

THE COSMIC-RAY PRECURSOR OF RELATIVISTIC COLLISIONLESS SHOCKS: A MISSING LINK IN GAMMA-RAY BURST AFTERGLOWS

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ABSTRACT

Collisionless shocks are commonly argued to be the sites of cosmic-ray (CR) acceleration. We study the influence of CRs on weakly magnetized relativistic collisionless shocks and apply our results to external shocks in gamma-ray burst (GRB) afterglows. The common view is that the transverse Weibel instability (TWI) generates a small-scale magnetic field that facilitates collisional coupling and thermalization in the shock transition. The TWI field is expected to decay rapidly, over a finite number of proton plasma skin depths from the transition. However, the synchrotron emission in GRB afterglows suggests that a strong and persistent magnetic field is present in the plasma that crosses the shock; the origin of this field is a key open question. Here we suggest that the common picture involving TWI demands revision. Namely, the CRs drive turbulence in the shock upstream on scales much larger than the skin depth. This turbulence generates a large-scale magnetic field that quenches TWI and produces a magnetized shock. The new field efficiently confines CRs and enhances the acceleration efficiency. The CRs modify the shocks in GRB afterglows at least while they remain relativistic. The origin of the magnetic field that gives rise to the synchrotron emission is plausibly in the CR-driven turbulence. We do not expect ultra-high-energy cosmic-ray production in external GRB shocks.

Subject headings: acceleration of particles — cosmic rays — gamma rays: bursts — plasmas — shock waves

1. INTRODUCTION

Collisionless shocks are observed on all astrophysical scales. The diffusive shock acceleration (DSA) mechanism is believed to accelerate cosmic rays (CRs) in these shocks (Bell 1978; Blandford & Ostriker 1978; Blandford & Eichler 1987). The CRs can carry a substantial fraction of the energy of the shock, and thus the CR pressure can influence the structure of the shock (Eichler 1979; Blandford 1980; Drury & Völk 1981; Ellison & Eichler 1984). This picture has recently received support from X-ray observations of supernova (SN) remnants (Warren et al. 2005). Recently, Bell (2004, 2005) has shown that CRs in SNe can drive turbulence and amplify magnetic fields in the shock upstream, and Dar & de Rújula (2005) speculate that relativistic jets could do the same.

The current lore is that weakly magnetized relativistic collisionless shocks are mediated by the transverse Weibel instability (TWI; Weibel 1959; Fried 1959). TWI produces a magnetic field near equipartition that provides collisional coupling in the shock transition layer (Gruzinov & Waxman 1999; Medvedev & Loeb 1999). TWI is universally observed in two- and three-dimensional particle-in-cell (PIC) simulations of unmagnetized colliding shells (e.g., Lee & Lampe 1973; Gruzinov 2001; Silva et al. 2003; Frederiksen et al. 2004; Jaroschek et al. 2004; Nishikawa et al. 2005; Medvedev et al. 2005; Kato 2005). The resulting field is small-scale, of the order of the proton plasma skin depth λ_s .

The role of TWI in particle acceleration and in the downstream thermodynamics is controversial. The field is expected to decay rapidly, within a few λ_s from the shock transition (Gruzinov 2001; Milosavljević & Nakar 2006). This decay is evident in three-dimensional PIC simulations (Spitkovsky 2005). Although there are claims that the decay saturates at distances of $\sim 100\lambda_s$ from the shock (Silva et al. 2003; Medvedev et al. 2005), survival of the field over larger distances has not been demonstrated.

According to the popular model of gamma-ray bursts (GRBs; e.g., Piran 2005a and references therein), the afterglow originates in a relativistic blast wave that propagates into an ambient e^-p plasma. The afterglow emission is ascribed to synchrotron radiation from nonthermal electrons that gyrate in the shock-generated magnetic field. Detailed studies of GRB spectra and light curves (Panaitescu & Kumar 2002; Yost et al. 2003) have shown that the magnetic energy density in the emitting region (the downstream shocked plasma) is a fraction $\epsilon_B \sim 10^{-2}$ to 10^{-3} of the internal energy density.³ This field must persist over at least a few percent of the width of the blast wave (Rossi & Rees 2003). Its origin has remained a key open question.

