MeV-GeV EMISSION FROM NEUTRON-LOADED SHORT GAMMA-RAY BURST JETS

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ABSTRACT

Recent discovery of the afterglow emission from short gamma-ray bursts suggests that binary neutron star or black hole–neutron star binary mergers are the likely progenitors of these short bursts. The accretion of neutron star material and its subsequent ejection by the central engine implies a neutron-rich outflow. We consider here a neutron-rich relativistic jet model of short bursts and investigate the high-energy neutrino and photon emission as neutrons and protons decouple from each other. We find that upcoming neutrino telescopes are unlikely to detect the 50 GeV neutrinos expected in this model. For bursts at $z \sim 0.1$, we find that the Gamma-Ray Large Area Space Telescope (GLAST) and ground-based Cerenkov telescopes should be able to detect prompt 100 MeV and 100 GeV photon signatures, respectively, which may help test the neutron star merger progenitor identification.

Subject headings: gamma rays: bursts — gamma rays: theory — ISM: jets and outflows — radiation mechanisms: nonthermal — stars: neutron

1. INTRODUCTION

The nature and origin of short gamma-ray bursts (SGRBs), with $t_c \lesssim 2$ s duration, were largely conjectural until recently. This changed dramatically with the detection of the afterglow and the identification of the host galaxy of GRB 050509b at a redshift $z = 0.226$ (Gehrels et al. 2005), as well as subsequent observations of other short-burst afterglows, e.g., GRB 050709 (Fox et al. 2005; Villasenor et al. 2005), GRB 050724 (Barthelmy et al. 2005), etc. These observations revealed that the short bursts are located mainly in elliptical hosts, while some are in spirals or irregulars, as expected from an old population such as neutron star–neutron star (NS-NS) or neutron star–black hole (NS-BH) mergers. They also provided, for the first time, details about the X-ray, and in some cases optical and radio afterglow light curves, including in two cases evidence for a jet break, qualitatively similar to the relativistic jet afterglows of long GRBs (see, e.g., Zhang & Mészáros 2004) for reviews of long GRBs). The average redshifts are lower than for long bursts, indicating an isotropic equivalent energy of SGRBs approximately 2 orders of magnitude lower than for long ones.

We investigate here a model of an SGRB jet that is initially loaded with neutrons and fewer protons, as expected from a neutron star composition. High-energy emission in this model is then expected from the photosphere of the relativistic jet outflow, including both thermal and nonthermal components, in addition to the radiation expected outside of it. We describe our jet model in § 2, high-energy emission processes in § 3, and expected neutrino and photon fluxes in § § 4 and 5, respectively. We summarize and discuss our results in § 6.

2. NEUTRON-LOADED JET DYNAMICS

We assume the total isotropic equivalent energy outflow in the SGRB jet of $L = 10^{50} L_{50}$ ergs s$^{-1}$. The jet is loaded with protons and neutrons with an initial neutron-to-proton number density ratio $n_n/n_p = 10\xi_{o,1}$ in the comoving frame. The total mass outflow rate in such a neutron-rich jet is $\dot{M} = 4\pi r^2 c (1 + \xi_o) n_n n_p$, initially, since $m_p \approx m_n \gg m_e$, and the thermal energy is negligible. A useful quantity is the total energy to mass flow ratio $\eta = L/Mc^2$, characterizing the dimensionless entropy or the baryon loading of the jet. In a typical long GRB ($t_c \gtrsim 2$ s), the final bulk Lorentz factor of the gamma-ray-emitting fireball is equivalent to $\eta$. In the case of SGRBs, we assume $\eta = 316/\eta_{2.5}$, similar to that of the long GRBs, motivated by observation.

The outflow starts at a radius $R_0 = 10^6 R_{6,0} \text{ cm}$ which is a few times the Schwarzschild radius $r_g = 2GM_{BH}/c^2$ of a solar-mass black hole created by binary mergers. The initial temperature of the plasma outflow is

$$T_o' = \left( \frac{L}{4\pi R_0^2 c^4} \right)^{1/4} \approx 1.2 \left( \frac{L_{50}}{R_{6,0}^2} \right)^{1/4} \text{ MeV}, \quad (1)$$

