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Generation of Magnetic Fluctuations Near a Shock Front in a Partially Ionized Medium

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Abstract—We investigate the generation mechanism of long-wavelength Alfvénic disturbances near the front of a collisionless shock that propagates in a partially ionized plasma. The wave generation and dissipation rates are calculated in the linear approximation. The instability is attributable to a current of energetic particles upstream of the shock front. The generation of long-wavelength magnetic fluctuations is most pronounced for strong shocks, but the effect is retained for shocks with a moderate particle acceleration efficiency without any noticeable modification of the shock structure by the pressure of accelerated particles. The mode generation time for supernova remnants in a partially ionized interstellar medium is shown to be shorter than their age. Long-wavelength magnetic disturbances determine the limiting energies of the particles accelerated at a shock by the Fermi mechanism. We discuss the application of the mechanism under consideration to explaining the observed properties of the SN 1006 remnant. *© 2005 Pleiades Publishing, Inc.*

Key words: *supernovae and supernova remnants, MHD instability, accelerated particles, magnetic field.*

INTRODUCTION

The acceleration of particles near collisionless shock fronts is an efficient conversion mechanism of the free energy of supersonic magnetized plasma flows into the energy of nonthermal particles. Popular models for the origin of cosmic rays (CRs) consider young supernova remnants as the main sources of high-energy CRs (Berezinskiĭ *et al.* 1990; Ptuskin and Zirakashvili 2003, 2005; Hillas 2005) and shock acceleration as the main generation mechanism of relativistic particles. In models for the collective CR acceleration by multiple interactions with supernova remnants and strong winds from massive early-type stars, ensembles of shocks with various intensities also play a crucial role (Bykov and Toptygin 2001).

The X-ray observations of supernova remnants performed by the Chandra telescope with an angular resolution of ∼1 arcsec (see, e.g., Vink 2004) point to an efficient generation mechanism of magnetic fields in the preshock region of the Cas A remnant.

Anisotropic CR distributions can lead to the generation of MHD waves (for a review, see Berezinskiĭ *et al.* 1990). The resonant generation of Alfvén waves by relativistic particles was considered as a possible mechanism of turbulence formation in the Galaxy; it provides the diffusion of CRs with energies up 100 GeV and the multiple scattering of accelerated particles near shock fronts (for a review, see Blandford and Eichler 1987). The nonresonant generation mechanisms of MHD waves can also be efficient near shock fronts. The instability of magnetosonic disturbances propagating in the upstream region of a strong shock modified by a CR pressure gradient is a possible nonresonant generation mechanism of a random magnetic field. The various cases of this instability were considered by Drury (1984), Zank and McKenzie (1987), Berezhko (1986), Chalov (1988), and Zank *et al.* (1990). Recently, Bell (2004) pointed out the possibility of an efficient generation of small-scale Alfvén modes (with scales smaller than the gyroradii of nonthermal particles) in a strong shock propagating in a completely ionized plasma.

In many cases, shocks propagate in a partially ionized medium. For supernova remnants interacting with molecular clouds (e.g., IC 443), the presence of a neutral component upstream of the shock front can lead to peculiar features in the regime of high-energy particle acceleration (see, e.g., Drury *et al.* 1996) and affects significantly the radiation spectra of such remnants (Bykov *et al.* 2000). A certain fraction of neutral hydrogen, helium, and metal atoms reach the

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shock front even in the case of a supernova remnant in a tenuous medium. These are observed in the optical and ultraviolet spectra of the shock front as a superposition of broad and narrow lines (in particular, for the $H\alpha$ line) in the remnants of SN 1006, Kepler, Tycho, RCW 86, the Cygnus Loop, etc. and are an efficient tool for estimating the shock velocity (Chevalier and Raymond 1978; Raymond 2001).

Below, we calculate the linear growth rate of long-wavelength Alfvénic oscillations excited by accelerated particles with allowance made for the dissipation of MHD oscillations in a medium containing not only a plasma, but also a neutral gas. The latter prevents fast background plasma neutralization of the electric current produced by relativistic particles and leads to a renormalization of the magnetic viscosity, which makes the growth of MHD oscillations possible.

