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EARLY PHOTON-SHOCK INTERACTION IN A STELLAR WIND: A SUB-GeV PHOTON FLASH AND HIGH-ENERGY NEUTRINO EMISSION FROM LONG GAMMA-RAY BURSTS

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ABSTRACT

For gamma-ray bursts (GRBs) born in a stellar wind, as the reverse shock crosses the ejecta, usually the shocked regions are still precipitated by the prompt MeV γ-ray emission. Because of the tight overlapping of the MeV photon flow with the shocked regions, the optical depth for the GeV photons produced in the shocks is very large. These high-energy photons are absorbed by the MeV photon flow and generate relativistic $e^{\pm}$ pairs. These pairs rescatter the soft X-ray photons from the forward shock as well as the prompt γ-ray photons and power detectable high-energy emission, a significant part of which is in the sub-GeV energy range. Since the total energy contained in the forward shock region and the reverse shock region are comparable, the predicted sub-GeV emission is independent of whether the GRB ejecta are magnetized (in which case the reverse shock inverse Compton and synchrotron self-Compton emission is suppressed). As a result, a sub-GeV flash is a generic signature for the GRB wind model, and it should be usually detectable by the future *Gamma-Ray Large Area Space Telescope* (*GLAST*). Overlapping also influences neutrino emission. Besides the 10^{15}–10^{17} eV neutrino emission powered by the interaction of the shock-accelerated protons with the synchrotron photons in both the forward and reverse shock regions, there comes another 10^{14} eV neutrino emission component powered by protons interacting with the MeV photon flow. This last component has a similar spectrum to that generated in the internal shock phase, but the typical energy is slightly lower.

Subject headings: acceleration of particles — elementary particles — gamma rays: bursts

1. INTRODUCTION

The leading model for long gamma-ray bursts (GRBs) invokes a relativistic jet emerging from a collapsar (e.g., Woosley 1993; Paczyński 1998). It has been widely expected that GRBs associated with supernovae should occur in a preburst stellar wind environment with particle number density $n \propto R^{-2}$ (e.g., Dai & Lu 1998; Mészáros et al. 1998; Chevalier & Li 2000). As the jet expands into a dense stellar wind, a forward shock (FS) and a reverse shock (RS) form (Chevalier & Li 2000; Dai & Lu 2001, hereafter DL01), whose synchrotron emission peaks in the soft X-ray band and in the ultraviolet-to-optical band, respectively. Photomeson interaction in the RS region also results in neutrino emission in the energy range 3×10^{15}–3×10^{17} eV (DL01).

However, until very recently, in nearly all the works on RS emission, possible overlapping of the initial prompt γ-ray emission (MeV photon flow) with the RS region has been ignored. Beloborodov (2005) pointed out that when overlapping is important, the RS emission would be suppressed significantly, since the electron cooling is dominated by inverse Compton (IC) scattering off the MeV photon flow. Correspondingly, the upscattering would power a strong GeV–TeV photon flash. Fan et al. (2005a) suggested that when overlapping is substantial, similar processes are also relevant to the early FS emission.

In this work, we show that overlapping is a common feature for GRBs born in a stellar wind (§ 2), so that the early afterglow wind model needs substantial revision. We then systematically study the consequence of this overlapping effect in the wind model. In § 3 we show that such an early photon-shock interaction inevitably leads to the prediction of a sub-GeV photon flash in the wind model, which is generally detectable by the *Gamma-Ray Large Area Space Telescope* (*GLAST*). In § 4, we discuss various high-energy neutrino emission processes in the early afterglow stage as the result of photon-shock interaction. We summarize our conclusions in § 5 with some discussions.

2. TIGHT OVERLAPPING OF THE MeV PHOTON FLOW WITH THE SHOCKED REGIONS

The overlapping of the MeV photon flow with the fireball ejecta is important if the RS-crossing radius satisfies

$$R_\times \ll 2\Gamma_0^2\Delta, \qquad (1)$$

where Γ_0 is the initial Lorentz factor of the ejecta (at the internal shock phase), $\Delta = cT_{90}/(1+z)$, $T_{90} \sim 20$ s is the observed duration of the long GRB, and $z \sim 1$ is the redshift. The shock-crossing radius $R_\times$ can be written in a general form

$$R_\times = \max\left(R_\gamma,\ \Gamma^2\Delta\right), \qquad (2)$$

where R_γ is the radius where the mass of the medium collected by the fireball is equal to $1/\Gamma_0$ of the fireball mass (which corresponds to $\Gamma \sim \Gamma_0/2$) and Γ is the Lorentz factor of the shocked ejecta. In equation (2), the first term dominates if the RS is nonrelativistic, while the second term dominates if the RS is relativistic. For the interstellar medium (ISM) case with a typical number density $n_{\rm ISM} \sim 1$ cm^{-3}, the RS is nonrelativistic, and $R_\times\,({\rm ISM}) = R_\gamma \approx 8.8\times10^{16}E_{53.6}^{1/3}n_{{\rm ISM},0}^{-1/3}\Gamma_{0,2.5}^{-2/3}$, which is usually larger than (or at least comparable to) $2\Gamma_0^2\Delta = (5.4\times10^{16}\ {\rm cm})\Gamma_{0,2.5}^2\Delta_{11.5}$. Throughout this paper, we adopt the convention $Q_x = Q/10^x$ to express physical parameters in cgs

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units. So, in the ISM case, the initial γ-rays cannot change the RS emission significantly. In the stellar wind case, on the other hand, one has $n = 3.0 \times 10^{35} A_* R^{-2}$ cm^{-3}, where $A_* = [\dot{M}/(10^{-5}\ M_\odot\ \mathrm{yr}^{-1})][v_w/(10^8\ \mathrm{cm\ s^{-1}})]^{-1}$, $\dot{M}$ is the mass-loss rate of the progenitor, and v_w is the wind velocity (Chevalier & Li 2000). With such a dense medium, the RS is relativistic (Chevalier & Li 2000; DL01); $\Gamma \approx \bar{\zeta}^{1/4}\Gamma_0^{1/2}/\sqrt{2}$, where $\bar{\zeta} \approx 2600 E_{53.6}/(A_* \Delta_{11.5} \Gamma_{0,2.5}^2)$ is the ratio of the proper mass density of the unshocked ejecta and the medium density and $E \sim 4 \times 10^{53}$ ergs is the isotropic energy of the ejecta, which is calculated by adopting $T_{90} = 20$ s, $z = 1$, and by assuming that the observed γ-ray luminosity is $L_\gamma \simeq 10^{52}$ ergs s^{-1} and that the radiation efficiency of the internal shocks is $\eta \simeq 0.2$. The total wind luminosity is therefore $L_{\mathrm{tot}} = L_\gamma/\eta = 5 \times 10^{52}$ ergs s^{-1}, and the afterglow kinetic energy luminosity is $L \simeq L_{\mathrm{tot}}(1-\eta) = 4 \times 10^{52}$ ergs s^{-1} in the beginning of the afterglow phase. The resulting bulk Lorentz factor in the shocked region is therefore

$$\Gamma \approx 87 L_{52.6}^{1/4} A_*^{-1/4}. \tag{3}$$

The RS crosses the ejecta at the radius

$$R_\times\,(\mathrm{wind}) \simeq (2.7 \times 10^{15}\ \mathrm{cm}) L_{52.6}^{1/2} A_*^{-1/2} \Delta_{11.5}. \tag{4}$$

