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THE UPSTREAM MAGNETIC FIELD OF COLLISIONLESS GRB SHOCKS

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ABSTRACT

Gamma-ray burst (GRB) afterglow emission is believed to be produced by synchrotron emission of electrons accelerated to high energy by a relativistic collisionless shock propagating into a weakly magnetized plasma. Afterglow observations have been used to constrain the postshock magnetic field and structure, as well as the accelerated electron energy distribution. Here we show that X-ray afterglow observations on day timescales constrain the preshock magnetic field to satisfy $B > 0.2(n/1\ \mathrm{cm}^{-3})^{5/8}$ mG, where n is the preshock density. This suggests that either the shock propagates into a highly magnetized fast, $v \sim 10^3$ km s^{-1}, wind, or the preshock magnetic field is strongly amplified, most likely by the streaming of high-energy shock-accelerated particles. More stringent constraints may be obtained by afterglow observations at high photon energy at late, >1 day, times.

Subject headings: acceleration of particles — gamma rays: bursts — magnetic fields — shock waves

1. INTRODUCTION

Diffusive (Fermi) acceleration of charged particles in collisionless shocks is believed to be the mechanism responsible for the production of cosmic rays, as well as for the nonthermal emission from a wide variety of high-energy astrophysical sources (for reviews see, e.g., Drury 1983; Blandford & Eichler 1987; Axford 1994). A theory of collisionless shocks based on first principles is, however, lacking. One of the major issues that is not understood is magnetic field amplification. Postshock (downstream) magnetic field strengths derived from recent observations of supernova remnant shocks are significantly higher than those expected from shock compression of the preshock (upstream) plasma, suggesting a significant amplification of the preshock magnetic field beyond compression (e.g., Völk et al. 2005). This amplification is most likely intimately related to the process of particle acceleration and is not understood theoretically.

GRB afterglow shocks present an extreme example of magnetic field amplification. Phenomenological considerations suggest that the afterglow is produced by synchrotron emission by electrons accelerated to high energy in a relativistic collisionless shock driven by the GRB explosion into the medium surrounding the progenitor (see, e.g., Zhang & Mészáros 2004; Piran 2005, for reviews). The postshock magnetic field is inferred to be near equipartition. If the shock propagates into an interstellar medium with characteristic field amplitude of $\simeq 1\ \mu$G, this implies amplification of the field energy density (beyond compression) by $\sim$7 orders of magnitude (Gruzinov & Waxman 1999; Gruzinov 2001). While the amplification of the field to near equipartition by electromagnetic plasma instabilities appears likely (e.g., Gruzinov & Waxman 1999; Medvedev & Loeb 1999), such instabilities tend to create a field varying on a plasma skin depth scale, c/ω_p, which is expected to decay in the postshock plasma at a distance of a few skin depths away from the shock. Observations indicate, however, that the field survives over a scale many orders of magnitude larger than c/ω_p. This suggests that the characteristic scale of field variation grows to values much larger than c/ω_p, and the challenge is thus to explain the formation of equipartition field on a scale $\gg c/\omega_p$ (Gruzinov & Waxman 1999; Gruzinov 2001). Various groups have recently attempted to address this challenge using numerical plasma simulations (Silva et al. 2003; Frederiksen et al. 2004; Jaroschek et al. 2004; Nishikawa et al. 2006; Medvedev et al. 2005; Kato 2005; Spitkovsky 2006). Since the calculations are extremely demanding numerically, the simulation boxes are typically only a few tens of skin depths wide, and a clear picture of field length scale growth has not yet emerged.

GRB afterglows provide a unique opportunity for diagnosing collisionless shock physics, as they allow one to observe a rapid evolution of the synchrotron spectrum over a wide span of wavelengths and for a wide range of shock Lorentz factors. Afterglow observations were used to constrain the downstream field and the accelerated electron energy distribution, at both high (Waxman 1997a; Freedman & Waxman 2001) and low (Waxman 1997b; Eichler & Waxman 2005) energy. Here we point out that afterglow observations also provide constraints on the upstream field. A characterization of the upstream field provides information not only on the circumburst medium, but also on the collisionless shock physics. Amplification of magnetic field fluctuations in the preshock plasma are naturally expected within the frame work of diffusive (Fermi) shock acceleration of particles (Bell 1978; Blandford & Ostriker 1978), as the high-energy particles stream ahead of the shock. Some evidence for such enhancement has been obtained from radio observations of supernova remnants (e.g., Achterberg et al. 1994). Recently, it has been proposed that the streaming of high-energy particles in nonrelativistic collisionless shocks may significantly amplify not only the fluctuations in the magnetic field, but also the (overall) amplitude of upstream field strength (Bell 2004). If such amplification is indeed achieved, it would allow acceleration of particles to higher energy. For supernovae, this may allow acceleration of particles beyond 10^{15} eV, and for GRBs it may allow acceleration to ultrahigh ($>10^{19}$ eV) energy in the afterglow shock (and not only in the internal shocks see Waxman 2001, for review).