INSTABILITY OF THE PRESHOCK REGION IN A PARTIALLY IONIZED MEDIUM

Let us consider a shock propagating in a partially ionized medium with a magnetic field. The matter upstream of a fairly strong shock front is assumed to be cold, and the gas pressure and viscosity may be disregarded. However, we take into account the Joule dissipation and the pressure of nonthermal particles.

Basic Equations

Let us write the magnetohydrodynamic equations in the preshock region in the rest frame of the shock front:

$$\nabla \times \mathbf{B} = \frac{4\pi}{c}(\mathbf{j} + \mathbf{j}^{\rm cr}), \qquad \nabla \times \mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}, \quad (1)$$

$$\rho\left(\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u}\cdot\nabla)\mathbf{u}\right) \quad (2)$$
$$= -\nabla P_{\rm cr} - \frac{1}{4\pi}\mathbf{B}\times(\nabla\times\mathbf{B}),$$

where ρ and $\mathbf{u}$ are, respectively, the mass density and the total macroscopic velocity of the medium (including the nonthermal plasma component); $P_{\rm cr}$ is the pressure of the nonthermal component; $\mathbf{E}$ and $\mathbf{B}$ are the total electric and magnetic field strengths; $\mathbf{j}$ is the current excited by these fields in the background plasma; and $\mathbf{j}^{\rm cr}$ is the macroscopic current produced by the accelerated particles. The latter was averaged over small-scale fluctuations and should be considered as the current produced by an extraneous (to the background medium) source. Equation (2) is applicable to describing the motions of the medium with scale lengths exceeding the mean free paths of the nonthermal energetic particles that determine $P_{\rm cr}$ and $\mathbf{j}^{\rm cr}$. The presence of an extraneous current of accelerated particles contributing to the total magnetic field is a major factor of the long-wavelength instability considered below. This distinguishes our case from the previously considered long-wavelength instabilities associated with the CR pressure gradient. An important factor is also the presence of a neutral component in the preshock region (even at a fairly high degree of ionization of the medium). In our analyzed one-dimensional case where all quantities depend on one z coordinate, both currents are transverse relative to the normal to the front plane.

The relation between the current $\mathbf{j}$ and the field vectors should be added to the above equations. In a medium with a fraction of neutral matter, this relation largely determines the form of Ohm's generalized law and the dissipative effect (i.e., the magnetic viscosity). This question has been explored in detail in the past decades (see Pikel'ner 1964; Ruzmaĭkin *et al.* 1988). We use Ohm's generalized law in the form given in the monograph by Pikel'ner (1966):

$$\mathbf{E} + \frac{1}{c}\mathbf{u}\times\mathbf{B} = \frac{1}{\sigma}\mathbf{j} + \frac{1}{n_i e c}\mathbf{j}\times\mathbf{B} \quad (3)$$
$$+ \frac{F^2\tau_i}{n_i m_i c^2}\mathbf{B}\times(\mathbf{j}\times\mathbf{B}),$$

where $\mathbf{j}$ is the current of background particles of a partially ionized plasma, $\mathbf{E}$ and $\mathbf{B}$ are the vectors of the total field produced by the current $\mathbf{j}$ and other (external) sources, $n_i \approx n_e$ are the number densities of the charged components, and τ_i and τ_e are the mean times between collisions of ions and electrons, respectively, with other particles (in our case, mainly with neutral atoms and molecules). Finally, $\sigma = n_e e^2\tau_e/m_e$ is the collisional electric conductivity and F is the mass fraction of the neutral particles. In a warm medium and neutral clouds far from the shock front, the fraction $F \approx 1$, but the ionizing radiation of the gas heated by a strong shock decreases appreciably in the preshock region, where neutral helium atoms can play an important role.