At this radius, the rear of the MeV photon flow only leads the rear of the ejecta by a distance of $R_\times\,(\mathrm{wind})/2\Gamma_0^2 \ll \Delta$, so that *during most of the time when the RS crosses the ejecta, the shocked regions are still precipitated by the MeV photons, so that the electrons and e$^\pm$ pairs in the shocks lose energy mainly by scattering with the MeV photons*. As a result, the conventional synchrotron ultraviolet-optical emission of the electrons in the RS region would be greatly suppressed. Let us define a parameter $Y \equiv U_\gamma/U_B$, where $U_\gamma \simeq L_\gamma/4\pi R^2\Gamma^2 c$ is the MeV photon energy density in the rest frame of the shocked region and $U_B = B^2/8\pi \simeq (\Gamma_0/\Gamma)^2 \epsilon_B n_4 m_p c^2$ is the comoving magnetic energy density in the same region, where $n_4 = (1-\eta)L/4\pi R^2 \Gamma_0^2 m_p c^3$ is the comoving density of the unshocked ejecta. As usual, we introduce magnetic field and electron equipartition parameters $\epsilon_B \sim 0.001$–0.1 and $\epsilon_e \sim 0.1$ into the shocked region. The wide range of ϵ_B is adopted with the consideration of broadband afterglow fits (e.g., Panaitescu & Kumar 2001) and the possibility that the ejecta may be magnetized (e.g., Fan et al. 2002; Zhang et al. 2003). In this work, we normalize our expression by taking $\epsilon_B \sim 0.01$. After some simple algebra, we have

$$Y \approx \eta/[(1-\eta)\epsilon_B]. \tag{5}$$

For $\epsilon_B \sim 0.001$–0.1 and $\eta \approx 0.2$, we have $Y \sim 2.5$–$250 \gg 1$. The cooling frequency ν_c (see eq. [15] of DL01) would then be lowered by a factor of $1/(1+Y)^2$. For $\epsilon_{B,-2} \sim 0.1$–1, the flux of the soft ultraviolet-optical photons is significantly lower than that calculated in DL01, so that the energy loss fraction for each proton through pion production (i.e., f_π, defined as the ratio between the comoving expansion time t'_d and the comoving proton cooling time t'_π due to photomeson cooling) should also be lowered by roughly a factor of $1/(1+Y)$ (see eqs. [21] and [22] in DL01). Consequently, the resulting detectable neutrino flux should be somewhat lower.

It is worth pointing out that for some other processes accompanying prompt γ-ray emission, such as the prediction of the ultraviolet flash powered by neutron-rich internal shocks (Fan & Wei 2004), the MeV photon flow also plays an important role in the cooling of the accelerated electrons. By taking into account the overlapping effect, Fan et al. (2005b) have shown that with reasonable parameters, the near-IR, optical, and UV emission flux from the neutron-rich internal shocks can be as strong as 10 mJy. Comparing with the standard internal shock models, it also suffers weaker synchrotron self-absorption and seems to be able to account for prompt long-wavelength observation of GRB 041219a (e.g., Blake et al. 2005; Vestrand et al. 2005).

3. SUB-GeV PHOTON FLASH

Beloborodov (2005) suggested that electrons accelerated by the RS would scatter the MeV photon flow and power a detectable GeV–TeV flash. It is suggested that this mechanism may potentially account for the distinct high-energy spectral component observed in GRB 941017 (González et al. 2003).

One uncertainty is that if in the RS region $\epsilon_B^r \geq 0.1$, as might have been the case for GRB 990123 (Fan et al. 2002; Zhang et al. 2003), the *Y*-parameter (eq. [5]) is not large, and the IC effect may not be strong. (Hereafter the superscripts *r* and *f* represent the RS and FS, respectively. If the parameters in both regions are the same, no superscript is marked.) Below we show that by taking into account the photon-shock interaction in the FS region, the IC process is generic regardless of the magnetization of the ejecta. In addition, the typical energy of the emerged spectrum peaks in the sub-GeV range (see below).

At $R_\times$(wind), the energy contained in the RS region can be written as

$$E_{\mathrm{tot}}^r \approx \gamma_{34}\Gamma M_{\mathrm{ej}} c^2 \approx \Gamma_0 M_{\mathrm{ej}} c^2/2 \approx E_{\mathrm{tot}}^f, \tag{6}$$

where we have taken the Lorentz factor of the shocked ejecta relative to the unshocked ejecta as $\gamma_{34} \approx (\Gamma_0/\Gamma + \Gamma/\Gamma_0)/2 \simeq \Gamma_0/2\Gamma \gg 1$ (which is valid for the wind RS model). Here $M_{\mathrm{ej}} \simeq E/(\Gamma_0 c^2)$ is the rest mass of the ejecta. Equation (6) suggests that the energy contained in the RS region is roughly equal to that in the FS region. As a result, if a significant part of the FS energy is emitted in the sub-GeV range (we show below that this is the case), whether the sub-GeV emission from the RS region is suppressed or not is observationally unimportant. In view of this fact, below we adopt the same physical parameters ($\epsilon_e \sim 0.1$ and $\epsilon_B \sim 0.01$) for both the RS and the FS regions.

The total thermal energies in the RS and FS regions are $E_{\mathrm{th}}^r \simeq (\gamma_{34}-1)E_{\mathrm{tot}}^r/\gamma_{34}$ and $E_{\mathrm{th}}^f \simeq (\Gamma-1)E_{\mathrm{tot}}^f/\Gamma$, respectively. The ratio of them is $E_{\mathrm{th}}^r/E_{\mathrm{th}}^f \approx (\gamma_{34}-1)/\gamma_{34}$. For typical parameters, $\gamma_{34} \sim 2$, we have $E_{\mathrm{th}}^r/E_{\mathrm{th}}^f \approx 1/2$. This result suggests that for both the high-energy photon emission and the neutrino emission (as discussed in § 4), the component from the FS region is the dominant one.

As usual, we assume that the shocked electrons distribute as $dN_e/d\gamma_e \propto \gamma_e^{-p}$ ($p \sim 2.2$) for $\gamma_m < \gamma_e < \gamma_M$, where $\gamma_M \sim 10^8/B^{1/2}$ is the maximum Lorentz factor of the electrons accelerated by shocks and B is the magnetic field strength of the shock regions, which can be calculated by

$$B \approx (1.86 \times 10^3\ \mathrm{G}) \epsilon_{B,-2}^{1/2} A_*^{-1/4} L_{52.6}^{1/4} R_{15}^{-1}. \tag{7}$$

For the FS and RS, γ_m can be estimated by

$$\begin{aligned}
\gamma_m^f &= (\Gamma-1)\epsilon_e (m_p/m_e)[(p-2)/(p-1)] \\
&\approx 2600\epsilon_{e,-1} A_*^{-1/4} L_{52.6}^{1/4}, \\
\gamma_m^r &= (\gamma_{34}-1)\epsilon_e (m_p/m_e)[(p-2)/(p-1)] \\
&\approx 23\epsilon_{e,-1} A_*^{1/4} L_{52.6}^{-1/4} \Gamma_{0,2.5}.
\end{aligned} \tag{8}$$

Performing IC correction on equation (14) of DL01, we also get the cooling Lorentz factor $\gamma_c \leq 1$, which means that the shocked electrons all cool so rapidly that essentially all of them have $\gamma_e \sim 1$. In other words, all electrons are in the fast cooling phase.