Compressional amplification of the weak preexisting magnetic field of the interstellar medium (ISM) yields $\epsilon_B \sim 10^{-9}$ (Gruzinov 2001) and does not explain the field in the emitting region. If TWI does develop in the transition layer, it generates a strong field in the vicinity of the shock transition. However, as discussed above, there is no evidence that it can persist over the required $\sim 10^9\lambda_s$ from the shock (e.g., Piran 2005b and references therein).

X-ray observations of GRB afterglows are modeled as an optically thin synchrotron spectrum requiring nonthermal electrons with Lorentz factors as high as $\sim 10^6$. The observations indicate that electrons, and thus protons as well, are efficiently accelerated in the shock to produce a hard power-law spectrum. DSA can achieve such Lorentz factors if the circumburst medium is magnetized at $\sim 1 \mu\text{G}$ level, as expected in the ISM. The weak magnetic field, however, does not directly affect the structure of the shock (e.g., TWI can still develop in the shock transition).

It seems that a self-consistent picture of relativistic astrophysical shocks includes high-energy particles (CRs) and at least a weakly magnetized upstream. Here we explore the interaction between these components. In § 2, we derive the conditions under which the CRs drive turbulence in the upstream; this turbulence generates a large-scale magnetic field that increases the acceleration

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³ Under some circumstances ϵ_B as low as $\sim 10^{-5}$ can fit the data (Eichler & Waxman 2005).

efficiency.⁴ In § 3, we argue that this mechanism is a candidate solution of the origin of the field inferred from the afterglow emission in external GRB shocks. In § 4, we discuss implications for TWI and for the production of ultra-high-energy cosmic rays (UHECRs).

2. COSMIC RAY–MAGNETIC FIELD INTERACTION

2.1. Cosmic-Ray Trajectories and the Return Current

Consider a relativistic shock wave with Lorentz factor Γ that accelerates particles to Lorentz factors $\gamma \gg \Gamma$. The accelerated particles are roughly isotropic in the downstream frame; a fraction of them are moving ahead of the shock. In the upstream frame, the velocities of the accelerated particles are directed within an angle $\theta \sim \sqrt{2}/\Gamma$ of the shock normal.⁵ If the upstream contains a magnetic field \mathbf{B} , the field deflects charged particles moving ahead of the shock. When the velocity of a particle is deflected to an angle $\sim 1/\Gamma$ from the shock normal, the shock catches up and reabsorbs the particle.

The orbit of a particle in the upstream field can be quasi-circular or diffusive, depending on how the radial distance X the particle traverses ahead of the shock compares to the relevant magnetic correlation length λ . The typical angle by which the magnetic field deflects the particle is $\Delta\theta \sim X/r_g$ when the orbit is quasi-circular and $\Delta\theta \sim (\lambda X)^{1/2}/r_g$ when the orbit is diffusive, where

$$r_g = \frac{\gamma mc^2}{eB(\lambda)} \quad (1)$$

is the gyroradius of the particle expressed in terms of the magnetic field $B(\lambda)$ on scales λ . Reabsorption in the shock occurs when $\Delta\theta \sim 1/\Gamma$, and thus the distance traversed by the particle in the quasi-circular and diffusive regime is

$$X \sim \begin{cases} r_g/\Gamma, & X \lesssim \lambda, \\ (r_g/\Gamma)^2/\lambda, & X \gg \lambda. \end{cases} \quad (2)$$

In spherical blast waves, a particle can be accelerated by repeated shock crossing if the time separating its emission and reabsorption in the shock $\sim X/c$ is smaller than the expansion time of the shock $\sim R/c$, where R is the shock radius.⁶ Therefore, $X \leq R$ is required for acceleration. The maximum distance that the particle reaches from the moving shock transition is substantially smaller,

$$\Delta \sim \frac{X}{\Gamma^2}, \quad (3)$$

because the particle and the shock are both moving at about the speed of light. Another requirement for acceleration is that the cooling time be longer than the acceleration time (see § 3).

In general, protons and electrons are both emitted into the shock upstream. However, the synchrotron and self-Compton cooling inhibit the acceleration of electrons, so the protons, which we refer to as the CR precursor of the shock, can reach higher energies and travel farther from the shock into the upstream. The upstream fluid sees a net positive CR current $|\mathbf{J}_{\text{CR}}| \sim ev_s n_{\text{CR}}$ in the region where

⁴ Gruzinov (2001) and Waxman (2004) speculate that CRs could amplify large-scale magnetic fields in external GRB shocks.

⁵ Unless otherwise noted, quantities are evaluated in the upstream frame.