For a magnetized plasma with an admixture of neutrals with $\omega_B\tau \gg 1$ (here, $\omega_B = eB_0/(mc)$), the three terms on the right-hand side of Eq. (3) differ by many orders of magnitude. The ratio of the third and second terms yields $\tau_i eB_0/(m_i c) = \omega_{Bi}\tau_i \gg 1$. The ratio of the third and first terms is of the order of $\tau_i B_0^2\sigma/(n_i m_i c^2) = (\omega_{Bi}\tau_i)(\omega_{Be}\tau_e) \gg 1$. These estimates become invalid only if the transverse (relative to $\mathbf{B}$) current component is anomalously small. However, the current is transverse in the case under consideration. Let us retain only the last term on the right-hand side of Eq. (3) and substitute $\mathbf{B} = \mathbf{B}_0 + \mathbf{b}$, $\mathbf{j} = (c/4\pi)\nabla\times\mathbf{b} - \mathbf{j}^{\rm cr}$ into it. This yields the following after linearization: $\mathbf{B}\times(\mathbf{j}\times\mathbf{B}) = (cB_0^2/(4\pi))\nabla\times\mathbf{B} - B_0^2\mathbf{j}^{\rm cr}$. Eliminating the electric field $\mathbf{E}$ from (3) and (1), we obtain a linearized equation for the magnetic field in the standard form,

$$\frac{\partial \mathbf{b}}{\partial t} + (\mathbf{u} \cdot \nabla)\mathbf{b} = (\mathbf{B}_0 \cdot \nabla)\mathbf{u}_\perp + \nu_\mathrm{m} \triangle \mathbf{b} + \frac{4\pi\nu_\mathrm{m}}{c} \nabla \times \mathbf{j}^\mathrm{cr}, \quad (4)$$

but with an essentially renormalized magnetic viscosity, which we express in terms of the effective electric conductivity σ_ef:

$$\nu_\mathrm{m} = \frac{c^2}{4\pi\sigma_\mathrm{ef}} = \frac{F^2 \tau_i B_0^2}{4\pi n_i m_i}.$$

The linearized equation of motion for the medium (2) takes the form

$$\rho \left(\frac{\partial \mathbf{u}_\perp}{\partial t} + (\mathbf{u} \cdot \nabla)\mathbf{u}_\perp \right) = -\frac{1}{4\pi} \mathbf{B}_0 \times (\nabla \times \mathbf{b}). \quad (5)$$

The pressure gradient of nonthermal particles does not appear in the equation for the transverse velocity. Now the constant velocity component of the preshock medium perpendicular to the shock front is denoted in Eqs. (4) and (5) by $\mathbf{u}$ without any subscript. The small velocity component parallel to the shock front has the subscript $\perp$. In what follows, we will consider a purely hydrogen medium with $m_i = m_p$ and $\omega_{Bi} = \omega_{Bp}$.

Calculating the Current of Accelerated Particles

Let us calculate the electric current that is generated by nonthermal particles under the electric and magnetic fields of an MHD oscillation. Denote these fields by $\mathbf{b}$ and $\mathbf{E}$ and consider the case with the simplest geometry where the external field $\mathbf{B}_0$ is uniform and directed along the normal to a plane boundless front, while the MHD oscillation field is perpendicular to the external field ($\mathbf{b} \perp \mathbf{B}_0$) and is a plane wave propagating along $\mathbf{B}_0$: $\mathbf{b} = \mathbf{b}_0 e^{i(\mathbf{k}\cdot\mathbf{r} - \omega t)}$, $\mathbf{k} = k\mathbf{e}_{||}$. The electric field of the wave is related to the magnetic field by the electromagnetic induction equation

$$\mathbf{k} \times \mathbf{E} = \frac{\omega}{c}\mathbf{b}, \qquad \mathbf{E} = -\frac{\omega}{ck}\mathbf{e}_\| \times \mathbf{b}.$$

We perform our analysis in the rest frame of the shock front.

The distribution function $f(\mathbf{r}, p_\perp, p_\|, \phi, t)$ of relativistic particles in the geometry of the problem under consideration satisfies the equation

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \frac{\partial f}{\partial \mathbf{r}} + e\mathbf{E} \cdot \frac{\partial f}{\partial \mathbf{p}} - \frac{ec}{\varepsilon}(\mathbf{B}_0 + \mathbf{b}) \cdot \mathcal{O} f = I[f], \quad (6)$$

where

$$\mathcal{O} = \mathbf{p} \times \frac{\partial}{\partial \mathbf{p}}$$

is the momentum rotation operator, ε is the particle energy, and $I[f]$ is the collision integral averaged over the turbulent spectrum that includes the interaction of relativistic particles with MHD turbulence. Specific expressions for this quantity in the lowest-order approximation in turbulence amplitude can be found in a monograph by Toptygin (1985), while nonlinear corrections are given in our previous review (Bykov and Toptygin 1993).