The electrons in the shocked regions scatter with the MeV photons and produce high-energy photons. Given the observed prompt GRB peak energy $\epsilon^b_{\gamma,\rm obs} = 300$ keV (throughout the paper, the subscript "obs" denotes the parameters measured in the observer frame), the electron energy above which the IC process is suppressed by the Klein-Nishina limit is given by (see also Beloborodov 2005)

$$\gamma_e^* = \frac{\Gamma m_e c^2}{(1+z)\epsilon^b_{\gamma,\rm obs}} \sim \Gamma. \tag{9}$$

In the RS region, one has $\gamma_e^* > \gamma_m^r$. The peak energy of the upscattered photons can be then estimated as

$$\epsilon^{r,\rm IC}_{m,\rm obs} \approx \gamma_m^{r2}\epsilon^b_{\gamma,\rm obs} \simeq (0.16\ {\rm GeV})\epsilon^2_{e,-1}A_*^{1/2}\Gamma^2_{0,2.5}L_{52.6}^{-1/2}\left(\frac{\epsilon^b_{\gamma,\rm obs}}{300\ {\rm keV}}\right). \tag{10}$$

In the FS region, on the other hand, one has $\gamma_m^f \gg \gamma_e^*$, so that the IC cooling by the MeV photons is suppressed significantly. The fresh energetic electrons lose energy mainly by synchrotron self-Compton (SSC) effects; the corresponding Compton parameter can be estimated by $Y_{\rm SSC} \simeq (\epsilon_e/\epsilon_B)^{1/2} \gg 1$. The typical energy of SSC photons reads

$$\epsilon^{f,\rm SSC}_{m,\rm obs} \approx h(\gamma_m^f)^4 e\Gamma B/[2\pi m_e c(1+z)] \approx (86\ {\rm GeV})(1+z)^{-1}\epsilon^4_{e,-1}\epsilon^{1/2}_{B,-2}A_*^{-3/2}L_{52.6}^{3/2}R_{15}^{-1}, \tag{11}$$

where h is Planck's constant.

Apparently, equations (10) and (11) are quite different. However, as shown below, photons more energetic than ∼1 GeV will be absorbed by the MeV photons near the shock-crossing radius, and most energy will be reprocessed and end up in the sub-GeV range. This is valid for both the RS and FS. The pair production optical depth for photons with energy ∼1 GeV [which will be absorbed by the soft photons with energy $\epsilon_{a,\rm obs} \simeq 2(\Gamma m_e c^2)^2/(1+z)^2$ GeV ∼ 1 MeV] is roughly (e.g., Svensson 1987)

$$\tau_{\gamma\gamma}(1\ {\rm GeV}) \approx \frac{11\sigma_{\rm T}N_{>\epsilon_{a,\rm obs}}}{720\pi R^2} \approx 20E_{\gamma,53}\left(\frac{1+z}{2}\right)L_{52.6}^{-1/2}A_*^{1/2}R_{15}^{-2}, \tag{12}$$

where $N_{>\epsilon_{a,\rm obs}}$ is the total number of photons satisfying $\epsilon_{\rm obs} > \epsilon_{a,\rm obs}$, which can be estimated by $N_{>\epsilon_{a,\rm obs}} \approx 0.2E_\gamma/[(1+z)\epsilon_{a,\rm obs}]$ for $\epsilon_{a,\rm obs} > \epsilon^b_{\gamma,\rm obs}$, where the coefficient 0.2 is a rough spectrum correction factor on the photon number calculation derived from the integration of the GRB spectrum. In this paper, we take a typical broken power-law GRB spectrum for MeV photons, i.e., $n(\epsilon_\gamma) \propto \epsilon_\gamma^{-1}$ for $\epsilon^b_\gamma/50 < \epsilon_\gamma < \epsilon^b_\gamma$ and $n(\epsilon_\gamma) \propto \epsilon_\gamma^{-2}$ for $\epsilon^b_\gamma < \epsilon_\gamma < 50\epsilon^b_\gamma$.

With typical parameters, we have $\tau_{\gamma\gamma}(1\ {\rm GeV}) \simeq 2$ at $R_\times$(wind). This means that photons more energetic than 1 GeV are absorbed by the MeV photons and produce $e^\pm$ pairs. The random Lorentz factor of these secondary pairs can be estimated as $\gamma_\pm \simeq (1+z)\epsilon_{\gamma,\rm obs}/2\Gamma m_e c^2 = 11(1+z)[\epsilon_{\gamma,\rm obs}/(1\ {\rm GeV})]$. For $\epsilon_{\gamma,\rm obs} = \epsilon^{f,\rm SSC}_{\gamma,\rm obs} \sim 32(1+z)^{-1}$ GeV, we have $\gamma_\pm^{f,\rm SSC} \simeq 350$. These pairs will scatter both the MeV photons and the X-ray photons produced in the forward shock region. Since $\gamma_\pm^{f,\rm SSC}$ is significantly larger than γ_e^*, the Klein-Nishina correction is still important for the pair MeV photon IC scattering.[4] The cross section can be approximated as $\sigma \approx A(x)\sigma_{\rm T}$, where $x \simeq \gamma_\pm^{f,\rm SSC}/\gamma_e^* \sim \gamma_\pm^{f,\rm SSC}/\Gamma \approx 4$ and $A(x) \approx (3/8x)(\ln 2x + 1/2) \approx 0.2$. The energy density of the soft X-rays (FS synchrotron emission) is about $g = [(1-\eta)/\eta][\epsilon_e/(1+Y_{\rm SSC})] \sim 0.1$ times of that of the MeV photons, so the fraction of energy of the secondary pairs that goes to scattering with the soft X-ray FS emission (with a typical frequency defined by eq. [15]) can be estimated as $[4g/A(4)]/[1+4g/A(4)] \sim 2/3$. The typical energy of the upscattered soft X-rays can be estimated as [at $R_\times$(wind)]

$$\epsilon^{f,\rm IC}_{m,\rm obs} \simeq \left(\gamma_\pm^{f,\rm SSC}\right)^2\epsilon^f_{m,\rm obs} \simeq (0.29\ {\rm GeV})\left(\frac{\gamma_\pm^{f,\rm SSC}}{350}\right)^2 \simeq (0.29\ {\rm GeV})\left(\frac{\gamma_\pm^{f,\rm SSC}}{350}\right)^2\left(\frac{1+z}{2}\right)^{-1}\epsilon^2_{e,-1}\epsilon^{1/2}_{B,-2}A_*^{-1}L_{52.6}. \tag{13}$$

The pairs also scatter the initial MeV photons. However, the upscattered photons are still too energetic to escape. A pair photon cascade likely develops until the photon energy is low enough to be transparent. For example, in the extreme Klein-Nishina limit, the typical energy of the upscattered MeV photons can be estimated by $(1+z)\epsilon^{\rm IC}_{\rm obs} \sim \Gamma\gamma_\pm^{f,\rm SSC}m_e c^2 \sim (15.6\ {\rm GeV})L_{52.6}^{1/4}A_*^{-1/4}(\gamma_\pm^{f,\rm SSC}/350)$. After interacting with the MeV photons again, tertiary pairs will be produced with a typical Lorentz factor of ∼180. They then rescatter various photon fields again. The detailed investigation of this cascade process has to be modeled numerically and is beyond the scope of this paper. Nonetheless, a generic picture is that at least $\frac{2}{3}$ of the total energy of the reprocessed photons is radiated in the sub-GeV range, making these GRBs interesting targets for *GLAST* observations.