⁶ It is also implicitly assumed that some fraction of accelerated particles can diffuse back to the shock after they find themselves in the downstream.

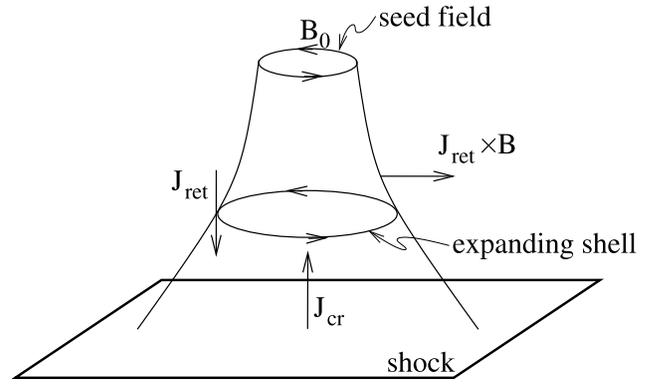


FIG. 1.—Schematic representation of the reaction of the upstream plasma to the CR streaming ahead of the shock. The figure is in the frame of the shock so that the upstream plasma moves vertically downward. An initially weak magnetic field loop is stretched by the Ampère force $(\mathbf{J}_{\text{ret}} \times \mathbf{B})/c$ and plasma is accelerated sideways as it flows toward the shock front.

it overlaps with the precursor. Here v_s and n_{CR} are the CR velocity and density. In the upstream region that is not accelerated to relativistic velocities by the CR pressure, the CR velocity is about the speed of light, $v_s \sim c$; thus, we have

$$J_{\text{CR}} \sim ecn_{\text{CR}}. \quad (4)$$

To maintain overall charge neutrality, the upstream responds to the CR current by supporting an opposite “return current” \mathbf{J}_{ret} that short-circuits the electric field associated with the charge separation in the CR precursor.⁷ If n_{CR} is smaller than the upstream density, the return current approximately balances the CR current,

$$J_{\text{ret}} \sim J_{\text{CR}}. \quad (5)$$

If the upstream contains a weak magnetic field, the return current couples with this field; next we explore the consequences of this coupling.

2.2. The Driving of Turbulence and Field Amplification

Here we argue that the return current drives turbulence in the upstream and suggest that this turbulence leads to magnetic field growth on scales much larger than the proton plasma skin depth. One type of interaction between the CR precursor and the magnetic field of the upstream, explored by Bell (2004, 2005) in the context of Newtonian shocks, is the Ampère force

$$\mathbf{F} = \frac{\mathbf{J}_{\text{ret}} \times \mathbf{B}}{c}, \quad (6)$$

which accelerates the upstream perpendicular to the shock velocity. We assume that the initial weak magnetic field has power on all scales, as expected in interstellar turbulence, and that the power on relevant scales is roughly scale-independent.

Consider a loop of radius λ of a weak initial magnetic field \mathbf{B}_0 parallel to the shock transition that is directed clockwise as seen from the shock, as shown in Figure 1, and assume that the upstream is initially stationary. The Ampère force accelerates the fluid away from the center of the loop. Bell (2004, 2005) carried out magnetohydrodynamic (MHD) simulations of this process; we base our estimates on his results.

⁷ The return current is established on the plasma time $\lesssim 10$ ms, which is instantaneous compared to the light crossing time of the CR precursor.

The dynamics of the upstream fluid can be approximated using MHD, and thus flux freezing during the expansion of the loop implies that

$$\frac{B_\theta}{\rho r} = \text{constant}, \quad (7)$$

where B_θ and ρ are the azimuthal magnetic field and the density at radius r , respectively. If the expansion is nonrelativistic and if the thermal and magnetic pressure are ignored so that only the Ampère force influences the motion, the radius of the loop accelerates according to

$$\frac{d^2 r}{dt^2} \sim \frac{J_{\text{ret}} B_\theta}{\rho c} \sim \frac{J_{\text{CR}} B_0}{\lambda \rho_0 c} r. \quad (8)$$

Therefore, the loop expands exponentially $r(t) \sim \lambda e^{\sigma t}$ at a rate

$$\sigma \sim \left(\frac{J_{\text{CR}} B_0}{\lambda \rho_0 c} \right)^{1/2}, \quad (9)$$

implying a velocity of expansion of $dr/dt \sim r\sigma$.