Let us linearize kinetic equation (6) by assuming the external field to be weak and by separating out the small part δf attributable to this field from the distribution function $f = f_0 + \delta f$:

$$\frac{\partial \delta f}{\partial t} + \mathbf{v} \cdot \frac{\partial \delta f}{\partial \mathbf{r}} - \frac{ec}{\varepsilon}\mathbf{B}_0 \cdot \mathcal{O}\delta f = -e\mathbf{E} \cdot \frac{\partial f_0}{\partial \mathbf{p}} + \frac{ec}{\varepsilon}\mathbf{b} \cdot \mathcal{O} f_0 \equiv Q(\mathbf{p}, z, t). \quad (7)$$

Here, f_0 is the stationary (in the rest frame of the shock front) distribution function of the accelerated particles in the absence of the MHD wave under consideration. The accelerated particles undergo strong scattering and have a weakly anisotropic distribution function that can be written as (Toptygin 1985)

$$f_0(z, \mathbf{p}) = \frac{1}{4\pi}\left[N(p, z) + \frac{3}{pv}\mathbf{p} \cdot \mathbf{J}(p, z)\right], \quad J \ll vN, \quad (8)$$

where

$$J_\alpha = -\kappa_{\alpha\beta}\frac{\partial N}{\partial x_\beta} - \frac{p}{3}\frac{\partial N}{\partial p}u_\alpha \quad (9)$$

is the differential flux density of the accelerated particles, $\kappa_{\alpha\beta}$ is their diffusion tensor, and $\mathbf{u}$ is the velocity of the medium. The applicability of relations (8) and (9) is limited by the condition $u \ll v$.

The isotropic part $N(p, z)$ of the distribution function in the preshock region can be easily calculated in the stationary case:

$$N(p, z) = (\alpha - 3)N_0 \frac{p_0^{\alpha - 3}}{p^\alpha} \times \exp\left[\int\limits_0^z \frac{u dz'}{\kappa_\|(p, z')}\right], \quad z \le 0, \quad p_0 \le p \le p_m. \quad (10)$$

Here, N_0 is the number density of relativistic particles with all energies, $\alpha = 3u/\Delta u$ is the spectral index, and $\Delta u > 0$ is the jump in velocity at the shock front. Solution (10) corresponds to the acceleration of test particles where a small part of the shock energy flux is transferred to them. In this case, the shock front modification by the accelerated particles may be disregarded and the velocity of the medium in the preshock region may be assumed to be approximately constant, $u \approx$ const. The spectral index is $\alpha > 4$ for a moderately strong shock front and $\alpha \le 4$ for a strong shock (see, e.g., Toptygin 1997). The specific value of α depends not only on the Mach number of the wave, but also on the rate of particle injection into the acceleration process. At $\alpha < 4$, the bulk of the energy belongs to the most energetic particles with $\varepsilon \lesssim \varepsilon_{\rm m} = cp_{\rm m}$, and several tens of percent of the total flow energy is spent on particle acceleration. In this calculation, we restrict ourselves to spectral indices $\alpha \ge 4$, suggesting a moderate acceleration rate at which the total kinetic energy of the accelerated particles generally does not exceed 10% of the system's total energy, although the maximum particle momentum $p_{\rm m}$ can be much larger than the injection momentum: $p_{\rm m} \gg p_0$. At such spectral indices, the total energy of the accelerated particles at the shock front ($z = 0$) depends on p_m very weakly not more strongly than the logarithmic law:

$$w_{\rm cr} \approx \int\limits_{p_0}^{p_m} cpN(p,0)p^2dp = \frac{\alpha-3}{\alpha-4}N_0 m_p c^2, \quad (11)$$

$$\alpha > 4; \quad \frac{\alpha-3}{\alpha-4} \to \ln\frac{p_m}{p_0} \quad \text{at} \quad \alpha \to 4.$$