For typical parameters taken in this work ($z = 1$ and $Q_x = 1$), the fluence of the sub-GeV flash from the very early FS can be estimated as

$$S \sim \frac{2}{3}\epsilon_e\Gamma(\Gamma-1)m_p c^2(1+z)\int_0^{R_\times({\rm wind})} R^2 n\, dR/D_L^2 \approx 2\times10^{-6}\ {\rm ergs\ cm^{-2}}, \tag{14}$$

where D_L is the luminosity distance, which is $\approx 2.2\times10^{28}$ cm for $z = 1$. The fluence for the RS component may be comparable to or less than this value. For the Energetic Gamma-Ray Experiment Telescope (EGRET), the estimated fluence threshold in the low integration time ($t_{\rm int}$) regime ($t_{\rm int} < 1.7\times10^3$ s) is $\sim 2.1\times10^{-6}$ ergs cm^{-2} for a typical photon energy ∼400 MeV (e.g., Zhang & Mészáros 2001). Our predicted fluence $S \sim 2\times10^{-6}$ ergs cm^{-2} at a sub-GeV energy range is below the EGRET detection threshold, which explains the nondetection of such flashes during the EGRET era. The fluence threshold of *GLAST*

[4] In the Klein-Nishina regime, for one electron, the energy loss in one scattering is $\sim\gamma_e m_e c^2$, and the energy loss rate is $P_{\rm IC} \sim \gamma_e m_e c^2 A(x)\sigma_{\rm T}n_\gamma c$, where $n_\gamma \sim U_\gamma(\epsilon^b_\gamma/\Gamma)$ is the number density of γ-ray photons. The synchrotron radiation power satisfies $P_{\rm syn} = (4/3)\sigma_{\rm T}\gamma_e^2\beta_e^2 U_B c$. The Compoton parameter can be estimated by $Y_{\rm KN} = P_{\rm IC}/P_{\rm syn} \sim A(x)U_\gamma/(xU_B)$, where we have taken $\gamma_e = \gamma_\pm^{f,\rm SSC}$. Since U_γ/U_B is just the Compton parameter Y in the Thomson regime, we have $Y_{\rm KN} \sim A(x)Y/x$ for $x \gg 1$, i.e., in the extreme Klein-Nishina regime the IC cooling is inefficient.

is roughly $\sim 4\times10^{-7}$ ergs cm^{-2} for $t_{\rm int} < 10^5$ s (e.g., Zhang & Mészáros 2001). We therefore expect that many sub-GeV photon flashes could be detected during the *GLAST* era.

4. NEUTRINO EMISSION

4.1. *Neutrinos from Photomeson Interaction with the Synchrotron Photons in RS and FS*

The protons accelerated from both the FS and the RS would produce high-energy neutrinos via photomeson interaction, mainly through Δ-resonance. In the RS region, DL01 have found that the typical neutrino energy is 3×10^{15}–3×10^{17} eV.

We now estimate the neutrino emission from the FS region. As mentioned in § 3 (see the paragraph below eq. [10]), the fresh electrons accelerated by the FS lose energy mainly by SSC cooling rather than by IC cooling. As a result, the flux above ν_m^f should be lowered by a factor of $1/(1+Y_{\rm SSC})$. The typical synchrotron radiation frequency (ν_m^f) of the FS reads

$$\nu_m^f \approx \left(3.1\times10^{18}\ {\rm Hz}\right)\epsilon_{e,-1}^2\epsilon_{B,-2}^{1/2}L_{52.6}A_*^{-1}R_{15}^{-1}. \quad (15)$$

The protons that interact with the $h\nu_m^f$ photons at the Δ-resonance have an observed energy of

$$\begin{aligned}\epsilon_{p,{\rm obs}}^b(1) &\approx (0.3\ {\rm GeV})^2\Gamma^2/\left[(1+z)h\nu_m^f\right]\\ &\approx \left(1.8\times10^{17}\ {\rm eV}\right)(1+z)^{-1}\epsilon_{e,-1}^{-2}\epsilon_{B,-2}^{-1/2}L_{52.6}^{-1/2}A_*^{1/2}R_{15}.\end{aligned} \quad (16)$$

The photomeson interactions include (1) production of π mesons, $p\gamma\rightarrow p+\pi^0$ and $p\gamma\rightarrow n+\pi^+$, and (2) decay of π mesons, $\pi^0\rightarrow 2\gamma$ and $\pi^+\rightarrow\mu^++\nu_\mu\rightarrow e^++\nu_e+\bar{\nu}_\mu+\nu_\mu$. When protons interact with the FS photons with a typical energy $h\nu_m^f$, these processes produce neutrinos with energy $\epsilon_{\nu,{\rm obs}}^b(1)\sim 0.05\epsilon_{p,{\rm obs}}^b(1)$ (e.g., Waxman & Bahcall 1997), i.e.,

$$\epsilon_{\nu,{\rm obs}}^b(1) \simeq \left(9\times10^{15}\ {\rm eV}\right)(1+z)^{-1}\epsilon_{B,-2}^{-1/2}\epsilon_{e,-1}^{-2}L_{52.6}^{-1/2}A_*^{1/2}R_{15}. \quad (17)$$

In principle, γ-rays of similar energies are produced by π^0-decay. But as shown in § 3, the MeV photon flow is optically thick for these high-energy photons, and a $e^\pm$ pair cascade develops. The bulk γ-rays emerge at much lower energies ($\sim$GeV), which adds to the sub-GeV emission discussed in § 3.

