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ARE GAMMA-RAY BURST SHOCKS MEDIATED BY THE WEIBEL INSTABILITY?

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ABSTRACT

It is estimated that the Weibel instability is not generally an effective mechanism for generating ultrarelativistic astrophysical shocks. Even if the upstream magnetic field is as low as in the interstellar medium, the shock is mediated not by the Weibel instability but by the Larmor rotation of protons in the background magnetic field. Future simulations should be able to verify or falsify our conclusion.

Subject headings: gamma rays: bursts — instabilities — magnetic fields — plasmas — shock waves

1. INTRODUCTION

There is much literature on gamma-ray burst (GRB) afterglows based on the assumption that the X-ray, optical, and radio afterglows are the synchrotron emission from relativistic electrons Fermi accelerated at the forward shock of the blast wave (see recent reviews by Mészáros 2002; Zhang & Mészáros 2004). A long-standing difficulty with this assumption has been that the inferred magnetic field needed to fit the afterglow data typically requires that the magnetic energy density exceeds by many orders of magnitude the magnetic energy density that would be expected from the shock compression of the interstellar magnetic field of the host galaxy (Gruzinov & Waxman 1999; Gruzinov 2001). Many authors have therefore assumed that the shock somehow manufactures field energy to meet this requirement, but no convincing mechanism has been proposed to date. One mechanism discussed is the Weibel instability (Medvedev & Loeb 1999; Silva et al. 2003; Frederiksen et al. 2004; Jaroschek et al. 2005; Medvedev et al. 2005; Kato 2005; Nishikawa et al. 2005), which has the fastest growth rate and produces relatively strong small-scale magnetic field even in an initially nonmagnetized plasma. It is expected that the thermalization of the upstream flow could occur via scattering of particles on the magnetic fluctuations. In the electron-positron plasma, the instability does generate the magnetic field at about 10% of the equipartition level and does provide the shock transition at the scale of a dozen electron inertial lengths (Spitkovsky 2005). However, simulations of colliding electron-proton flows show that while the electrons are readily isotropized, the protons acquire only small scattering in angles after passing the simulation box (Frederiksen et al. 2004). How long the field persists after the shock is also an important question (Gruzinov 2001), but here we discuss whether the Weibel instability can even cause the shock in the first place.

Moiseev & Sagdeev (1963; see also Sagdeev 1966) analyzed the structure of the nonrelativistic Weibel-driven shock and found that the width of the shock transition should be very large because electrons easily screen proton currents, thus suppressing development of the instability. Failure of the Weibel instability to preempt other shock mechanisms, except for very large Alfvén Mach numbers, has been discussed in the context of nonrelativistic shocks by Blandford & Eichler (1987). Here we estimate the width of the Weibel-driven shock in the ultrarelativistic electron-proton plasma. Although based on a number of physical assumptions about the behavior of the plasma parameters at the nonlinear stage of the Weibel instability, such analytical scalings are necessary in any case because, for evident reasons, simulations of plasmas are possible only with artificially low proton-to-electron mass ratios (e.g., Frederiksen et al. [2004] took $m_p/m_e = 16$). In this paper, we present the parameters in physically motivated dimensionless form, and we believe that our assumptions could be checked by numerical simulations. Only by combining numerical simulations with analytical scalings can we achieve reliable conclusions about the properties of real shocks.

As a model for the shock formation, we consider collision of two oppositely directed plasma flows. Eventually, two diverging shocks should be formed with plasma at rest between them. However, at the initial stage, the two flows interpenetrate each other exciting turbulent electromagnetic fields. Particles are eventually thermalized by scattering off these turbulent fields. As electrons are thermalized relatively rapidly, we consider development of the Weibel instability in two proton beams propagating through relativistically hot isotropic electron gas. We estimate the proton isotropization length in such a system and conclude that the Weibel-mediated shocks are so wide that even in the interstellar medium, the shock should be formed at the scale of the Larmor radius of the proton in the background magnetic field.

This paper is organized as follows. In § 2, we find the growth rate of the proton Weibel instability. Saturation of the instability is considered in § 3. In § 4, we exploit the obtained results in order to estimate the width of the Weibel-mediated shock transition. Section 5 contains the discussion.

