ON PARTICLE ACCELERATION RATE IN GAMMA-RAY BURST AFTERGLOWS

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ABSTRACT

It is well known that collisionless shocks are major sites of particle acceleration in the universe, but the details of the acceleration process are still not well understood. The particle acceleration rate, which can shed light on the acceleration process, is rarely measured in astrophysical environments. Here, we use observations of gamma-ray burst (GRB) afterglows, which are weakly magnetized relativistic collisionless shocks in ion–electron plasma, to constrain the rate of particle acceleration in such shocks. We find, based on X-ray and GeV afterglows, an acceleration rate that is most likely very fast, approaching the Bohm limit, when the shock Lorentz factor is in the range of $\Gamma \sim 10–100$. In that case X-ray observations may be consistent with no amplification of the magnetic field in the shock upstream region. We examine the X-ray afterglow of GRB 060729, which is observed for 642 days showing a sharp decay in the flux starting about 400 days after the burst, when the shock Lorentz factor is $\sim 5$. We find that inability to accelerate X-ray-emitting electrons at late time provides a natural explanation for the sharp decay, and that also in that case acceleration must be rather fast, and cannot be more than a 100 times slower than the Bohm limit. We conclude that particle acceleration is most likely fast in GRB afterglows, at least as long as the blast wave is ultrarelativistic.

Key words: acceleration of particles – gamma-ray burst: general – shock waves

1. INTRODUCTION

Astrophysical collisionless shocks are efficient particle accelerators. The signature of ultrarelativistic particles that are accelerated in these shocks is seen in a variety of astrophysical phenomena and over a wide range of environments. Nevertheless, despite an extensive study, the acceleration processes are still largely unknown. One of the leading candidates is the diffusive shock acceleration (DSA; e.g., Bell 1978; Blandford & Ostriker 1978; Blandford & Eichler 1987), where charged particles are accelerated by crossing the shock back and forth. The acceleration time in DSA depends on the duration that it takes a particle to close a single cycle, i.e., to cross the shock back and forth one time. This time increases with the particle energy, and it sets the maximal Lorentz factor, $\gamma_{\text{max}}$, that a particle can achieve. Thus, measuring $\gamma_{\text{max}}$ provides a direct information about the acceleration process and about the physical conditions in the acceleration site.

The reflection of particles back and forth through the shock is believed to be done by scattering on fluctuating magnetic fields. This process is typically approximated as a diffusion in the direction of the particle velocity, in which case the duration of the acceleration depends on the particle mean free path, $\lambda$. It is reasonable to assume that typically the shortest possible mean free path is the Larmor radius $r_l$ (Bohm limit), and thus to parameterize the diffusion by $\lambda = \eta r_l$ where $\eta \gg 1$ is generally expected, although this is not a hard lower limit. The value of $\eta$, which measures the diffusion efficiency and how fast the acceleration, was constrained only in a small number of systems. Probably the best estimate is obtained from the gamma-ray spectrum of the Crab nebula (de Jager & Harding 1992). The spectrum shows two components, where the lower energy component, which is most likely dominated by synchrotron, shows a cutoff around 100 MeV. The fact that synchrotron emission reaches these energies, given the rapid synchrotron cooling, implies $\eta \approx 1$. Higher energy synchrotron emission during flares suggest that maybe even $\eta < 1$ is required (Abdo et al. 2011). Another system where $\eta$ was claimed to be measured is the supernova remnant SNR RXJ1713.72-3946, where Uchiyama et al. (2007) find that the observed X-ray variability on a year timescale indicates on $\eta \sim 1$. Thus, two very different acceleration sites, one relativistic, possibly highly magnetized shock in pair plasma and the other Newtonian shock in ion–electron plasma suggest that particle acceleration, if dominated by DSA, is extremely fast. Calculating $\eta$ from first principles is impossible at this point, since it depends on the unknown shock structure, and most importantly on the properties of the upstream and downstream magnetic fields. Calculations of $\eta$ in relativistic shocks were done only by assuming the magnetic field structure. For example, Lemoine & Pelletier (2003) and Lemoine & Revenu (2006) find that when a Kolmogorov magnetic turbulence spectrum is assumed in the upstream region, then $\eta \approx 10$.

