SPECTRUM AND DURATION OF DELAYED MeV–GeV EMISSION OF GAMMA-RAY BURSTS IN COSMIC BACKGROUND RADIATION FIELDS

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ABSTRACT

We generally analyze prompt high-energy emission above a few hundred GeV due to synchrotron self-Compton scattering in internal shocks. However, such photons cannot be detected because they may collide with cosmic infrared background photons, leading to electron/positron pair production. Inverse Compton scattering of the resulting electron/positron pairs off cosmic microwave background photons will produce delayed MeV–GeV emission, which may be much stronger than a typical high-energy afterglow in the external shock model. We expand on the model of Cheng & Cheng by deriving the emission spectrum and duration in the standard fireball shock model. A typical duration of the emission is \( \sim 10^3 \) s, and the time-integrated scattered photon spectrum is \( \nu^{-\frac{p+6}{4}} \), where \( p \) is the index of the electron energy distribution behind internal shocks. This is slightly harder than the synchrotron photon spectrum, \( \nu^{-\frac{p+2}{2}} \). The lower energy property of the scattered photon spectrum is dependent on the spectral energy distribution of the cosmic infrared background radiation. Therefore, future observations on such delayed MeV–GeV emission and the higher energy spectral cutoff by the *Gamma-Ray Large Area Space Telescope* (GLAST) would provide a probe of the cosmic infrared background radiation.

Subject headings: diffuse radiation — gamma rays: bursts — gamma rays: theory

1. INTRODUCTION

Gamma-ray bursts (GRBs) are the brightest electromagnetic phenomena in the universe. Their radiation features have been well understood even though their origin has been unknown. In the standard fireball shock model (Piran 1999; van Paradijs, Kouveliotou, & Wijers 2000; Mészáros 2002), GRBs are explained to be caused by the dissipation of kinetic energy of an expanding fireball with an average Lorentz factor more than 100 in the internal shocks produced by collisions between different shells in the fireball. Afterglows are also considered to be due to the dissipation of kinetic energy in the external shocks generated by collision of the fireball with its surrounding medium. Such a model has been supported by a variety of observations on GRBs and afterglows (e.g., Wijers, Rees, & Mészáros 1997; Waxman 1997; Vietri 1997). Theoretically, it has also predicted that inverse Compton scattering of the resulting electron/positron pairs off CMB photons will produce delayed MeV–GeV emission by deriving the time-integrated spectrum of such emission and by estimating the emission duration. In § 3 we summarize our findings and discuss their implications. Cheng & Cheng (1996) have discussed similar delayed MeV–GeV emission, and here we further calculate the emission spectrum and estimate the emission duration in the standard fireball shock model.

2. HIGH-ENERGY RADIATION FROM INTERNAL SHOCKS

Let us consider a compact source that produces a wind (i.e., a variable fireball) with an average luminosity of \( L \) and with an average mass-loss rate of \( \dot{M} = L/\eta c^2 \). Initially, the bulk Lorentz factor of the wind, \( \Gamma \), increases linearly with radius, until most of the wind energy is converted to kinetic energy and \( \Gamma \) saturates at \( \Gamma \sim \eta \). After this time, fluctuations of \( \Gamma \) due to variabilities in \( L \) and/or \( \dot{M} \) on timescale \( \Delta t \) could lead to internal shocks in the wind at radius \( R \sim \Gamma^2 c \Delta t \). We also assume that the internal shock Lorentz factors \( \Gamma_i \) are approximately a few in the wind rest frame, so that the internal shocks could reconvert a significant fraction of the kinetic energy to internal energy, which is then emitted as \( \gamma \)-rays by synchrotron and inverse Compton (IC) radiation of the electrons accelerated by the shocks.