Similar to Halzen & Hooper (2002) and Guetta et al. (2004), at $\epsilon_{\nu,{\rm obs}}^b(1)$, the fraction of the energy loss of protons through photomeson interaction with the $h\nu_m^f$ photons, $f_\pi(1)$, can be parameterized as (in the calculation, we have taken $Y_{\rm SSC}\approx 3.2\epsilon_{e,-1}^{1/2}\epsilon_{B,-2}^{-1/2}$)

$$\begin{aligned}f_\pi(1) &\sim \sigma_\Delta n_\gamma(1)\Delta R'\langle x_{p\rightarrow\pi}\rangle\\ &\approx \frac{\epsilon_e E}{2(1+Y_{\rm SSC})\left(h\nu_m^f\right)4\pi RR_\times}\sigma_\Delta\langle x_{p\rightarrow\pi}\rangle\\ &\approx 0.66\epsilon_{e,-1}^{-3/2}L_{52.6}^{-1/2}A_*^{3/2},\end{aligned} \quad (18)$$

where $\sigma_\Delta\approx 5\times10^{-28}$ cm^2 is the cross section of the Δ-resonance, $\langle x_{p\rightarrow\pi}\rangle\simeq 0.2$ is the average fraction of the energy transferred from the initial proton to the produced pion, $\Delta R'$ is the width of the shock, and $n_\gamma(1)\approx (R/R_\times)[\epsilon_e E/(2h\nu_m^f)]/[4\pi(1+Y_{\rm SSC})R^2\Delta R']$ (for $R\leq R_\times$) is the photon number density, where $E\sim 4\times10^{53}$ ergs is the isotropic energy of the ejecta at the beginning of the afterglow, as defined in § 2. It is interesting to see that $f_\pi(1)$ is independent of t, since $\nu_m^f\propto R^{-1}$.

Below we estimate the maximum proton energy $\left(\epsilon_p^M\right)$ accelerated by the shocks. In general, $\epsilon_p^M=\min[\epsilon_p^M(1),\epsilon_p^M(2),\epsilon_p^M(3)]$ satisfies three constraints.

1. The comoving shock acceleration time $t_a'\sim\epsilon_p/\Gamma eBc$ should be smaller than the comoving wind expansion time $t_d'\sim R/\Gamma c$, which yields $\epsilon_p^M(1)\sim eBR$. The numerical value reads

$$\epsilon_p^M(1)\simeq\left(5\times10^{20}\ {\rm eV}\right)\epsilon_{B,-2}^{1/2}L_{52.6}^{1/4}A_*^{-1/4}. \quad (19)$$

In this work, the superscript $'$ represents the parameter measured in the comoving frame.

2. The comoving proton synchrotron cooling timescale $t_{\rm cool}'=(6\pi m_p^4c^3/\sigma_{\rm T}m_e^2)\Gamma\epsilon_p^{-1}B^{-2}$ should be longer than the comoving acceleration timescale t_a' (e.g., Li et al. 2002), which results in

$$\epsilon_p^M(2)\simeq\left(4\times10^{20}\ {\rm eV}\right)\epsilon_{B,-2}^{-1/4}L_{52.6}^{1/8}A_*^{-1/8}R_{15}^{1/2}. \quad (20)$$

3. The comoving proton cooling timescale due to photomeson interaction t_π' should also be longer than the comoving acceleration timescale t_a'. Here t_π' can be derived from $\bar{f}_\pi\simeq t_d'/t_\pi'$ (e.g., Waxman & Bahcall 1997); i.e., $t_\pi'\simeq(R/\Gamma c)\bar{f}_\pi^{-1}$, where $\bar{f}_\pi=f_\pi(1)(\epsilon_p/\epsilon_p^b)^{1/2}$ for $\epsilon_p>\epsilon_p^b(1)$. Requiring $t_\pi'=t_a'$ yields $\epsilon_p/(\Gamma eBc)=R/\left(\bar{f}_\pi\Gamma c\right)$, which can be simplified as $\epsilon_p^M(3)\simeq\bar{f}_\pi^{-1}\epsilon_p^M(1)\simeq f_\pi(1)[\epsilon_p^M(3)/\epsilon_p^b]^{-1/2}\epsilon_p^M(1)$. We then have

$$\begin{aligned}\epsilon_p^M(3) &\approx f_\pi(1)^{-2/3}\left[\epsilon_p^M(1)\right]^{2/3}\left(\epsilon_p^b\right)^{1/3}\\ &\simeq\left(4\times10^{19}\ {\rm eV}\right)\epsilon_{e,-1}^{1/3}\epsilon_{B,-2}^{1/6}L_{52.6}^{1/3}A_*^{-1}R_{15}^{1/3}.\end{aligned} \quad (21)$$

Since $\epsilon_{p,{\rm obs}}^b\ll\epsilon_p^M=\epsilon_p^M(3)\sim 4\times10^{19}$ eV, the predicted flux of the neutrino emission is insensitive to the actual value of ϵ_p^M.

After the pions (muons) are generated, the high-energy pions (muons) may lose energy via synchrotron emission before decaying, thus reducing the energy of the decay neutrinos (e.g., Guetta et al. 2004). For pions, this effect should be taken into account when the pion lifetime $\tau_\pi'\approx 2.6\times10^{-8}\epsilon_\pi'/(m_\pi c^2)$ s is comparable to the synchrotron loss time $t_{\pi,{\rm syn}}'=6\pi m_\pi^4c^3/(\sigma_{\rm T}m_e^2\epsilon_\pi'B^2)$. We can then define a critical pion energy $\epsilon_\pi'^c$ by requiring $\tau_\pi'=t_{\pi,{\rm syn}}'$, which gives

$$\epsilon_\pi'^c\approx\left(1.1\times10^{17}\ {\rm eV}\right)\epsilon_{B,-2}^{-1/2}L_{52.6}^{-1/4}A_*^{1/4}R_{15}, \quad (22)$$

above which the neutrino emission is suppressed significantly. Correspondingly, we have a critical observed energy for ν_μ (since $\epsilon_\pi'^c\approx 4\epsilon_{\nu_\mu}'^c$),

$$\epsilon_{\nu_\mu,{\rm obs}}^c\approx\frac{\Gamma\epsilon_\pi'^c}{4(1+z)}\approx\left(1.2\times10^{18}\ {\rm eV}\right)\left(\frac{1+z}{2}\right)^{-1}\epsilon_{B,-2}^{-1/2}R_{15}, \quad (23)$$

above which the slope of the ν_μ spectrum steepens by 2, since $t_{\pi,{\rm syn}}'/\tau_\pi'\propto\epsilon_\pi'^{-2}$.

Muons have a lifetime $\sim$100 times longer than that of pions, so that the energy cutoff of $\bar{\nu}_\mu$ and ν_e will therefore be 10 times smaller than $\epsilon_{\nu_\mu,{\rm obs}}^c$ (e.g., Guetta et al. 2004). With equation (23) we have

$$\epsilon_{(\bar{\nu}_\mu,\nu_e),{\rm obs}}^c\approx\frac{\epsilon_{\nu_\mu,{\rm obs}}^c}{10}\approx\left(1.2\times10^{17}\ {\rm eV}\right)\left(\frac{1+z}{2}\right)^{-1}\epsilon_{B,-2}^{-1/2}R_{15}. \quad (24)$$

