θ_∞ is the angle between **B** and the z axis at large z. This may be understood physically, since the temperature gradient parallel to **B** equals $\cos\Theta dT/dz$, and the component of the heat flux parallel to z varies as $\cos^2\Theta dT/dz$.

Since the magnetic fields within diffuse clouds (determined from the Zeeman effect of 21-cm lines) are apparently about the same as those in the warm ionized medium (Heiles 1987), a model of straight, parallel lines of force extending into the hot gas may, perhaps, provide a reasonable approximation. On this basis the thickening of the conduction front with time, resulting from the one-dimensional character of the heat flow, should be an important effect, probably more so than the dependence of the flow on Θ_∞.

One attribute of these various conductive evaporating envelopes, with or without magnetic fields, is that they contain highly ionized atoms such as O^{+5}, whose OVI absorption features have been observed along numerous lines of sight through the interstellar gas. At a kinetic electron temperature above 10^5K oxygen atoms will be highly ionized by electron collisions. If collisional ionization equilibrium is assumed in conductive envelopes, the fraction of oxygen in O^{+5} ions is greatest at $T = 3\times10^5$K.

In these envelopes the temperature of each fluid element in the moving gas is changing with time, and hence collisional ionization equilibrium will not be fully reached. The relative numbers of atoms in different stages of ionization must be computed from the relevant differential equations, including the rates of ionization and recombination. Such computations (Ballet et al. 1986) for an outwardly expanding conductive envelope give a substantial increase in the total number of O^{+5} ions in the envelope. We denote by F the ratio of this number to its value in collisional equilibrium; this factor F depends on $n_{ef}a_{pc}$, where again n_{ef} is the asymptotic electron density, far from the cloud, and a_{pc} is the cloud radius in parsecs. Detailed calculations show that for an asymptotic temperature T_f equal to 10^6K, F increases from 2.5 at $n_{ef}a_{pc} = 0.10$ pc/cm^3 to 40 at $n_{ef}a_{pc} = 0.01$ pc/cm^3.

These results, combined with an assumed cosmic composition, indicate that $<n(O^{+5})>$, the mean particle density of O^{+5} ions in the galactic disc, is about 2×10^{-7}cm^{-3} for a = 5 pc, $n_{ef} = 0.01$ cm^{-3} and $T_f = 10^6$K; the filling factors f_c and f_h for the cold clouds and the hot gas in the galactic disc are set equal to 0.02 and 1, respectively. If n_{ef} is decreased to 0.001 cm^{-3}, the computed O^{+5} density rises to 5×10^{-7}cm^{-3}. These densities exceed by an order of magnitude or somewhat more the observed mean value of about 2×10^{-8}cm^{-3} (Jenkins 1987).

Later computations for the ionization distribution in expanding conductive envelopes (Böhringer and Hartquist 1987) give similar results. These were used to compute $N(O^{+5})$, the column density of O^{+5} ions, integrated over dr from r = a to r = 10a. The resultant values were found to be clustered around 10^{13}cm^{-2}, with a relatively slow dependence on $n_{ef}a$; as this parameter increases, the increased mass loss rate is offset by a decrease in F. This theoretical column density of 10^{13}cm^{-2} agrees with the observed mean value (Jenkins 1978b) for two-thirds of the O^{+5} gas. (The remaining components have larger column densities and were not included in the value of $<n(O^{+5})>$ cited above.) However, the observed values should exceed these theoretical ones by a geometrical factor. For a line of sight through the cloud center the column density is twice the theoretical one. For lines passing tangentially through the shell of highest $n(O^{+5})$ the increase will be somewhat greater; an average increase by a factor two should provide a rough approximation.

While one would not expect such idealized models to correspond closely with reality, it is of interest to note that plausible changes in the assumed parameters can bring theory and observation into rough agreement. It is known (York et al. 1983) that oxygen is depleted in interstellar clouds, with a depletion factor δ_O between 0.4 and 0.7. If we set $\delta_O = 1/2$, this factor compensates for the geometrical factor discussed above, leaving the observed and computed column densities in agreement. If also f_h, the filling factor of the hot gas, is set equal to 0.2, the theoretical value for $<n(O^{+5})>$ is then reduced by an order of magnitude and agrees with the observations to within the many uncertainties involved. About this same value of f_h is needed to give the observed average number (Jenkins 1978b) of about six O^{+5} conductive envelopes (each with a radius typically of about 2a) in the line of sight per kiloparsec.