The comoving electron number density of the unshocked wind matter at radius \( R \) is \( n_e' = L/(4\pi \Gamma^8 \Delta t^2 m_p c^5) \), where \( m_p \) is the proton mass. The internal energy density of the shocked wind matter is thus approximated by \( e' \approx \frac{1}{2} \Gamma_i (\Gamma_i - 1) n_e' m_p c^2 \) (Blandford & McKee 1976). The electron energy distribution behind the shocks is usually a
power law: \( dn_e/d\gamma_e \propto \gamma_e^{-\gamma} \) for \( \gamma_e \geq \gamma_m \). Assuming that \( \epsilon_e \) and \( \epsilon_B \) are constant fractions of the internal energy density going into the electrons and the magnetic field, respectively, we obtain the minimum electron Lorentz factor \( \gamma_m = ((p-2)/(p-1))(m_p/m_e)c_0^2/(\epsilon_e) \) (where \( m_p \) is the electron mass) and the magnetic field strength of the shocked wind matter \( B_0 = (8\pi \rho_e e^3)^{1/2} = 6.2 \times 10^{13} \text{G} \) where \( \epsilon_e = \epsilon_{1014} \times 10^{-13} \times 10^{-2} \times \Gamma_{500}^{1/2} \Delta \tau_{-2}^{-1/2} \text{G} \), with \( \Gamma_{500} = \Gamma_{500}/500 \), \( L_{52} = L/10^{52} \) ergs s\(^{-1}\), and \( \Delta \tau_{-2} = \Delta t/10^{-2} \text{s} \). Thus, we estimate the characteristic frequency of synchrotron radiation from the shock-accelerated electrons,

\[
\nu_m = \Gamma_{500}^{2} \epsilon_B \frac{eB}{2\pi m_e c} = 3.0 \times 10^{19} \left( \frac{p-2}{p-1} \right)^{2} \epsilon_{0.5}^{0.5} \frac{1}{\Gamma_{500}^{1/2} \Delta \tau_{-2}^{-1/2}} \times \left( \Gamma_{500}^{1/2} \Gamma_{500}^{2} \Delta \tau_{-2}^{-1/2} \text{Hz} \right),
\]

(1)

where \( \epsilon_{0.5} = \epsilon_e/0.5 \). According to Sari, Piran, & Narayan (1998), the cooling Lorentz factor \( \gamma_c \ll \gamma_m \) for typical parameters of the internal shock model, implying that the shock-accelerated electrons are in the fast cooling regime. Therefore, the synchrotron spectral luminosity is

\[
L_{\nu} \equiv \frac{\hbar \nu N_{\nu}}{\Delta t} = \left\{ \begin{array}{ll}
L_{\nu m} (\nu/\nu_m)^{-1/2} & \text{if } \nu < \nu_m , \\
L_{\nu m} (\nu/\nu_m)^{-p/2} & \text{if } \nu > \nu_m ,
\end{array} \right.
\]

(2)

where \( \hbar \) is Planck’s constant and \( N_{\nu} \) is the synchrotron photon number per unit frequency in \( \Delta t \).

Although GRBs are optically thin to electron scattering, some synchrotron photons will Compton scatter on the shock-accelerated electrons, producing an additional IC component at higher energies (Papathanassiou & Mészáros 1996; Pilla & Loeb 1998; Panaitescu & Mészáros 2000). The ratio of the IC to synchrotron radiation luminosity can be estimated by

\[
Y \equiv \frac{L_{IC}}{L_{syn}} = \frac{L_{IC}}{L_{\nu m} \nu_{m} \nu_{m}^{-1/2}} = \left( \frac{\epsilon_e}{\epsilon_B} \right)^{1/2} \gamma_m^{2},
\]

(3)

where we have assumed that \( \epsilon_e \gg \epsilon_B \) and the radiating electrons are in the fast cooling regime (Panaitescu & Kumar 2000; Sari & Esin 2001; Zhang & Mészáros 2001). Since the characteristic IC frequency

\[
\nu_{m}^{IC} \simeq \gamma_m^{2} \nu_{m} = 1.0 \times 10^{16} \left( \frac{p-2}{p-1} \right)^{4} \epsilon_{0.5}^{0.5} \Gamma_{500}^{1/2} \Delta \tau_{-2}^{-1/2} \left( \Gamma_{500}^{1/2} \right)^{2} \Delta \tau_{-2}^{-1/2} \text{Hz},
\]

(4)

the ratio of the IC spectral luminosity at \( \nu_{m}^{IC} \) to synchrotron spectral luminosity at \( \nu_{m} \) is

\[
\frac{L_{IC}}{L_{\nu m}} = \gamma_m^{2} Y,
\]