We show in § 2 that X-ray afterglow observations may be used to put constraints on the upstream field amplitude. In § 3 we apply the results derived in § 2 to afterglow observations and derive lower limits on the upstream magnetic field for several GRBs. In § 4 we summarize our results and discuss their implications, including the implications for numerical simulations and for recent discussions of upstream magnetic fields in relativistic GRB afterglow shocks (Lyubarsky & Eichler 2006; Milosavljević & Nakar 2005).

2. SHOCK ACCELERATION AND MAXIMUM SYNCHROTRON ENERGY

Within the diffusive shock acceleration framework, high-energy particles cross the shock front multiple times, and gradually gain energy as in each crossing (from upstream to downstream or vice versa) they are scattered by a macroscopic flow *approaching* them. We first derive in § 2.1 a lower limit to the acceleration time for a given upstream magnetic field. In § 2.2 we estimate the energy-loss time of high-energy electrons due to inverse-Compton (IC) scattering of afterglow photons. The maximum energy to which electrons can be accelerated and the maximum energy of emitted synchrotron photons is then estimated for a given upstream field by comparing the acceleration and loss times.

2.1. *Acceleration Time*

Consider a relativistic shock expanding into the circumburst medium, with postshock fluid of Lorentz factor $\Gamma \gg 1$ relative to the upstream medium. The high-energy electrons that cross the shock from the downstream region to the upstream region are confined (due to the Lorentz boost) in the upstream fluid frame to a narrow cone around the shock normal of opening angle $\sim 1/\Gamma$. The residence time of such electrons in the upstream, i.e., the time that the electrons spend in the upstream before being scattered back into the downstream, is approximately given by the time it takes for the electrons to be deflected by an angle $\sim 1/\Gamma$. Once the high-energy electrons travel at an angle, with respect to the shock normal, that is larger than $1/\Gamma$, they are overtaken by the shock. Thus, the upstream residence time may be estimated as

$$t_u \sim \gamma_e m_e c/\Gamma e B_u = \gamma_e' m_e c/eB_u. \tag{1}$$

Here B_u is the characteristic upstream field amplitude, γ_e is the electron Lorentz factor measured in the upstream fluid frame, and γ_e' is the electron Lorentz factor measured in the downstream fluid frame. From here on, primed variables denote parameter values measured in the downstream frame.

The acceleration time to energy $\gamma_e m_e c^2$ is given by the sum of the residence times in the upstream and downstream over the many cycles of upstream/downstream scattering the electron undergoes as its energy is increased. Nevertheless, the acceleration time is not much larger than the upstream residence time, t_u, for two reasons. First, under the assumption that the preshock magnetic field is not amplified and its amplitude is that present in the preshock plasma, the much stronger, near-equipartition magnetic field in the downstream deflects the electron much faster than it is being deflected in the upstream. This implies that a single "cycle time," i.e., the time for back and forth crossing of the shock, is dominated by the upstream residence time. Second, since the electron's energy is increased by some (energy-independent) factor in each cycle (e.g., Gallant & Achterberg 1999), the fact that the upstream residence time increases linearly with the electron's energy implies that the acceleration time is dominated by the last cycle time. Thus, the acceleration time may be written as

$$t_a = g\frac{\gamma_e' m_e c}{eB_u}, \tag{2}$$

where g is a correction factor with weak (logarithmic) dependence on γ_e. The exact value of g depends on the detailed assumptions regarding the structure of the field and the scattering process. For reasonable assumptions on the magnetic field, values as small as $g \approx 10$ may be obtained (e.g., Lemoine & Pelletier 2003; Lemoine & Revenu 2006). We conservatively adopt this value in what follows.