The O^{+5} velocity distribution computed for the conductive model appears to be not inconsistent with the observations. For the thermal velocity spread of these ions within a single envelope the model yields a value between 14 and 18 km/s "in most cases" (Böhringer and Hartquist 1987), corresponding to temperatures between 4×10^5 and 6×10^5K. Because of overlapping components the observed line profiles can give only the minimum values of v_m, the rms velocity dispersion. The envelope expansion velocity can increase v_m appreciably for some lines of sight, but will have little effect on the minimum values, corresponding to lines of sight passing tangentially through the outer layers. Thus one would expect from the theory that for envelopes in general the minimum v_m should be about 14 km/s.

While a number of the observed OVI profiles (Jenkins 1978a) have values of v_m less than 14 km/s, ranging down to 10 km/s, Jenkins (1978b) points out that because of observational errors one can infer only that values as low as 14 km/s are "indeed rather common." More precise observations would be required to determine whether the actual distribution of v_m values is consistent with the conductive envelope theory.

The observed dispersion of radial velocities for the system of OVI absorbing regions is about 26 km/s (Jenkins 1978b), higher than the 6 km/s observed for most interstellar clouds but perhaps consistent with the velocities anticipated for clouds which have recently been compressed and accelerated by a supernova shock wave. Along individual lines of sight the velocity structure of the OVI lines appears correlated with that for lines of the less highly ionized species, SiIII and NII (Cowie et al 1979), supporting a common origin in the same set of conductive envelopes for at least some components of these different lines.

Two additional problems must be considered before the theory of expanding conductive envelopes could be accepted as an approximate quantitative explanation of the OVI observations. The first is the part played by the warm gas (WNM). According to the theory, $\langle n(O^{+5}) \rangle$ is proportional to the filling factor of the gas which is evaporating, which for the WNM is at least an order of magnitude greater than for the cold diffuse clouds. Perhaps evaporation of warm gas does not contribute many O^{+5} ions because this low-density gas evaporates so quickly when it is in contact with hot gas. The second problem is the effect of a magnetic field, which as we have seen above is likely to produce a marked thickening of the conductive envelope with time, as a result of the resultant one-dimensional flow. The theoretical calculation of O^{+5} densities should be repeated, taking this effect into account.

Some absorption by O^{+5} ions should also be produced in the inner regions of an idealized spherical supernova remnant, when some of the remnant gas cools through temperatures of some 3×10^5K. Since the initial temperature of the remnant decreases outwards, the outer layers cool first and a radiative cooling front gradually eats its way into the inner hot bubble, surrounded by the cold shell of the snowplow stage. As time goes on, the outwards velocity of the gas which is cooling will become less, and by the time one remnant overlaps with others this velocity may be sufficiently low to be consistent with the OVI observations. If a compressed transverse magnetic field is present, the restoring force on the gas may cause a more abrupt drop in the expansion velocities. The inner layers of hot gas which are cooling down to 3×10^5K and below have been proposed (Cox 1986) as sites for the observed OVI absorption. If effects of initial inhomogeneities and of thermal conduction were included, these cooling layers might have a somewhat similar appearance, though a different course of development, as the conductive envelopes between cold clouds and hot gas discussed above.

The development of a supernova remnant must be strongly affected by the various interactions with clouds, especially by cloud compression and evaporation. A detailed numerical calculation of this dependence was carried out some years ago (Cowie et al 1981), based on the clouds present in the galactic disc generally. Recently a different hydrodynamic technique has been applied (Wolff and Durisen 1987) to this same situation with general overall agreement. In particular, vigorous evaporation from clouds keeps the density nearly independent of r. As a result, the cold shell which is formed by radiative cooling, at a remnant age of typically some 10^5 years, first appears well inside the remnant rather than immediately behind the shock layer.