where $a = \pi^2 k^4/15(hc)^3 = 7.6 \times 10^{15}$ ergs cm$^{-3}$ K$^{-4}$ is the radiation density constant. The optically thick, hot plasma expands adiabatically due to the radiation pressure. The comoving temperature drops as $T'(r) \propto r/R_0$, and the plasma expands with an increasing bulk Lorentz factor $\Gamma(r) \propto r/R_0$ with the radius $r$ following the adiabatic law. Electrons (both $e^+$ and $\nu_e \approx n_\gamma$), which are coupled to the photons via Compton scattering, and protons, which are coupled to the electrons via Coulomb interaction, are held together in the expanding, optically thick plasma. Neutrons, however, are somewhat loosely coupled to the protons by elastic nuclear scattering with a cross section $\sigma_{\nu p} \approx 3 \times 10^{-26}$ cm$^2$, which is roughly 1/20 of the Thomson cross section ($\sigma_{\text{Th}} \approx 6.65 \times 10^{-25}$ cm$^2$). This considerably changes the dynamics of a neutron-rich jet (e.g., Derishev et al. 1999a; Bahcall & Mészáros 2000; Beloborodov 2003) as we discuss next.

The protons and neutrons are coupled together with a common bulk Lorentz factor $\Gamma_{\nu p}(r) \approx \Gamma_{\nu p}(r) \approx \Gamma(r)$ as long as the $n-p$ collision time $t_{np} \approx (n_n \sigma_{np} c)^{-1}$ is shorter than the plasma expansion time $t_{exp} \approx r/c\Gamma(r)$. The critical dimensionless entropy for which $n-p$ decoupling happens, from the condition $t_{np} = t_{exp}$ is

$$\eta_{np} \approx \left[ \frac{L_{np}}{4\pi R_0^2 c^3 (1 + \xi_o)} \right]^{1/4} \approx 150 \left( \frac{L_{50}}{R_{6,0}^2 (1 + \xi_{o,1})} \right)^{1/4}. \quad (2)$$

1 We use primed variables in the relativistic plasma frame and unprimed variables in the observer or laboratory frame.
As a result, the neutrons and protons in our SGRB jet model ($\eta \approx \eta_{np}$) decouple at a radius where the nuclear scattering optical depth $\tau_{np} \approx n'_p \sigma_{np} R_{np}/\Gamma(R_{np}) = 1$ as

$$R_{np} \approx \left[ \frac{L \sigma_{np}}{4\pi R_0 n_{np} m_p c^3 (1 + \xi_o)} \right]^{1/3} R_0 = \frac{\eta_{np}}{\eta}^{1/3} \approx 1.2 \times 10^9 \left[ \frac{L_{50} R_{2.5}}{n_{np} \eta_{np} (1 + \xi_{o_1})} \right]^{1/3} \text{ cm.}$$

The bulk Lorentz factor of the neutrons does not increase further and saturates to a final value of

$$\Gamma_{n,f} = \frac{R_{np}}{R_0} \approx 115 \left[ \frac{L_{50}}{n_{np} \eta_{np} (1 + \xi_{o_1})} \right]^{1/3},$$

at $r = R_{np}$. The neutrons start to lag behind the protons at $r \approx R_{np}$, and the total energy outflow due to protons is

$$\dot{L} = L - \Gamma_{n,f} \dot{M} c^2 \xi_o = L \left( 1 - \frac{\Gamma_{n,f}}{\eta} \frac{\xi_o}{1 + \xi_o} \right) \approx 7 \times 10^{49} \text{ ergs s}^{-1},$$

which is close to the initial value of $L$, since $\Gamma_{n,f} < \eta$. Here $\dot{M} \approx M/(1 + \xi_o)$ is the proton mass outflow rate after decoupling. The outflow is still optically thick at $r = R_{np}$, since $\sigma_{th} \approx 20 \sigma_{np}$. The decoupling of radiation from plasma electrons takes place for a critical value for the dimensionless entropy, for which the Thomson scattering timescale $\hat{\tau}_{th} \approx (n'_e \sigma_{th} c)^{-1}$ is equal to the plasma expansion time, given by

$$\frac{\dot{\eta}}{\dot{\eta}_{had}} \approx \frac{\dot{L}/\sigma_{th} \eta_{np}}{4\pi R_0 n_{np} m_p c^3} \approx 5 \times 10^9 \left( \frac{L_{49.8} R_{28.6}}{R_{50.8}} \right)^{1/4} \frac{\eta}{\eta_{np}} \approx 4.4.$$  