(5)

which is derived from equation (3), and the ratio of the IC to synchrotron photon number per unit frequency in \( \Delta t \) is further given by

\[
\frac{N_{IC}}{N_{\nu}} = \gamma_m^{2} Y, \text{ for } \nu \geq \nu_{m}^{IC}.
\]

(6)

The intrinsic cutoff energy in the IC component is defined by \( \epsilon_{cut}^{IC} = \min(\epsilon_{KN}^{IC}, \epsilon_{0}^{IC}) \), where \( \epsilon_{KN}^{IC} \) is the Klein-Nishina limit and \( \epsilon_{0}^{IC} \) is the energy at which a photon may be attenuated because of pair production through interactions with softer photons (also radiated from the internal shocks) whose frequency is equal to or larger than \( \nu_{cut} = (\Gamma_{500} e^3 c^2)/(\hbar \epsilon_{0}^{IC}) \). We first estimate \( \epsilon_{KN}^{IC} \). There is the maximum Lorentz factor \( \gamma_m \) in the electron energy distribution behind the shocks due to the fact that the shock-acceleration time cannot exceed the cooling time. This gives \( \gamma_m \sim \{3 \epsilon_e/\sigma_T(B(1 + Y))\}^{1/2} \), where \( \sigma_T \) is the Thomson cross section. Thus, we obtain the Klein-Nishina limit \( \epsilon_{KN}^{IC} = \gamma_m \Gamma_{500} e^3 c^2 \simeq 560 \epsilon_{0.5}^{1/2} (\Gamma_{500}^{1/2} \Delta \tau_{-2}^{-1/2} \text{TeV} \). To compute \( \epsilon_{0}^{IC} \), we adopt an analytical approach proposed by Lithwick & Sari (2001). Integrating the synchrotron spectral luminosity over frequency should lead to the total synchrotron luminosity, which further gives the synchrotron photon number per unit frequency at \( \nu_{cut} \): \( N_{\nu_{cut}} = \{(p-2)/(2[p(p-1)])\} \epsilon_e(1+Y)^{-1} (\hbar \nu_{cut}^{3})^{-1} L \Delta t \). The total synchrotron photon number at frequencies lower than \( \nu_{cut} \) can be estimated by \( N_{\nu_{cut}} \int_{\nu_{cut}}^{\infty} N_{\nu}d\nu = \int_{\nu_{cut}}^{\infty} N_{\nu m}(\nu/\nu_{m})^{-p/2}d\nu = ((p-2)/(p-1)) \epsilon_e(1+Y)^{-1} (\hbar \nu_{cut}^{3})^{-1} L \Delta t \), where the cutoff energy is given by equation (8).

3. Pair Production and Relaeradation in Cosmic Background Photon Fields

High-energy \( \gamma \)-rays emitted from the internal shocks may not only be intrinsically attenuated due to pair production through interactions with softer photons from the same radiation regions but may also be absorbed in the external background radiation fields when these \( \gamma \)-rays travel toward the observer. In the latter case, the observed cutoff energy, \( \epsilon_{cut} \), is determined by the redshift \( z \) and the comoving energy, \( \epsilon_{c}(z) \), of the background photons that dominate the pair production optical depth \( \tau_{\gamma}^{opt} \) because the cross section peaks when \( \epsilon_{cut}^{opt} \epsilon_{c}(z) \sim (m_e e^2)^2/(1+z) \). In general, that at low redshifts the spectral energy distribution of the starlight background radiation peaks at infrared wavelengths and thus \( \epsilon_{cut}^{opt} \) is in the TeV energy range (Salamon & Stecker 1998). Detailed predictions of \( \epsilon_{cut}^{opt} \) differ from model to model, depending on the theoretical treatment adopted for the stellar emissivity. For example, the observed cutoff energy when \( \tau_{\gamma}^{opt} = 1 \) for \( z = 1 \) is \( 50 \text{GeV} \leq \epsilon_{cut}^{opt} \leq 80 \text{GeV} \) for different models used in Salamon & Stecker (1998).