IV. *Effects far from the Galactic Plane*

The possibility that a hot gas might form a galactic corona at kiloparsec distances from the galactic plane (Spitzer 1956) has given particular interest to the possible role of supernova remnants in this connection, since the hot gases in such remnants provide an obvious heat source. While the remnant from a single stellar explosion can provide heating far from the galactic plane if the supernova is located in the halo or if the ambient particle density is $10^{-2}cm^{-3}$ or less, more energetic events provide a better source. It has long been clear (Chevalier and Gardner 1974) that adjacent explosions of several supernovae, as might be expected in young stellar groups, would more easily break out of the gaseous galactic disc and pervade the halo. Recent studies have emphasized such sequential explosions of supernovae and the "superbubbles" which they produce in this and other galaxies.

A primary reason for this emphasis has been the accumulating observational evidence (McCray and Snow 1979; Heiles 1979 and 1987a; Tomisaka et al. 1981; McCray and Kafatos 1987) for the existence of superbubbles not only in our own Galaxy but in other Local Group spiral and irregular systems. Extensive arcs and filaments seen in 21-cm emission reveal "supershells" of neutral H. X rays are detected from the hot gas within some of these supershells, and surveys of Balmer emission lines confirm such extended structures. The observed radii range from a hundred to a thousand parsecs; the total energies range up to 10^{53} erg or more, corresponding to about a hundred supernovae.

A second reason for the recent focus on superbubbles is the strong theoretical expectation that such concentrations of supernovae in time and space should in fact exist. Some supernovae (most of the Type I class) originate in old stars, which in the galactic disc show virtually no concentration in groups. However, the Type II supernovae and a few of Type I (Wheeler and Levreault 1985) result from core collapse in young, massive stars, which are apparently formed to a large extent in groups, including clusters and associations. Some of these escape as runaways. According to a recent survey (Gies 1987) about 70 percent of the O stars are now in groups and most of the supernovae produced when these massive stars die will be found within a radius of a few times 10 pc and a time interval of several times 10^7 years (McCray and Kafatos 1987). If several stellar groups are formed within the same cloud complex, the various superbubbles produced may overlap, producing an even more spectacular explosion.

Theoretical models of these superbubbles may be constructed, based on most of the assumptions made for a one-supernova remnant. For a superbubble, however, the disturbance can spread so far from the galactic plane that all physical quantities must be functions of two spatial dimensions, cylindrical radius r and vertical height z. Offsetting somewhat this complication is the initial homogeneity which may characterize the ambient interstellar gas, thanks to the processing of this medium by energetic stellar winds and intense ultraviolet stellar radiation emitted from the massive bright stars before these die and explode.

The nature of blast waves and outgoing winds under these conditions has received much study (see Schiano 1985 for a review). An important aspect of such disturbances in an exponential atmosphere, for example, is that they can lead to "blowout," in which the region of the outgoing shock at the greatest height starts to accelerate, and attains a much increased velocity in the high layers of very low density.

Numerical computations have shown the development of galactic superbubbles produced by consecutive supernovae (Tomisaka and Ikeuchi 1986; Mac Low and McCray 1988). While some details are still controversial, the general outline seems clear. With realistic assumptions for the ambient density as a function of z, blowout does not occur unless the center of the superbubble is about 100 pc or more from the galactic plane. In any case, an energetic superbubble expands rapidly in z and rises far above the galactic plane. In one example (Mac Low and McCray 1988), 80 supernovae occur uniformly during 10^7 years at $z = 0$; the ambient particle density at $z = 0$ equals 1.0 cm^{-3} (0.19 cm^{-3} of warm gas and 0.81 cm^{-3} of cold; I am indebted to R. H. McCray for providing me with the values actually used in the computations) and for large z varies as $0.19 \times \exp(-z/H)$ cm^{-3}, where $H = 500$ pc. After 10^7 years the hot expanding gas has reached $z \approx 500$ pc, as compared with $r \approx 300$ pc reached at $z = 0$. The temperature of the rising gas is of order 10^6K.

In a second example (Tomisaka and Ikeuchi 1986), with 50 supernovae in 10^7 years, again centered at $z = 0$, the total mass of the ambient gas in a column 1 cm^2 in cross-section is less by an order of magnitude; the interstellar particle density for large z equals $0.035 \times \exp(-z/H)$ cm^{-3}, where $H = 250$ pc. In this case the superbubble expands more rapidly, reaching $z \approx 1000$ pc in 10^7 years, as compared with $r \approx 400$ pc reached at $z = 0$.