This is close to $\dot{\eta}/\dot{\eta}_{had}$, so the outflow becomes transparent during the acceleration phase $\Gamma \propto r$, at a radius $r = R_\gamma$, where the Thomson optical depth $\tau_{th} \approx n'_e R_\gamma / \Gamma(R_\gamma)$ is 1, and

$$\dot{r} \approx R_\gamma \dot{\eta}_{had} \frac{\dot{\eta}_{had}}{\dot{\eta}} \approx 3.2 \times 10^9 \text{ cm},$$

which is close to $R_{np}$. The final Lorentz factor of the (proton) outflow is given by (Rossi et al. 2006)

$$\Gamma_{p,f} \approx \dot{\eta}_{had} \left( \frac{\dot{\eta}}{\dot{\eta}_{had}} \right)^{1/9} \approx 624,$$

which is achieved at a radius $r \sim R_\gamma$ for this $\eta$. This suggests a possible way to generate the high bulk Lorentz factor estimated...
in short GRBs without any spectral lags in the data (Norris & Bonnell 2006).

Note that for $\eta < \eta_{np}$, one has $\Gamma_{n,f} = \Gamma_{p,f} = \eta$ for a baryon-loaded jet outflow satisfying the condition $\eta < \eta_{had}$. Here $\eta_{had}$ is the initial critical radiation entropy for combined proton and neutron outflow, defined with an equation similar to equation (6) as

$$\eta_{had} = \left[ \frac{L_{\text{th}}}{4\pi R_{\text{th}} m_p c^3 (1 + \xi_o)} \right]^{1/4} \approx 322 \left[ \frac{L_{50}}{R_{60} (1 + \xi_{o1})} \right]^{1/4}. \quad (10)$$

We have plotted the final proton and neutron bulk Lorentz factors as functions of the baryon-loading parameter in Figure 1. For a given total energy outflow rate $L$, the maximum allowed values of $\eta$, from equation (10), are shown as three vertical solid lines for three values of initial neutron-to-proton ratio $\xi_o = 1, 5, 10$. Note that for $\eta \sim 316$, a typical value used for long GRB modeling, the neutron and proton components always decouple (see eq. [2]) in the case of SGRBs because of a lower $L$. On the contrary, $n$-p decouple only for $\xi_o \gtrsim 4$ in the case of a long GRB with $\eta \sim 316$, $L = 10^{52} L_{52}$ erg s$^{-1}$, and $R_{50} = 10^{10} R_{10}$ cm. Also note the high $\Gamma_{p,f}$ values of 519, 572, and 624 for $\xi_o = 1, 5,$ and 10, respectively, in Figure 1 from equation (9).

3. INELASTIC NEUTRON-PROTON SCATTERING

The elastic component of the total np cross section, which dominates at lower energy, drops rapidly above the pion production threshold of $\approx 140$ MeV, where the inelastic np scattering cross section is $\sigma_{np} \approx \sigma_{pp}$. Hence, the condition $\eta \gtrsim \eta_{np}$ (and subsequently $\hat{\eta} \gtrsim \eta_{had}$) for $n$-p decoupling implies that there will be pion production by the threshold processes, with roughly equal (total unit) probability to produce a $\pi^\pm$ pair and/or $2\eta^0$.

$$pn \rightarrow \begin{cases} nn \pi^+ & \pi^+ \rightarrow \mu^- \nu_\mu \rightarrow e^- \nu_e \bar{\nu}_\mu \nu_\mu, \\ pp \pi^- & \pi^- \rightarrow \mu^- \bar{\nu}_\mu \rightarrow e^- \bar{\nu}_e \nu_\mu, \\ pn \pi^0 & \pi^0 \rightarrow \gamma \gamma. \end{cases} \quad (11)$$

The secondary pions produced by the inelastic np scattering are at rest in the center-of-mass (c.m.) frame of interaction. However, the neutrons are not cold in the comoving outflow frame as we discuss next, and the c.m. frame acquires a velocity relative to this frame.