2. THE WEIBEL INSTABILITY

Let us consider two proton beams of equal strength propagating in opposite directions along the z-axis. For the sake of simplicity, let us adopt the waterbag distribution function,

$$F_p(\boldsymbol{p}) = \frac{1}{2\pi p_{\perp 0}^2}\left[\delta\left(p_z - p_{\parallel 0}\right) + \delta\left(p_z + p_{\parallel 0}\right)\right]\Theta\left(p_{\perp 0}^2 - p_\perp^2\right), \tag{1}$$

where $\Theta(x)$ is the Heaviside step function and $p_\perp = \left(p_x^2 + p_y^2\right)^{1/2}$. Assume that the beams propagate through an isotropic relativistically hot electron plasma with the distribution function $F_e(p)$.

This configuration is known to be unstable because a small transverse magnetic field $\boldsymbol{B} = B\hat{\boldsymbol{x}}\exp(-i\omega + iky)$ would drive the oppositely moving protons into current layers of opposite sign, which reinforce the initial field (see, e.g., Fig. 1 in Medvedev & Loeb 1999). Evolution of the electromagnetic field is governed by Maxwell's equations,

$$\nabla \times \boldsymbol{E} = -\frac{\partial \boldsymbol{B}}{\partial t}, \quad \nabla \times \boldsymbol{B} = \frac{\partial \boldsymbol{E}}{\partial t} + 4\pi \boldsymbol{j}, \tag{2}$$

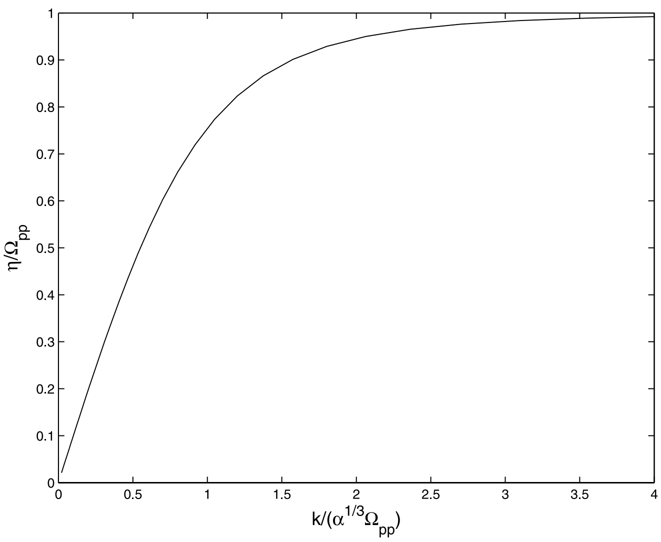

FIG. 1.—Growth rate of the Weibel instability of cold proton beams.

which are written in Fourier space as

$$kE = \omega B, \quad i(\omega E - kB) = 4\pi j. \tag{3}$$

Note that only z-components of $\boldsymbol{E}$ and $\boldsymbol{j}$ are present in this configuration. The current density, j, can be found from a solution to the linearized Vlasov equation

$$i(\omega - kv_y)\delta F_{p,e} = \pm e\left[(E - v_y B)\frac{\partial F_{p,e}}{\partial p_z} + v_z B\frac{\partial F_{p,e}}{\partial p_y}\right] \tag{4}$$

as

$$j = en\int v_z(\delta F_p - \delta F_e)\, d\boldsymbol{p}. \tag{5}$$

As usual (e.g., Krall & Trivelpiece 1973), the condition for the set of equations (3)–(5) to have a nonzero solution yields the dispersion equation

$$k^2 = \omega^2\left[1 + \chi_e(\omega, k) + \chi_p(\omega, k)\right], \tag{6}$$

where

$$\chi_\alpha = \frac{4\pi e^2 n}{\omega^2}\int v_z\left(\frac{\partial F_\alpha}{\partial p_z} + \frac{kv_z}{\omega + i0 - kv_y}\frac{\partial F_\alpha}{\partial p_y}\right) d\boldsymbol{p} \tag{7}$$

is the susceptibility of the plasma species α.