Here, we examine the constraints that can be obtained on $\eta$ from observations of cutoff, or the lack thereof, in gamma-ray burst (GRB) afterglow light curves and spectra. These afterglows are almost certainly generated by relativistic blast waves that propagate into a weakly magnetized ion–electron plasma (for reviews see Piran 2004; Mészáros 2006; Nakar 2007). Previous studies of particle acceleration in GRB afterglows assumed $\eta = 1–10$ and used the lack of spectral cutoff in the observed X-ray and GeV emission to constrain the magnetic field upstream and/or downstream of the shock (Li & Waxman 2006; Piran & Nakar 2010, hereafter PN10; Barniol Duran & Kumar 2011a; Li & Zhao 2011). Here, we take a different approach asking how well $\eta$ can be constrained. Moreover, we find that the fast decay observed in the extraordinarily year long X-ray afterglow of GRB 060729 is naturally explained by the inability of the shock to accelerate X-ray-emitting electrons, providing a measurement of $\eta$ in that case.

In Section 2, we describe the various limits on $\gamma_{\text{max}}$. The resulting limits on $\eta$, for various circumburst density profiles,
are derived in Section 3. The special case of the GRB with the longest duration X-ray afterglow, GRB 060729, is discussed in Section 4.

2. LIMITS ON THE MAXIMAL LORENTZ FACTOR

The two main factors that limit the acceleration of particles in a decelerating relativistic blast wave are confinement and cooling (e.g., Li & Waxman 2006; PN10; Barniol Duran & Kumar 2011a). Below we shortly discuss these limits (see PN10 for details). Observations indicate that the magnetic field in the downstream region is amplified in GRB external shocks well beyond the effect of compression. Thus, most of the emission take place in the downstream region while a particle spends most of its acceleration time in the upstream region. Throughout the paper we assume that $\eta$ is similar in the shock downstream and upstream, but we highlight which of the observations constrain $\eta$ in the downstream region and which constrain $\eta$ in the upstream region.

Confinement is limited by the ability of the accelerated particle that is moving in the upstream region to cross the shock back into the downstream region. Thus, confinement is limited by $\eta$ in the upstream region. Its limit on the maximal Lorentz factor is set by the requirement that the particle complete a turn of $180^\circ$, as seen in the shock frame, while the shock propagates a distance $f_u R$, where $f_u$ accounts for the shock deceleration (see PN10). Thus,

$$\gamma_{\text{conf}} \approx \frac{eB_u}{\eta m_e c^2} f_u R,$$

where $\ldots$ denotes quantities in the shock rest frame, $R$ is the shock radius, $e$ and $m_e$ are the electron charge and mass, and $c$ is the speed of light, and $B_u$ is the rest-frame upstream magnetic field.

The two processes that dominate cooling are synchrotron and inverse Compton (IC). In both cases, the maximal Lorentz factor is found by equating the acceleration time (i.e., the time to complete a Fermi cycle in relativistic shocks) to the relevant cooling time. In the case of synchrotron

$$\gamma_{\text{synch}} \approx \left( \frac{6\pi e}{\eta B_d \sigma_T} \right)^{1/2},$$

where $\sigma_T$ is the Thomson cross-section and $B_d$ is the downstream magnetic field. Since observations indicate that the downstream magnetic field is amplified by the shock, synchrotron cooling in the downstream region is more limiting than in the upstream region and it sets a limit on $\eta$ in the downstream region.