2.2. *Energy-Loss Time and Maximum Synchrotron Energy*

The maximum energy of accelerated electrons is limited by several factors. First, it is limited by the time available for acceleration, which is comparable to the shock expansion time. Balancing t_a given by equation (2) with the expansion time $\Gamma^2 t$, where t is the time measured by a distant observer, sets an upper limit to the electron Lorentz factor,

$$\gamma'_{\max} \simeq \frac{eB_u \Gamma^2 t}{g m_e c} \sim 10^7 g_1^{-1} \Gamma_2^2 t_3 B_{\mu\mathrm{G}}. \tag{3}$$

Here $B_u = 1B_{\mu\mathrm{G}}$ μG, $g = 10g_1$, $\Gamma = 10^2\Gamma_2$, and $t = 10^3 t_3$ s.

A more stringent constraint is obtained by considering the electrons' energy loss. Under the assumption that the upstream magnetic field is not significantly amplified, synchrotron cooling in the upstream is negligible. Electrons in the upstream lose energy primarily by IC scattering of afterglow synchrotron photons emitted in the downstream region. Since afterglow modeling typically implies that the energy density in radiation is similar (or larger) than the downstream magnetic field energy density, the large residence time of electrons in the upstream implies that IC losses of electrons in the upstream region is the dominant energy loss. This conclusion is valid, of course, provided IC scattering is not deep in the Klein-Nishina (KN) regime.

A comparison of the acceleration and energy-loss time is most easily carried out in the downstream frame. In this frame, the acceleration time is $t_a' \simeq t_a/\Gamma$. The energy-loss time may be estimated as follows. The energy density in synchrotron radiation is similar in the upstream and in the downstream regions, and both can be denoted as U'_{ph}. We show below that the KN effect is not important. Neglecting the KN effect, the cooling time due to IC scattering of (downstream-emitted) afterglow synchrotron photons is

$$t_c' \simeq \frac{3m_e c}{4\sigma_{\mathrm{T}} U'_{\mathrm{ph}} \gamma_e'}. \tag{4}$$

It is instructive to note that a similar estimate of t_c may be obtained by considering the energy-loss rate in the upstream frame. In this frame, the angular distributions of both photon and electron momenta are concentrated within a narrow cone of opening angle $1/\Gamma$ around the shock normal, while the photon number density and (individual) photon and electron energies are larger by a factor Γ than their downstream values, $n_{\mathrm{ph}} \simeq \Gamma n'_{\mathrm{ph}}$, $\nu \simeq \Gamma\nu'$, and $\gamma_e \simeq \Gamma\gamma_e'$. The energy-loss rate in the upstream frame is approximately given by

$$\dot{E} \sim n_{\mathrm{ph}}\sigma_{\mathrm{T}} c[1-\cos(1/\Gamma)]\Gamma\gamma_e'^2 h\nu' \sim n'_{\mathrm{ph}}\sigma_{\mathrm{T}} c \gamma_e'^2 h\nu',$$

where the collision rate is corrected by a factor $1-\cos(1/\Gamma) \approx 1/2\Gamma^2$, since the characteristic angle between photon and electron momenta is $1/\Gamma$, and $\Gamma\gamma_e'^2 h\nu'$ is the characteristic (upstream) energy of a scattered photon. Thus, the energy-loss rates in both upstream and downstream frames are similar, $\dot{E} \sim \dot{E}' \sim \gamma_e'^2 \sigma_{\mathrm{T}} c U'_{\mathrm{ph}}$, which implies $t_c \sim \Gamma t_c' \sim \Gamma m_e c/\gamma_e' \sigma_{\mathrm{T}} U'_{\mathrm{ph}}$ as obtained in equation (4).

Comparing t'_c and $t'_a \simeq t_a/\Gamma$ given by equation (2), we find

$$\gamma'_{\max} \simeq \left(\frac{3eB_u\Gamma}{4g\sigma_T U'_{\rm ph}}\right)^{1/2} \simeq 8.5\times10^4\left(\frac{B_{\mu G}\Gamma_2}{g_1 U'_{\rm ph,0}}\right)^{1/2}, \quad (5)$$

where $U'_{\rm ph} = 10^0 U'_{\rm ph,0}$ ergs cm^{-3}. Denoting by f the fraction of shock-accelerated electron energy converted to synchrotron radiation, we have $U'_{\rm ph} = 6f\epsilon_{e,-1}\Gamma_2^2 n_0$ ergs cm^{-3}, where $n = 10^0 n_0$ cm^{-3} is the proper density of the upstream plasma and $\epsilon_e = 10^{-1}\epsilon_{e,-1}$ is the fraction of downstream thermal energy carried by electrons. As we show below, f is of order unity, and hence $U'_{\rm ph,0} \sim 1$.