Substituting the proton distribution function equation (1) into equation (7) yields

$$\chi_p(\omega, k) = -\frac{2\Omega_{pp}^2}{\omega^2}\left[\frac{1}{m_p\gamma + \sqrt{m_p^2\gamma^2 - p_{\perp 0}^2}} \times\left(m_p\gamma - \frac{p_{\parallel 0}^2}{\sqrt{m_p^2\gamma^2 - p_{\perp 0}^2}}\right) + \frac{p_{\parallel 0}^2}{p_{\perp 0}^2}\left(\frac{\omega m_p\gamma}{\sqrt{\omega^2 m_p^2\gamma^2 - k^2 p_{\perp 0}^2}} - 1\right)\right], \tag{8}$$

where $\Omega_{pp} \equiv (4\pi e^2/m_p\gamma)^{1/2}$ is the relativistic proton plasma frequency and $\gamma = [1 + (p_{\perp 0}^2 + p_{\parallel 0}^2)/m_p^2]^{1/2}$. Below only the strongly anisotropic, highly relativistic case is considered, $p_\perp \ll p_\parallel$ and $\gamma \gg 1$. Then one can neglect the first term in the square brackets.

The susceptibility of isotropic electrons is written as

$$\begin{aligned}\chi_e(\omega, k) &= \frac{4\pi e^2 n}{\omega}\int \frac{v_z}{\omega + i0 - kv_y}\frac{\partial F_e}{\partial p_z}\, d\boldsymbol{p} \\ &= \frac{4\pi e^2 n}{\omega}\int \frac{v\sin^2\theta\cos^2\varphi}{\omega + i0 - kv\cos\theta}\frac{dF_e}{dp}p^2\, dp\, d\cos\theta\, d\varphi \\ &= \frac{4\pi^2 e^2 n}{\omega}\left[\wp\int\frac{v(1 - x^2)}{\omega - kvx}\frac{dF_e}{dp}p^2\, dx\, dp\right. \\ &\qquad \left. - \frac{\pi i}{k}\int\left(1 - \frac{\omega^2}{k^2 v^2}\right)p^2\frac{dF_e}{dp}\, dp\right].\end{aligned} \tag{9}$$

Here we used the spherical coordinates in the momentum space and the Plemeli formula. The Weibel instability operates in the low-frequency limit, $\omega \ll k$. In this case the imaginary part of χ_e dominates because at $\omega \to 0$, the principal value of the integral in x goes to zero. The physical reason is that only electrons with small v_y contribute to the susceptibility because other electrons "see" a rapidly oscillating field as they move in the y-direction over a distance larger than $1/k$ for the time $\sim 1/\omega$. Now one can write

$$\chi_e(\omega, k) = i\frac{\pi\Omega_{pe}^2}{4k\omega}, \tag{10}$$

where $\Omega_{pe}^2 = 4\pi e^2 n/T$, and $1/T \equiv 8\pi\int F_e p\, dp$. The parameter T is equal to the electron temperature if the electron spectrum is Maxwellian and $T \gg m_e$. Note that in their analysis of the proton Weibel instability, Wiersma & Achterberg (2004) erroneously used the expression $\chi_e = -\Omega_{pe}^2/\omega^2$, which is valid only in the high-frequency limit, $\omega \gg k$, and therefore is irrelevant to the case of interest.

Now one can write the dispersion equation (6) in the low-frequency limit:

$$k^2 + \frac{2\Omega_{pp}^2 p_{\parallel 0}^2}{p_{\perp 0}^2}\left(\frac{\omega m_p\gamma}{\sqrt{\omega^2 m_p^2\gamma^2 - k^2 p_{\perp 0}^2}} - 1\right) - i\frac{\pi\Omega_{pe}^2\omega}{4k} = 0. \tag{11}$$

In the limit of negligible angular spread of the proton beams, $p_{\perp 0} \to 0$, the dispersion equation is reduced to a simple cubic equation,

$$\frac{\alpha}{\kappa^3}x^3 + x^2 - 1 = 0, \tag{12}$$

where $\alpha \equiv \pi\Omega_{pp}^2/4\Omega_{pe}^2 = \pi m_p\gamma/4T$, $\kappa \equiv k/\Omega_{pp}$, and $x \equiv -i\omega/\Omega_{pp}$. The system is unstable provided $\mathrm{Re}\,(x) > 0$. In the small- and long-wavelength limits, the growth rate $\eta \equiv \mathrm{Im}\,(\omega) = x\Omega_{pp}$ is found as

$$\eta = \begin{cases} k\alpha^{-1/3}, & k \ll \alpha^{1/3}\Omega_{pp}, \\ \Omega_{pp}, & k \gg \alpha^{1/3}\Omega_{pp}. \end{cases} \tag{13}$$

