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SIMULATIONS OF COLLISIONLESS SHOCKS

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ABSTRACT

A problem of critical importance to space and astrophysics is the existence and properties of high-Mach-number (HMN) shocks. In this letter we present the results of simulations of perpendicular shocks with Alfvén Mach number 22. We show that the shock structure is a sensitive function of resistivity, becoming turbulent when the resistivity is too low. We discuss the problem of electron heating, and the extension of our results to higher Mach numbers.

INTRODUCTION

A shock is a nonlinear supersonic compressive wave across which ordered flow energy is converted into disordered thermal energy and magnetic energy. In a collisionless shock, the mean free path of ion-electron and electron-electron collisions is very long, so dissipation is provided by anamolous wave-particle scattering and ion reflection. The best known example of a collisionless shock is the earth's bow shock, formed when the supersonic, super-Alfénic solar wind runs into the earth's magnetosphere.[1] Other astrophysical examples are the low Mach number shocks produced when an early-type star first turns on,[2] and the high Mach number shocks driven by supernova remnants.[3]

In recent years, the ISEE satellite program has stimulated an intense theoretical and observational program in collisionless shock physics. By combining theory, simulations, and observations, a good understanding of the Earth's bow shock has emerged but there remain, of course, many unanswered questions. One of the most important of these is the nature of high-Mach-number (HMN) shocks. At present, the fastest shock observed within the solar system was at Jupiter, and had a magnetosonic (fast mode) Mach number of 12.[4] The lack of higher Mach numbers have led to the suggestions that (1) there is a "2nd critical" Mach number, above which the shock physics changes radically, or (2) that the solar wind, which generates the planetary bow shocks, rarely exceeds magnetosonic Mach numbers above 10.

The Jovian bow shock observed by Russell et al.[4] is displayed in Figure 1. The magnetic field (in units of $\gamma = 10^{-5}$ Gauss) is plotted as a function of time (hours and minutes). Upstream of the shock (right), the magnetic field strength is less than one γ, the plasma β (ratio of plasma to magnetic pressure) is 2.9, and the shock-normal angle θ_{Bn} is 78°. The magnetic field at the shock is observed to rise (overshoot) well above its mean downstream value.

Particle observations,[5] theory,[6] and simulations[7] have shown that this overshoot is a consequence of a reflected ion beam which is

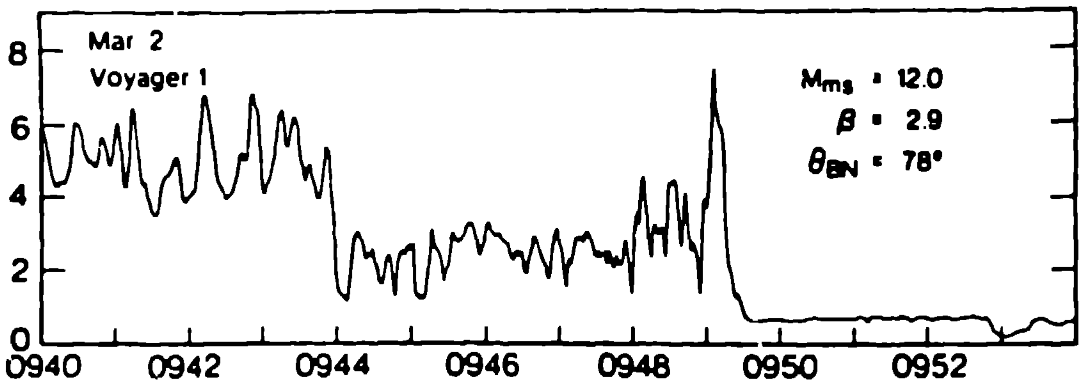

Fig 1. Magnetic field (10^{-5} G) vs. time (hours and minutes) for an outbound Jovian bow shock crossing (from Ref. 4).

the primary source of dissipation for collisionless shocks at these Mach numbers.

The structure of HMN shocks is also a topic of interest to astrophysics. Observations of the x-ray emissions from young super nova remnants show the existance of a well defined spherical shell of hot electrons.[8] This shell is believed to be the downstream wake of an outward propagating shock, with Alfvén Mach number as high as 1000.[9] If true, then there is no apparent upper limit to the speed of an astrophysical shock, and additionally, such shocks strongly heat the electrons. Predicting how such heating occurs, and what the relative ion to electron temperature ratio is downstream of the shock, is an unresolved issue of plasma astrophysics.

In this paper we present numerical simulations of HMN shocks. We review the results of Quest,[10] who showed that in the absence of resistivity HMN shocks are unsteady, and compare these results with shock simulations including electron dissipation.