The Lorentz factor of the neutrons relative to the outflowing protons, at the radius $r \sim R_e \sim R_{np}$, is given by

$$\Gamma_{\text{rel}} = \Gamma_p \Gamma_n (1 - \beta_p \beta_n) \approx \frac{1}{2} \left( \frac{\Gamma_{p,f}}{\Gamma_{n,f}} + \frac{\Gamma_{n,f}}{\Gamma_{p,f}} \right) \approx 2.8. \quad (12)$$

Thus, in the (proton) outflow rest frame the neutron energy is $E_n = \Gamma_{\text{rel}} m_p c^2$ and the Lorentz factor of the c.m. is given by $\Gamma_{c.m.} = (\Gamma_{p,n} m_p c^2)/(m_p^2 c^4 + m_e^2 c^4 + 2 \mu_v m_p c^2)^{1/2} \approx 1.4$. The energies of the decay neutrinos are $\nu_\mu \approx 30$ MeV from $\pi^+$ and $\nu_e \approx 30$ MeV from $\pi^-$, and $\bar{\nu}_e \approx 30$ MeV and $\bar{\nu}_\mu \approx 50$ MeV from $\mu^\pm$. The $\pi^0$-decay photon energy is $\epsilon_\gamma \approx 70$ MeV. The observed energies of the decay products would be

$$\epsilon = \epsilon' \Gamma_{p,f} = \begin{cases} 26 \text{ GeV} & \nu_\mu \text{ and } \nu_e \text{ from } \pi^+ \text{ and } \mu^+, \\ 43 \text{ GeV} & \nu_\mu \text{ from } \mu^+, \\ 60 \text{ GeV} & \gamma \end{cases} \times \Gamma_{p,2.8}. \quad (13)$$

Here $\Gamma_{p,2.8} = \Gamma_{p,f}/624$.

Note that additional pions may be produced via $pp \rightarrow p\pi^+$ and $nn \rightarrow np^{-}$ inelastic interactions with cross sections similar to the inelastic np interaction. The $nn$ process in particular may be important for neutron destruction in the outflow (Rossi et al. 2006). Neutrinos and photons produced from $nn$ and $pp$ processes, however, are of much lower energy, as the relative velocities between the n-n or p-p components (e.g., arising from a variable outflow) may not exceed that between n-p in equation (12).

4. NEUTRINO EMISSION

The inelastic n-p scattering opacity is $\tau_2 \approx \tau_{np} \approx 1$ at $r \sim R_{np} \sim R_e$ and each proton produces, on average, a $\pi^\pm$ pair half of the time. The decay neutrinos escape the plasma outflow as soon as they are created. The neutrino and antineutrino flux rate from an SGRB at $z \approx 0.1$ (corresponding to a luminosity distance $D_L \approx 10^{27} D_{27}$ cm for a standard cosmology with Hubble constant $H_0 \approx 70$ km s$^{-1}$ Mpc$^{-1}$) is

$$\dot{N}_{\nu_e} \approx \dot{N}_{\bar{\nu}_e} \approx 2N_{\nu_e} \approx 2N_{\bar{\nu}_e} \approx \frac{\dot{L}}{4\pi D_L^2 \Gamma_{p,f} m_p c^2} \approx 5.7 \times 10^{-6} \frac{\dot{L}_{49.8}}{D_{27}^2 \Gamma_{p,2.8}} \text{ cm}^{-2} \text{ s}^{-1}. \quad (14)$$

The neutrino-nucleon ($\nu/N$) total cross section for detection is $\sigma_{\nu/N} \approx 6 \times 10^{-37}$ cm$^2$ at $E_\nu \approx 60$ GeV (see, e.g., Table 1 in Gandhi et al. 1998). Upcoming kilometer-scale water Cerenkov detectors such as ANTARES and KM3NeT are designed to have sensitivity for detecting neutrinos of energy $\gtrsim 50$ GeV. The total number of nucleons in a gigaton detector is $N_N \approx 1.2 \times 10^{39}$, which folded with the cross section and muon neutrino flux from a typical SGRB at $z \sim 0.1$ yields an event rate of

$$\dot{N}(\nu_\mu + \bar{\nu}_e \text{ event}) \approx 2N_{\nu_e} \sigma_{\nu/N} N_N \approx 0.01 \frac{\dot{L}_{49.8} D_{27}^2}{\Gamma_{p,2.8}} \text{ s}^{-1} \text{ Gton}^{-1}. \quad (15)$$

per burst. Hence, it is unlikely that upcoming neutrino telescopes would be able to detect these neutrinos from individual SGRBs, based on our model as discussed here. The number of events from diffuse flux can be calculated assuming a burst rate of $\sim 300/(4\pi) \text{ yr}^{-1}$ sr$^{-1}$ (approximately 30% of the long GRB rate) distributed over the whole sky and an optimistic $2\pi$ sr angular sensitivity of the detector as $300/(4\pi) \times 2\pi \times \dot{N}_{\nu_e} (\mu + \bar{\nu}_e \text{ event}) \approx 1 \text{ yr}^{-1}$. Here, we have assumed $t \sim 1 \text{ s}$ as the typical duration of the SGRBs.