The full solution to equation (12) is shown in Figure 1. One can see that if $T \sim \gamma m_e$, as one can expect within the shock structure

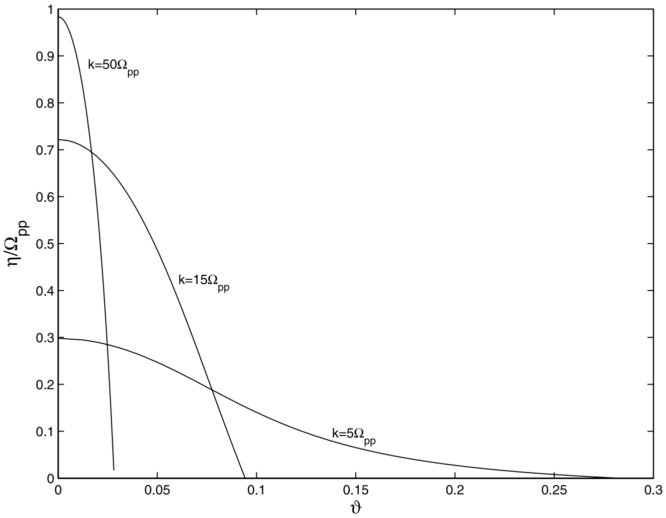

FIG. 2.—Dependence of the growth rate on the angular spread of the proton beams $\vartheta \equiv p_{\perp 0}/p_{\parallel 0}$ at $T = m_e\gamma/3$.

(see discussion in § 4), the most unstable are short-wave perturbations, $k \gtrsim (m_p/m_e)^{1/3}\Omega_{pp}$.

When the angular spread of the beams increases, the growth rate of the instability decreases (Fig. 2). The threshold of the instability can be easily found by substituting $\omega = 0$ into equation (11); this yields

$$\frac{p_{\perp 0}}{p_{\parallel 0}} = \frac{\sqrt{2}\Omega_{pp}}{k}. \tag{14}$$

So at small wavelengths, where the growth rate is maximal, the instability stops after a small spread in the angular velocity distribution is achieved.

3. STABILIZATION OF THE WEIBEL INSTABILITY

The dispersion relation equation (6) was found assuming that the particles are nonmagnetized; i.e., their trajectories are nearly straight. The instability saturates when this condition is violated either for protons or for electrons. As a result of the instability, current filaments are formed along the direction of the proton motion. The magnetic field forms a sort of cocoon around these filaments and eventually traps protons within the filaments; then the instability stops (Yang et al. 1994; Wiersma & Achterberg 2004). The quiver motion of a proton within the current filament can be described by the linearized equation

$$m_p\gamma\frac{d^2\xi}{dt^2} = ev_zB(\xi, t), \tag{15}$$

where ξ is the proton displacement in the transverse direction. Near the axis of the filament, the magnetic field can be written as $B \approx k\xi B_0$, where $B_0 \propto \exp(\eta t)$ is the amplitude of the perturbation. Then equation (15) describes oscillations in the transverse direction with the frequency $\omega_0 = \left(eB_0k/m_p\gamma\right)^{1/2}$ and the (growing) amplitude $\xi = eB_0/(\gamma m_p\eta^2)$. The proton is trapped within the filament when the oscillation amplitude becomes less than $1/k$ or, which is the same, when the frequency ω_0 exceeds the growth rate of the instability, η. This occurs when the magnetic field reaches the value

$$B_{\rm trap} = \frac{m_p\gamma\eta^2}{ek} = \frac{m_p\gamma}{e}\times\begin{cases} k\alpha^{-2/3}, & k \ll \alpha^{1/3}\Omega_{pp}, \\ \Omega_{pp}^2/k, & k \gg \alpha^{1/3}\Omega_{pp}. \end{cases} \tag{16}$$

One can consider electrons as nonmagnetized if their Larmor radius exceeds the characteristic scale of the unstable perturbation (Moiseev & Sagdeev 1963; Sagdeev 1966). This condition is violated when the field reaches the value

$$B_{\rm fr} = \frac{Tk}{e}. \tag{17}$$

Then the magnetic field becomes frozen into the electrons and the magnetic flux ceases growing.