Equation (18) is valid only for protons at the break energy $\epsilon^b_{p,\rm obs}(1)$. Considering emission from all the protons, the resulting neutrino spectrum would trace the photon spectrum. Generally one has

$$f_\pi \approx \begin{cases} f_\pi(1)\left[\dfrac{\epsilon_{p,\rm obs}}{\epsilon^b_{p,\rm obs}(1)}\right]^{1.1}, & \epsilon_{p,\rm obs} < \epsilon^b_{p,\rm obs}(1), \\ f_\pi(1)\left[\dfrac{\epsilon_{p,\rm obs}}{\epsilon^b_{p,\rm obs}(1)}\right]^{0.5}, & \epsilon^b_{p,\rm obs}(1) < \epsilon_{p,\rm obs} < \epsilon^0_{p,\rm obs}, \\ 1, & \epsilon^0_{p,\rm obs} < \epsilon_{p,\rm obs} < \epsilon^M_{p,\rm obs}, \end{cases} \tag{25}$$

where $\epsilon^0_{p,\rm obs}$ is defined as $(\epsilon^0_{p,\rm obs}/\epsilon^b_{p,\rm obs})^{0.5} = 1/f_\pi(1)$; $\epsilon^M_{p,\rm obs} = \epsilon^M_p/(1+z)$. The neutrino spectrum finally reads

$$\frac{\epsilon^2_{\nu,\rm obs} d^2N_\nu}{d\epsilon_{\nu,\rm obs}dt} \simeq \frac{1}{8}\frac{(1+z)L}{4\pi D_L^2} \times \begin{cases} f_\pi(1)\left[\dfrac{\epsilon_{\nu,\rm obs}}{\epsilon^b_{\nu,\rm obs}(1)}\right]^{1.1}, & \epsilon_{\nu,\rm obs} < \epsilon^b_{\nu,\rm obs}(1), \\ f_\pi(1)\left[\dfrac{\epsilon_{\nu,\rm obs}}{\epsilon^b_{\nu,\rm obs}(1)}\right]^{0.5}, & \epsilon^b_{\nu,\rm obs}(1) < \epsilon_{\nu,\rm obs} < \epsilon^0_{\nu,\rm obs}, \\ 1, & \epsilon^0_{\nu,\rm obs} < \epsilon_{\nu,\rm obs} < \epsilon^c_{\nu,\rm obs}, \\ \left(\dfrac{\epsilon_{\nu,\rm obs}}{\epsilon^c_{\nu,\rm obs}}\right)^{-2}, & \epsilon^c_{\nu,\rm obs} < \epsilon_{\nu,\rm obs} < \epsilon^M_{\nu,\rm obs}, \end{cases} \tag{26}$$

where the factor of $\frac{1}{8}$ takes into account the fact that charged and neutral pions are produced with roughly equal probabilities and that each neutrino carries $\sim\frac{1}{4}$ of the pion energy, i.e., $\epsilon^0_{\nu,\rm obs} \approx 0.05\epsilon^0_{p,\rm obs}$. In equation (26), $\epsilon^c_{\nu,\rm obs} < \epsilon^M_{\nu,\rm obs}$ is assumed. In the case of $\epsilon^c_{\nu,\rm obs} > \epsilon^M_{\nu,\rm obs}$, the neutrino flux is 0 for $\epsilon_{\nu,\rm obs} > \epsilon^M_{\nu,\rm obs}$.

As in many publications, we assume that the protons in the FS are distributed as $dn/d\gamma_p \propto \gamma_p^{-2}$ for $\gamma_{p,m} < \gamma_p < \gamma_{p,M}$, where $\gamma_{p,m} = (1-\epsilon_e)(\Gamma-1)/\ln(\gamma_{p,M}/\gamma_{p,m})$ and $\gamma_{p,M} = \epsilon^M_p/\Gamma m_p c^2$. For a neutrino detector with a surface area of $S_{\rm det} \sim 1$ km^2, the detectable events for one burst can be estimated by [taking $R = 0.5R_\times$ (wind)]

$$\begin{aligned} N_{\rm events}(1) &\approx \frac{1}{8}\frac{(1-\epsilon_e)E}{\epsilon^b_{\nu,\rm obs}(1)}\frac{1}{4\pi D_L^2}S_{\rm det}P_{\nu\to\mu} \\ &\quad\times \frac{\ln\left[\epsilon^0_{\nu,\rm obs}/\epsilon^b_{\nu,\rm obs}(1)\right]f_\pi(1) + \ln\left(\epsilon^c_{\nu,\rm obs}/\epsilon^0_{\nu,\rm obs}\right)}{\ln\left(\gamma_{p,M}/\gamma_{p,m}\right)} \\ &\sim 0.002\,E_{53.6}D^{-2}_{L,28.34}, \end{aligned} \tag{27}$$

where $P_{\nu\to\mu} \simeq 6\times10^{-4}\left[\epsilon_{\nu,\rm obs}/(3\times10^{15}\ \rm eV)\right]^{0.5}$ is the probability that a neutrino produces a detectable high-energy muon for $\epsilon_{\nu,\rm obs} > 10^3$ TeV (Gaisser et al. 1995). The factor of $\frac{1}{8}$ takes into account the fact that there are three neutrinos produced, with each carrying about $\frac{1}{8}$ of the energy loss of the proton, and that due to neutrino flavor oscillation, only $\frac{1}{3}$ of the neutrinos are the detectable muon neutrinos. For a nearby GRB with $z \sim 0.1$ ($D_L \sim 1.4\times10^{27}$ cm) and $E \sim 4\times10^{53}$ ergs, the number of resulting detectable events is $\sim$0.4. For such a nearby energetic GRB, the prompt γ-ray energy fluence would be $F_\gamma = (1+z)E_\gamma/(4\pi D_L^2) \sim (10^{-3}$ ergs cm$^{-2})E_{\gamma,53}D^{-2}_{L,27.15}$. This means that only for those very bright GRBs can the predicted neutrinos be detected. Similar conclusions have been also drawn in Dermer & Atoyan (2003) and Guetta et al. (2004). We note that low-redshift long GRBs are rare. So far we have detected at least three of them, but neither of them is energetic enough for the above purpose.

Another method to estimate the detectability of the neutrino early afterglow emission is to assume that GRBs are the dominant sources of high-energy cosmic rays (from the internal or external shocks; Waxman 1995; Vietri 1995). This leads to a detectability about an order of magnitude higher (see, e.g., DL01) than the value presented above. One possible reason for the discrepancy is that this second method is based on the assumption that ultra–high-energy cosmic rays are mostly produced by GRBs, which might not be the case.

The atmospheric neutrino background flux is $\phi_{\nu,\rm bkg} \sim 10^{-12}\left[\epsilon_{\nu,\rm obs}/(10^{14}\ \rm eV)\right]^{-5/2}$ cm^2 s^{-1} sr^{-1}. The typical angular resolution of the planned neutrino telescopes at TeV energies is about 1°. For a typical GRB with $T_{90} \sim 20$ s, $E \sim 4\times10^{53}$ ergs, and $z = 1$, at $\epsilon_{\nu,\rm obs} \sim 5\times10^{15}$ eV, the background neutrino number in this time span is $\sim5\times10^{-9}$, which is $\ll$0.002 (eq. [27]). If a $\sim5\times10^{15}$ eV neutrino that is well correlated with one GRB both in time and coordinates has been detected, the chance probability of being an atmospheric neutrino is only $\sim3\times10^{-5}$. Therefore, these GRB neutrinos can be detected above the background (see also Mészáros & Waxman 2001).

4.2. *Neutrinos from Photomeson Interaction with the MeV Photon Flow:* $\sim10^{14}$ eV

Because of the significant overlapping of the MeV photon flow and the shocked regions, high-energy protons with energy larger than 2×10^{15} eV accelerated from both the RS and the FS would interact with the MeV flow effectively. We now study the neutrino emission process in this case.