The upstream accelerates until the pressure overtakes the Ampère force (we assume that the expansion remains subluminal at this point). In a realistic environment where the magnetic field fluctuates on all scales, we expect the expansion on each scale to saturate when neighboring magnetic “rings” on similar scales collide, namely, when $r \sim 2\lambda$. This happens when

$$\sigma t \sim \frac{\sigma \Delta}{c} \sim 1. \quad (10)$$

Compressional amplification of the magnetic field at this point is of the order of unity. However, it is commonly argued (Kulsrud 2005 and references therein) that a turbulent dynamo produces magnetic energy near equipartition with the turbulent energy. Therefore, the upstream fluid that is exposed to a current J_{CR} over a time t will develop turbulent motion on all scales

$$\lambda \lesssim \lambda_{\text{max}} \sim \frac{J_{\text{CR}} B_0 t^2}{\rho_0 c}. \quad (11)$$

Each eddy amplifies the magnetic energy on its own scale on an e -folding time $1/\sigma(\lambda)$. The bulk of the kinetic energy initially stored in expanding shells $\sim \frac{1}{2}(\sigma \lambda_{\text{max}})^2$ turns into turbulent motion and the magnetic field on scales $\sim \lambda_{\text{max}}$. The field in equipartition with the energy in turbulent eddies equals

$$B_1 \sim \left(\frac{4\pi \lambda J_{\text{CR}} B_0}{c \rho_0} \right)^{1/2}. \quad (12)$$

This is the minimum field strength generated in the shock upstream on scales $\lambda \leq \lambda_{\text{max}}$. The interaction between the CRs and the generated field could continue to accelerate the fluid after the shells have collided once. This could result in additional turbulent driving and field growth. Then a saturation occurs when the Ampère force balances the magnetic tension

$$|\mathbf{F}| \sim \frac{B^2}{4\pi \lambda}, \quad (13)$$

which implies a limit on the final field of

$$B_2 \sim \frac{4\pi \lambda J_{\text{ret}}}{c} \sim \frac{4\pi \lambda J_{\text{CR}}}{c}. \quad (14)$$

Bell (2004), who considered Newtonian shocks, argues that such a field is generated as a consequence of CR streaming. Therefore, we can limit the generated magnetic field to lie between B_1 and B_2 .

For some λ , the above values of B formally imply superluminal motion. We do not analyze the evolution of the magnetic field after its energy approaches equipartition with the rest energy of the upstream.⁸ Our treatment is applicable at distances from the shock where the magnetic field does not reach equipartition, $B < (4\pi \rho_{\text{ISM}})^{1/2} c$.

2.3. The Cosmic-Ray Spectrum and the Quasi-Steady State

The generated magnetic field can modify the spatial profile of J_{CR} by confining CRs closer to the shock than the preexisting field. The maximum coherence length of the new field cannot exceed the maximum distance CRs reach from the shock $\sim R/\Gamma^2$. Therefore, the propagation of the most energetic CRs traversing a distance $X \sim R$ between emission and reabsorption is always diffusive in the new field. If the CR motion is dominated by the new field,⁹ the shock settles in a quasi-steady state in which to every distance $\Delta \leq R/\Gamma^2$ from the shock corresponds a γ such that CRs with Lorentz factor γ are confined within a distance Δ from the shock by the magnetic field generated by the streaming of these same CRs.

The density of the CRs at a distance Δ is

$$n_{\text{CR}}(\gamma) \sim \frac{1}{4\pi R^2 \Delta(\gamma)} \frac{dN}{d \ln \gamma}, \quad (15)$$

where $dN/d\gamma$ is the CR energy spectrum. The CR spectrum produced by DSA is typically a power law

$$\frac{dN}{d\gamma} \propto \gamma^{-p}, \quad (16)$$

in a range $\gamma_{\text{min}} \leq \gamma \leq \gamma_{\text{max}}$. We take the minimum Lorentz factor of CRs to equal (Achterberg et al. 2001 and references therein)

$$\gamma_{\text{min}} \sim \Gamma^2. \quad (17)$$

Numerical simulations of acceleration in ultrarelativistic shocks (Achterberg et al. 2001; Ellison & Double 2002; Lemoine & Pelletier 2003) agree with the spectral index $p = 20/9 \approx 2.22$ derived analytically by Keshet & Waxman (2005) for isotropic diffusion. While the true spectral index will depend on the detailed interaction between CRs and magnetic turbulence (e.g., Ellison & Double 2004; Lemoine & Revenu 2006; Niemiec & Ostrowski 2006; Lemoine et al. 2006) and may differ from this derived value, we adopt

$$p = \frac{20}{9} \quad (18)$$

⁸ Magnetic field exceeding equipartition with the rest energy of the upstream would affect the hydrodynamic profile of the shock and accelerate the upstream in the direction of shock propagation, thereby reducing v_s and J_{CR} .