5. PHOTON EMISSION

The $\gamma$-ray emission from an SGRB, however, is more promising. Although the plasma outflow becomes transparent to Compton scattering at a radius $r \sim R_e \sim R_{np} \sim 10^9$ cm, not all the high-energy photons due to $\pi^0$ decay can escape. Only those created below a skin depth of $R_e$ escape with a probability (Bahcall & Mészáros 2000)

$$P_{\pi^0} \approx \tau_{\pi^0} (R_e) \approx R_{np}/R_e \approx 0.4. \quad (16)$$

The corresponding number flux of 60 GeV $\pi^0$-decay photons at Earth is

$$\dot{N}_{\gamma,x0} \approx \frac{P_{\pi^0} \dot{L}}{4\pi D_L^2 \Gamma_{p,f} m_p c^2} \approx 2.0 \times 10^{-6} \frac{\dot{L}_{49.8}}{D_{27}^2 \Gamma_{p,2.8}} \text{ cm}^{-2} \text{ s}^{-1}. \quad (17)$$
which is below the detection threshold of the GLAST “Large Area Telescope” (LAT), whose effective area is $\sim 10^5$ cm$^2$ in this energy range (Gehrels & Michelson 1999). However, ground-based Cerenkov telescopes with $\lesssim 100$ GeV threshold may detect these photons. In particular, Milagro has an effective area of $\sim 5 \times 10^5$ cm$^2$ for a zenith angle $\lesssim 15^\circ$ in this energy range (Dingus 2004). For a yearly SGRB rate of $\sim 300$ yr$^{-1}$, this implies a possible detection rate of $\lesssim 5$ yr$^{-1}$, assuming a 90% duty cycle.

The high energy photons from $\pi^0$ decay created below $R_i$ interact with thermal photons of energy $\epsilon_{\gamma,\text{ph}} \approx T_{\text{ph}} R_i$ which are carried along with the plasma. The number density of these photons is high, and all $\pi^0$-decay photons of energy $\epsilon_{\gamma,\pi^0} \approx 100$ MeV originating below the $T_{\text{ph}} \approx 1$ skin depth produce $e^\pm$ pairs satisfying the condition $\epsilon_{\gamma,\pi^0} \gtrsim m_e c^2$. Assuming equal energy distribution among final particles, charged-pion decay $e^\pm$ and $\gamma \gamma \rightarrow e^\pm$ have a characteristic energy $\sim 35 m_e$ MeV and initiate electromagnetic cascades by inverse-Compton (IC) scattering of thermal photons, creating new generations of pairs, which interact again with the thermal photons. Contrary to the optically thick case of $T_{\text{Th}} \gg 1$, where all upscattered photons are absorbed by thermal photons producing $e^\pm$ pairs, these cascades are not saturated (Derishev et al. 1999a). The upscattered photons eventually escape from the photosphere with characteristic energy $\epsilon_{\gamma,\pi^0} \approx m_e c^2$, below the $e^\pm$ pair production threshold energy.

The fraction of the $\pi^\pm$ and $\pi^0$ energy carried by the escaping $\gamma$-rays depends on the details. An order of magnitude estimate for this fraction is $\xi_{\gamma} = 0.1 \xi_{\gamma,\pi^0}$ (Rossi et al. 2006). Thus, the number flux of the escaping $\gamma$-rays of energy

$$\epsilon_{\gamma,e} \approx \Gamma_{p,f} m_e c^2 \approx 320 \Gamma_{p,2.8} \text{ MeV}$$

at Earth is

$$N_{\gamma,e} \approx \frac{4L}{4\pi D_1^2 \Gamma_{p,f} m_e c^2} \left( \frac{\gamma_{\text{ph}} m_e c^2}{m_e c^2} \right)$$

$$\approx 8.6 \times 10^{-4} \frac{\hat{L}_{49.8} \xi_{\gamma,\text{Th}}^{-1}}{D_1^2 \Gamma_{p,2.8}} \text{ cm}^{-2} \text{ s}^{-1}.$$