IC cooling is more efficient in the upstream region, since the radiation field is similar in both sides of shock, but a particle spends more time in the upstream region, where the magnetic field is lower. Thus, $\gamma_{\text{IC}}$ is limited by $\eta$ in the upstream region:

$$\gamma_{\text{IC}} \approx \left( \frac{3eB'_u}{4\pi\sigma_T U_{\text{rad}}(\ll \nu_{\text{KN}})} \right)^{1/2} \gamma_{\text{synch}} \left( \frac{B'_u}{Y(\gamma_{\text{IC}}) B_d} \right)^{1/2},$$

where $B'_u$ is the magnetic field in the upstream region as measured in the shock frame (related to the rest-frame magnetic field by $B'_u \approx \Gamma B_u$). $U_{\text{rad}}(\ll \nu_{\text{KN}})$ is the shock frame radiation energy density at frequencies smaller than $\nu_{\text{KN}}(\gamma_{\text{IC}}) = \frac{m_e c^2}{h \gamma_{\text{IC}}}$,

where $\hbar$ is the Planck constant. Equation (3) also gives the relation between $\gamma_{\text{IC}}$ and $\gamma_{\text{synch}}$ using the ratio between IC and the synchrotron cooling rate in the downstream region,

$$Y(\gamma_{\text{IC}}) = U_{\text{rad}}(\ll \nu_{\text{KN}})/(B_d^2/8\pi).$$

3. MAXIMAL OBSERVED FREQUENCY AND LIMITS ON $\eta$

Below we derive the constraints that X-ray and GeV afterglows set on $\eta$. We assume that the observed emission is synchrotron radiation generated by a quasi-spherical decelerating adiabatic blast wave. This is almost certainly the case in many X-ray afterglows, at least during the first day. The origin of the observed long-lasting GeV emission, which is seen up to $\sim 1000$ s after some bursts, is still unclear, although observations suggest that it is also synchrotron emission from the decelerating blast wave (Kumar & Barniol Duran 2009, 2010; Ghisellini et al. 2010). Since afterglow observations suggest that the circumburst density profile varies from one GRB to another, we consider here two typical external density profiles, one constant as expected for the interstellar medium (ISM) and one $\propto R^{-2}$ as expected for a stellar wind.

3.1. ISM

Under the assumption of spherical expansion in a constant density $n$, the radius and Lorentz factor of an adiabatic blast wave with energy $E$ at an observer time $t$ are (e.g., Sari et al. 1998)

$$R \approx 6 \times 10^{17} \text{cm} \left( \frac{E_{53}}{n} \right)^{1/4} \left( \frac{1 + \frac{z}{3}}{\frac{3}{5}} \right)^{-1/4} t_5^{1/4},$$

$$\Gamma \approx 12 \left( \frac{E_{53}}{n} \right)^{1/8} \left( \frac{1 + \frac{z}{3}}{\frac{3}{5}} \right)^{3/8} t_5^{-3/8},$$

where $z$ is the burst redshift and $q_5$ denotes the value of $q/10^5$ in c.g.s. units. We assume that the downstream magnetic field is a constant fraction, $B_{B,d}$, of the internal energy behind the shock so $B_d \approx \left( 32\pi e_B m_p c^3/3m_e \right)^{1/2}$, where $m_p$ is the proton mass. If the magnetic field in the upstream is not amplified by a precursor to the shock then it is expected to be constant and of the order of $10 \mu$G. If it is amplified then it may be significantly larger. Finally, $f_u = 1/3$ in ISM (PN10). Using the synchrotron emission from the downstream region, $h \nu = \Gamma^{\gamma-2} e B_d^2/(2\pi m_e c)$, we obtain the maximal frequency that is dictated by the limits discussed above:

$$h\nu_{\text{conf}} \approx 2 \times 10^{10} \text{eV} \eta^{-2} E_{53}^{3/4} \epsilon_{B,d}^{1/2} B_d^2 \nu_0^{-1/4} t_5^{-1/4},$$

$$h\nu_{\text{sync}} \approx 2.5 \times 10^8 \text{eV} \eta^{-1/4} t_5^{3/8},$$

$$h\nu_{\text{IC}} \approx 1.5 \times 10^5 \text{eV} \eta^{-1/4} B_d^2 \nu_0^{-5/8} t_5^{-1/4},$$

where here, and throughout the paper, we derive values for the typical $z = 2$. We also ignore dependencies on parameters that are raised to the power of 1/8 since these cannot affect the result by an order of magnitude, which is the accuracy of our calculation to begin with. The constraints that we derive for confinement and synchrotron are similar to those of PN10, which assumed $\eta = 1$, and the constraint on the IC cooling is similar to the one derived in Li & Waxman (2006), which used a canonical value of $\eta = 10$.