SIMULATIONS

The numerical model we will use is a one dimensional (in x) electromagnetic hybrid code which follows the individual particle orbits of the ions and treats the electrons as a resistive fluid. Because the electrons are massless, plasma oscillations are suppressed, and the neglect of the displacement current in Ampere's Law eliminates light waves. As a consequence, large spatial and temporal steps are possible, which allows following the evolution of the shock for several ion gyroperiods. This model has been described in detail previously.[11] Plasma is continuously injected from the left-hand boundary (x=0) and moves in the positive x-direction (see Fig. 2a). When the plasma hits the right-hand boundary (x=L) it is reflected and a shock is launched (see Fig.

2b). As the shock continues to propagate through the box, its separation distance from the wall becomes greater than the downstream ion gyroradius, effectively separating piston and shock heated plasma (Fig. 2c). The simulation run is continued until the

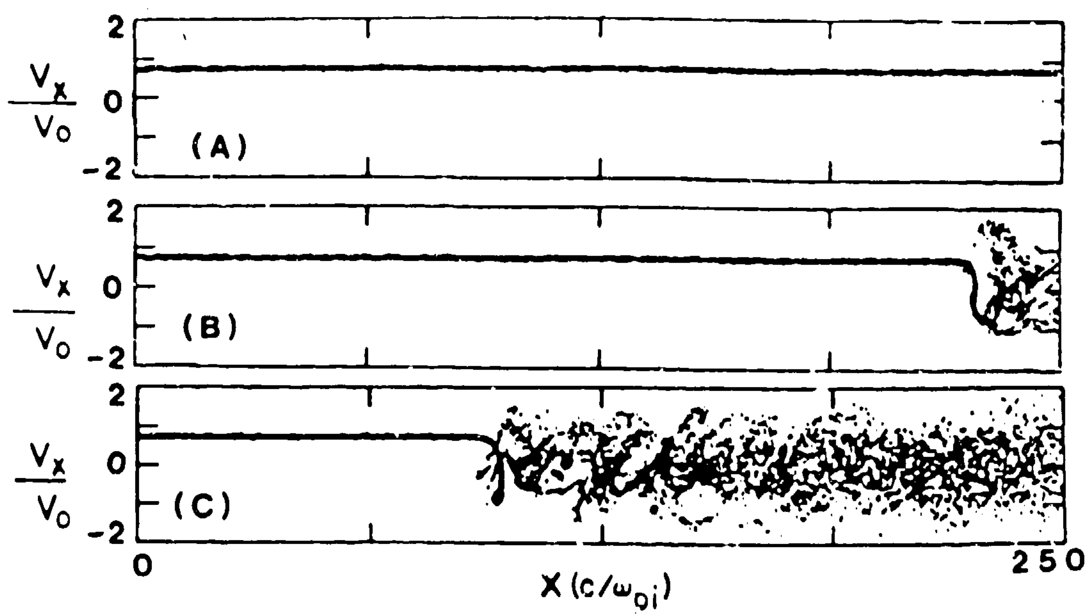

Fig 2. V_x - x phase space initially (A), and at later times (B-C) in the simulation run.

shock generated magnetic turbulence decays before reaching the right boundary, insuring that piston effects do not govern the shock.

The shock simulation we will examine[10] in perpendicular ($\underline{B} = B_z \hat{z}$), propagates at an Alfvén Mach number of 22, and has an upstream electron and ion β of 0.5 respectively, where β is the ratio of thermal pressure to magnetic pressure. The upstream ratio of ion plasma frequency (ω_{pi}/ω_{ci}) is 2×10^4 and the resistivity η is between $1.5 - 12 \times 10^{-4} \omega_{pi}^{-1}$. A check of the average downstream values show that $n_d/n_u \cong 3$, $B_d/B_u \cong 3$ and $T_{id}/(1/2 M_i V_o^2) \cong 0.5$ where u denotes upstream, d downstream, and V_o the shock speed. These results are consistent with the 2-dimensional Rankine-Hugoniot relations, with most of the shock energy being deposited in the ions.

Because the shock speed is well above the critical Alfvén Mach number, (approximately 3 for these upstream conditions) dissipation by electron heating alone is insufficient to stop shock steepening, and ion reflection results. Simulations of resistive perpendicular shocks with $3 < M_A < 10$ and $\beta = 1$ have shown that, in this Mach range, the shock structure is reasonable steady.[7] A fraction of the incoming ions is reflected by a potential barrier and magnetic ramp at the shock front. These ions gyrate in front of the shock, gaining energy from the $\underline{E} \times \underline{B}$ electric field, and are carried downstream. After thermalizing with the directly transmitted ions, a heated downstream population results. As the Mach number is increased, the reflection process continues to be the dominant source of dissipation, but can be highly oscillatory, depending on the magniture of the resistivity.

