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A MECHANISM FOR STRONG SHOCK ELECTRON HEATING IN SUPERNOVA REMNANTS

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ABSTRACT

It is shown that collisionless shock waves propagating away from a supernova may be directly responsible for the 10 keV X-ray emission seen in supernova remnants. A sequence of plasma instabilities (Buneman and ion acoustic) between the reflected and/or transmitted ions and the background electrons at the foot of the shock front can give rise to rapid anomalous heating of electrons. Hybrid simulations of a perpendicular collisionless shock are presented to demonstrate that this heating can arise within a self-consistently computed shock structure.

Subject headings: nebulae: supernova remnants — plasmas — shock waves — stars: supernovae

I. INTRODUCTION

While it has been well known for many years that supernova remnants are a strong source of X-rays (e.g., Danziger and Gorenstein 1983), the electron heating process required for this emission has been an unsolved question. It is commonly assumed (e.g., Chevalier 1981; Low 1982) that when the magnetic field and the plasma in a stellar envelope move out into the local interstellar medium during a supernova, a piston is formed which subsequently gives rise to a pair of forward and reverse shock waves (McKee 1974). Heating at these shocks is then considered to be responsible for producing the X-ray emission in the supernova remnants. In this *Letter*, we assume that the shocks are quasi-perpendicular, so that the shock will propagate normal to the ambient magnetic field (Low 1982). The difficulty in understanding the heating mechanism arises because the shocks will be collisionless and so knowledge of the microscale plasma processes in such shocks is essential (McKee and Hollenbach 1980).

Until recently our understanding of moderate ($M_A > 2$) and strong ($M_A > 10$), quasi-perpendicular collisionless shocks was poor. (Here M_A is the Alfvén Mach number defined as $M_A = U_1/V_A$, where U_1 is the upstream velocity in the shock frame and V_A is the upstream Alfvén speed.) Recent studies (Leroy *et al.* 1982) showed that the reflection of upstream ions plays a dominant role in the structure of quasi-perpendicular shocks. Twenty percent or so of the incident ions were reflected, forming a cross-field ion beam in the immediate upstream region; these ions subsequently gyrated in the magnetic field and moved downstream. No means of heating electrons strongly at such shocks was found (Wu *et al.* 1984). However, Papadopoulos (1988, hereafter P88) noted that above certain Mach numbers, strong electron heating could take place in the shock foot due to plasma instabilities involving the reflected ion beams, and the suggestion was made that the results could be of relevance to SNR electron heating. The purpose of this *Letter* is to present a more self-consistent calculation than that of P88.

II. SEQUENCE OF ION BEAM INSTABILITIES

The reflected ion beam mentioned above moves counter to the ambient incoming flow with a relative velocity (U_R) of between 1.5 and $2U_1$. The requirement of current neutrality produces a shift in the electron velocity moment below U_1 and the incoming electrons and ions are also separated in velocity space because the cross-shock electric potential affects the ions and electrons differently (Leroy *et al.* 1982). The ion beam is susceptible to the Buneman instability (Buneman 1958) between the ambient electrons and the reflected ions if

$$2U_1 > V_e\,, \quad \text{or} \quad M_A > 20\sqrt{\beta_e}\,, \tag{1}$$

where $\beta_e = 8\pi n T_e/B^2$ is the electron plasma beta, $V_e = (2kT_e/m_e)^{1/2}$ is the electron thermal speed, and we assume that $U_R \approx 2U_1$ so that condition (1) is satisfied easily for a SNR shock with β_e of order 1 and $M_A \approx 100$. The Buneman instability gives rise to very rapid electron heating ($10\omega_e{}^{-1}$ s, where ω_e is the electron plasma frequency: Davidson *et al.* 1970) which saturates at

$$\frac{T_e}{T_{e0}} \approx 1 + 2\left(\frac{m_e}{m_i}\right)\frac{m_i U_1{}^2}{T_{e0}} \tag{2}$$

(P88), where T_{e0} is the upstream temperature. This rapid increase in T_e now permits the possibility of an ion acoustic instability between either the incoming or reflected ions and the ambient electrons. The condition for the ion acoustic instability is

$$\frac{V_D}{V_e} > \sqrt{\alpha}\sqrt{\frac{m_e}{m_i}} + (1-\alpha)^{3/2}\left(\frac{T_e}{T_i}\right)^{3/2}\exp\left(-\frac{\alpha T_e}{2T_i}\right), \tag{3}$$

where α is the density ratio of the reflected (or transmitted) ion beam to the electron density and V_D is the relative drift between the ions and electrons (Ichimaru 1973). P88 anticipated that the ion acoustic instability would operate between only the transmitted ions and the electrons whereas our simulations presented in the following section suggest that instabilities between both reflected and incident ions beams could be significant.