⁹ The preexisting magnetic field could have power on the largest scales $\sim R$, and thus it could dominate the CR motion even if the new field is stronger.

in what follows. If E_{CR} is the total energy in CRs in the shock wave, then the CR spectrum is normalized such that

$$\int_{\gamma_{\min}}^{\gamma_{\max}} \gamma m_p c^2 \frac{dN}{d\gamma} d\gamma = E_{\text{CR}}. \quad (19)$$

We relate the Lorentz factor of the shock to the total energy in the blast wave (Blandford & McKee 1976)

$$\Gamma \approx \left(\frac{E_{\text{tot}}}{n_0 m_p c^2 R^3} \right)^{1/2}, \quad (20)$$

where $n_0 \approx \rho_0/m_p$. We assume further that a fraction

$$\epsilon_{\text{CR}} \equiv \frac{E_{\text{CR}}}{E_{\text{tot}}} \quad (21)$$

of the blast wave is stored in the CRs.

We proceed to outline the procedure by which the shock-accelerated CR Lorentz factor, the shock-generated magnetic field coherence length, and the energy density in the shock-generated magnetic field are calculated self-consistently as a function of distance from the shock Δ . In § 3, we apply this procedure to external GRB shocks.

2.4. The Self-Consistent Solution

At any distance Δ from the shock, the CR current is dominated by the least energetic CRs that reach this distance. CRs with smaller γ are confined to shorter distances from the shock. The CR current $J_{\text{CR}}(\Delta)$ is larger at smaller Δ . Note the time over which an upstream element is exposed to $J_{\text{CR}}(\Delta)$ is proportional to $\Delta \ll R$; i.e., since the cosmic rays are confined so close to the shock front in an ultrarelativistic shock, the time available for generation of the magnetic field is short. The field is generated on scales λ for which $\sigma(\lambda, \Delta)\Delta/c \gtrsim 1$.¹⁰ In calculating σ , the distance CRs reach from the shock and the CR number and current densities are all calculated using the generated field B , rather than the preexisting field B_0 . A self-consistent solution for the CR and magnetic field profile ahead of the shock shows that CRs with smaller γ generate a stronger magnetic field (because of their larger J_{CR}) on smaller scales λ (because of the shorter exposure time) and are confined by their self-generated field. The maximum CR Lorentz factor γ_{\max} and the maximum field correlation length λ_{\max} are obtained for $\Delta \sim R/\Gamma^2$ (see § 2.1).

To solve for $\gamma(\Delta)$, $\lambda(\Delta)$, and $B(\Delta)$, we proceed as follows. First, we eliminate r_g and X from equations (1), (2) (the diffusive case), and (3) to obtain

$$\lambda \Delta \sim \frac{\gamma^2 m_p^2 c^4}{\Gamma^4 e^2 B^2}. \quad (22)$$

Next, we eliminate n_{CR} , γ_{\min} , and E_{CR} from equations (4) and (15)–(21) to obtain

$$J_{\text{CR}} = \frac{\Gamma^{4/9} \epsilon_{\text{CR}} E_{\text{tot}} e}{18\pi \gamma^{11/9} m_p c R^2 \Delta}. \quad (23)$$

Next, we eliminate σ from equations (9) and (10) to obtain

$$\left(\frac{J_{\text{CR}} B_0}{\lambda \rho_0 c} \right)^{1/2} \frac{\Delta}{c} \sim 1. \quad (24)$$

Finally, we set the strength of the saturated magnetic field to B_2 where the Ampère force balances magnetic tension (see § 2.2; the results for saturation at B_1 are qualitatively the same)

$$B \sim \frac{4\pi \lambda J_{\text{CR}}}{c}. \quad (25)$$

We substitute Γ from equation (20) in equations (22) and (23). Then we substitute J_{CR} from equation (23) in equations (24) and (25). We finally solve for γ , λ , and B as a function of Δ and the parameters of the blast wave. It is convenient to express the magnetic field strength in terms of

$$\epsilon_B \equiv \frac{B^2}{4\pi \rho c^2} \quad (26)$$

and to express distance from the shock in terms of the dimensionless parameter

$$\tilde{\Delta} \equiv \frac{\Delta \Gamma^2}{R} \leq 1. \quad (27)$$

The maximum γ and λ correspond to CRs that reach farthest from the shock, i.e., to $\tilde{\Delta} = 1$.