The obtained limits on the magnetic field are shown in Figure 3. The instability develops until the magnetic field $B_{\rm sat} = \min(B_{\rm fr}, B_{\rm trap})$ is reached. One can see from Figure 3 that the maximal field is achieved in perturbations with the wavenumber

$$k_0 = \alpha^{1/2}\Omega_{pp} = \sqrt{\frac{\pi}{4}}\Omega_{pe}; \tag{18}$$

the corresponding wavelength is of the order of the inertial length of electrons. The energy of the generated field scales as the energy of electrons:

$$\frac{B^2}{8\pi} = \frac{1}{8\pi\alpha}\left(\frac{m_p\gamma\Omega_{pp}}{e}\right)^2 = \frac{2}{\pi}nT. \tag{19}$$

Within the shock structure, one can conveniently normalize the electron temperature as

$$T = \frac{\tau m_e\gamma}{3}, \tag{20}$$

where τ is a dimensionless parameter; $\tau = 1$ means that the average electron energy remains the same as upstream of the shock. Now one can estimate the fraction of the upstream kinetic energy transformed into the energy of the magnetic field as

$$\epsilon_B = \frac{2\tau m_e}{3\pi m_p}. \tag{21}$$

4. THE WIDTH OF THE SHOCK WAVE

When two oppositely directed plasma streams collide, the Weibel instability generates small-scale magnetic fields; particle scattering off these magnetic fluctuations provides an isotropization mechanism necessary for the shock transition to form. The electron streaming is halted easily, whereas protons still plow on through an

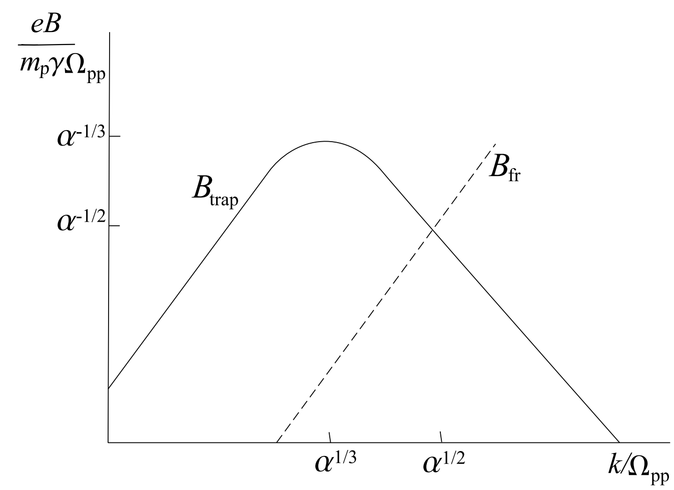

FIG. 3.—Saturation magnetic field as a function of the wavenumber. The limit $B_{\rm trap}$ due to the proton trapping is shown by the solid line; the limit $B_{\rm fr}$ when the field becomes frozen into electrons is shown by the dashed line.

isotropic electron gas. The shock is formed on the scale defined by slow diffusion of protons in the momentum space (Moiseev & Sagdeev 1963; Sagdeev 1966).

Let us denote the transverse scale of the magnetic inhomogeneities by d and the amplitude of magnetic fluctuations by B. The proton is scattered over a characteristic correlation length by the angle

$$\delta\theta = \frac{eBd}{\sin(\theta) m_p \gamma}, \tag{22}$$

where θ is the angle the proton makes with the flow direction. Here we take into account the fact that the Weibel instability generates strongly elongated current filaments, so the proton passes the distance $l = d/\sin\theta$ within the same filament (of course l and $\delta\theta$ remain finite at $\theta \to 0$; however, we are interested in the isotropization scale, which is determined by $\theta \sim 1$). The angular diffusion coefficient is estimated as

$$D = \frac{(\delta\theta)^2}{l} = \frac{e^2 B^2 d}{\sin(\theta) m_p^2 \gamma^2}, \tag{23}$$