At an arbitrary spectral index $\alpha \ge 4$, the particle distribution function and the current take the form

$$f_0(z,\mathbf{p}) = \frac{(\alpha-3)N_0 p_0^{\alpha-3}}{4\pi p^\alpha} \times \left[1 + \frac{u}{v}(\alpha-3)\cos\theta\right] e^{uz/\kappa_\parallel},$$

$$Q(\phi) = (\mathbf{b}\cdot\mathbf{e}_\phi)\frac{(\alpha-3)eN_0 p_0^{\alpha-3}}{4\pi\Omega p^{\alpha+1}} \times \left[(\alpha-3)\frac{u}{c} - \frac{\omega}{ck}\left(\alpha + \frac{u}{v}(\alpha+1)(\alpha-3)\cos\theta\right)\right] \times e^{uz/\kappa_\parallel}\sin\theta,$$

$$\delta f = \int\limits_{\pm\infty}^{\phi} Q(\phi')e^{a(\phi-\phi')}d\phi' = (\mathbf{b}\cdot\mathbf{e}_\phi)\frac{(\alpha-3)eN_0 p_0^{\alpha-3}}{4\pi\Omega p^{\alpha+1}} \times \left[(\alpha-3)\frac{u}{c} - \frac{\omega}{ck}\left(\alpha + \frac{u}{v}(\alpha+1)(\alpha-3)\cos\theta\right)\right] \times \frac{\mathbf{b}\cdot\mathbf{e}_\perp - a\mathbf{b}\cdot\mathbf{e}_\phi}{1+a^2}e^{uz/\kappa_\parallel}\sin\theta,$$

where

$$a = \frac{1}{\Omega}[v_\parallel u/\kappa_\parallel - i(\omega - kv_\parallel)].$$

The expression for the current of relativistic particles can be reduced to

$$\mathbf{j}^{\rm cr} = \int e\mathbf{v}\delta f(p,\theta,\phi)p^2dp\sin\theta d\theta d\phi = \int p^2dp\sin^3\theta d\theta\frac{(\alpha-3)ve^2N_0p_0^{\alpha-3}}{4\Omega p^{\alpha+1}} \times \left[(\alpha-3)\frac{u}{c} - \frac{\omega}{ck}\left(\alpha + \frac{u}{v}(\alpha+1)(\alpha-3)\cos\theta\right)\right] \times \frac{\mathbf{b} + a\mathbf{e}_\parallel \times \mathbf{b}}{1+a^2}e^{uz/\kappa_\parallel}.$$

In the previous formulas, except formula (11), there was no substitution $v \to c$; these are also valid for nonrelativistic energies, but at $v \gg u$. Below, we consider the relativistic case, $v \approx c$, $p_0 = m_pc$.

At present, there is no consistent theory to calculate the turbulence spectrum in the vicinity of a shock. The turbulence-determined diffusion coefficient of energetic particles has to be specified from model considerations. The Bohm diffusion model is most popular (see, e.g., the reviews by Jones and Ellison 1991 and Malkov and Drury 2001). This model assumes strong turbulence at which the local transport mean free path $\Lambda(p)$ of a particle is of the order of its gyroradius:

$$\Lambda(p) \approx \eta r_{\rm g}(p) = \frac{cp}{eB},$$

$$\tilde{\kappa}_\parallel = \frac{c\Lambda}{3}, \quad p_0 \le p \le p_{\rm m},$$

where the parameter $\eta \gtrsim 1$ and p_0 and $p_{\rm m}$ bound the range of momenta under consideration; the case of $p_{\rm m} \gg p_0$ is of considerable interest. The turbulent and regular fields are assumed to be of the same order of magnitude: $B \approx B_0$. The latter condition is consistent with the assumption of a minor fraction of the energy being transferred to the accelerated particles, since the mechanical energy density in a strong shock under typical astrophysical conditions is several orders of magnitude higher than the energy density of the primary magnetic field. The pattern of particle transport depends on the relationship between the local diffusion coefficient $\tilde{\kappa}_\parallel$ and the parameters of the turbulent plasma motions with scale lengths exceeding Λ. If the fluctuation amplitude of the macroscopic velocity of the medium $\delta u(l)$ with a scale length $l \gg \Lambda$ is so large that $\delta ul \gtrsim \tilde{\kappa}_\parallel$, then the turbulent particle transport will play a major role on these scales. In this model, the diffusion coefficient does not depend on the particle energy over a wide energy range (Bykov and Toptygin 1993).