The ion acoustic instability gives rise to electron heating at a rate given by the Sagdeev anomalous collision frequency (Sagdeev and Galeev 1969):

$$\nu^* \approx \frac{\omega_e}{32\pi}\,. \tag{4}$$

Using this pair of results, P88 demonstrated that 12% of the upstream kinetic energy could be converted into electron thermal energy provided the ratio of the ion plasma frequency (ω_i) to the ion gyrofrequency (Ω_i) based on the upstream quantities exceeds 100, a condition trivially satisfied in almost every space plasma. This condition arises from the necessity that the instability length scale be less than the shock foot length. For a shock velocity of 2000–10,000 km s^{-1} (McKee and Hollenbach 1980), temperatures of 2–50 keV are obtained from this simple model.

However, the ideas of P88 ignored the self-consistent shock ion dynamics which are known to be important (Leroy *et al.* 1982). It is therefore crucial to see whether a simulated high Mach number shock can give rise to the sequence of instabilities discussed above. We address that issue next.

III. ONE-DIMENSIONAL HYBRID SIMULATIONS OF STRONG SHOCKS

Over the past five years or so, one-dimensional (along the shock normal) hybrid simulations have proved to be a powerful yet economical way of studying the structure of collisionless shock waves (Leroy *et al.* 1982). Hybrid simulations treat ions as individual particles to be followed in self-consistently computed electromagnetic fields, whereas the electrons are treated as a massless, resistive fluid. The resistivity is treated by the use of phenomenological transport coefficients, computed using our knowledge of the nonlinear evolution of plasma microinstabilities. A resistive drag is applied to the ions such as to reduce the relative electron-ion streaming of each particle. Hybrid codes also incorporate an electron energy equation with an ohmic heating term such that energy is conserved between electrons and ions during the anomalous collision process.

In this *Letter*, we present two simulations, one in which the Buneman and ion acoustic instabilities are assumed not to occur, and one in which they do. We concentrate on a $M_A = 50$ shock with the upstream ion and electron plasma betas both set to 0.5. Lengths are normalized to the upstream ion inertial length (c/ω_i) and times are normalized to ω_i^{-1}. The computational box is $200c/\omega_i$ long with 800 cells, so that a time step of $8\omega_i^{-1}$ is needed. A total of 40,000 particles are used. In the first run (A), the anomalous resistivity η^* is small so that minimal electron heating results. In the second simulation (B), we use the following resistivity prescription:

$$\eta^* = 0.1\omega_i^{-1},\ V_D > V_e\ , \tag{5}$$
$$\eta^* = 10^{-3}\omega_i^{-1},\ V_D > V_{IA}\ ,$$

where V_D is the relative electron-ion drift and V_{IA} is the threshold velocity for triggering the ion acoustic instability (see eq. [4]). The value of η^* for the Buneman instability is not especially important since the system evolves very rapidly to marginal stability. We also choose the ion acoustic resistivity to be somewhat smaller than given in equation (5) so that, given the uncertainties in deriving equation (5), any simulated heating rate will be on the conservative side.

Within the realms of current computer resources, it is difficult to simulate a SNR shock with M_A between 10^2 and 10^3. Instead, we simulate a shock at a lower M_A where electron heating can still take place and extrapolate our results to higher M_A. Justification of this extrapolation can also be found in P88, where it is shown that the local magnetosonic Mach number falls to a value of around 40 at the edge of the foot where the Buneman instability occurs, independently of the original value of M_A.

To carry out our studies, we first simulate a $M_A = 50$ shock with $\eta^* = 10^{-5}\omega_i^{-1}$ for several thousand time steps to allow a quasi-steady shock structure to be set up. This resistivity is typical of that produced by, for example, lower hybrid turbulence (Papadopoulos 1985). Then, in run B we switch on the tests for triggering the Buneman and ion acoustic instabilities and study their effect on the shock structure. Figure 1 shows a comparison of the two simulations. Figures 1*a*, 1*b*, and 1*c* show (v_x, x), (v_y, x) phase space and the electron temperature, respectively, for run A, whereas Figures 1*d*, 1*e*, and 1*f* show the same quantities for run B. Figures 1*a*, 1*b*, and 1*c* are shown at $5.2\Omega_i^{-1}$, whereas Figures 1*d*, 1*e*, and 1*f* are at $9\Omega_i^{-1}$, where Ω_i is based upon the upstream magnetic field strength. The important feature to note is the behavior of the electon temperature. Figure 1*f* shows that in run B, the upstream electrons are heated by a factor of around 1000, whereas in run A, the heating ratio is of order 10–30. For the values of M_A and β_e used, the electron heating corresponds to a transfer of very approximately 18% of the initial upstream kinetic energy into electron thermal energy, a number close to that proposed in an ad hoc manner in P88.

