tom of Fig. 3(a). These satellites are shifted by approximately 70 to 100 kHz from the oscillator frequency. Near the region of maximum refractive index the amplitude of the satellites grows and becomes nearly equal to that of the central peak. Figure 3(b) shows simultaneous values of the index of refraction, the wave amplitude, and the satellite amplitude as a function of radial position.

There are oscillations in this plasma which modulate the density by about 10% and which appear to be drift waves. The dependence of amplitude on radius and magnetic field and the dependence of frequency on magnetic field and ion mass indicate that these low-frequency oscillations are drift waves. We have observed that the frequency of these oscillations is always exactly equal to the frequency separation between the satellites and central peak for the driven high-frequency waves. These data indicate a strong nonlinear interaction between the driven waves propagating near the lower hybrid resonance and the drift waves which are always present in the plasma column. Furthermore, as the amplitude of the high-frequency waves increases, the low-frequency oscillation amplitude is affected and usually decreased. It is possible that this effect may be caused by a change in the plasma equilibrium induced by the driven oscillations rather than dynamic stabilization. However, the relatively intense satellites indicate a strong interaction between the low-frequency and high-frequency waves that could be of considerable interest to those workers concerned with dynamic or feedback stabilization.

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Ion Heating in High–Mach-Number, Oblique, Collisionless Shock Waves

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A mechanism for strong ion heating in oblique, collisionless shock waves is found in numerical-simulation experiments and identified with an essentially electrostatic two-ion-beam instability excited nonlinearly by the potential oscillations accompanying the whistler precursor.

In the theory of collisionless shock waves in a magnetized plasma two regimes are conveniently distinguished: (a) the resistive regime, valid for low Mach numbers, $M_A \lesssim 3$, where the main dissipation process is electron heating by anomalous resistivity; and (b) the viscous regime, valid for higher Mach numbers, where anomalous ion heating prevails. The resistive regime is now quite well understood, since the instability mechanism has been identified.[1] As it produces a microturbulence with wavelength $\lambda \sim \lambda_D$ which is much smaller than typical magnetic scales, $c/\omega_{pe} - c/\omega_{pi}$, a fluid description in terms of a phenomenological resistivity is adequate.

The so-called viscous or supercritical regime is more complicated to understand, and several different ideas on how ion heating takes place have been proposed.[2-3] The case of a perpendicular (with respect to the ambient magnetic field B) shock wave has been elucidated recently by computer-simulation experiments.[4-5] It was found that some dissipation occurs by trapping of ions in the magnetosonic wave train behind the shock front, but that the beam of reflected ions, which is the main source of energy to be dissipated, flows upstream nearly unperturbed (in those computations the ions are assumed to be unmagnetized). Thus, ion gyration effects must play a major role in the final thermalization process, which is in agreement with the conclusion drawn in Ref. 3.

For shock waves propagating obliquely with respect to the magnetic field, a different behavior is possible. It is well known that outside a narrow angular range about the plane perpendicular to $\vec{B}$ the trailing magnetosonic wave vanishes and the whistler precursor appears. Oblique shock waves play an important role in astrophysics, the best known example being Earth's bow shock. The predominant feature here is the strong magnetic turbulence, which suggests that the dissipation process is connected with the excitation of magnetic oscillations. In this Letter we show the existence of a strong dissipation mechanism which is absent in the perpendicular case. We first present results of numerical-simulation experiments and then give a qualitative theoretical explanation of the nonlinear processes observed.

The numerical model is one-dimensional in space and three-dimensional in velocity space and describes both electrons and ions by a number of simulated particles $[N=(2-8)\times 10^4$ per species$]$ including their electrostatic and electromagnetic interactions (the displacement current is neglected which is consistent in one dimension). Thus the time scales ω_{pe}^{-1}, Ω_e^{-1}, and Ω_i^{-1} and the spatial scales λ_D, c/ω_{pe}, and c/ω_{pi} appear, which for practical computations imposes some restrictions on the mass ratio. In the runs discussed below we use $m_i/m_e=64$ and 128. The units chosen are c/ω_{pi}, c_A, and $B^2/4\pi n$ (B is the unperturbed magnetic field) so that we have temperature $T_{e,i}=\frac{1}{2}\beta_{e,i}$, where $\beta=\beta_e+\beta_i$ is the ratio of kinetic to magnetic pressure. The shock wave is produced by an electric field E_{0y} which is induced at the edge of the system and drives a magnetic piston into the plasma (the coordinate system is such that the shock propagates in the x direction; the unperturbed magnetic field is in the x,z plane).