The actual ambient gas density at high z is probably between the values assumed in these two examples. Hence one may conclude that some hot gas from superbubbles reaches z values of at least 500 pc and probably substantially more. This conclusion is strengthened by calculations of superbubbles centered initially at z = 100 pc. Such computations were made for each of the two examples cited above, and even for the case with the higher ambient density show hot gas rising to some 1500 pc, where blowout begins. In view of the many uncertainties affecting the dymanical calculations and also the overall energy budget of superbubbles (Heiles 1987a), definitive models are not yet to be expected.

If the superbubble gas rising to kiloparsec heights is mostly confined to the Galaxy, the material will recirculate, falling towards the galactic plane as cooler gas, presumably concentrated in clouds formed through thermal instabilities. Such models have been called galactic fountains; the superbubbles, spewing out hot gas at great altitudes, have also been likened to smoking chimneys (Ikeuchi 1987).

The distribution of such falling clouds in velocity and in space when they reach the base of the corona has been computed with an idealized model (Bregman 1980). First the formation of clouds was analyzed, using two-dimensional hydrodynamic calculations for the uprushing hot gas; condensation into cold clouds was assumed whenever the gas temperature fell below 10^4K as a result of adiabatic expansion and radiative cooling. As boundary conditions, in a thin layer at z = 0 and at all times after t = 0, the gas temperature was taken to be about 10^6K, with a particle density of about 10^{-3}cm^{-3}, varying slowly with distance from the galactic center. Since the initial density at high z was negligibly small, the boundary layer generates outward winds, expanding into a vacuum; such models can be regarded as exploding gaseous discs, which approach a quasi-steady state as a result of cloud condensation.

After the clouds formed, these were assumed to move ballistically, independently of the ambient pressure and density, falling back to z = 0. The final distribution of these clouds in velocity and in galactocentric distance is in general agreement with the observed high-velocity clouds, seen mostly in 21-cm radiation. While this model is too idealized to permit definitive conclusions, the general agreement with available 21-cm data on cloud velocities seems impressive. In a related dynamical investigation (Corbelli and Salpeter 1988) the downward cooling flow from superbubbles is analyzed and its effects in extending or compressing the galactic HI disc are discussed.

One may infer from these discussions that hot supernova remnants ejected into the halo, especially from successive explosions within a young stellar group, provide a substantial source of thermal energy to the halo and can possibly maintain the coronal gas at a high temperature. However, the observational evidence for the existence of such a corona at high z is not yet conclusive (see below). While such a corona provides a natural explanation for a large scale height of the halo gas, implied by the observations (Savage 1987; Jenkins 1987), alternative methods of supporting the gas, including magnetic fields and cosmic rays, must also be considered; these topics are treated in the following section.

In an exploration of alternatives to a hot coronal gas, the high state of ionization found for C and Si atoms at high z may be attributed at least in part to photoionization rather than to collisional ionization (Jenkins 1987; Savage 1987). A detailed computation (Bregman and Harrington 1986) has led to the conclusion that these observed C^{+3} and Si^{+3} ions could be produced by photons from hot O-type stars and from the central stars of planetary nebulae. This analysis assumed that 20 percent of the energetic photons from such stars could escape from the Galaxy, either because the radiating stars are at high z, above the absorbing layer of the galactic disc, or because there are gaps in the distribution of neutral H. This escape fraction is highly uncertain. Looking outward from the Sun one finds (for $\delta > -40°$) a minimum HI column density of about $5 \times 10^{19} cm^{-2}$ (Lockman et al 1986b). For this column density the optical thickness of neutral H for a 48-eV photon, just capable of ionizing C^{+2}, is about 9. If ionizing stellar radiation cannot escape from the galactic disc, photons from active galaxies and quasars can, perhaps, account for the observed C^{+3}, if n_e is less than about 2×10^{-3} cm^{-3} (Fransson and Chevalier 1985).

In any case the N^{+4} ions observed along a few lines of sight at high z are difficult to account for by photoionization. It is primarily for this reason that the presence of hot gas in the halo seems likely.