The solution of kinetic equation (7) for a constant diffusion coefficient has the simplest form that we consider below, although there are no fundamental difficulties in considering systems with $\kappa_{\parallel}(p)$ dependent on the momentum p. For the CR current after calculating the integrals, we obtain

$$\mathbf{j}^{\rm cr} = -(\sigma'_{\rm cr} + i\sigma''_{\rm cr})\mathbf{e}_{\parallel} \times \mathbf{b} + (g' + ig'')\mathbf{b}, \qquad (12)$$

where the following notation is used:

$$\sigma'_{\rm cr} = \frac{(\alpha-3)^2\omega_0^2\omega u}{8(\alpha-1)\omega_{Bp}^2 c}\left(\frac{ck}{\omega_{Bp}}\right)^{\alpha-4},$$

$$\sigma''_{\rm cr} = \frac{(\alpha+1)(\alpha-3)^2\omega_0^2\omega u}{60\pi(\alpha-4)\omega_{Bp}^2 c},$$

$$g' = -\frac{\omega_0^2}{12\pi\omega_{Bp}}\left(\frac{3u}{c} + \alpha\frac{\omega - uk}{ck}\right),$$

$$g'' = \frac{(\alpha-3)\omega_0^2}{8\alpha(\alpha-2)\omega_{Bp}}\left(\frac{ck}{\omega_{Bp}}\right)^{\alpha-3}\left(\frac{3u}{c} + \alpha\frac{\omega - uk}{ck}\right),$$

$$\omega_0^2 = \frac{4\pi e^2 N_0}{m_p}, \qquad \xi_0 = \frac{\omega_{Bp}}{ck}.$$

When passing to $\alpha \to 4$, we should make the substitution $(\alpha-4)^{-1} \to \ln\xi_0 + 8/15$.

The current can also be represented in a conventional form similar to Ohm's law in a gyrotropic medium, $\mathbf{j}^{\rm cr} = \sigma_{\rm ef}\mathbf{E} + \sigma_{\rm H}\mathbf{e}_{\parallel} \times \mathbf{E}$, but representation (12) in the MHD approximation is more convenient.

The Growth Rate of MHD Waves

Equations (4) and (6) together with the extraneous current of relativistic particles (12) allow us to derive a dispersion relation for the waves that can be excited in a weakly ionized medium with an admixture of nonthermal energetic particles. We seek a solution of these equations for $|z| \ll \kappa_{\parallel}/u$ in the form $\mathbf{b}$, $\mathbf{u}_{\perp} \propto e^{i(kz-\omega t)}$. Equation (5) allows the transverse velocity to be expressed in terms of the magnetic field:

$$\mathbf{u}_{\perp} = \frac{kB_0}{4\pi\rho\omega'}\mathbf{b},$$

where $\omega' = \omega - uk$ is the oscillation frequency in the frame of reference comoving with the medium. Using this relation, we derive a system of equations for the magnetic field components from Eqs. (4) and (12):

$$\left\{\omega'^2 - (u_{\rm A}k)^2 + \frac{4\pi}{c}\nu_{\rm m}\sigma'_{\rm cr}k\omega' + i\nu_{\rm m}k^2\omega' \times \left[1 + \frac{4\pi}{ck}\sigma''_{\rm cr}\right]\right\}b_x - \frac{4\pi}{c}\nu_{\rm m}(g' + ig'')k\omega' b_y = 0,$$

$$\frac{4\pi}{c}\nu_{\rm m}(g' + ig'')k\omega' b_x + \left\{\omega'^2 - (u_{\rm A}k)^2 + \frac{4\pi}{c}\nu_{\rm m}\sigma'_{\rm cr}k\omega' + i\nu_{\rm m}k^2\omega'\left[1 + \frac{4\pi}{ck}\sigma''_{\rm cr}\right]\right\}b_y = 0,$$

where $u_{\rm A}$ is the Alfvén velocity. Equating the determinant of this system to zero yields a dispersion relation that defines several oscillation branches:

$$\omega'^2 - (u_{\rm A}k)^2 + \frac{4\pi}{c}\nu_{\rm m}(\sigma'_{\rm cr} \mp g')k\omega' + i\nu_{\rm m}k^2\omega'\left[1 + \frac{4\pi}{ck}(\sigma''_{\rm cr} \pm g')\right] = 0. \qquad (13)$$