If the resistive diffusion length (proportional to the resistivity and inversely proportional to the upstream flow speed) is set much smaller than a spatial cell size, then it is not possible to stop shock ramp steepening by resistive dissipation. Under these conditions we find that the shock exhibits a periodicity (1/3 of an upstream gyroperiod) in which the shock steepens, breaks and overturns, and then steepens again. During this cycle, roughly all of the ions are transmitted through the shock, followed by a brief period of total reflection.

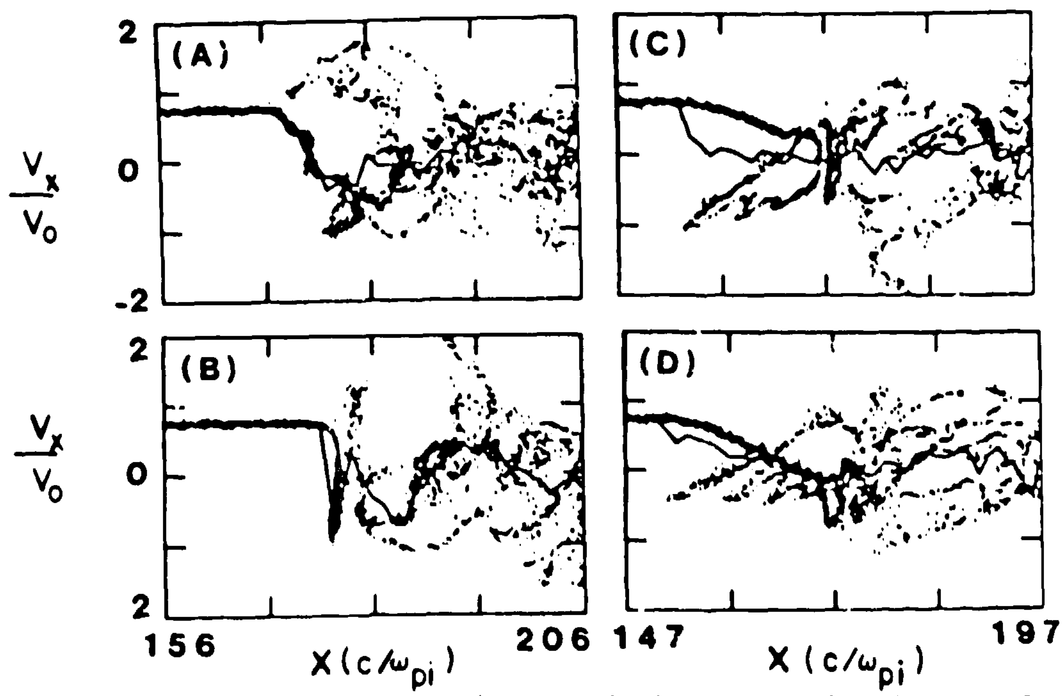

Fig 3. V_x - x phase space at 4 times during a wave breaking cycle. Resistivity η is set to 0 for this run.

In Fig. 3a we show a close-up of the V_x - x phase space after the shock has advanced roughly 1/3 of the way into the simulation box. The solid line is the average value of V_x. At this time the shock transition consists of a smooth ramp, with an energetic ion population downstream. The energetic ions are the result of the previous reflection cycle. In Fig. 3b we see the ramp has steepened to its minimum thickness and is starting to reflect the incoming ions. In Fig. 3c the ions have been reflected, travel upstream some distance, then turn around and head downstream. The ramp thickness is now very broad and in completing the cycle will steepen because of an E field which accelerates particles in the negative x direction. This returns us to 3a.

The behavior of the shock is very turbulent, and reminiscent of earlier shock studies in which a periodic formation and destruction of the shock was observed.[12] An important difference in our results, however, is that the minimum shock thickness (just before breaking) is numerically determined by the cell size. There is no resistive ion scattering or electron inertia in the code, so wave steepening cannot be balanced by dispersion or ion diffusion. Thus, our results demonstrate the process by which the shock will heat ions downstream in the absence of anomalous scattering (by wave breaking), but we are unable to predict details of the structure such as the magnitude of the turbulent magnetic overshoot. Such specifics will require 2-dimensional particle codes, which include both cross-field instabilities and electron inertia.

