EARLY PHOTON-SHOCK INTERACTION IN A STELLAR WIND: A SUB-GEV PHOTON FLASH AND HIGH ENERGY NEUTRINO EMISSION FROM LONG GRBS

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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⁺e⁻ pairs. These pairs re-scatter the soft X-ray photons from the forward shock as well as the prompt γ-ray photons and power detectable high energy emission, 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 on whether the GRB ejecta are magnetized (in which case the reverse shock IC 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 typically detectable by the future Gamma-Ray Large Area Telescope (GLAST). Overlapping also influence neutrino emission. Besides the $10^{15} \sim 10^{17}$ 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 the one generated in the internal shock phase, but the typical energy is slightly lower.

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

1. INTRODUCTION

The leading model for long gamma-ray bursts (GRBs) invokes a relativistic jet emerging from a collapsar (e.g. Woosley 1993; Paczynski 1998). It has been widely expected that GRBs associated with supernovae should occur in a pre-burst stellar wind environment with particle number density $n \propto R^{-2}$ (e.g. Dai & Lu 1998; Mészáros, Rees & Wijers 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 of $3 \times 10^{15}$ eV - $3 \times 10^{17}$ 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. Fin 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 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

$$R_x \ll 2\Gamma_0^2 \Delta,$$

where $\Gamma_0$ is the initial Lorentz factor of the ejecta (at the internal shock phase), $\Delta = cT_90/(1+z)$, $T_90 \sim 20s$ is the observed duration of the long GRB, and $z \sim 1$ is the redshift. The shock crossing radius $R_x$ can be written in a general form

$$R_x = \max(R_\gamma, \Gamma^2 \Delta),$$

where $R_\gamma$ 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 $\Gamma$ is the Lorentz factor of the shocked ejecta. In equation 2, the first term dominates if the RS is non-relativistic, while the second term dominates if the RS is relativistic. For the ISM case with a typical number density $n_{\text{ISM}} \sim 1\text{cm}^{-3}$, the RS is non-relativistic and one has $R_x(\text{ISM}) = R_\gamma \approx 8.8 \times 10^{16} E_{53.6}^{-1/3} \Gamma^{-2/3} \text{cm}$. 


which is usually larger than (or at least comparable to) $2F_0^2\Delta = 5.4 \times 10^{16}\text{cm} \Gamma_0.25 \Delta_{11.5}$. Throughout this paper, we adopt the convention $Q_x = Q/10^x$ to express physical parameters in cgs units. So, in the ISM case, the initial $\gamma$-rays cannot change the RS emission significantly. In the stellar wind case, on the other hand, one has $n = 3.0 \times 10^{35} A, R^{-2}\text{cm}^{-3}$, where $A = (M/10^{-5} M_\odot \text{yr}^{-1})(w/10^6 \text{cm} \text{s}^{-1})^{-1}, M$ is the mass loss rate of the progenitor, and $w$ is the wind velocity (Chevalier & Li 2000). With such a dense medium, the RS is relativistic (Chevalier & Li 2000; DL01): $\Gamma \approx \zeta_{1/2}^{1/2} / \sqrt{2},$ where $\zeta \approx 2600E_{53,6}/(A_7 \Delta_{11.5} \Gamma_2 2.5)$ is the ratio of proper mass density of the un-shocked ejecta and the medium density, $E \sim 4 \times 10^{53}$ ergs is the isotropic energy of the ejecta, which is calculated by adopting $T_90 = 20\text{s}, z = 1$, and by assuming that the observed $\gamma$-ray luminosity is $L_y \approx 10^{52}$ ergs s$^{-1}$ and that the radiation efficiency of the internal shocks is $\eta \approx 0.2$. The total wind luminosity is therefore $L_{\text{tot}} = L_y/\eta = 5 \times 10^{52}$ ergs s$^{-1}$, and the afterglow kinetic energy luminosity is $L \approx L_{\text{tot}}(1 - \eta) = 4 \times 10^{52}$ ergs s$^{-1}$ in the beginning of the afterglow phase. The resulting bulk Lorentz factor in the shocked region is therefore

$$\Gamma \approx 87 L_{52.6} A_{-1/4}^{1/4}. \quad (3)$$

The RS crosses the ejecta at the radius

$$R_\times (\text{wind}) \approx 2.7 \times 10^{15} \text{cm} L_{52.6}^{1/2} A_{-1/2}^{1/2} \Delta_{11.5}. \quad (4)$$