In the rest frame of the shocked regions, γ-ray photons with the observed typical energy $\epsilon^b_{\gamma,\rm obs} \approx 300$ keV have an energy $\epsilon^b_\gamma \approx (1+z)\epsilon^b_{\gamma,\rm obs}/\Gamma$. The protons that interact with the photons with energy ϵ^b_γ at the Δ-resonance have an observed energy of

$$\begin{aligned} \epsilon^b_{p,\rm obs}(2) &\approx (0.3\ \rm GeV)^2\Gamma^2/(1+z)^2\epsilon^b_{\gamma,\rm obs} \\ &= \left(2\times10^{15}\ \rm eV\right)L^{1/2}_{52.6}A_*^{-1/2}\left(\frac{1+z}{2}\right)^{-2}\left(\frac{\epsilon^b_{\gamma,\rm obs}}{300\ \rm keV}\right)^{-1}. \end{aligned} \tag{28}$$

This energy is well below the observed maximum proton energy ϵ^M_p (eqs. [19]–[21]). The generated neutrinos have a typical energy

$$\epsilon^b_{\nu,\rm obs}(2) \sim \left(10^{14}\ \rm eV\right)L^{1/2}_{52.6}A_*^{-1/2}\left(\frac{1+z}{2}\right)^{-2}\left(\frac{\epsilon^b_{\gamma,\rm obs}}{300\ \rm keV}\right)^{-1}. \tag{29}$$

Similar to § 4.1, at $\epsilon^b_{p,\rm obs}(2)$, the fraction of the energy loss of protons through photomeson interaction with the MeV photons can be parameterized as

$$f_\pi(2) \approx 0.18L_{\gamma,52}L^{-1/2}_{52.6}A_*^{1/2}\left[\frac{(1+z)\epsilon^b_{\gamma,\rm obs}}{600\ \rm keV}\right]^{-1}R^{-1}_{15}. \tag{30}$$

The MeV flow photons have a broken power-law spectrum, i.e., $n(\epsilon_\gamma) \propto \epsilon_\gamma^{-1}$ below $\epsilon_{\gamma,\rm obs}$ and $n(\epsilon_\gamma) \propto \epsilon_\gamma^{-2}$ above $\epsilon^b_{\gamma,\rm obs}$. So the general form of the fraction of the energy loss of protons through photomeson interaction with the MeV photons is

$$f_\pi \approx f_\pi(2) \begin{cases} 1, & \epsilon_{p,\mathrm{obs}} > \epsilon^b_{p,\mathrm{obs}}(2), \\ \dfrac{\epsilon_{p,\mathrm{obs}}}{\epsilon^b_{p,\mathrm{obs}}(2)}, & \epsilon_{p,\mathrm{obs}} \le \epsilon^b_{p,\mathrm{obs}}(2). \end{cases} \tag{31}$$

As a result, the differential neutrino spectrum is also a broken power law, i.e., $n(\epsilon_\nu) \propto \epsilon_\nu^{-1}$ below $\epsilon^b_{\nu,\mathrm{obs}}(2)$ and $n(\epsilon_\nu) \propto \epsilon_\nu^{-2}$ above $\epsilon^b_{\nu,\mathrm{obs}}(2)$ (see also Waxman & Bahcall 1997). However, for $\epsilon_{\nu,\mathrm{obs}} > \epsilon^0_{\nu,\mathrm{obs}}$, the energy of the corresponding protons has been lost mainly by interacting with the keV soft photons from the FS region (see § 4.1 for details). So the neutrino spectrum cannot extend to $\epsilon_{\nu,\mathrm{obs}} > \epsilon^0_{\nu,\mathrm{obs}}$.

Similar to equation (27), we now estimate the detectability of such 10^{14} eV neutrinos in the early afterglow phase. For one typical GRB, taking $R = 0.5R_\times$ (wind), one has

$$\begin{aligned} N_{\mathrm{events}}(2) &\approx \frac{1}{8}\frac{(1-\epsilon_e)E}{\epsilon^b_{\nu,\mathrm{obs}}(2)}\frac{1}{4\pi D_L^2} S_{\mathrm{det}} P_{\nu\to\mu} \\ &\times \frac{\ln\left[\epsilon^0_{\nu,\mathrm{obs}}/\epsilon^b_{\nu,\mathrm{obs}}(2)\right] f_\pi(2)}{\ln(\gamma_{p,M}/\gamma_{p,m})} \\ &\sim 0.001 E_{53.6} D_{L,28.34}^{-2}, \end{aligned} \tag{32}$$

where $P_{\nu\to\mu} \simeq 1.7\times10^{-6}[\epsilon_{\nu,\mathrm{obs}}/(1\ \mathrm{TeV})]^{0.8}$ (e.g., Halzen & Hooper 2002) for 1 TeV $< \epsilon_{\nu,\mathrm{obs}} < 10^3$ TeV. The resulting detectability is comparable with that of 5×10^{15} eV (see eq. [27]). These neutrinos can be only marginally detectable above the atmospheric neutrino background. For a typical GRB with $T_{90} \sim 20$ s, $E \sim 4\times10^{53}$ ergs, and $z = 1$, the 10^{14} eV neutrino background event number is $\sim 9\times10^{-5}$ in the time span. This is <0.001 (eq. [32]), but not significantly. For a detected $\sim 10^{14}$ eV neutrino well correlated with one GRB, the chance probability of being an atmospheric neutrino is $\sim$10%.

The total detectable events for one typical GRB born in the wind environment ($z = 1$) is $\sim$0.003; this rate can also be enhanced significantly when fluctuations in distance and energy are considered (e.g., Halzen & Hooper 2002). Therefore, nearby GRBs such as GRB 030329 are of great interest (e.g., Razzaque et al. 2004). Thanks to the correlation of the neutrinos with GRBs, both in time and coordinates, these GRB neutrinos are likely to be detected with the planned cubic kilometer high-energy neutrino detectors such as IceCube, Astronomy with a Neutrino Telescope and Abyss Environmental Research (ANTARES), and the Neutrino Extended Submarine Telescope with Oceanographic Research (NESTOR).

One caveat is that the $\sim 10^{14}$ eV neutrinos would also overlap with the $\sim 5\times10^{14}$ eV neutrinos generated from the internal shocks (e.g., Waxman & Bahcall 1997), so that it might not be easy to distinguish both species. We note, however, that if GRB ejecta are strongly magnetized so that the prompt emission is not powered by internal shocks, the 10^{14} eV neutrinos discussed here are still produced.

5. DISCUSSION AND CONCLUSION

Massive stars are likely the progenitors of long, soft GRBs. For an ultrarelativistic fireball emerging from a massive star and expanding into a dense stellar wind medium, a strong relativistic RS develops so early that both the RS and FS regions overlap in time and space with the prompt MeV photon flow. This overlapping leads to significant modification of the early RS and FS emission, since the dominant cooling process for electrons is likely IC with the MeV photons (especially for the RS region [Beloborodov 2005]; see also Fan et al. [2005a] for a more detailed calculation on the modification of the very early *R*-band light curve).