V. *Structure of the Halo Gas*

The injection of a superbubble into the halo will certainly have dramatic effects on the properties of the local gas, with consequences that are difficult to predict. After such an eruption is over, the local halo will presumably relax to some physical state that may even remain somewhat constant, at least in a statistical sense, until the next great explosion nearby. Thus one may reasonably ask what sort of steady state may be physically possible, subject to what little we know about conditions in the halo. The simplest assumption is that in such a steady state all quantities are functions only of z, the distance from the galactic plane. We discuss here two recent such one-dimensional models of halo gas.

A basic element in any such model is the assumed topography of **B,** the magnetic field. The magnetic pressure p_B, equal to $B^2/8\pi$, makes a major contribution to the pressure in the galactic disc provided that $B_z = 0$. On the other hand, cosmic rays can effectively stream only parallel to the magnetic lines of force, and hence the escape of these energetic particles from the Galaxy is most simply explained if **B** has an appreciable z component. The first model is a hydrostatic equilibrium configuration which is based on the magnetic pressure, with **B** assumed everywhere parallel to the galactic plane; this approach is consistent with the observed direction of **B** in the galactic disc, as determined both from pulsar rotation measures and from the optical polarization of starlight, resulting from alignment of interstellar grains. The second model is a cosmic-ray supported configuration which is based on the escape of cosmic rays, with B_z the only non-vanishing component of **B**; this approach is consistent with the escape of relativistic particles, which seems required by the observed composition of cosmic rays.

The most detailed discussion of models in hydrostatic equilibrium with **B** parallel to the galactic plane includes (Bloemen 1987) all the components of the interstellar gas, cold, warm and hot. A central feature of this discussion is a condition for hydromagnetic stability, particularly against instabilities of Parker type (Parker 1966), in which each line of magnetic force is raised in some regions and depressed in others, with the gas flowing along **B** into the depressed regions. This stability condition (Lachièze-Rey et al. 1980) requires that p_G, the total gas pressure (including the turbulent pressure which is produced by cloud motions), exceed $(dp_{TOT}/dz)/(\gamma d\ln\varrho/dz)$, where p_{TOT} is the sum of all the pressures acting on the gas, – p_G, p_B and the cosmic-ray pressure p_R; as usual, γ is the ratio $\delta\ln p_G/\delta\ln\varrho$ during perturbations of a gaseous element. This criterion, which is derived only for the idealized hydrostatic equilibrium model, may not be a sufficient condition for stability of the model.

One of the significant features included is the large scale height of the warm HI gas, which is set equal to 400 pc in accordance with recent observations (Lockman et al. 1986a). Since the gravitational acceleration increases with z out to at least 1 kpc, the total weight of this component is relatively large, requiring (Cox 1986; Bloemen 1987) a greater value than previously assumed for $p_{TOT}(0)$, the value of p_{TOT} at $z = 0$.

Since $p_G(0)$ and $p_R(0)$ are reasonably well known, this increase in $p_{TOT}(0)$ requires increasing $p_B(0)$, giving an rms B field of about 6μG. While so large a field agrees with a variety of estimates (Bloemen 1987; Heiles 1987b), its consistency with the precise measures of B from pulsar data is a primary requirement. These data give a mean field B_m in the solar neighbourhood which ranges from about 1.6 to 3.5μG depending on the choice of region averaged and on other features of the data analysis (Heiles 1987b). We set B_m equal to 2.5μG. While the rms dispersion of measured fields is less than B_m, the rms random field B_r may considerably exceed B_m if the scale size is substantially less than the pulsar distances.

Such a random field will have dynamical consequences, producing oscillations of the lines of force, together with the attached interstellar clouds. The mean magnetic energy of such oscillations, equal to $B_r^2/8\pi$, should be roughly equal to the mean kinetic energy density $\varrho v^2/2$; if for ϱ we take the smoothed density in cold diffuse clouds (corresponding to $n_H = 0.7 cm^{-3}$) and for v the two-dimensional rms cloud velocity of 8.4 km/s, then $B_r = 4$μG. This value is consistent with the observed dispersion of pulsar measures if the scale size for field variations is in the range from 100 to 200 pc (Thomson and Nelson 1980), depending on the relative importance of fluctuations in n_e and B. The quadratic sum of B_m and B_r then gives about 5μG, not far below the value obtained from $p_{TOT}(0)$. One inference from this discussion is that above about 100 pc, where the density of the cold gas is much reduced and the rms gas velocity is not much greater, B_r^2 may be significantly below its value in the galactic disc.