3. GAMMA-RAY BURST AFTERGLOWS

External shocks in GRBs are the best astrophysical candidates of weakly magnetized relativistic collisionless shocks. They have Lorentz factors $\Gamma \gtrsim 100$ when the blast wave is at a radius $R \sim 10^{17}$ cm (Piran 2005a and references therein). The blast wave decelerates and becomes Newtonian at $R \sim 10^{18}$ cm. The external shocks are initially beamed in the form of a jet with a typical opening angle $\theta_j \sim 0.1$ rad. The angular size of the causally connected region in the shock is $1/\Gamma$. At early times when $\Gamma > 1/\theta_j$, the blast wave has isotropic equivalent energy $E_{\text{tot}} = 10^{52}$ – 10^{54} ergs and evolves as a spherical fragment with a Lorentz factor given by equation (20). The evolution at late times when $\Gamma \lesssim 1/\theta_j$ is poorly understood: the opening angle is expected to increase, Γ decays faster (perhaps exponentially) with R , and the isotropic equivalent energy approaches $\sim 10^{51}$ ergs.

Here we explore the interaction between the CRs and the preexisting magnetic field in the circumburst medium while the evolution of the blast wave is quasi-spherical ($\Gamma > 1/\theta_j$). We expect that our results are qualitatively applicable also when $\Gamma < 1/\theta_j$ by taking $E_{\text{tot}} \sim 10^{51}$ ergs and $R \sim 10^{18}$ cm. We derive the maximum Lorentz factor γ of CRs that reach a distance Δ from the shock, the maximum magnetic field correlation length λ at this distance, and the fraction ϵ_B of the energy density in the magnetic fields on scales λ .

Following the solution outlined in § 2.4 we obtain

$$\begin{aligned} \gamma &\sim 5 \times 10^7 \tilde{\Delta}^{0.25} \epsilon_{-1}^{0.62} B_{-6}^{0.37} E_{53}^{0.75} R_{17.5}^{-1.4} n_0^{-0.37}, \\ \lambda &\sim 10^{11} \tilde{\Delta}^{0.70} \epsilon_{-1}^{0.25} B_{-6}^{0.55} E_{53}^{-0.7} R_{17.5}^{3.04} n_0^{-0.55} \text{ cm}, \\ \epsilon_B &\sim 10^{-2} \tilde{\Delta}^{-1.21} \epsilon_{-1}^{0.99} B_{-6}^{0.19} E_{53}^{1.21} R_{17.5}^{-3.84} n_0^{-1.19}, \end{aligned} \quad (28)$$

where $\tilde{\Delta} = 1$ corresponds to the highest energy CRs that are accelerated in the shock. Here $\epsilon_{\text{CR}} = 0.1 \epsilon_{-1}$ is the fraction of the energy E_{tot} carried by CRs ahead of the shock, $E_{\text{tot}} = 10^{53} E_{53}$ ergs, $B_0 = B_{-6} \mu\text{G}$, $n_{\text{ISM}} = n_0 \text{ cm}^{-3}$, and $R = 10^{17.5} R_{17.5}$ cm. We have used $p = 20/9$, but the dependence on p is weak. Evidently, in the shocks considered here, the field

¹⁰ In reality, field saturation will require multiple e -foldings, $\sigma \Delta/c \sim$ a few.

generated in the precursor has the capacity to confine CRs with higher energies than a preexisting microgauss field.¹¹

Note that $\lambda_{\max} \gg \lambda_s \sim 2 \times 10^7 n_0^{-1/2}$ cm. The fractional magnetic energy density ϵ_B should be approximately preserved as the fluid passes the shock transition. The dependence of ϵ_B on the radius and the field correlation length is $\epsilon_B \propto R^{1.41} \lambda^{-1.73}$ (it is independent of E_{tot} if ϵ_{CR} is constant). Note that ϵ_B is larger on smaller scales, since it is generated at smaller Δ by the larger J_{CR} that is dominated by CRs with smaller γ .