At this radius, the rear of the MeV photon flow only leads the rear of the ejecta by a distance of $R_\times (\text{wind})/2T_9^2 \ll \Delta$, so that during most of the time when 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 $Y \equiv U_r/U_b$, where $U_r \approx L_y/4\pi R^2 T^2 e$ is the MeV photon energy density in the rest frame of the shocked region, and $U_b = B^2/8\pi \approx (\Gamma_0/T)^2 e^{-1} B_n m_p c^2$ is the comoving magnetic energy density in the same region, where $n_4 = (1 - \eta) L/4\pi R^2 T^2 m_p c^3$ is the comoving density of the un-shocked ejecta. As usual, we introduce magnetic field and electron equipartition parameters $\epsilon_B \sim 0.001 - 0.1$ and $\epsilon_e \sim 0.1$ in the shocked region. The wide range of $\epsilon_B$ 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

$$Y \approx \eta/(1 - \eta) \epsilon_B. \quad (5)$$

For $\epsilon_B \sim 0.001 - 0.1$ and $\eta \sim 0.2$, we have $Y \sim 2.5 - 250 \gg 1$. The cooling frequency $\nu_c$ (see Eq. (15) of DL01) would then be lowered by a factor $1/(1 + Y)^2$. For $\epsilon_B, \sim 0.1 - 1$, the flux of the soft ultraviolet-optical photons is significantly lower than the one calculated in DL01, so that the energy loss fraction for each proton through pion production (i.e. $f_\pi$, defined as the ratio between the comoving expansion time $t_4^\pi$ and the comoving proton cooling time $t_4^\pi$ due to photonsem cooling) should be also lowered by roughly a factor of $1/(1 + Y)$ (see equations [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 $\gamma$--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 on 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 model, 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 GRB990123 (Fan et al. 2002; Zhang et al. 2003), the $Y$ parameter (equation (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 will 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. Also, the typical energy of the emerged spectrum peaks in the sub-GeV range (see below).

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

$$E_r^{\text{tot}} \approx \gamma_{34} \Gamma M_{ej} c^2 \approx \Gamma_0 M_{ej} c^2 / 2 \approx E_f^{\text{tot}}, \quad (6)$$

where we have taken the Lorentz factor of the shocked ejecta relative to the unshocked ejecta as $\gamma_{34} \approx (\Gamma_0/\Gamma + \Gamma_0/\Gamma_0)/2 \approx \Gamma_0/2 \Gamma \approx 1$ (which is valid for the wind RS model). Here $M_{ej} \approx E/(\Gamma \rho c^2)$ 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'll show below that it 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 $E_{th}^r \approx (\gamma_{34} - 1) E_f^{\text{tot}} / \gamma_{34}$ and $E_{th}^f \approx (1 - \Gamma) E_f^{\text{tot}} / \Gamma$, respectively. The ratio of them is $E_{th}^r/E_{th}^f \approx (\gamma_{34} - 1) / \gamma_{34}$. For typical parameters, $\gamma_{34} \sim 2$, we have $E_{th}^r : E_{th}^f \approx 1 : 2$. 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 $dN_e/d\gamma_e \propto \gamma_e^{-\pi}$ ($p \sim 2.2$) for $\gamma_m < \gamma_e < \gamma_M$, where $\approx \gamma_{34} - 1 \approx \gamma_{34}$.
where $\gamma_m \sim 10^8/B^{1/2}$ is the maximum Lorentz factor of the electrons accelerated by shocks, $B$ is the magnetic field strength of the shock regions, which can be calculated by

$$B \approx 1.86 \times 10^9 G \frac{1/2}{\epsilon_{b,-2} A_{\ast}^{-1/4}} L_{52.6}^{1/4} R_{15}^{-1}$$

(7)

For the FS and RS, $\gamma_m$ can be estimated by

$$\gamma_m = \langle 1 + z \rangle \epsilon_{A,\gamma,\gamma,\gamma} \frac{\langle p - 2 \rangle}{(p - 1)}$$

$$\approx 2600 \epsilon_{A,\gamma,\gamma,\gamma} \frac{1}{L_{52.6}}$$

$$\gamma_m = (\tau_{\gamma,\gamma,\gamma} - 1) \epsilon_{A,\gamma,\gamma,\gamma} \frac{\langle p - 2 \rangle}{(p - 1)}$$

$$\approx 23 \epsilon_{A,\gamma,\gamma,\gamma} L_{52.6}^{-1/4} R_{15}^{-1/4} G_{2.5}$$