While the numerical calculations ignore a number of effects (for example, the gradient of $p_B + p_{CR}$ usually assumed to support the cold gas in the galactic disc), the detailed discussion indicates that models which include an increased $p_{TOT}(0)$ together with the hydromagnetic stability criterion promise to fit together various aspects of the interstellar medium at high z. To provide sufficient p_G to satisfy the stability criterion a gas of high pressure but low density apparently suffices. For an assumed halo exponential scale height H_h of 6 kpc the halo temperature needed at high z is about 10^6K, with lower values possible at about a kiloparsec. The value of $n_h(0)$, defined as the atomic particle density of the hot gas extrapolated to $z = 0$, is determined from the added constraint which the observed synchrotron radio emission places on p_{CR} and p_B; it turns out that $n_h(0)$ varies about as H_h^{-2}. If H_h is assumed to equal 6 kpc, a value of 5×10^{-3} provides a "best estimate" for $n_h(0)$.

How such a physical state might be established and maintained raises several physical problems. In particular, how can the temperature of the coronal gas be held in a steady state, with heating balanced by radiative losses? A gas in such radiative equilibrium at $T = 2 \times 10^5$K, for example, is thermally unstable and tends to heat up or cool down. A transient situation, with the gas either heating or cooling through these high temperatures, can provide a more likely explanation for the highly ionized atoms observed in the halo. For example, with plausible assumptions for the flow of initially hot gas through the galactic corona, a calculation of effective recombination rates gives N^{+4} column densities vertically through the corona (Edgar and Chevalier 1986) in good agreement with observed values.

Another problem is how do cosmic rays escape from the Galaxy if **B** is parallel to the galactic plane. Streaming along spiral arms to large distances and then spraying out is conceptually possible but may require diffusion over too great a distance. Diffusion of cosmic-ray particles transverse to **B** can be produced by small-scale magnetic turbulence (Cesarsky 1980). However, this turbulence is probably weak at high z, where $\varrho v^2/2$ is small compared to $B^2/8\pi$; as a result, any turbulent diffusion is likely to be confined to the galactic disc.

We turn to the second theoretical model, in which the escape of cosmic rays is a basic element of the picture, with **B** taken parallel to z. This model (Hartquist and Morfill 1986) takes into account the detailed physical processes involved when energetic ions stream outwards along lines of magnetic force, exciting Alfvén waves. These waves scatter the ions, which drift along **B** by a combination of diffusion and convection. The model calculations assume hydrostatic equilibrium, with dp_R/dz balancing $-g\varrho$, and two energy equations, one for p_R (which includes convection, diffusion and conversion of cosmic-ray energy into wave energy) and one for the wave energy density. In this latter equation the wave damping normally provided by collisions of ions with neutral atoms is ineffective because of low density and high ionization. Instead damping is attributed to interactions between upward travelling and downward travelling Alfvén waves. It is not clear what physical process will produce downward waves of the necessary strength. The stability of such an equilibrium model is also unclear.

Apart from these problems, the model presents a full and self-consistent picture for the particular magnetic topography assumed. Since $p_R(0)$ is only a fraction of $p_{TOT}(0)$, this model applies only to the halo gas, at values of z where the particle density does not exceed about $10^{-3}cm^{-3}$. The energy input from the diffusing cosmic rays into Alfvén waves and into thermal energy via wave-wave damping provides a heat source for the gas. Depending on the parameters assumed, such models would be consistent either with equilibrium kinetic temperatures somewhat below 10^5K and photon ionization as a source of C^{+3}, or with higher temperatures and collisional ionization. Since thermal instability is probable, a transient solution, with the temperature of a fluid element rising or falling, may again be needed to explain the observed CIV and SiIV line strengths.