Equation (6) implies that for canonical GRB parameters confinement does not play an important role when the afterglow blast wave propagates into ISM, with the possible exception
of very early time GeV emission.\(^1\) The limit provided by synchrotron cooling is the most robust as it is independent of almost anything,\(^2\) except for the time since the explosion, which is typically well measured:

\[
\eta \lesssim 3 \left( \frac{h \nu_{\text{obs}}}{1 \text{ GeV}} \right)^{-1} \left( \frac{t}{100 \text{ s}} \right)^{-3/8}.
\]  

(7)

Several GeV photons are seen \(\sim 100 \text{ s}\) after the burst starts in a number of the Fermi-LAT GRBs (e.g., Abdo et al. 2009a, 2009b; Ackermann et al. 2011) while \(>100 \text{ MeV}\) photons are seen in large numbers up to \(\sim 1000 \text{ s}\) after the burst in many Fermi-LAT GRBs. Therefore, if the GeV emission is emitted by synchrotron process in the external shock, as suggested by several authors (Kumar & Barniol Duran 2009, 2010; Ghisellini et al. 2010) then the acceleration process in ultrarelativistic (\(\Gamma \sim 100\)), weakly magnetized, shock must be extremely fast\(^3\) with \(\eta \lesssim 1\) in the shock downstream.

The IC constraint depends on the value of the \(Y\) parameter. \(Y_{\text{GeV}},\) the \(Y\) parameter of GeV emitting electrons, varies by many orders of magnitude across the relevant phase space (Nakar et al. 2009; Li & Zhao 2011; Barniol Duran & Kumar 2011b). As a result the synchrotron cooling limit is more stringent in part of the phase space. Therefore, considering the robustness and tightness of the synchrotron constraint, we do not attempt to cover here the possible IC limits on the GeV emission. However, when considering X-ray emission the synchrotron limit is very loose. Therefore, we consider the IC limit in that case, for which we need to evaluate \(Y_x\). This is not trivial due to Klein–Nishina (KN) effects that play different roles over various areas of the phase space. An upper limit on \(Y_x\) can be easily obtained by assuming that the electrons are in the fast cooling regime and that KN effects are negligible. In that case \(Y_x = \sqrt{\varepsilon_e / \varepsilon_{\text{B,d}}}\) if \(\varepsilon_e > \varepsilon_{\text{B,d}}\) (e.g., Sar & Esin 2001), where \(\varepsilon_e\) is the fraction of the internal energy behind the shock that goes into accelerated electrons. For typical parameters the electrons are cooling slowly at \(t > 10^8 \text{ s}\), implying that \(Y_x\) is smaller for electron distribution with a power-law index \(p > 2\): \(Y_x = \sqrt{\varepsilon_e / \varepsilon_{\text{B,d}}(\gamma / \gamma_m)^{2-p}/7}\), where \(\gamma_m\) is the typical (also minimal) Lorentz factor of accelerated electrons and \(\gamma_c\) is the Lorentz factor of electrons that are cooling over dynamical timescale. In addition, over a large range of the parameter space KN suppression can be important, reducing the value of \(Y_x\) further. To account for these effects we use Equations (46), (59), (60), and (63) of Nakar et al. (2009), which take consideration of the KN effects and their feedback on the electron distribution, to scan the phase space for the value of \(\varepsilon_{\text{B,d}} / Y_x\) (which appears in the IC limit of Equation (6)). We scan the parameter phase space and find that if the fraction of downstream region internal energy that goes to electrons is \(\varepsilon_e = 0.1\), the electron distribution power-law index is in the range \(p = 2-2.8\) (Curran et al. 2010) and \(\varepsilon_{\text{B,d}} > 10^{-3}\), then the value of \(\varepsilon_{\text{B,d}} / Y_x\) is typically in the range of 0.3–3 and its dependence on the other parameters, \(\eta\), \(E\), and \(t\) is rather weak (most of the dependence in this range is on \(p\) due to the fraction of energy that is in fast cooling electrons, while KN suppression is rather mild). When \(\varepsilon_{\text{B,d}} \ll 10^{-3}\) and/or \(\varepsilon_e \ll 0.1\) KN effects significantly suppress \(Y_x\) and \(\varepsilon_{\text{B,d}} / Y_x \ll 1\). We therefore conclude that for the canonical values of \(\varepsilon_{\text{B,d}} > 10^{-3},\varepsilon_e = 0.1, n_0 \approx 1\), and \(p = 2-2.8\),