The pair production optical depth, \( \tau_{\gamma}^{opt} \), strongly depends on the \( \gamma \)-ray energy \( (\epsilon_{cut}^{opt}) \). Salamon & Stecker (1998) numerically calculated \( \tau_{\gamma}^{opt} \) as a function of the photon energy for several fixed redshifts by considering the stellar emissivity with and without metallicity correction, shown in their Figures 6 and 7, respectively. It can be seen from their Figure 6 that if \( z = 1 \), then \( \tau_{\gamma}^{opt} \simeq 1 \) for \( \epsilon_{cut}^{opt} = 50 \text{GeV} \), but \( \tau_{\gamma}^{opt} \simeq 10 \) for \( \epsilon_{cut}^{opt} = 300 \text{GeV} \). In this model of Salamon & Stecker (1998), therefore, a photon with energy of \( \epsilon_{cut}^{opt} \approx 300 \text{GeV} \) must have a pair production optical depth
of \( \tau_{\gamma\gamma} \gg 1 \). This implies that such high-energy \( \gamma \)-rays may be locally attenuated once they are radiated from internal shocks. For simplicity, we will neglect any redshift correction for low-redshift sources discussed in this paper.

The resulting electron/positron pairs have Lorentz factors of \( \gamma_e \equiv \delta_e^0 / (2m_e c^2) \geq \gamma_{\min} \equiv \delta_0^0 / (2m_e c^2) = 3 \times 10^3 \delta_{0,300}^0 \), where \( \delta_0^0 = \delta_{0,300}^0 \times 300 \) GeV is the minimum energy of photons that are locally attenuated in the external infrared background radiation field. The pairs will Compton scatter the CMB photons. As a result, the initial energy of a microwave photon, \( h \nu_0 \), is boosted by IC scattering up to an average value \( \sim \gamma_e^2 h \nu_0 \sim 57(\delta_{0,300}^0)^2 \) MeV, where \( h \nu_0 = 2.7kT \) is the mean energy of the CMB photons with \( T \simeq 2.73 \) K and \( k \) is the Boltzmann’s constant. The IC lifetime (in the local rest frame) of an electron with a Lorentz factor of \( \gamma_e \) reads

\[
\tau(\gamma_e) = \frac{3m_e c}{4\gamma_e e \sigma_T u_{\text{cmb}}} \leq 2.4 \times 10^{14} (\delta_{0,300}^0)^{-1} \text{s}, \tag{9}
\]

where \( u_{\text{cmb}} = aT^4 \) is the CMB energy density and \( a \) is the radiation constant. It is not difficult to find that the typical length in which most of the electron energy is lost, \( \sim c \tau(\gamma_e) \), is much less than the distance from the source to the observer, implying that energy loss of the electron due to IC scattering is also local. Beyond this length, energy loss of the electron and its emission become insignificant in the absence of any acceleration.

The typical duration of the scattered photons after the GRB trigger in the observer’s frame is estimated by \( \tau_{\text{obs}} = \max(\tau_{\text{obs}}^1, \tau_{\text{obs}}^2) \), where \( \tau_{\text{obs}}^1 \) is the observed IC cooling lifetime,

\[
\tau_{\text{obs}}^1 \simeq \frac{\tau(\gamma_e)}{2\gamma_e^2} \leq 1.3 \times 10^3 (\delta_{0,300}^0)^{-3} \text{s}, \tag{10}
\]

and \( \tau_{\text{obs}}^2 \simeq R_{\text{pair}} / (2\gamma_e c) \) is the angular timescale (Piran 1999). Here \( R_{\text{pair}} = (0.26 \sigma_T \Gamma T^2)^{-1} \simeq 5.8 \times 10^{24} (n_{\text{IR}} / 1 \text{ cm}^{-3})^{-1} \text{cm} \) is the typical pair production radius (where \( n_{\text{IR}} \sim 1 \text{ cm}^{-3} \) is the cosmic infrared photon number density; Protheroe & Stanee 1993). Thus, we have

\[
\tau_{\text{obs}}^2 \leq 1.0 \times 10^3 (n_{\text{IR}} / 1 \text{ cm}^{-3})^{-1} (\delta_{0,300}^0)^{-2} \text{s}. \tag{11}
\]

Therefore, most of the scattered photons will reach the observer in \( \sim 10^3 \) s. This time is much longer than a typical GRB duration.