A time-dependent solution would also be helpful in reconciling the magnetic field perpendicular to the galactic plane, which is suggested by the cosmic ray data, with the parallel field, which is observed in the solar neighbourhood and which seems needed to support the cold and warm components of the interstellar gas. Such a situation is provided, of course, by the occasional eruption of superbubbles. If some fraction of the rising hot gas escapes from the Galaxy, the magnetic lines of force will be stretched far out, with some reconnection perhaps recurring. As pointed out by several researchers (Cesarsky 1980), such a time-dependent sequence might permit the intermittent escape of cosmic-ray particles.

Evidently progress in understanding the hot gas takes place through a succession of idealized models, a familiar state of affairs in much of modern astrophysics. Many of these models have been proposed, some have been developed in considerable detail, and a few have been described here. For the most part they are not very realistic, but they may provide a tentative picture of the hot gas, how it arises and what range of effects it produces. In any case these models are helpful in suggesting new theoretical questions and new observational programs.

Acknowledgements

The preparation of this paper has been greatly aided by comments and suggestions from several astronomers, including J.B.G.M. Bloemen, J.N. Bregman, B.T. Draine, L.L. Cowie, T.W. Hartquist, E.B. Jenkins, R. McCray, C.F. McKee, J.P. Ostriker, B.D. Savage and M. Wardle. Several very detailed discussions with D.P. Cox have been particularly helpful.

Appendix: Oblique Isothermal Shocks

We summarize some relevant information (Landau et al. 1984; Kantrowitz and Petschek 1966) on a plane shock travelling at some oblique angle across a magnetic field. The shock front is stationary and in the preferred reference frame **v,** the local fluid velocity, is parallel to **B**. We denote by θ the angle between **B** and **n**, the unit vector normal to the shock front. Subscripts 1 and 2 denote preshock and postshock quantities; a subscript n denotes a component normal to the shock front.

It is straightforward to write all the shock jump conditions and to solve these for the various postshock quantities. We assume an isothermal shock, with $T_2 = T_1$. The resultant cubic equation for ϱ_2/ϱ_1 gives values which in Figure 1 are plotted against θ_1, the preshock angle between **n** and $\mathbf{B}_1$. The values of V_A (for $B = B_1$) and of C_s taken in the computations are those given above for the warm neutral medium, with $V_s = |v_{1n}|$ set equal to 40 km/s. The values of θ_2 corresponding to various points are indicated in the Figure. The analysis is not strictly applicable for $\theta_1 = 90°$, since for this transverse shock there is no reference frame in which $\mathbf{v}_1$ is parallel to $\mathbf{B}_1$. However, the limiting results as θ_1 approaches 90° agree with those found separately for **n** perpendicular to **B.**

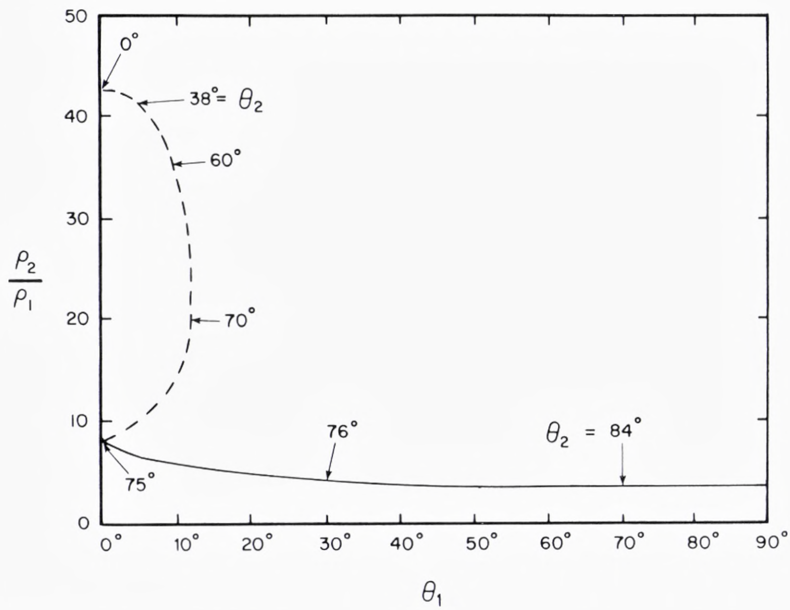

Figure 1. Compression factor across an oblique hydromagnetic shock. The density ratio ϱ_2/ϱ_1, determined from the shock jump conditions, is plotted against θ_1, the preshock angle between the magnetic field and the shock-front normal. Values of the postshock θ_2 are indicated at various points. The dashed curve represents extraneous solutions, not realizable in a single shock.