\[
\eta \lesssim 15 \frac{10 \text{keV}}{h \nu_{\text{obs}}} B_{\mu s}^{-5} s_{10}^{-3/8},
\]  

(8)

in the shock upstream region. If, however, \(\varepsilon_{\text{B,d}} \ll 10^{-3}\) or \(\varepsilon_e \ll 0.1\) or \(n_0 \ll 1\), then X-ray emission does not provide strong constraints on \(\eta\).

Many afterglows show X-ray emission (0.2–10 keV) that is bright for days and in some cases weeks, without showing a clear sign of spectral softening (Liang et al. 2008; Racusin et al. 2009). Thus, since afterglow modeling typically implies \(\varepsilon_{\text{B,d}} > 10^{-3}\) or \(\varepsilon_e \approx 0.1\) and an ISM circumburst environment, observations of X-ray afterglows suggest that the acceleration mancinism in relativistic shocks (\(\Gamma \sim 10-50\)) is fast. As evident from Equation (8), this limit depends on various parameters. Some are constrained rather well, e.g., \(\varepsilon_e\), while others are less constrained, e.g., \(n_0\). Most important is the dependence on \(B_{\mu s}\).

There is a viable possibility that the interaction of accelerated particles that run ahead of the shock significantly amplifies the upstream magnetic field (Blandford & Eichler 1987; Bell 2004; Milosavljević & Nakar 2006). If this is the case then the limits provided by X-ray observations are rather loose. In fact, Li & Waxman (2006) concluded, based on X-ray observations, that the upstream magnetic field must be amplified at least up to \(\frac{1}{5} B_{\mu s}\). This conclusion was based on the assumption that \(\eta = 10\). They also take as a canonical value \(\varepsilon_{\text{B,d}} / Y_x \approx 0.1\), which they calculate by considering only the part of the phase space where KN effects are negligible and by taking \(p = 2\), for which, \(Y_x\) is not suppressed by the slow cooling of most of the electrons. Our results show that a more careful estimate of \(\varepsilon_{\text{B,d}} / Y_x\), reduces the Li & Waxman (2006) limit by at least a factor of a few. In addition, if the acceleration is as fast as suggested by the recently detected GeV emission, and \(\eta \sim 1\) also in the upstream region then the limits on the upstream field drop to \(\mu G\) level, implying that current X-ray observations do not provide strong evidence for magnetic field amplification in GRB afterglows.

### 3.2. Wind

The mass density profile in a wind from massive stars is \(\rho = AR^{-2}\). Under the assumption of spherical expansion the radius and Lorentz factor of an adiabatic blast wave with energy \(E\) at an observer time \(t\) are (Chevalier & Li 2000)

\[
R = 3 \times 10^{17} \text{ cm} \left( \frac{E_{53}}{A_{\nu}} \right)^{1/2} \left( \frac{1 + z}{3} \right)^{1/2} t_{10}^{1/2},
\]  

\[
\Gamma = 12 \left( \frac{E_{53}}{A_{\nu}} \right)^{1/4} \left( \frac{1 + z}{3} \right)^{1/4} t_{10}^{1/4},
\]  