In the absence of energetic nonthermal particles ($\sigma_{\rm cr} = g = 0$), we obtain a dispersion relation that describes damped Alfén waves:

$$\omega' = \pm u_{\rm A}k - i\gamma, \qquad \gamma = \frac{1}{2}\nu_{\rm m}k^2$$

(in the approximation of $\gamma \ll u_{\rm A}k$). In the presence of energetic nonthermal particles, the term in square brackets is added to the imaginary part in Eq. (13); this term can have different signs, since $\sigma''_{\rm cr}/g' \approx u\ln\xi_0/c \ll 1$. The oscillation branch that corresponds to the plus in front of g' at fairly low values of

$$k < k_{\rm c} = \frac{4\pi}{c}|g'|$$

will grow, $\mathbf{b} \propto e^{\gamma t}$, at a rate

$$\gamma \approx \frac{1}{2}\nu_{\rm m}k^2\left(\frac{4\pi|g'|}{ck} - 1\right). \qquad (14)$$

Significantly, the nonresonant generation mechanism of long-wavelength magnetic fluctuations in a partially ionized medium requires no pressure gradient (or other inhomogeneity) in the preshock region. This distinguishes it from the previously mentioned instabilities of an inhomogeneous preshock region modified by the CR pressure. However, one might expect growth rate (14) to remain valid an order of magnitude up to $k \gtrsim 2\pi/l_{\rm m} \approx 2\pi u/\kappa_{\parallel}$ in the presence of an inhomogeneous preshock region as well. The front inhomogeneity can be consistently taken into account in the geometrical optics approximation in a way similar to that used by Chalov (1988). A certain renormalization of the real part of the frequency also takes place. The presence of a neutral component (even with a relatively low mass fraction F) in the case of a magnetized medium with $\omega_{Bi}\tau \gg 1$ changes radically Ohm's law and gives rise to long-wavelength instability. For a shock with an Alfvén Mach number $M_{\rm A} > 4$, the estimate of growth rate (14) is

$$\gamma \sim \frac{F^2}{6}(\omega_{Bi}\tau_i)\frac{N_0}{n_i}ku. \qquad (15)$$

THE GENERATION OF MAGNETIC FIELD FLUCTUATIONS IN THE SHOCKS OF SUPERNOVA REMNANTS

Let us consider the possible applications of the Alfvén wave generation mechanism in the shock of the SN 1006 remnant considered above. This is one of the young supernova remnants known from ancient historical chronicles (Lozinskaya 1986), and it probably belongs to type-Ia supernovae. The X-ray emission from the SN 1006 remnant is characterized by bright thin segments located in the northeast (NE) and southwest (SW) parts of a roughly spherical shell $\sim 30'$ in diameter. The bright NE part has recently been studied in detail in the Chandra X-ray observatory by Long *et al.* (2003) and Bamba *et al.* (2003). The X-ray spectrum of the thin bright NE segment is dominated by a nonthermal continuum that is commonly interpreted as the synchrotron radiation of electrons with energies of $\sim$10–100 TeV in the vicinity of a shock. The high angular resolution ($\sim 1''$) of the ACIS CCD detector in the Chandra observatory allowed Long *et al.* (2003) to detect a sharp jump in radiation intensity. They established that the intensity of the radiation at energies above 1.2 keV immediately upstream of the shock front does not exceed 1.5% of the maximum brightness immediately downstream of the shock front. The bright NE segment of the X-ray synchrotron radiation is $\approx 10''$ in width ($1'' = 3.3 \times 10^{16}$ cm at an estimated distance to SN 1006 of 2.2 kpc). The presence of a weak radio halo in remnant supernovae and estimates of the diffusion coefficients for relativistic electrons were previously discussed by Achterberg *et al.* (1994), but the upper limit on the brightness of the synchrotron halo set by Long *et al.* (2003) is most stringent (see Ballet 2005).