which yields the isotropization scale

$$L = \frac{1}{d}\left(\frac{m_p \gamma}{eB}\right)^2. \tag{24}$$

Motivated by the estimates expressed in equations (18), (19), and (20), we normalize the characteristic inhomogeneity scale, d, and turbulent magnetic field amplitude, B, as

$$B = 4\xi\sqrt{\frac{\tau m_e \gamma}{3}}, \quad d = \frac{\zeta}{\alpha^{1/2}\Omega_{pp}}, \tag{25}$$

where τ, ξ, and ζ are dimensionless parameters. Now the shock width can be expressed as

$$L = \left(\frac{3\pi m_p}{4\tau m_e}\right)^{3/2} \frac{1}{\xi^2 \zeta \Omega_{pp}}. \tag{26}$$

The estimated shock width, L, is thus very large compared to the proton inertial length $1/\Omega_{pp}$ assuming that the electron temperature as well as the scale and amplitude of the generated magnetic field do not significantly exceed their fiducial values, i.e., $\tau \sim \xi \sim \zeta \sim 1$. We now explain why this is expected to be the case.

If $\tau \sim m_p/m_e \gg 1$, then although the shock width could then be brought down to the proton skin depth, this would beg the question of how the electrons are heated, which is merely passing along the question of a shock mechanism. Hededal (2005) found in his 2.5-dimensional (2.5D) simulations that the electrons are heated almost to equipartition with the ions, but he simulated collision of the mildly relativistic flows (the relative Lorentz factor was only 3). In the non- or mildly relativistic case, the longitudinal two-stream instability excites electrostatic oscillations that could result in significant plasma heating. It is not evident even in this case whether the electron temperature could rise about 1000 times, which is necessary in order to reach equipartition in the real plasma (in Hededal's simulations, $m_p/m_e = 16$, so the growth of the electron temperature by a factor of a few is sufficient to achieve equipartition). We do not consider this heating mechanism because we are interested in the highly relativistic case in which the transverse Weibel instability dominates the electrostatic one (e.g., Jaroschek et al. [2005] simulated colliding streams in a wide range of Lorentz factors and demonstrated that in the highly relativistic case, the instability becomes purely transverse). Such an instability generates quasi-static magnetic fluctuations, and we believe that the electron scattering off these fluctuations could not provide significant heating. Hededal et al. (2004) found in their 3D simulations that a radial electric field appears around the ion filaments and that this results in a specific mechanism of acceleration of electrons. However, the spectrum of the accelerated electrons is steep (the spectral slope is 2.7), so this process does not affect most of electrons and does not influence the characteristic electron "temperature."

Similarly, we expect $\xi \lesssim 1$ because we know of no reason to expect that the magnetic field would grow above the limit (19). On the contrary, such a small-scale magnetic field should decay (Gruzinov 2001). There is evidence for hierarchical merging of current filaments (Silva et al. 2003; Frederiksen et al. 2004; Medvedev et al. 2005; Kato 2005), so ζ might exceed unity. However, at the scale larger than the electron Larmor radius, which is of the order of $\sim 1/(\alpha^{1/2}\Omega_{pp})$, the field is already frozen into the electron gas; therefore, ζ could hardly ever grow significantly beyond unity. One should also take into account that the current filaments are unstable to a kink-like mode (Milosavljević & Nakar 2006; such an instability is actually observed in simulations by Frederiksen et al. (2004) and Hededal (2005), which also stimulates the field decay. Therefore, we believe that there is no reason for the Weibel-driven shock transition to be significantly narrower than equation (26) predicts. There is, however, reason to suspect that the transition is even more gradual than predicted by equation (26); this is decay of the small-scale magnetic field.

The highest resolution published simulations of shell collisions with $m_p/m_e = 16$ (Frederiksen et al. 2004) do show that while electrons are readily isotropized, the proton beams achieve only a small angular spread when passing the simulation box of the length $40/\Omega_{pp}$. The experiment duration, $120/\Omega_{pp}$, was 3 times larger than the particle crossing time; by the end of the simulations the spatial wavelength of the magnetic fluctuations achieved one-half of the width of the simulation box. In physical units, this is written as $\lambda = 20/\Omega_{pe}$, $k = 0.31\Omega_{pe}$, so that $\zeta = 3$. Although one cannot exclude the possibility that the pattern growth was frustrated by the finite size of the simulation box, we believe, for the reasons outlined above, that ζ would not grow significantly in any case.