We next derive the scattered photon spectrum. From equations (2), (3), and (6), we first write the spectrum of the electrons and positrons as follows:

\[
\frac{dN_e}{d\gamma_e} = 2\gamma_e^{p-2} Y_{\nu_m} N_{\nu_0} \left( \frac{h \nu_m}{2m_e c^2} \right)^{p/2} \gamma_e^{-(p+2)/2},
\]

if \( \gamma_{\min} \leq \gamma_e \leq \gamma_{\max} \), \tag{12}

where \( \gamma_{\max} = \delta_{0,300}^0 / (2m_e c^2) \). The scattered photon number per unit frequency is given by (Blumenthal & Gould 1970, hereafter BG70)

\[
N_{\nu}^\text{sc} = \frac{dN_{\nu}}{d\gamma_e} \frac{dN_{\gamma_e}}{dt} \tau(\gamma_e) d\gamma_e, \tag{13}
\]

where \( dN_{\nu_e} / dt \; d\nu \), expressed by equation (2.42) of BG70 in the Thomson limit, is the spectrum of photons scattered by an electron with Lorentz factor of \( \gamma_e \) from a segment of the CMB photon gas of differential number density \( n(\epsilon) d\epsilon \), and \( t \) is the time measured in the local rest frame. Please note that the scattered photon number per unit frequency given by equation (13) is time-integrated by multiplying \( \tau(\gamma_e) \) in equation (2.61) of BG70. After integrating over \( \gamma_e \) and \( \epsilon \), we obtain the ratio of the scattered to synchrotron photon number per unit frequency,

\[
\frac{N_{\nu}^\text{sc}}{N_{\nu}} = \frac{9(m_e c^2)^{-3(p-2)/2}}{2^{p+8}/2^2 \pi^2 \left( (hc)^2 \right) aT^4} \gamma_e^{p-2} \frac{Y_F(q)(kT)^{(p+14)/4} (h \nu)^{(p-2)/4}}{F(3) \left( (m_e c^2)^2 \right)} , \tag{14}
\]

where \( F(q) \) is a function of \( q = (p+4)/2 \),

\[
F(q) = \frac{2^{q+1} (q+4 q+11)}{(q+3)^2 (q+1)(q+5)} \Gamma[(q+5)/2] \zeta[(q+5)/2], \tag{15}
\]

with the Gamma function and Riemann zeta function. Therefore, the scattered photon spectrum becomes

\[
N_{\nu}^\text{sc} \propto \nu^{-(p+6)/4} . \tag{16}
\]

The lower energy limit of the scattered photons is

\[
h \nu_1 \sim \gamma_{\min}^2 (2.7kT) = 57(\delta_{0,300}^0)^2 \text{ MeV}, \tag{17}
\]

while the upper energy limit is

\[
h \nu_2 \sim \gamma_{\max}^2 (2.7kT) \simeq 16 \epsilon_{0.5}^{-2(2p-3)/p} \epsilon_{B,-2}^{-1} (\Gamma_1/2)^{-5(p-2)/p} \Gamma^{(8(p+1))/p} L_{52}^{-(p+2)/p} A_{-2}^2 \text{ MeV}, \tag{18}
\]

which is derived from equation (8).

We now consider a simple example. The typical parameters of the internal shock model are taken: \( \epsilon_{e,0.5} = \epsilon_{B,-2} = L_{52} = A_{-2} = \Gamma_1/2 = 1, \; \Gamma_{500} = 1.4, \) and \( p = 2.2 \). Furthermore, \( \delta_{0,300}^0 = 1 \) is assumed. Therefore, we obtain \( h \nu_1 \sim 57 \) MeV and \( h \nu_2 \sim 0.8 \) GeV. From equation (14), we further find \( N_{\nu}^\text{sc} / N_{\nu} \) increases from \( \sim 5.0 \) to \( \sim 5.8 \) when the observed energy of the scattered photons increases from \( h \nu_1 \) to \( h \nu_2 \). This leads to a delayed MeV–GeV emission component. It is interesting to note that the total energy release of such emission is not only comparable to that of a typical GRB, but also much larger than the total MeV–GeV energy release in an early afterglow as calculated by Dermer, Chiang, & Mitman (2000) and Zhang & Mészáros (2001).