The escaping–Comptonized-photon spectrum of the cascade is $dN_{\gamma,e}/d\epsilon_{\gamma,e} \propto \epsilon_{\gamma,e}^{-q}$, where $q \approx 5/3$ for $\epsilon_{\gamma} < \epsilon_{\gamma,e}$ and $q \approx 2.2$ for $\epsilon_{\gamma,e} < \epsilon_{\gamma} < \epsilon_{\gamma,\text{max}}$ (Belyanin & Derishev 2001). The maximum photon energy is $\epsilon_{\gamma,\text{max}} \approx \gamma_{\text{ph}}^2 \epsilon_{\gamma,\text{Th}} \Gamma_{p,f} \approx 20$ GeV. The fluence threshold for GLAST is $\sim 4 \times 10^{-8}$ ergs cm$^{-2}$ for a short integration time (Gehrels & Michelson 1999). Thus, the Comptonized-photon fluence $\epsilon_{\gamma,e} F_{\gamma,e} \approx \epsilon_{\gamma,e} N_{\gamma,e} \xi_{\gamma,e} \approx 4.4 \times 10^{-7}$ ergs cm$^{-2}$ for a typical SGRB of $t \sim 1$ s duration at $z \approx 0.1$ should be detectable by GLAST (see Fig. 2).

Apart from the $\pi^\pm$-decay photons of energy 60 GeV and $e^\pm$ pair cascade photons of energy 320 MeV, the spectrum of radiation from the photosphere of an SGRB would also contain thermal photons of peak energy $\epsilon_{\gamma,b} \approx 2.8 \epsilon_{\gamma,\text{Th}} \Gamma_{p,f} \approx 6.4 \Gamma_{p,2.8} \text{ MeV}$.

**Fig. 2.**—Photon spectra of the expected photospheric signals (thermal and IC cascade by pion-decay electrons and photons) and the synchrotron spectrum from internal shocks, compared to the GLAST LAT sensitivity. The dashed, dotted, and dotted-dash curves correspond to the initial neutron-to-proton ratio $\xi_0 = n_n/n_p$ = 1, 5, and 10, respectively, as in Fig. 1. The other parameters used are $L = 10^{55} L_{18}$ ergs s$^{-1}$, $R_i = 10^4$ cm, and $n_i = 316$ for all curves. To calculate the internal shock synchrotron spectra, we used a fixed variability time $t_i = 10^{-3}$ s. The peaks of the spectra are correlated with the final bulk Lorentz factor of the proton outflow: $\Gamma_{p,f} = 519, 572,$ and 624, respectively, for $\xi_0 = 1, 5,$ and 10 (see main text for more details).
photospheric photons of peak energy is 
\[ \hat{N}_{\gamma,b} \approx \frac{L}{4\pi D^2 I_{\gamma,b}} \approx 0.5 \frac{L_{49.8}}{D_{27}^2 I_{p,2.8}} \text{ cm}^{-2} \text{ s}^{-1}. \]  
(20)

The photon fluxes at various energies discussed above are associated with the jet photosphere and last as long as the usual prompt \( \gamma \)-ray burst emission. Hence, these emissions would be contemporaneous with the emission at \( \sim \text{MeV} \) associated with internal shocks or other dissipation processes above the photosphere, aside from a small time lag between the onset of the two events of the order of a millisecond. We discuss emission from internal shocks next.

5.1. Comparison with Internal Shock Emission

The internal-shock model of plasma shells colliding at some radius above the photosphere is the most commonly discussed mechanism to produce the observed \( \gamma \)-rays in the fireball-shock model. The average bulk Lorentz factor of the colliding shells is \( \Gamma_I \approx \Gamma_{p,f} \approx 624 \Gamma_{12.8} \). With a millisecond variability time \( t_v = 10^{-3} \text{ s} \), the internal shocks take place at a radius \( r_I \approx 2L_I^2 t_v \approx 2.3 \times 10^{17} \text{ cm} \). Note that this is much smaller than the \( \beta \)-decay radius \( \sim 3.4 \times 10^{18} \text{ cm} \). Hence, most neutrons have not had time to decay until well beyond the internal shock radius.