Since the upstream is turbulent on scales $\lesssim \lambda_{\max}$, and since this turbulence is accompanied by density inhomogeneities, the shock transition and the fluid passing the transition are also expected to be turbulent on the same scales. The evolution of the magnetic field after it passes the transition is described by the laws governing MHD turbulence, and we do not attempt to describe it here. The value of ϵ_B in the downstream region will be that reached in the shock upstream, modified by any evolution that the turbulence experiences as it travels downstream past the shock transition.

The main premise of our treatment is that the precursor is dominated by protons and thus carries a net positive current. To compete with the protons in the distance they reach ahead of the shock, the electrons would have to have Lorentz factors $\gamma_e > \gamma_{\min} m_p/m_e \sim \Gamma^2 m_p/m_e$. Electrons with $\gamma_e \sim 10^8$ – 10^{10} in the rest frame of the upstream are cooled by the synchrotron mechanism on timescales much shorter than the age of the shock, and thus their acceleration to even higher energies is suppressed. Therefore, indeed, the shock precursor at proton Lorentz factors $\gamma \gtrsim 10^5$ is completely dominated by protons and carries a net positive current.

4. DISCUSSION

The picture presented here casts doubt on the role of TWI in collisionless shocks. It is commonly argued that TWI generates the magnetic field that facilitates collisional coupling in the shock transition. However, if the upstream is weakly magnetized, the shock accelerates CRs, and these CRs drive upstream turbulence and magnetic field generation. When the upstream fluid reaches the transition—where TWI is assumed to take place—the large-scale field generated in the precursor will plausibly quench TWI. Hededal & Nishikawa (2005) find that TWI is quenched for ratios of the electron plasma to cyclotron frequency $\omega_{pe}/\omega_{ce} < 5$. For nonrelativistic upstream electrons this implies quenching for $\epsilon_B \gtrsim 2 \times 10^{-5}$, which is expected from our analysis.¹²

Recently, a question has been raised whether external GRB shocks are sites of UHECR acceleration ($\gamma \sim 10^{10}$ – 10^{11} ; Dermer 2002; Waxman 2004). The idea that UHECRs are produced in

GRBs (Waxman 1995; Vietri 1995; Milgrom & Usov 1995) is motivated by the observation that the cosmic energy densities in UHECRs and GRB fireballs are comparable. The presumed absence of strong upstream magnetic fields in external GRB shocks has focused attention on UHECR production in shocks propagating into relativistic ejecta (Waxman 2004; Gialis & Pelletier 2005). Our analysis of external shocks indicates that the maximum Lorentz factor of CRs produced there is well below the UHECR range. Even under the most optimistic assumptions ($\epsilon_B \sim 1$ on scales $\lambda \sim R/\Gamma^2$), the resulting Lorentz factors are

$$\gamma < 3 \times 10^{10} E_{\text{tot},51}^{1/3} n_0^{1/6}, \quad (29)$$

where $E_{\text{tot}} = 10^{51} E_{\text{tot},51}$ ergs is the total, beaming-corrected energy of the blast wave. Therefore, we do not expect UHECR acceleration in external GRB shocks.

5. CONCLUSIONS

CRs accelerated in relativistic collisionless shocks excite large-scale turbulence and magnetic field generation in the shock upstream. The field generated in the shock precursor has power on scales much larger than the proton plasma skin depth. The propagation of CRs in the generated field is diffusive. In external GRB shocks, CRs are accelerated to higher energies in the generated field than in the preexisting field of the ISM. The generated field reaches equipartition with the energy density in the fluid. The shock transition is turbulent with a hydrodynamic profile dominated by CR pressure. The commonly invoked TWI is probably quenched in relativistic collisionless shocks by the magnetic field and the turbulence generated in the shock precursor. External GRB shocks do not accelerate UHECRs. PIC simulations of collisionless shocks must include a weak upstream magnetic field and simulate a spatial domain as large as the CR gyroradius in the upstream to observe the acceleration of particles beyond equipartition.

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¹¹ The maximum Lorentz factor to which CRs are accelerated in the preexisting field of the ISM is $\gamma_{\max}(B_0) \sim 4.6 \times 10^6 E_{33}^{1/2} B_{-6} n_0^{-1/2} R_{17.5}^{-1/2}$.

¹² Because of the finite transverse velocities of the CRs ($\sim c/\Gamma$ in the frame of the ISM), the CRs themselves will not excite TWI in the upstream.

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