Figure 1 illustrates a run with $m_i/m_e=128$, $c/c_A=64$, $T_e=\frac{1}{4}$, and $T_i=\frac{1}{16}$, such that $c/\omega_{pi}=128\lambda_D$. The driving field is $E_{0y}=20$ where E is in units of Bc_A/c. The angle between $\vec{B}$ and the shock normal is 45°. The main result is that in contrast to the perpendicular case there is strong interaction between the upstream plasma and the reflected ion beam. Ions are trapped in the potential wells associated with the whistler wave. To demonstrate that gyration effects of the ions play no role in this process, the same run was repeated with the Lorentz force in the equation of motion of the ions switched off, which shows essentially the same picture. The Mach number of the leading edge of the wave train is $M_A\simeq 4.9$. The ion phase-space vortices are soon filled up, so that the density reaches its downstream value within a few oscillation periods. While electron heating is moderate, since the effect of anomalous resistivity cannot be taken into account in our model, the ion temperature becomes much larger than the electron temperature because of the thermalization of the reflected ions.

Figure 2 shows an x, v_x ion phase-space plot for a run with $m_i/m_e=64$, all other parameters being unchanged. To give a clearer picture of the electrostatic interaction, this case is shown in the version with unmagnetized ions. Comparing Fig. 2 with Fig. 1, it is seen that the ion-beam interaction is much weaker for the smaller mass ratio. The main part of the reflected ions form a nearly unperturbed beam flowing upstream. As in the case of a perpendicular shock wave these ions are turned around into the shocked plasma only by gyration effects.[6]

We also performed a run, which is not shown here, with $m_i/m_e=128$, but $E_{0y}=32$, producing a faster shock wave, $M_A\simeq 5.8$. Here, too, the electrostatic beam interaction is seen to be somewhat reduced as compared with the lower Mach-number case $M_A\simeq 4.9$ in Fig. 1. However, the rapidly increasing strength of the ion-beam interaction with increasing mass ratio as seen from a comparison of Figs. 1 and 2 allows for the conclusion that for the hydrogen mass ratio this interaction should play a dominant role up to very high Mach numbers.

The strong ion dissipation process observed in the shock front suggests the presence of a counterstreaming ion-ion instability. For the case of two ion beams streaming perpendicularly to a magnetic field, the parameter range for instability was investigated in Ref. 3 by solving the linear dispersion relation. It can be seen that the linear instability is absent when the relative velocity V of the ion beams exceeds a certain