The equipartition magnetic field in the shock region is \( B = (2\varepsilon_B L/r_I) \frac{1}{2} \approx 9.1 \times 10^{15} \text{ G} \) with an equipartition fraction \( \varepsilon_B = 0.1 \varepsilon_{B,1} \). The shock-accelerated electrons are assumed to follow a power-law energy distribution \( dN_e/d\gamma_e \sim \gamma_e^{-\delta \gamma_e} \) for \( \gamma_e \geq \gamma_{\text{min}} \approx \varepsilon_e m_p/m_e \) with \( \alpha \approx 2.4 \), which results in synchrotron radiation with a peak energy
\[ \epsilon_{\gamma,m} = \frac{\hbar c \Gamma_I}{2m_e c^2} \left( 3 \frac{\gamma_e^{\delta \gamma_e}}{\epsilon_{\gamma,\text{min}}^{\delta \gamma_e}} \right)^{1/2} \approx 330 \left( \frac{\epsilon_{\gamma,\text{min}}^{\delta \gamma_e}}{\Gamma_{12.8}^{\delta \gamma_e} t_v^{\delta \gamma_e}} \right)^{1/2} \text{ keV}. \]  
(21)

Here \( \varepsilon_e = 0.1 \varepsilon_{e,-1} \) is the fraction of the fireball’s kinetic energy transferred to the electrons by shocks. The corresponding number flux at Earth of these peak synchrotron photons is
\[ \hat{N}_{\gamma,m} \approx \frac{\hat{L}_{\gamma,\text{e}}}{4\pi D^2 I_{\gamma,m}} \approx 1.0 \left( \frac{L_{49.8} k^2 I_{\gamma,m}^{1/2}}{D_{27}^2 I_{\gamma,m}^{1/2}} \right)^{1/2} \text{ cm}^{-2} \text{ s}^{-1}. \]  
(22)

The typical synchrotron radiation spectrum from power-law–distributed electron energy is \( dN_e/d\epsilon_e \sim \epsilon_e^{-(\delta \gamma_e + 2)/2} \sim \epsilon_e^{-2.2} \) for \( \epsilon_e > \epsilon_{\text{min}} \) (see Fig. 2). The observed typical GRB spectra, \( dN_e/d\epsilon_e \sim \epsilon_e^{-1} \), is slightly different from the theoretical synchrotron spectrum with an index \(-3/2\).

We have plotted the energy spectra of synchrotron radiation from the internal shocks, and thermal and IC cascades from the jet photosphere in Figure 2. We used \( L = 10^{50} L_{50} \) ergs s\(^{-1}\), \( R_p = 10^6 \) cm, and \( \eta = 316 \) for all curves and \( \xi_0 = 1, 5, 10 \) for the dashed, dotted, and dashed-dotted curves, respectively. The peak energy of thermal photons \( (\epsilon_{\gamma,b}) \) decreases from \( 6.4 \text{ keV} \) for \( \xi_0 = 10 \) to \( 3 \text{ keV} \) for \( \xi_0 = 1 \) as the photosphere radius \( (R_p) \) in equation (8) increases with decreasing \( \xi_0 \). The variation of the peak energy \( (\epsilon_{\gamma,b}) \) of the photospheric cascade is directly proportional to the variation of \( \Gamma_{p,f} \) with \( \xi_0 \) and increases with \( \xi_0 \). We have normalized these spectra by integrating the corresponding differential number fluxes from \( \epsilon_{\gamma,e} \) up to the respective maximum energy \( (\epsilon_{\gamma,max}) \) and equating the fraction of the total energy carried by the pions to the respective total number flux in equation (19) with fixed \( \xi_0 = 0.1 \).

The variation of the synchrotron peak energy \( (\epsilon_{\gamma,m}) \), in the range \( 300-400 \text{ keV} \) with \( \xi_0 \), can be explained similarly. We have normalized the synchrotron spectra by integrating the corresponding differential number fluxes from \( \epsilon_{\gamma,m} \) up to the energy \( 10^{-4} \text{ MeV} \) for \( \xi_0 \). The typical energy range for the BATSE and BAT detectors, and equating to the respective total number flux in equation (22). We have plotted an extension of the synchrotron spectra up to \( 1 \text{ GeV} \), in order to compare them with other spectral components.