In what follows we provide a detailed derivation of f, and hence of $U'_{\rm ph}$. It is important to note here that when the KN correction is taken into account, $U'_{\rm ph}$ should be replaced by $U'_{\rm ph}(<\nu_{\rm KN})$, where $h\nu_{\rm KN} = \Gamma mc^2/\gamma'_{\max}$ is the photon energy above which IC cooling becomes less efficient due to the KN effect. For the derivation that follows, it is useful to define the downstream Compton $Y_{\max}$ parameter for electrons with Lorentz factor $\gamma'_{\max}$ as $Y_{\max} = [U'_{\rm ph}(<\nu_{\rm KN})/U'_B]_d$, and replace $U'_{\rm ph}$ in equation (5) with $U'_{\rm ph} = (U'_B)_d Y_{\max}$. The characteristic frequency of synchrotron photons emitted by electrons with $\gamma'_{\max}$, $\nu_{\max} \simeq 0.29\Gamma\gamma'^2_{\max}eB'_d/2\pi m_e c$, is then given by equation (5) to be

$$\begin{aligned}\nu_{\max} &\simeq 2.9\times10^{18}\frac{B_{\mu G}}{g_1(\epsilon_{B,-2}n_0)^{1/2}}\frac{\Gamma_2}{Y_{\max}}\ \text{Hz}\\ &\simeq 2.7\times10^{17}\frac{B_{\mu G}E_{53}^{1/8}}{g_1\epsilon_{B,-2}^{1/2}n_0^{5/8}}Y_{\max}^{-1}t_d^{-3/8}\ \text{Hz}. \quad (6)\end{aligned}$$

Here $t = 1t_d$ days, $\epsilon_B = 10^{-2}\epsilon_{B,-2}$ is the fraction of downstream thermal energy carried by the downstream magnetic field, and $E = 10^{53}E_{53}$ ergs is the (isotropic equivalent) explosion energy. The second equality holds for the case of the uniform density circumburst medium, for which the Lorentz factor drops with radius R following the Blandford-McKee self-similar solution, $\Gamma = (17E/16\pi nm_p c^2)^{1/2}R^{-3/2}$ (Blandford & McKee 1976), and the relation between Γ and the observer time t is $t = R/4\Gamma^2 c$ (Waxman 1997c). Note also that we have assumed an isotropic electron distribution and fully tangled magnetic field downstream and taken into account the fact that the synchrotron radiation peaks at 0.29 times the gyration frequency of the relevant electrons.

Let us finally derive the value of $Y_{\max}$ with consideration of the KN effect. In order to estimate $Y_{\max}$, we first discuss the energy distribution of afterglow photons, which is determined by the energy distribution of shock-accelerated electrons. Afterglow observations imply that the postshock electron energy distribution follows a power law, $dn_e/d\gamma_e \propto \gamma_e^{-p}$ for $\gamma_m \le \gamma_e \le \gamma_{\max}$, with a spectral index $p \approx 2.2$ (Waxman 1997a; Freedman & Waxman 2001, and references therein; Wu et al. 2004). This energy distribution is consistent with the theoretical value derived for isotropic diffusion of accelerated particles (in the test particle limit) in both numerical calculations (Bednarz & Ostrowski 1998; Kirk et al. 2000; Achterberg et al. 2001) and analytic analyses (Keshet & Waxman 2005). The synchrotron spectrum of such an electron distribution follows a power law, with power per unit frequency $f_\nu \propto \nu^{-(p-1)/2}$, up to the energy where the radiative (synchrotron and IC) losses of the electrons become important. This occurs at electron energy for which the radiative loss time is shorter than the adiabatic cooling time, i.e., than the energy-loss time due to the expansion of the postshock plasma. The adiabatic loss time, $t'_{\rm ad} \simeq 6R/13c\Gamma$ (Gruzinov & Waxman 1999) with $R = 4\Gamma^2 ct$ (for uniform circumburst density), is longer than the radiative cooling time for electrons with Lorentz factors exceeding $\gamma'_c = 3mc/4\sigma_T(U'_B)_d(1+Y_C)t'_{\rm ad}$. Here Y_C is the Compton Y parameter for electrons with $\gamma'_e = \gamma'_c$. The characteristic synchrotron frequency of photons emitted by electrons with $\gamma'_e = \gamma'_c$, $\nu_c \simeq 0.29\Gamma\gamma'^2_c eB'_d/2\pi m_e c$, is (for uniform circumburst density)