Figure 1 shows that large compressions are possible only for a narrow range of θ_1 ($\theta_1 \leqslant 12.8°$). The two solutions which at $\theta_1 = 0$ have low compression ($\varrho_2/\varrho_1 = 8.2$) are "switch-on shocks," with $\mathbf{B}_1$ parallel to $\mathbf{n}$, while $\mathbf{B}_2$ is steeply inclined, at an angle θ_2 which here equals 75°. For these shocks the azimuthal angle of $\mathbf{B}_2$ cannot be determined from the shock jump equations but depends on the boundary conditions.

Two additional features must be considered in connection with these results. In the first place, only the lower curve in Figure 1 is physically realizable. The shocks represented by points on the upper dashed line, including all those with $\varrho_2/\varrho_1 > V_s^2/V_A^2 = 8.2$, would, if formed, immediately disintegrate; this disruption is much more rapid than an exponential growth of small perturbations and cannot be regarded as an instability. Such shocks are termed "non-evolutionary" (Landau et al. 1984) or "extraneous" (Kantrowitz and Petschek 1966). The fact that θ_2 is in the opposite direction from θ_1 for these higher-compression solutions of the jump conditions is characteristic of many such extraneous shocks.

In the second place, the end result which would be produced by an extraneous

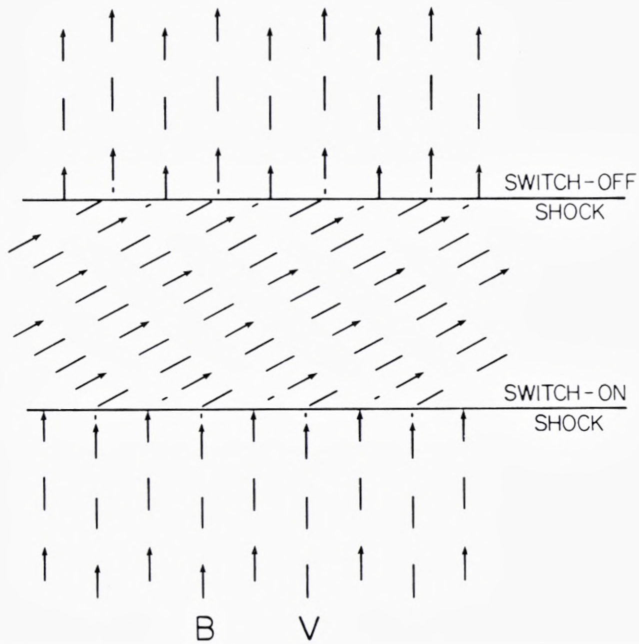

Figure 2. Direction of magnetic field and velocity in a double shock. The leading switch-on shock produces a transverse magnetic field component, which is eliminated in the subsequent switch-off shock. The distance between the two shocks is arbitrary but constant. The overall compression ratio $\varrho_3/\varrho_1 = v_1/v_3$ is that produced by a single shock with the same v_1 but with $B = 0$.

shock, if this could exist, can under some conditions be produced by two successive shocks. This possibility has been pointed out (Kantrowitz and Petschek 1966) in particular for the high-compression shock travelling parallel to $\mathbf{B}_1$; i.e., with $\theta_1 = 0$. If V_A exceeds C_s and if, for an isothermal shock, V_s has any value greater than V_A, this shock is extraneous and cannot exist. However, the switch-on shock which appears instead can be followed by a switch-off shock, which travels at a fixed separation from the first shock and which bends the magnetic field back to its original direction, normal to the two shock fronts. The resultant geometry is indicated in Figure 2. The second shock produces an additional compression by a factor 5.2 for the parameters assumed here, giving an overall compression factor of 43, the value found above for $B = 0$. What is possible for $\theta_1 = 0$ should also be possible for some other values of θ_1, especially if the two shocks are allowed to separate gradually and if the final direction of $\mathbf{B}$ is allowed to differ somewhat from the original direction. Further analysis is needed to indicate under what conditions a supernova remnant in the snowplow stage can strongly compress the warm interstellar medium.

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