The synchrotron spectrum does not extend very far above \( \epsilon_{\gamma,m} \), for the parameters used here, from a calculation of the maximum electron Lorentz factor. The energy at which the Rayleigh-Jeans portion of the thermal photospheric starts to dominate over the synchrotron spectrum is \( \epsilon_{\gamma} \approx 2 \text{ MeV} \) for \( \xi_0 = 5 \) and \( 10 \) curves in Figure 1. For \( \xi_0 = 1 \), the thermal component is small over the whole energy range. Above this, besides the Comptonized photospheric cascade spectrum, there could also be an IC scattered component of the synchrotron photons (SSC), whose importance is characterized by the parameter \( Y = (1+1 (1+4\epsilon_e/e_B)^{1/2})/2 \approx 0.6 \) for our present model. The corresponding minimum IC peak is \( \gamma_e^{\text{min}} \sim 10^4-10^5 \) higher than the synchrotron peak, at a photon energy \( \epsilon_{\gamma,m,\text{IC}} \sim 10 \text{ GeV} \). Hence, an SSC component is not important in this case \( (Y < 1) \). Thus, at energies above which the synchrotron spectrum cuts off, the Comptonized photospheric cascade photon spectrum would be clearly distinguishable by GLAST (see Fig. 1). Although for some choice of parameters an SSC contribution could become significant at energies \( > 10 \text{ GeV} \), this is expected to be at energies much higher than the peak of the Comptonized cascade peak \( \epsilon_{\gamma,c} \approx 300 \text{ MeV} \). The maximum SSC photon energy, in the Thomson limit \( \gamma_e^{\text{min}} \ll m_e c^2/\epsilon_{\gamma,m} \), would be \( \epsilon_{\gamma,\text{IC}} \ll \epsilon_{\gamma,c} \approx m_e c^2 / \epsilon_{\gamma,c} \approx 300 \text{ GeV} \). At energies well above \( \sim 10 \text{ GeV} \), the two components (Comptonized photospheric cascade and SSC) may be hard to distinguish, but in the range of \( \sim 5 \text{ MeV} \rightarrow 10 \text{ GeV} \), the Comptonized photospheric cascade should dominate. Below the peak \( \epsilon_{\gamma} \sim \epsilon_{\gamma,c} \), the Comptonized photospheric cascade component is lower than the thermal photospheric peak, and at even lower energies it is also lower than the synchrotron component.

6. DISCUSSION

If SGRBs are produced as a result of double neutron star or neutron star–black hole binary mergers, the neutron-rich outflow may result naturally in a high-proton Lorentz factor \( \Gamma_{p,f} \approx 600 \). Besides the usual nonthermal (e.g., internal shock) synchrotron spectrum, this implies the presence of a thermal photospheric component peaking around \( \sim 10 \text{ MeV} \). The \( n-p \) decoupling will lead to higher energy photons (\( \gtrsim 0.1 \text{ GeV} \)) from neutral and charged pion decay cascades and neutrinos (\( \gtrsim 30 \text{ GeV} \)) directly from pion and muon decays. There may be additional X-ray emission when the neutron-decay proton shell with initial lower Lorentz factor impacts with the decelerating proton shell initially moving with a high Lorentz factor (Dermer & Atiyon 2006). The all-particle light curves can also be affected by the neutron decay (Derishev et al. 1999b; Rossi et al. 2006). Here we have concentrated on the prompt (\( t \sim 2 \text{ s} \)) high-energy (MeV to tens of GeV) signatures of neutron-rich outflows. For short bursts at \( z \sim 0.1 \), the \( \sim 50 \text{ GeV} \) neutrinos are unlikely to be detectable with currently planned gigaton Cerenkov detectors. However, the pion decay photons at \( \sim 60 \text{ GeV} \) escaping from within a \( r_{Th} \sim 1 \) skin depth of the photosphere should be detectable with
high-area, high-duty-cycle Cerenkov detectors such as Milagro, at an estimated rate of 5 bursts per year. The pion decay photons created below the photosphere will result in a Comptonized cascade spectrum, which, peaking in the range $\gtrsim 0.3$ GeV, dominates the nonthermal shock synchrotron-IC spectrum, and should be detectable with GLAST. Both the prompt GeV photon components discussed here are essentially contemporaneous with the usual MeV range short-GRB prompt emission, $t_\gamma \lesssim 2$ s. Their detection or lack thereof could constrain the neutron content of the outflow and the nature of the short-burst progenitors.

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