$$\begin{aligned}\nu_c &\simeq \frac{2.5\times10^{13}}{(\epsilon_{B,-2}n_0)^{3/2}(1+Y_C)^2\Gamma_2^4 t_3^2}\ \text{Hz}\\ &\simeq \frac{7.5\times10^{13}}{\epsilon_{B,-2}^{3/2}E_{53}^{1/2}n_0}\frac{t_d^{-1/2}}{(1+Y_C)^2}\ \text{Hz}. \quad (7)\end{aligned}$$

At higher frequency, the photon spectrum steepens from $f_\nu \propto \nu^{-(p-1)/2} \approx \nu^{-1/2}$ to $f_\nu \propto \nu^{-p/2} \approx \nu^{-1}$.

In order to complete the description of the synchrotron photon spectrum, let us estimate the lowest energy of the accelerated electrons. The fraction of downstream thermal energy carried by electrons, ϵ_e, is typically $\epsilon_e \gtrsim 0.1$ (Freedman & Waxman 2001). This implies a minimum Lorentz factor of accelerated electrons of $\gamma'_m \simeq \epsilon_e(m_p/m_e)\Gamma$, for which the characteristic frequency of synchrotron emission is $\nu_m \simeq 10^{13}\epsilon_{e,-1}^2(\epsilon_{B,-2}n_0)^{1/2}\Gamma_1^4$ Hz. For $t > 1$ hr the afterglow is typically in the "slow cooling regime," with $\gamma'_m < \gamma'_c$. The synchrotron spectrum is therefore well approximated by a broken power law, with $f_\nu \propto \nu^{-(p-1)/2} \approx \nu^{-1/2}$ for $\nu_m < \nu < \nu_c$ and $f_\nu \propto \nu^{-p/2} \approx \nu^{-1}$ for $\nu > \nu_c$. In this case, the energy density of synchrotron photons is mainly concentrated around ν_c, i.e., $U_{\rm syn} \simeq U_{\rm ph}(<\nu_c)$, and we may approximate

$$Y_{\max} = Y_{\rm syn}\min[1,\ (\nu_{\rm KN}/\nu_c)^{1/2}], \quad (8)$$

where $Y_{\rm syn} \equiv (U'_{\rm syn}/U'_B)_d$ is the downstream energy density ratio between the synchrotron radiation and the magnetic field, and $\nu_{\rm KN}$ is the frequency of photons for which IC scattering of electrons with $\gamma'_e = \gamma'_{\max}$ is in the KN regime. Here $\nu_{\rm KN} = \Gamma m_e c^2/h\gamma'_{\max}$, i.e.,

$$\begin{aligned}\nu_{\rm KN} &= 1.3\times10^{17}\left(\frac{g_1\epsilon_{B,-2}n_0}{B_{\mu G}}\right)^{1/2}\Gamma_2^{3/2}Y_{\max}^{1/2}\ \text{Hz}\\ &\simeq 3.1\times10^{15}\left(\frac{g_1\epsilon_{B,-2}}{B_{\mu G}}\right)^{1/2}E_{53}^{3/16}n_0^{5/16}t_d^{-9/16}Y_{\max}^{1/2}\ \text{Hz}. \quad (9)\end{aligned}$$

Here too, the second equality holds for the uniform density circumburst medium.

Using the equations for $\nu_{\rm KN}$ and ν_c, we have

$$\frac{\nu_{\rm KN}}{\nu_c} = 41\frac{g_1^{1/2}\epsilon_{B,-2}^2 E_{53}^{11/16}n_0^{21/16}}{B_{\mu G}}\frac{(1+Y_C)^2 Y_{\max}^{1/2}}{t_d^{1/16}}. \quad (10)$$