According to estimate (26), the full shock transition is too wide to be simulated numerically even with a moderate mass ratio $m_p/m_e > 10$. On the other hand, scaling (26) may hardly ever be applied to the case $m_p/m_e \lesssim 10$ because it was obtained under the assumption that the proton and electron scales are well separated, i.e., that $(m_p/m_e)^{1/2} \gg 1$. Therefore, a direct numerical check of this scaling is very difficult. Nevertheless, it would be very useful to follow the behavior of the parameters τ, ξ, and ζ in numerical simulations even with a low mass ratio. Even 2.5D simulations of the proton Weibel instability in the isotropic electron gas would clarify the behavior of these parameters in the highly nonlinear regime.

5. DISCUSSION

Estimate (26) was obtained under the assumption that there is no magnetic field in the upstream flow. If the flow is magnetized, a shock transition may be formed at the scale of the proton Larmor radius; therefore, the above estimates are valid only if $eB_0L/m_p\gamma < 1$, where B_0 is the magnetic field in the upstream flow. One can conveniently characterize the magnetization of the flow by the parameter $\sigma = B_0^2/(4\pi m_p \gamma n)$. Making use of equation (26), one finds that the shock may be driven by the Weibel instability if

$$\sigma < \xi^4 \zeta^2 \left(\frac{4\tau m_e}{3\pi m_p}\right)^3 = 1.5 \times 10^{-11} \xi^4 \tau^3 \zeta^2. \qquad (27)$$

If the shock propagates through the interstellar medium, the magnetization exceeds the right-hand side of equation (27) by factor of about 30: $\sigma = 5 \times 10^{-10} B_{-5.5}^2 n^{-1}$, where $B = 10^{-5.5} B_{-5.5}$ G is the interstellar magnetic field and n cm^{-3} the gas number density. This suggests that forward shocks that presumably produce GRB afterglows are mediated not by the Weibel instability but by the Larmor rotation of protons in the background magnetic field. This does not mean that the Weibel instability does not work at all. On the contrary, it does develop and may well create some small-scale magnetic field that is much stronger than the background field. The scattering off these magnetic fluctuations results in diffusion of the protons in angles. However, because the scattering does not manage to isotropize the protons at the scale less than the Larmor radius, it is hard to see how the Weibel instability could convert the kinetic energy to another form. The strong dependence of estimate (21) on the parameters ξ, τ, and ζ makes accurate determination of their values, presumably by simulations, crucial to solidify this conclusion.

Estimate (27) shows that a fraction $\epsilon_B \sim 10^{-4}$ of the total energy is converted into magnetic energy unless the electrons are heated additionally within the shock structure. The generated small-scale field should decay (Gruzinov 2001; Milosavljević & Nakar 2006), so ϵ_B could be even lower. According to Panaitescu & Kumar (2002) and Yost et al. (2003), the observed spectra and light curves of the GRB afterglows imply $\epsilon_B \sim 10^{-3}$–10^{-1} in most cases. Eichler & Waxman (2005) demonstrated that the above estimates can be rescaled such that the observations are fitted with values of ϵ_B that are smaller by an arbitrary factor f, $m_e/m_p < f < 1$. Taking this into account one can see that the Weibel instability could provide the necessary field unless the field decay is too strong. On the other hand, as the global structure of the shock transition is dictated by the Larmor rotation of the protons in the background field, some new physics could come into play.

A presumably important physical mechanism is the synchrotron maser instability at the shock front. This instability generates semicoherent, low-frequency electromagnetic waves (Gallant et al. 1992; Hoshino et al. 1992; Y. Lyubarsky 2006, in preparation). In low-magnetized flows, the amplitude of these waves exceeds the amplitude of the shock compressed background field. In this case, relativistic particles radiate in the field of the waves via nonlinear Compton scattering (e.g., Melrose 1980, p. 136). The power and characteristic frequencies of this emission are similar to those for synchrotron emission in a magnetic field of the same strength; therefore, the observed GRB afterglows could be attributed to the nonlinear Compton scattering off the electromagnetic waves generated by the synchrotron maser instability at the shock front. It is beyond the scope of this paper to redo afterglow theory with the spectrum of low-frequency electromagnetic waves that is expected from this instability.

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