Due to the tight overlapping of the MeV photon flow and the shocked regions (both FS and RS), the optical depth of the multi-GeV photons is larger than unity. These high-energy photons (SSC and/or IC emission components from the FS and RS regions) are absorbed by the MeV photon flow and therefore generate relativistic $e^\pm$ pairs. These $e^\pm$ pairs rescatter the soft X-rays in the FS region and the prompt γ-rays and power detectable high-energy emission. A significant portion of the reprocessed photon energy is distributed in the sub-GeV band, making wind-interaction GRBs interesting targets for *GLAST.* In the wind model, the total thermal energy contained in the FS region is larger than or at least comparable with that contained in RS region. The above effect is therefore generic, regardless of the possible magnetization of the ejecta (e.g., Fan et al. 2002; Zhang et al. 2003; in which case the IC emission as well as the SSC emission of the RS are suppressed). The predicted fluence of these sub-GeV flashes is high enough to be detected by *GLAST.*

Equally interesting is that the early photon-shock interaction leads to interesting neutrino emission signatures. Besides the conventional neutrino component powered by the photomeson interaction in the RS region (DL01), in this paper we studied two new neutrino emission components, i.e., a $\sim 5\times10^{15}$ eV component powered by the photomeson interaction in the FS region and a $\sim 10^{14}$ eV component powered by photomeson interaction with the MeV photon flow. The detectability of the two components is comparable. The spectrum of a 10^{14} eV neutrino is similar to that of the neutrinos generated in the internal shock phase, since the seed photons are nearly the same, although the typical observed energy is softened by a factor of $\sim(\Gamma/\Gamma_0)^2$. The spectrum of the 5×10^{15} eV neutrinos takes a more complicated form (see eq. [26]). A caveat is that these neutrinos overlap with the $\sim 5\times10^{14}$ eV internal shock neutrino component (Waxman & Bahcall 1997) and hence are difficult to distinguish. On the other hand, they are generic regardless of whether the MeV emission is powered by the internal shocks or not.

In this work, it is assumed that the wind medium is at rest at any radius. This may not be the case if the radiation front of the GRB is taken into account. As shown in Beloborodov (2002), due to the large amount of $e^\pm$ pair loading (the pairs are created by the interaction of the prompt γ-rays with the backscattered γ-rays by the wind medium), the medium will be accelerated, and the ejecta will move in a cavity until it reaches a radius $R_{\mathrm{gap}} \sim 3.3\times10^{15}E^{1/2}_{\gamma,53}$ (e.g., Beloborodov 2002). As a result, the shock-crossing radius (see eq. [2]) should be $\sim R_{\mathrm{gap}} + R_\times$ (wind) $\sim 2R_\times$ (wind). The overlapping of the MeV photon flow and the shocked regions is still significant, since the ejecta essentially does not decelerate at $R < R_{\mathrm{gap}}$. However, $f_\pi(1)$ and $f_\pi(2)$ involved in the above calculation should be smaller by a factor of about $\frac{1}{4}$; so are the predicted neutrino fluxes. The optical depth for the multi-GeV photons (see eq. [12]) is $\sim$1, so that they are still absorbed by the MeV photon flow and generate relativistic $e^\pm$ pairs, which rescatter the initial MeV γ-rays to the sub-GeV energy range. Note that the introduction of the radiation front will soften the FS shock emission significantly, since the upstream of the FS is $e^\pm$-rich. If we assume that the number of pairs is *k* times that of the electrons associated with protons, γ_m^f (see eq. [8]) should be reduced to $\gamma_m^f/(1 + 2k)$. If k is as large as several tens, the electrons accelerated in the FS will be cooled mainly via IC with the MeV photon flow rather than synchrotron radiation and SSC emission. A significant part of that IC component may be in the sub-GeV energy range. Therefore, sub-GeV flashes are still expected when the radiation front effect is taken into account.

Only the wind model is discussed in this work. In the ISM model, if the GRB is long enough, the overlapping of the MeV photon flow and the shocked regions could also be tight. In this case, $R_\times$ (ISM) $\sim 10^{17}$ cm, and the optical depth of the GeV photons is very small (since $\tau_{\gamma\gamma} \propto R^{-2}$; see eq. [12]). In such a case, GeV–TeV flashes are expected, as suggested by Beloborodov (2005). Therefore, the sub-GeV photon flashes are a signature of GRBs born in a stellar wind.

We note that the overlapping effect increases the neutrino flux predicted from the collapsar GRB model to be close to that predicted in the supranova model (cf. Dermer & Atoyan 2003). This is because the overlapping MeV photons effectively provide a convenient target for photomeson interactions, similar to the external photons involved in the supranova model.

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REFERENCES

Beloborodov, A. M. 2002, ApJ, 565, 808
———. 2005, ApJ, 618, L13
Blake, C. H., et al. 2005, Nature, 435, 181
Chevalier, R. A., & Li, Z. Y. 2000, ApJ, 536, 195
Dai, Z. G., & Lu, T. 1998, MNRAS, 298, 87
———. 2001, ApJ, 551, 249 (DL01)
Dermer, C. D., & Atoyan, A. 2003, Phys. Rev. Lett., 91, 071102
Fan, Y. Z., Dai, Z. G., Huang, Y. F., & Lu, T. 2002, Chinese J. Astron. Astrophys., 2, 449
Fan, Y. Z., & Wei, D. M. 2004, ApJ, 615, L69
Fan, Y. Z., Zhang, B., & Wei, D. M. 2005a, ApJ, 628, 298
———. 2005b, ApJ, 628, L25
Gaisser, T. K., Halzen, F., & Stanev, T. 1995, Phys. Rep., 258, 173
González, M. M., Dingus, B. L., Kaneko, Y., Preece, R. D., Dermer, C. D., & Briggs, M. S. 2003, Nature, 424, 749
Guetta, D., Hooper, D., Aliarez-Muniz, J., Halzen, F., & Reuveni, E. 2004, Astropart. Phys., 20, 429
Halzen, F., & Hooper, D. 2002, Rep. Prog. Phys., 65, 1025
Li, Z., Dai, Z. G., & Lu, T. 2002, A&A, 396, 303
Mészáros, P., Rees, M. J., & Wijers, R. A. M. J. 1998, ApJ, 499, 301
Mészáros, P., & Waxman, E. 2001, Phys. Rev. Lett., 87, 171102
Paczyński, B. 1998, ApJ, 494, L45
Panaitescu, A., & Kumar, P. 2001, ApJ, 554, 667
Razzaque, S., Mészáros, P., & Waxman, E. 2004, Phys. Rev. D, 69, 023001
Svensson, R. 1987, MNRAS, 227, 403
Vestrand, W. T., et al. 2005, Nature, 435, 178
Vietri, M. 1995, ApJ, 453, 883
Waxman, E. 1995, ApJ, 452, L1
Waxman, E., & Bahcall, J. N. 1997, Phys. Rev. Lett., 78, 2292
Woosley, S. E. 1993, ApJ, 405, 273
Zhang, B., Kobayashi, S., & Mészáros, P. 2003, ApJ, 595, 950
Zhang, B., & Mészáros, P. 2001, ApJ, 559, 110