(9)
where $A_\ast = A/(5 \times 10^{11}\ g\ cm^{-1})$. Similar to the ISM case we assume that the downstream magnetic is a constant fraction, $\epsilon_{B,d,-2}$, of the internal energy behind the shock. Contrary to the ISM case, the upstream magnetic field is not constant. The magnetic field in the upstream region depends on the wind magnetization and flux freezing implies $B_u \propto R^{-1}$, assuming that upstream field is not amplified by the shock precursor. The normalization depends on the wind velocity and on the surface rotation velocity and magnetic field (Goldreich & Julian 1970), which are not tightly constrained. For typical parameters of a Wolf–Rayet wind a field of $\sim 10\, \mu G$ is expected at $R = 10^{18}\ cm$ (Eichler & Usov 1993), but it can be more than an order of magnitude larger or smaller. Therefore, we write the upstream magnetic field as $B_u = 1\, \mu G(R/10^{18}\ cm)^{-1}B_{u,\mu G,19}$. Using this parameterization and $f_a = 1/2$ (PN10) the various constraints on the maximal observed frequencies are

$$\begin{gathered}
h_{\text{conf}} = 3 \times 10^{11}\ eV\, \eta^{2/3}\, E_{53}^{1/4} A_{*,4}^{1/2} B_{u,\mu G,19}^{-1/4} \\
h_{\text{sync}} = 2 \times 10^{9}\ eV\, \eta^{2} E_{53}^{1/4} A_{*,4}^{1/4} \nu^{-1} \tau_{-1}^{-1/4} \\
h_{\text{IC}} = 3 \times 10^{9}\ eV\, \eta^{4/3} E_{53}^{1/4} B_{u,\mu G,19}^{2} \nu^{-2} A_{*,4}^{-3/4} \nu^{-1} \tau_{-1}^{-1/4}.
\end{gathered}$$

The most robust synchrotron limit is relevant only to the GeV emission:

$$\eta \lesssim \left( \frac{h_{\text{obs}}}{1\, \text{GeV}} \right)^{-1} \left( \frac{E_{53}}{A_{*,4}} \right)^{1/4} \left( \frac{\tau}{100\, \text{s}} \right)^{-1/4}. \quad (11)$$

Implying that if the observed GeV emission is produced by synchrotron from an external shock in a wind environment then the acceleration mechanism must be extremely fast. This limit is very similar to the one obtained in the case of an ISM density profile (Equation (7)) and is therefore general for any reasonable circumburst density profile.

The IC limit depends on the value of the $Y$ parameter. For the same reasons discussed in the ISM case, we consider here IC limits only on the X-ray emission. Unlike the ISM case the $Y$ parameter of X-ray-emitting electrons, $Y_\ast$, depends strongly on time. In a wind density profile the observed synchrotron cooling frequency, $\nu_c$, where most of the synchrotron energy is emitted, is increasing with time. As a result, KN effects become significantly more dominant with time, suppressing the IC cooling of X-ray-emitting electrons at late time. The standard afterglow theory in a wind (Chevalier & Li 2000) provides the value of $\nu_c(t)$ and of the Lorentz factor of X-ray-emitting electrons, $\gamma_e(t)$. Since at slow cooling most of the synchrotron luminosity is emitted at $\nu_c$ (for $p < 3$), KN effects are negligible for X-ray-emitting electrons as long as $\gamma_e\nu_e / \Gamma < m_e c^2$. Thus, the time at which this inequality becomes an equality provides a good approximation to the time at which KN effects on $Y_\ast$ become important:

$$t_{\gamma_e,\text{KN}} \sim 4 \times 10^6\, s A_{53}^{10/7} E_{19,d^{-2}E_{e,-1}}^{4/7}. \quad (12)$$

This approximation assumes slow cooling and that the cooling frequency is below the X-ray (if the latter is not satisfied then cooling is not the limiting factor anyway, see below). It also ignores the effect of IC cooling on $\nu_c$ which can only delay the time at which KN effects become important. Thus, at $t \sim 10^4-10^5\ s$, where these conditions are typically valid, KN effects are negligible and $E_{B,d^{-2}Y_\ast}^{1/2}$ is of order unity (for the same reasons discussed above Equation (8) in the context of ISM). Thus, observations of 10 keV photons during the first day imply

$$\eta \lesssim 30 \left( \frac{h_{\text{obs}}}{10\, \text{keV}} \right)^{-1} \left( \frac{E_{53}}{A_{*,4}} \right)^{1/4} B_{u,\mu G,19}^{-1/4}. \quad (13)$$