This implies that unless the upstream field is amplified to $B_{\mu G} \gg 1$, $\nu_c \lesssim \nu_{\rm KN}$. In this case, we may approximate $Y_{\max} \approx Y_{\rm syn}$. As for $Y_{\rm syn}$, if $Y_{\rm syn} \gtrsim 1$ and only single IC scattering is considered (multiple IC scattering is typically suppressed, as it is well within the KN regime), then $Y_{\rm syn} \simeq (\eta\epsilon_e/\epsilon_B)^{1/2}$, where $\eta = (\nu_c/\nu_m)^{-(p-2)/2}$ is the "radiative efficiency" of the postshock electrons (e.g., Sari & Esin 2001). For $p \approx 2$, we have $\eta \approx 1$, and since afterglow observations imply $\epsilon_e \gtrsim 0.1$, we find $\eta\epsilon_e/\epsilon_B \gtrsim 1$

and $Y_{\rm syn} \gtrsim 1$. To summarize, afterglow observations imply $Y_{\rm max} \approx Y_{\rm syn} \approx (\eta\epsilon_e/\epsilon_B)^{1/2} \approx (\epsilon_e/\epsilon_B)^{1/2} \sim$ a few.

With $Y_{\rm max} \approx (\eta\epsilon_e/\epsilon_B)^{1/2}$, equation (6) then implies a maximum afterglow synchrotron photon energy of

$$h\nu_{\rm max}^{\rm obs} \simeq 0.3 \frac{B_{\mu{\rm G}} E_{53}^{1/8}}{g_1(\eta\epsilon_{e,-1})^{1/2}[(1+z)n_0]^{5/8}} t_d^{-3/8} \text{ keV}, \qquad (11)$$

where we have taken into account the redshift z of the source. Note that $\nu_{\rm max}^{\rm obs}$ is independent of the poorly known ϵ_B. At photons energies larger than $h\nu_{\rm max}^{\rm obs}$, an exponential cutoff is expected in the afterglow spectrum. Since for $B_{\mu{\rm G}} \sim 1$ this cutoff is expected to take place in the soft X-ray band, X-ray afterglow observations provide an interesting constraint on the upstream field. If the synchrotron spectrum is observed to extend beyond an energy $h\nu^{\rm obs}$, then a lower limit to the upstream field is implied,

$$B_u > 50 \frac{g_1(\eta\epsilon_{e,-1})^{1/2}}{E_{53}^{1/8}} \left(n_0 \frac{1+z}{2}\right)^{5/8} t_d^{3/8} \frac{h\nu^{\rm obs}}{10 \text{ keV}} \ \mu\text{G}. \qquad (12)$$

3. LOWER LIMITS ON UPSTREAM FIELD STRENGTH

The *BeppoSAX* catalog of X-ray afterglows is presented in De Pasquale et al. (2006). X-ray data are typically available between ~0.3 and ~1 days. The 0.1–10 keV spectra are well fit by a simple, photoelectrically absorbed power law. As shown in Figure 1 of De Pasquale et al. (2006), most of the spectra have been well determined up to 10 keV, with no sign of a cutoff or steepening of the spectra at 10 keV. This implies, using equation (12),

$$B_u \gg 50 \frac{g_1(\eta\epsilon_{e,-1})^{1/2}}{E_{53}^{1/8}} \left(n_0 \frac{1+z}{2}\right)^{5/8} \mu\text{G}. \qquad (13)$$

This lower limit is only weakly dependent on E, but somewhat sensitive to ϵ_e. Values of $\epsilon_e \gtrsim 0.1$ are inferred from afterglow modeling and are consistent with the clustering of explosion energy (Frail et al. 2001) and X-ray afterglow luminosity (Freedman & Waxman 2001).

In several cases, stronger constraints on the magnetic field are obtained.

GRB 990123.—This is the only case in which the afterglow was detected up to a photon energy of 60 keV (using the *BeppoSAX* PDS instrument; Maiorano et al. 2005; Corsi et al. 2005), implying nonthermal emission up to this high energy at ≃0.5 days, and a limit on the magnetic field of $B_u > 0.2n_0^{5/8}$ mG.

GRB 030329.—In this case, there are very late-time detections of X-ray emission at 37, 61, and 258 days by *XMM-Newton* (Tiengo et al. 2003). Since at late times the expansion may not be highly relativistic, as assumed in the analysis of this paper, and since late-time emission could be affected by the associated supernova, we only consider here the observation at 37 days, with $h\nu_{\rm max}^{\rm obs} > 10$ keV. Equation (12) then implies $B_u > 0.2n_0^{5/8}$ mG. We note that this numerical value is an underestimate, since at this late time the shock expansion is not spherical (due to jet expansion), leading to a faster decrease of Γ than assumed in our analysis.