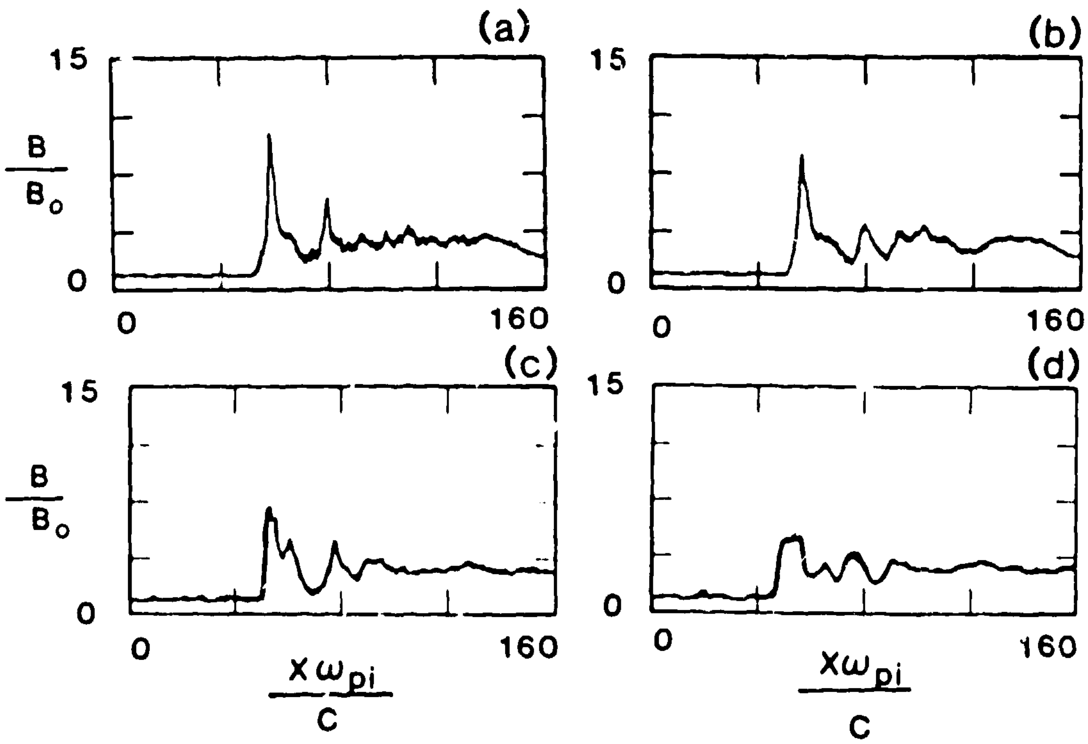

Fig 4. B - x plots for four different values of η = (A) 1.5×10^{-4}, (B) 3×10^{-4}, (C) 6×10^{-4}, and (D) 1.2×10^{-3}.

As the resistive diffusion length is increased to a magnitude greater than a cell size, the temporal behavior of the shock becomes much quieter. In figure 4a and 4b we show the magnetic field profile for runs with the resistivity set at 1.5 and 3×10^{-4} ω_{pi}^{-1}. While the magnetic overshoots are large (~10 times the upstream value in 4a), they are quite stationary, varying by less than 6% during the course of the run. These results indicate that there exists a range of finite resistivity over which stationary solutions, similar to those examined by Leroy et al., at lower Mach numbers. In fact, a simple set of jump conditions, with the fraction of reflected ions set as a free parameter, yield predictions quite similar to the above two runs. As the resistivity is increased further, the average magnetic overshoot and the fraction of reflected ions continue to decrease (see Fig. 4c, 4d), but the magnitude of RMS deviations increase strongly (13% for the magnetic overshoot). The problem is that increasing the resistivity decreases the fraction of the reflected ions. In order to maintain a steady state the additional dissipation must come from the heated electrons. There is a limit, however, to the amount of total electron heating (see for example, arguments in Leroy et al.[7]). When the mean number of reflected ions becomes too small, the shock structure oscillates.

CONCLUSIONS

Given the various classes of solutions for high-Mach number shocks, which will actually apply? The answer to that question will depend on the efficiency of wave-particle instabilities heating the shocked plasmas. Observationally, most shocks with plasma $\beta \approx 1$ appear quite stationary, even at the higher Mach numbers. By contrast, shocks with $\beta \gg 1$ are very unsteady.[13] It is tempting to speculate that in the former case ($\beta \approx 1$), the lower-hybrid drift instability acts to smooth the shock structure, while in the latter ($\beta \gg 1$), the mode is stabilized, resulting in cyclic shock steepening and wave breaking. An observational study is currently in progress to resolve these issues.[14] As the Mach number continues to increase, so does the amount of resistivity required to maintain a specified shock ramp thickness. It seems likely that for sufficiently fast shocks the resistivity will not keep up, and wave breaking will result. Another important point is that because of the one-dimensionality of the simulation and the suppression of short-wavelength oscillations the thermalization of the downstream shocked plasma is likely to be quite different from what has been presented here. Since the gyrational energy is perpendicular to the B field, anisotropy driven ion instabilities will generate large fluctuating fields. Current driven modes and beam modes could be destabilized, peaking at short wavelengths and driving the electrons resistive. For extremely high Mach numbers, the large fluctuating fields generated by the ion anisotropy could be absorbed by the resistive electrons, resulting in strong electron heating. Clearly, a great deal of work is required to clarify these and other issues raised by these simulations.

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