This result is similar to the one obtained in the ISM. It implies that the conclusion that the observed X-ray afterglows indicate on a fast acceleration is largely independent of the circumburst density profiles. The same is applicable to the conclusion that currently there is no strong indication for amplification of the magnetic field in the shock upstream. Note that this limit is valid only of the X-rays that are observed to be above the cooling frequency. If X-ray photons are not cooling over the system dynamical time (e.g., due to a very low value of $\epsilon_B$) then the confinement limit, which require that the X-ray-emitting electrons spend less time than the dynamical time in the upstream, is more constraining than the IC limit and should be used instead.

The confinement limit is unimportant at early time, but it becomes more stringent with time and may become the dominant limit at very late time, $t \sim 10^7\ s$ or even earlier if $\epsilon_B \ll 10^{-3}$. X-ray afterglows that are observed at such late time are very rare, but they do exist, as we discuss in the following section.

4. GRB 060729

GRB 060729 is the burst (at $z = 0.54$) with the latest X-ray detection, 642 days after the burst (Grupe et al. 2010, hereafter G10). The late time X-ray emission show a temporal break, from $F_\nu \propto t^{-(1.32 \pm 0.14)}$ to $F_\nu \propto t^{-(1.61 \pm 0.16)}$, roughly $10^6\ s$ after the explosion. At the same time the X-ray spectrum varies from $F_\nu \propto \nu^{-(1.18 \pm 0.11)}$ to $F_\nu \propto \nu^{-(0.89 \pm 0.11)}$. This simultaneous temporal break and spectral hardening fits very well (within 1σ) a passage of the cooling frequency, which increases with time, through the X-ray band (G10). This behavior of increasing cooling frequency is expected in a wind external medium. G10 find that a model of a spherical blast wave in a wind profile medium, where $E = 10^{53}\ erg$ (isotropic equivalent), $A_{53} = 0.1$, $\epsilon_B = 0.003$, and $\epsilon_e = 0.1$, fits the data well until $t \approx 4 \times 10^7\ s$, when a very sharp temporal break is observed. The spectral evolution during this late break is hard to constrain, due to the faintness of the signal, but it shows indications of softening.

The origin of the late temporal break is not well determined. G10 discuss two possible origins—a jet break or a break in the electron distribution. They find that it is hard to reconcile the late break with a jet origin, although they cannot rule it out. On the other hand, a spectral origin can provide a more consistent explanation. In that case the most natural source of the temporal break is the inability of the shock to accelerate X-ray-emitting particles. In that case these observations provide the first direct measurement (not only an upper limit) of $\eta$. Note that according to the model of G10, the blast wave is still relativistic even a year after the burst, $\Gamma \approx 5$, due to the large blast wave energy and low external density. At late time, when the cooling frequency is above the X-ray band, the limit on acceleration of X-ray-emitting electrons must be due to confinement. Thus if indeed the late break in the afterglow, at $t \approx 4 \times 10^7\ s$, is due to limited acceleration then

$$\eta \approx 100 A_{53}^{1/4} \left( \frac{E_{B,d}}{0.003} \right)^{1/4} B_{u,\mu G,19}. \quad (14)$$
Hence, unless the upstream field is significantly amplified by the shock precursor, acceleration cannot be very slow also when $\Gamma \approx 5$. Moreover, if $B_{\mu G,19} \lesssim 0.1$ then the acceleration must be very fast and the origin of the observed break is almost certainly due to limited acceleration. If $B_{\mu G,19} \gtrsim 1$, and the break is due to limited electron acceleration, then $\eta \gtrsim 100$ which is significantly larger than the value suggested by earlier X-ray ($\sim$day) and GeV ($\sim 10^3$ s) observations. This may suggest that the efficiency of particle acceleration is reduced when the shock approaches mildly relativistic velocities. Finally, even if the break is not related at all to electron acceleration then the equality in Equation (14) becomes an upper limit on $\eta$.