GRB 050904.—For this high-z burst ($z = 6.3$), the *Swift* XRT (0.2–10 keV) observation at 3×10^5 s shows emission up to hard X-rays, 1.4–73 keV in the source rest frame (Cusumano et al. 2006). Equation (12) then implies $B_u > 3(n/100 \text{ cm}^{-3})^{5/8}$ mG $= 0.2n_0^{5/8}$ mG. Here we have also given a value for large n, as the analysis of absorption lines in the optical afterglow implies $n_e = 10^{2.3\pm0.7}$ cm^{-3} (Kawai et al. 2006).

4. DISCUSSION

We have shown that late-time, $t \gtrsim 1$ day, X-ray observations of GRB afterglows provide interesting constraints on the upstream magnetic field. A lower limit to the magnetic field is obtained by requiring the acceleration time of electrons producing high-energy synchrotron photons to be shorter than their energy-loss time due to inverse-Compton scattering of afterglow photons. The lower limit for the magnetic field is given in equation (12) as a function of the energy of the observed high-energy photons, $h\nu$. We have then shown in § 3 that X-ray afterglow observations typically imply an upstream field strength of $B \gg 0.05n_0^{5/8}$ mG, where $n = 10^0 n_0$ cm^{-3} is the upstream plasma density. In several cases, $B > 0.2n_0^{5/8}$ mG is obtained.

The large magnetic fields inferred are unlikely to be present in the interstellar medium (ISM) of the galaxies hosting the GRBs. This is easy to see by assuming equipartition between the magnetic field energy density and the turbulent energy density in the ISM, $B^2/8\pi = nm_p v_T^2/2$, which implies that the turbulent velocity required to support the inferred $B > 0.2n_0^{5/8}$ mG field is $v_T \gtrsim 500n_0^{1/8}$ km s^{-1}. Indeed, mG magnetic fields are observed in starburst galaxies (Thompson et al. 2006), but the ISM density in such galaxies is high, $n_0 \gg 1$, implying that the constraint $B > 0.2n_0^{5/8}$ mG is not satisfied.

Given the above estimate of $v_T \gtrsim 500n_0^{1/8}$ km s^{-1}, it is clear that a sufficiently strong magnetic field may be present in a wind emitted by the GRB progenitor, and into which the shock propagates, provided that the wind is fast, $v_w \gtrsim 10^3$ km s^{-1}, and highly magnetized, i.e., provided that it carries magnetic field energy density that is not much smaller than its kinetic energy density. Such interpretation faces two challenges. Wolf-Rayet stars possess fast, $\sim 10^3$ km s^{-1}, winds and are the likely progenitors of SN of Type Ic, which are associated with GRBs (e.g., Zhang & Mészáros 2004; Piran 2005). However, such hot stars have radiative envelopes and are not expected therefore to have magnetically driven winds. Nevertheless, the possibility that magnetic fields play an important role in the winds of such hot stars cannot be ruled out (e.g., Poe et al. 1989; Ignace & Gayley 2003). The second challenge is that afterglow observations so far have not provided a clear indication for GRB shocks propagating into a circumburst wind (e.g., Zhang & Mészáros 2004). However, it should be noted that late-time, ≳1 day, afterglow observations do not allow us to clearly distinguish between propagation into a wind and propagation into the uniform density characteristic of the ISM (e.g., Livio & Waxman 2000), and that early-time afterglow observations provided by *Swift* are not yet properly understood (e.g., Fan & Piran 2006).

If the GRB shock is not propagating into a highly magnetized wind, then the large upstream magnetic fields inferred by our analysis require strong amplification of the field ahead of the shock, most likely by the streaming of high-energy particles. If this is the correct interpretation, then numerical simulations of relativistic, weakly magnetized collisionless shocks that properly describe the shock structure should present strong preshock magnetic field amplification, and thus also particle acceleration (note that the precursor magnetic field obtained in the simulation of Spitkovksy [2006] is a numerical effect, due to the reflection of particles from the edge of the simulation box).

Finally, it has recently been concluded by Lyubarsky & Eichler (2006) that the saturation of the Weibel instability implies that the shock is mediated by deflections of particles in the preexisting, ~ 1 μG upstream field, and Milosavljević & Nakar (2005) have postulated, in their phenomenological analysis interaction of the accelerated particles with the preexisting ~ 1 μG upstream field. The constraints derived by our analysis indicated that this conclusion (postulation) may not apply to collisionless GRB shocks.

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