The optical and ultraviolet spectra of SN 1006 obtained by Korreck *et al.* (2004) agree with the estimates of the neutral mass fraction $F \sim 0.1$ in the upstream region of a shock propagating with a velocity $v_{\rm sh} \sim 2300$ km s^{-1}. The gas density upstream of the front of the NE sector of the shock in SN 1006 was estimated to be $n_i \sim 0.1$ cm^{-3}. Using the hydrogen charge exchange reaction rate at a temperature of $\sim 10^4$ K from Kulsrud and Cesarsky (1971), we obtain the mean free path of a hydrogen atom relative to charge exchange equal to the minimum instability wavelength $\lambda_0 = 2\pi k_0^{-1} \sim 2 \times 10^{16}$ cm (since $k_0 \ll k_c$) and an estimate of the magnetization factor, $\omega_{Bi}\tau_i \gtrsim 10^7 B n_{-1}$, where B is measured in μG. Thus, using relation (15), we obtain a characteristic mode growth time scale of $\sim 6 \times 10^2 (N_0/n_i)^{-1}$ (s), which allows magnetic fields with scale lengths of the order of λ_0 to be amplified over the lifetime of SN 1006 if the rate of proton injection into the acceleration process at the shock admits of $N_0/n_i \gtrsim 10^{-7}$. Using relation (11), we can make sure that the energy density of the accelerated particles accounts for a few fractions of a percent of the onflow kinetic energy density. If we restrict our analysis to injection rates that admit of CR energy densities $w_{\rm cr}$ on the order of several percent of the flow kinetic energy density ($\sim m_p n_i v_{\rm sh}^2$), then we will get the possibility of the generation of magnetic field fluctuations upstream of the shock front with amplitudes of $\delta B \sim 30$ μG. A compression of the transverse field component at the discontinuity in a strong shock with $R \approx 4$ (here, we consider a single-fluid wave without an extended preshock region, since we investigate the case of low proton injection rates) can yield magnetic fields of ~ 100 μG in the postshock region. Magnetic fields in the postshock region of SN 1006 of ~ 100 μG allow the observed intensity distribution of the X-ray continuum emission in the NE segment of the shock (Long *et al.* 2003; Berezhko *et al.* 2003; Ballet 2005) to be explained in terms of the rapid synchrotron cooling of relativistic electrons downstream of the shock front. In the case of magnetic field generation in a partially ionized medium, no significant pressure of the CR nucleon component ($\sim m_p n_i v_{\rm sh}^2$), assumed in the model by Berezhko *et al.* (2003) in the preshock region, is required. The upper limit for the ratio of the synchrotron luminosity upstream of the shock front to the maximum luminosity in the transverse segment of the shock is $R^{-\Gamma}$, where Γ is the photon spectral index of the synchrotron radiation at energies above $\sim$1.5 keV. In our case of compression in a strong single-fluid shock with $R \approx 4$ and a synchrotron index $\Gamma \sim 3$, we obtain a ratio of the specific luminosities close to 1.5%, in agreement with the limit set by Long *et al.* (2003).

Another interesting application of the physical mechanism of magnetic field generation by a shock in a plasma medium with a neutral component may be radio filaments near the Galactic center (Morris and Serabyn 1996). The observations by Yusef-Zadeh *et al.* (2005) point to the possible association of some of the filaments with supernova remnants.

CONCLUSIONS

We investigated a new type of instability of a partially ionized plasma with relativistic particles accelerated at a strong shock front. Alfvénic MHD waves are excited by a nonresonant mechanism upstream of the shock front. Applying this mechanism to the remnant of SN 1006 allows us to explain the main features of the emission from this remnant and the possible magnetic field amplification downstream of the shock front (~ 100 μG) without assuming a significant nonlinear modification of the profile of the preshock region by the accelerated particles.

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