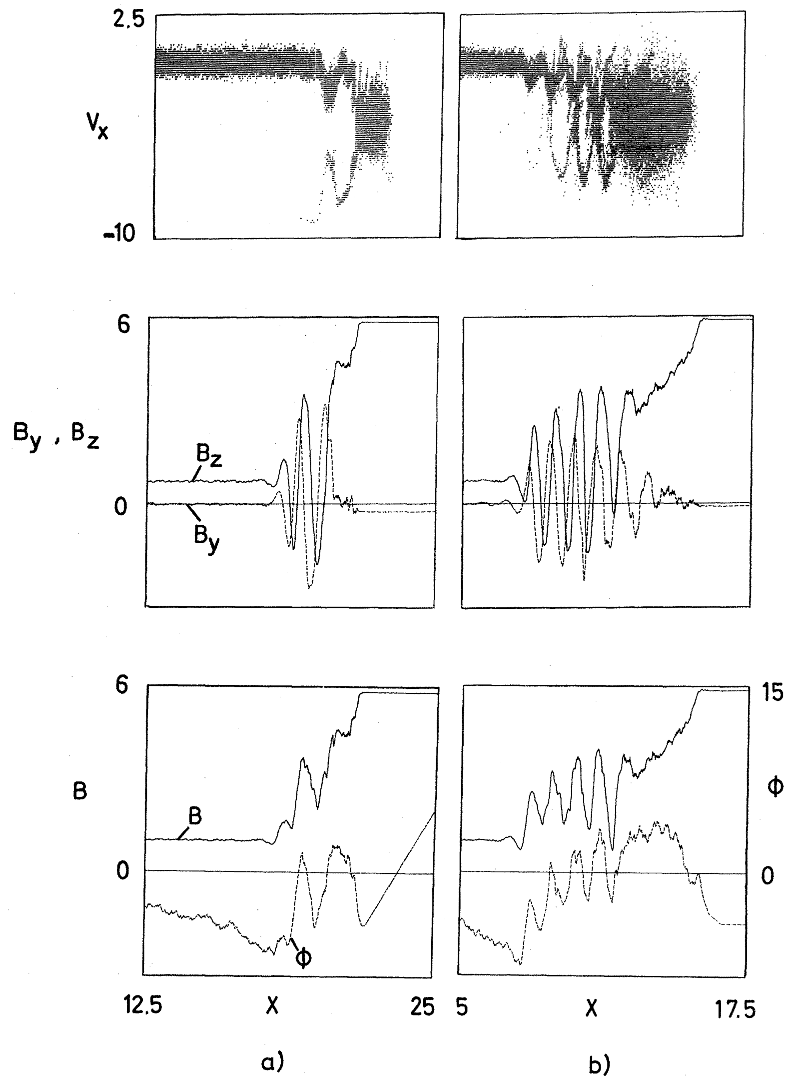

FIG. 1. Ion x, v_x phase-space plots, magnetic fields $B_y(x)$, $B_z(x)$, $B(x)$, and electric potential $\varphi(x)$ at times (a) $t = 1.25\Omega_i^{-1}$, (b) $t = 3.25\Omega_i^{-1}$. System length $L = 25c/\omega_{pi}$, mass ratio $m_i/m_e = 128$, $\lambda_D = 1/128c/\omega_{pi}$.

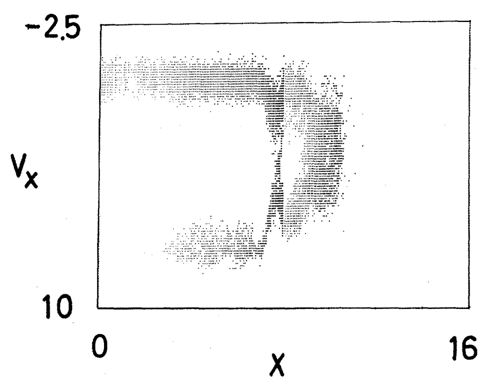

FIG. 2. Ion x, v_x phase-space plot at time $t = 2.00\Omega_i^{-1}$ of a run with $m_i/m_e = 64$ and unmagnetized ions; all other parameters as in Fig. 1.

value $V_0 \simeq 2c_A(1+\beta)^{1/2}$. This is always the case in supercritical perpendicular shock waves, which explains the observation made in Refs. 4 and 5 that the reflected ions do not interact with the incoming plasma.

The dispersion relation is quite different for two ion beams propagating obliquely at an angle θ with respect to B. For cold ion beams the dispersion relation was derived in Ref. 2. When evaluating it, we find instability for much higher relative velocities, $V < c_A(m_i/m_e)^{1/2}$, assuming $\Omega_e \sim \omega_{pi}$ and $kc/\omega_{pe} \sim \cos\theta \sim 1$ and some other conditions not given here. However, it can easily be seen that this "beam-whistler instability," as it was called in Ref. 2, is not an instability between the two ion beams but between one ion beam and the electron component with phase velocities near those of either ion beam. Numerical simulation of the development of this beam-whistler instability shows that it is quite ineffective. Hence no appropriate linear instability mechanism is available.

We claim that the instability seen in Fig. 1 is essentially the electrostatic two-ion-beam instability excited *nonlinearly*. The two-ion-beam configuration is only linearly stable, but nonlinearly unstable if the initial perturbation φ_0 is sufficiently large. Separate numerical simulations of the two-ion-beam system initially modulated by a large-amplitude wave confirm this picture. Thus the electrostatic dissipation of the reflected ions is due to the presence of the whistler precursor and its coupling to electrostatic oscillations. For perpendicular shock waves no whistler is excited and hence no electrostatic thermalization of the reflected ions occurs. The same is true of nearly parallel shock propagation, since here the whistler becomes purely electromagnetic.

We have presented a mechanism of anomalous ion dissipation in oblique collisionless shock waves. We still have to relate our results to experimental observations made, for instance, on the bow shock. The very coherent, laminar behavior seen in Fig. 1 is probably due to the one-dimensional character of the model (similar behavior was found in the perpendicular case[4] and for purely electrostatic shock waves[7]). In a higher-dimensional space, wave-coupling processes will probably disturb the whistler wave train and make it more turbulent. An essential feature of the dissipation process is that it is intimately connected with strong excitation of magnetic field oscillations. This explains the highly oscillatory behavior of the bow-shock magnetic field profiles, which did not appear in previous numerical-simulation computations,[6] where the only ion dissipation process was by gyration. In the present model anomalous resistivity was not taken into account, which could damp the wave train to a certain degree. However, the wild magnetic turbulence in the bow shock indicates that this damping is not very effective. It should be noted that the present nonlinear instability mechanism has no relation to an electrostatic shock model discussed earlier,[8] which implies a *linear* instability between the upstream and the downstream plasmas (instead of the reflected ion beam) with scale $\lambda \sim \lambda_D$ (instead of c/ω_{pe}) and hence no coupling to magnetic field oscillations.

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