The ion shock structures are fairly similar in each case. Ions are still being reflected [especially evident in (v_y, x)] phase space, there is still the characteristic magnetic field overshoot at the shock ramp (not shown: Leroy *et al.* 1982), and the anisotropic downstream ion distribution still arises so that inclusion of the Buneman/ion acoustic resistivity does not change the basic shock structure. The ion reflection is unsteady in this case, as is to be expected from strong, resistive shocks (Quest 1986). The shock also remains at roughly the same location in the computational box, indicating that pressure balance still holds.

The detailed physics of the resistive shock is as follows. As the ions reflected at the shock front move outward, they increase their net drift since they now experience a motional electric field in the *y*-direction. When this drift exeeds V_e, the Buneman resistivity is switched on and the electron temperature increases to reach marginal stability. This is the reason for the initial rapid heating in Figure 1*f*. The ion acoustic instability between either the reflected and transmitted ions and the beam can now take place giving further electron heating. However, the ion acoustic instability is curtailed rapidly as the system moves to marginal stability due to the resistive drag on the ions. The motional electric fields acting on each ion beam due to their net drift with respect to the total ion velocity moment continue to drive V_D so that the heating process is thus a classic case of the marginal stability approach to plasma microinstabilities (Lampe, Manheimer, and Papadopoulos 1975). As seen from Figure 1*f*, the electrons are continually heated in the shock foot by the ion acoustic instability, and by the time they reach the ramp (the site of the electric potential maximum), they have been heated by a factor of around 1000.

Having successfully demonstrated that the mechanism proposed by P88 is applicable for a $M_A = 50$ shock, we are faced with the question of how to extrapolate the results to higher Mach numbers. It is relatively easy to show that a simple scaling of the amount of Buneman and ion acoustic heating with M_A exists. The simulations of Davidson *et al.* (1970) suggest that the electron heating scales with U_1^2 as shown in equation (2). For ion acoustic heating, the spatial heating rate is given by