5. SUMMARY

In this paper, we examined the constraints that GRB afterglow observations place on the acceleration rate, within the DSA framework, in relativistic, weakly magnetized, collisionless shocks in ion–electron plasma. We examine shocks that propagate into a constant density medium (ISM) and into a decreasing density of massive stellar wind. We consider three major factors that limit the acceleration in such shocks: confinement, synchrotron cooling, and IC cooling. We find that at early times ($\sim 10^3$ s) the best limits are set by synchrotron cooling of GeV emitting electrons while at intermediate times ($\sim 10^7$ s) IC cooling of X-ray-emitting electrons provides the best constraints. These results are independent of the circumburst medium density profile. At very late time ($\gtrsim 10^8$ s) confinement may become the dominant factor in a wind environment while IC cooling remains the dominant factor that limits the acceleration in ISM.

Examining available observations, the tightest limits are obtained by GeV photons that are observed 100–1000 s after the burst, if these are synchrotron photons from the external shock. The origin of these photons is not determined yet, but they are seen long after the prompt emission fades and therefore most likely originate in the external shock. Various modelings of the GeV emission find that synchrotron emission can explain the observations well (Kumar & Barniol Duran 2009, 2010; Ghisellini et al. 2010). If this is true, then the observed GeV emission requires an extremely fast acceleration, at the Bohm limit or faster, i.e., $\eta \lesssim 1$. This limit is very robust since it is almost independent of any of the shock parameters such as energy, density, etc. On timescales of 100–1000 s the shock Lorentz factor is $\sim 100$

X-ray ($\sim 10$) keV photons are regularly observed on timescales of hours–days, where $\Gamma \sim 10$–50. The IC cooling of these photons also provides a tight limit on the acceleration rate: $n \lesssim 15 B_{\mu G,19}^{-5/8}$ in an ISM (cf. Li & Waxman 2006) and $n \lesssim 30 B_{\mu G,19} A_{-3/4}$ in a wind. On the one hand, these limits are less robust than those obtained by the GeV photons, due to the uncertainty in $B_{\mu G}$ and the external density, but on the other hand the certainty that these X-ray photons are emitted by synchrotron process in the external blast wave is much higher. Note that if, as suggested by the GeV data (and by the observations of other acceleration sites such as the Crab nebula; Abdo et al. 2011), acceleration can be as fast as $\eta \sim 1$, then the available X-ray observations may be consistent with no amplification of the magnetic field in the shock upstream region (contrary to previous conclusions of Li & Waxman 2006).

Finally, GRB 060729 is the burst with the longest duration X-ray afterglow observed to date. Its afterglow shows a sharp decline in the integrated X-ray flux $4 \times 10^7$ s after the burst (G10). This decline is most likely accompanied by a spectral softening. This decline can be explained naturally if the synchrotron frequency of the maximally accelerated electron is crossing the X-ray band at $t \approx 4 \times 10^7$ s. If this is the case then this is a direct measurement of $\eta$. The afterglow light curve is consistent with a wind circumburst density and a cooling frequency that crosses the X-ray band at $t \approx 10^6$ s, implying that at later time the X-ray emission can be limited only by confinement. Using the fit of G10 to the afterglow parameters, the shock Lorentz factor at the time of the fast decline is $\Gamma \approx 5$ and $\eta \sim 100 A_{-1/4}^{1/4} (\epsilon_{B,d}/0.003)^{1/4} B_{\mu G,19}$. If the fast decline at $t > 4 \times 10^7$ s is not due to shock acceleration limit then the equality becomes an upper limit. These results suggest that the acceleration rate remains rather fast also at lower Lorentz factors. To conclude, we find that GeV and X-ray afterglow observations provide independent limits on $\eta$. The combination of these limits strongly suggest that particle acceleration is fast in relativistic, weakly magnetized, collisionless shocks in ion–electron plasma. Namely, diffusion in the shock upstream and downstream regions takes place close to the Bohm limit at $\Gamma \approx 100$ and it remains fast during the shock deceleration, at least up to $\Gamma \approx 5$.

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