4. DISCUSSION AND CONCLUSIONS

We have discussed prompt high-energy emission from the internal shocks, which are produced by collisions between shells with different Lorentz factors in an ultrarelativistic wind ejected from the central engine of GRBs. Such emission may arise from synchrotron self-Compton scattering in the internal shocks. We derived the intrinsic cutoff energy due to pair production through interactions of high-energy photons with softer photons from the same radiation regions. Our intrinsic cutoff energy is consistent with the observed spectra of low-redshift GRBs such as GRB
940217 and GRB 970417a. However, another cutoff energy appears in the $\gamma$-ray spectra. This cutoff results from interactions of high-energy photons with external cosmic infrared background photons. The electron/positron pairs produced during such interactions own a significant fraction of the explosion energy, and they will Compton scatter the CMB photons. We expanded on the Cheng & Cheng (1996) model by deriving the emission spectrum and duration in the standard fireball shock model. The resulting emission will be able to reach the observer in $\sim 10^3$ s. The time-integrated, scattered photon spectrum is $N_\nu \propto \nu^{-(p+6)/4}$. This is slightly harder than the internal shock emission spectrum, $N_\nu \propto \nu^{-(p+2)/2}$. For typical parameters of the internal shock model, the observed energies of the scattered photons are in the MeV–GeV range, and the ratio of the scattered to synchrotron photon number per unit frequency, $N_\nu/N_\nu$, is typically a few.

Of course, whether or not there is such delayed emission depends on the parameters of the internal shock model. Some parameters (e.g., $L$, $\epsilon_e$, and $\epsilon_B$) adopted in this paper are consistent with detailed fits to the multi wavelength data of a few afterglows (Wijers & Galama 1999; Freedman & Waxman 2001; Panaitescu & Kumar 2002). Two other parameters (e.g., $p$ and $\Delta t$) are consistent with the BATSE observations (Preece et al. 2000). The bulk Lorentz factor of the wind, $\Gamma$, is a crucial parameter because $\delta^\text{in}_\gamma$ is strongly dependent on $\Gamma$ (see eq. [8]). The observed TeV emission from GRB 970417a requires that $\Gamma$ be equal to or larger than 700, provided that the other parameters are given in the simple example of § 3. Furthermore, some arguments on the optical flash of GRB 990123 in the reverse shock model (Wang, Dai, & Lu 2000; Soderberg & Ramirez-Ruiz 2002) show that $\Gamma$ may be of the order of $10^3$ (but also see Sari & Piran 1999, who suggested that $\Gamma$ in this burst could be about 200). If $\Gamma$ is close to or larger than 600, we can indeed see a delayed MeV–GeV emission spectrum. If $\Gamma < 600$, however, the intrinsic cutoff energy could be less than the external pair production minimum energy at which a photon is locally attenuated through interaction with the infrared background radiation, so that there could not be the delayed MeV–GeV emission discussed here.

It is seen from equations (6) and (14) that the Compton parameter $Y$ determines the TeV emission from synchrotron self-Compton scattering in internal shocks, and thus the spectral photon number of the delayed MeV–GeV emission. In addition, additional TeV photons produced by other mechanisms, e.g., photo-pion and inelastic proton-neutron collisions (Waxman & Bahcall 1997; Böttcher & Dermer 1998; Derishev, Kocharovsky, & Kocharovsky 1999; Bahcall & Mészáros 2000; Dai & Lu 2001), and synchrotron self-Compton scattering both in external reverse shocks (Wang, Dai, & Lu 2001a, 2001b) and in early-time forward shocks (Dermer, Chiang, & Mitman 2000; Zhang & Mészáros 2001), may be absorbed in the cosmic infrared photon field, leading to electron/positron pairs. The inverse Compton scattering of such pairs off CMB photons may have a non-negligible contribution to the delayed MeV–GeV emission.

The MeV–GeV emission studied in this paper will be detectable by next-generation $\gamma$-ray satellites such as the Gamma-Ray Large Area Space Telescope (GLAST). In the Swift-GLAST era, plenty of GRBs localized by the Swift satellite are detectable by GLAST. Such observations, particularly on the delayed MeV–GeV emission and higher energy spectral cutoff, may provide a probe of the cosmic infrared background radiation. This in turn may help to constrain some of the most fundamental uncertainties in physical models of the star formation.

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