$$U_1 \frac{dT_e}{dx} = \nu^* V_D \,|\, V_D - V_e \,|\, m_e \tag{6}$$

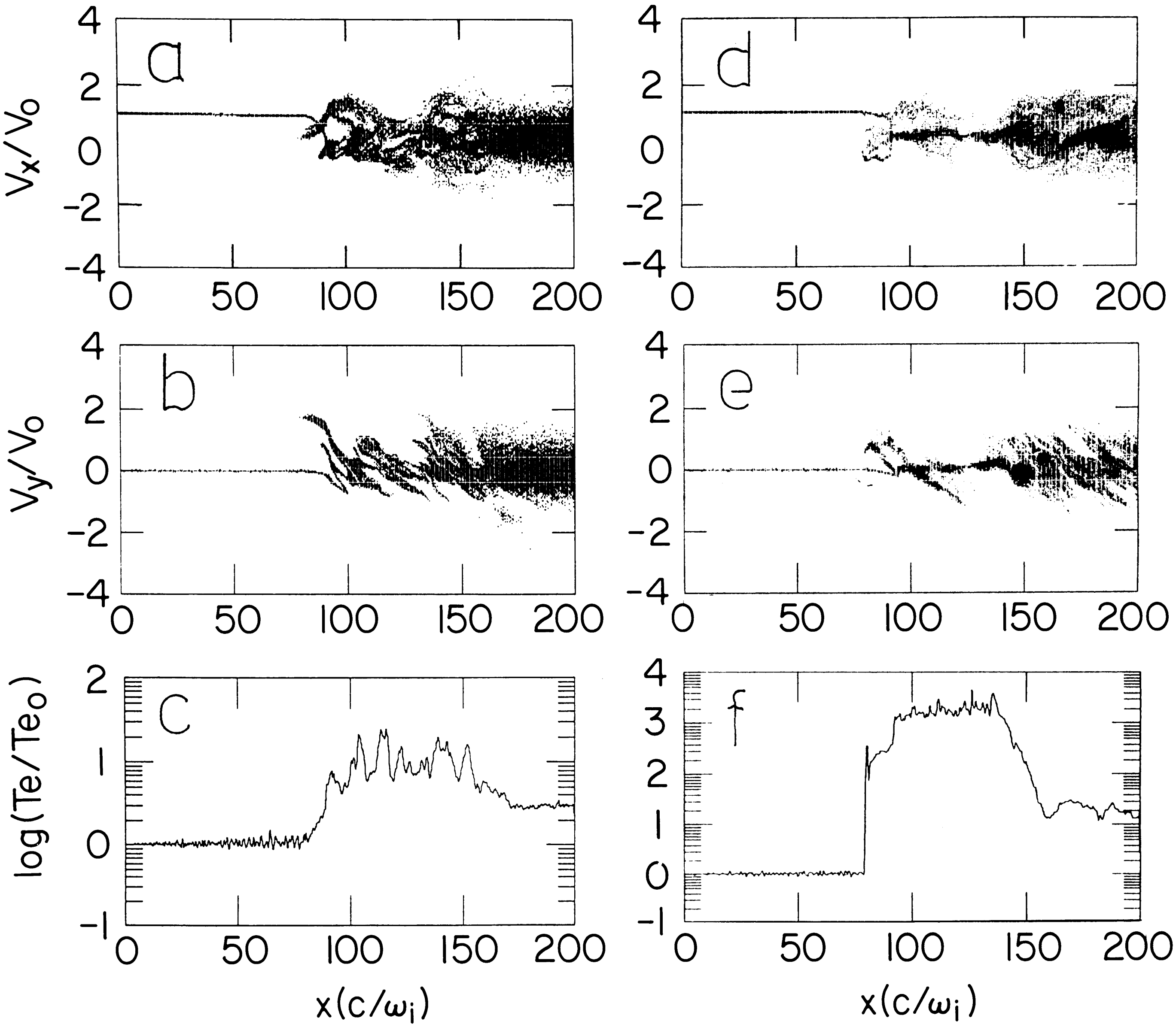

FIG. 1.—(v_x, x) phase space (1*a* and 1*d*), (v_y, x) phase space (1*b* and 1*e*), and the electron temperature (1*c* and 1*f*) for run A (1*a*, 1*b*, and 1*c*) and run B (1*d*, 1*e*, and 1*f*) as a function of distance in units of ion inertial lengths (c/ω_i). Run A has a small uniform resistivity, and run B has the resisitivity prescription given in eq. (5). Figs. 1*a*–1*c* are shown at a time $t = 5.2\Omega_i^{-1}$, and Figs. 1*d*–1*f* are shown $4\Omega_i^{-1}$ later. In the phase space plots, the velocity is normalized to the upstream incoming ambient flow speed. In Figs. 1*c* and 1*f*, the temperature is normalized to the ambient upstream temperature: note that Figs. 1*c* and 1*f* have different vertical scales.

(P88, eq. [12]). The shock thickness dx scales as U_1/Ω_i, so that using equation (4), it is obvious that the amount of electron heating scales with U_1^2 and the ratio of the final to initial electron temperature scales with M_A^2/β_{e0}. Applying this scaling to our results, we find that a $M_A = 500$ shock with $\beta_{e0} = 0.5$ gives an electron temperature ratio of 10^5 across the shock. For any reasonable upstream conditions, such a temperature jump should produce keV X-ray–emitting electrons.

IV. CONCLUSIONS

We have examined the question of how electrons are heated at collisionless shocks in supernova remnants. A model discussed earlier by Papadopoulos (1988) has been put on much firmer ground; by means of simple scaling arguments it was shown that it is relatively easy for electrons at SNR shocks to gain a great deal of energy (a factor of 10^3 for a $M_A = 50$ shock). The heating arises because at high Mach numbers the reflected ion beam is unstable to the Buneman instability. This gives rise to electron heating subsequently triggering the ion acoustic instability between the transmitted and/or reflected ions and the electrons. This results in further electron heating. Hybrid simulations of such a shock were performed, and it was shown that this mechanism could indeed arise in a self-consistently computed shock structure without significant modification to the shock structure. It was also noted that the amount of heating should scale as M_A^2, so that for a SNR shock with $M_A = 10^3$, the electron temperature should increase by a factor of 10^5.

We note that these results do not give information about details of electron spectra seen at SNRs since the simulations use fluid electrons. To predict such results at SNR shocks would require use of a full particle code with all the attendant complexities and assumptions involved. The model is hence not predictive in that important issue.

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