(8)

Performing IC correction on equation (14) of DL01, we also get the cooling Lorentz factor $\gamma_e \lesssim 1$, 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 $\epsilon_{A,\gamma,\gamma,\gamma} = 300$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)

$$\gamma_e \approx \frac{1 + z}{\epsilon_{A,\gamma,\gamma,\gamma} \eta} \sim \Gamma.$$  

(9)

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

$$\epsilon_{A,\gamma,\gamma,\gamma} = \gamma_{A,\gamma,\gamma,\gamma} \epsilon_{A,\gamma,\gamma,\gamma} \sim 0.16\text{GeV} \epsilon_{A,\gamma,\gamma,\gamma}^{1/2} L_{52.6}^{1/2} \left( \frac{1}{300\text{keV}} \right).$$

(10)

In the FS region, on the other hand, one has $\gamma_{A,\gamma,\gamma,\gamma} > \gamma_e$, 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 $\gamma_{SSC} \approx \epsilon_{A,\gamma,\gamma,\gamma} e^B \gg 1$. The typical energy of SSC photons reads

$$\epsilon_{SSC} \approx \eta \left[ \gamma_e^{1/4} \Gamma \right] \left[ 2 \pi m_e c (1 + z) \right]$$

$$= 860 \text{GeV} (1 + z)^{-1/4} \epsilon_{A,\gamma,\gamma,\gamma}^{1/2} A_{\gamma,\gamma,\gamma}^{-3/2} L_{52.6}^{1/2} R_{15}^{-1/2},$$

(11)

where $h$ is the Planck’s constant.

Apparently, equations (11) and (12) are quite different. However, as shown below, photons more energetic than $\sim 1$ GeV will be absorbed by the MeV photons near the shock crossing radius, and most energy will 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 $\sim 1$ GeV (which will be absorbed by the soft photons with energy $\epsilon_{A,\gamma,\gamma,\gamma} \approx 2(\Gamma m_e c^2)/(1 + z)^2 \text{GeV} \sim 1$MeV) is roughly (e.g. Svensson 1987)

$$\tau_{\gamma,\gamma}(1\text{GeV}) \approx \frac{110 \pi N_{\gamma,\gamma,\gamma,\gamma}}{270 \pi R^2}$$

$$\approx 20 E_{\gamma,\gamma,\gamma,\gamma} \left( \frac{1 + z}{2} \right) L_{52.6}^{-1/2} A_{\gamma,\gamma,\gamma,\gamma}^{1/2} R_{15}^{-2}.$$  

(12)

where $N_{\gamma,\gamma,\gamma,\gamma}$ is the total number of photons satisfying $\epsilon_{A,\gamma,\gamma,\gamma} > \epsilon_{A,\gamma,\gamma,\gamma}$, which can be estimated by $N_{\gamma,\gamma,\gamma,\gamma} \approx 0.2 E_{\gamma,\gamma,\gamma,\gamma}/[(1 + z) \epsilon_{A,\gamma,\gamma,\gamma}]$ for $\epsilon_{A,\gamma,\gamma,\gamma} > \epsilon_{A,\gamma,\gamma,\gamma}^{1/4} \gamma_{A,\gamma,\gamma,\gamma}$, 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., $\epsilon_{\gamma,\gamma,\gamma,\gamma}(1 \text{GeV}) \approx 2 \times 10^{-6} \text{ergs cm}^{-2}$.
where $D_L$ is the luminosity distance, which $\approx 2.2 \times 10^{24}$ cm for $z = 1$. The fluence for the RS component may be comparable or less than this value. For EGRET, the estimated fluence threshold in the low integration time ($t_{\text{int}} < 1.7 \times 10^3$ s) is $\sim 2.1 \times 10^{-6}$ ergs cm$^{-2}$ for a typical photon energy $\sim 400$ MeV (e.g. Zhang & Mészáros 2001). Our predicted fluence $S \sim 10^{-6}$ ergs cm$^{-2}$ at sub-GeV energy range is below the EGRET detection threshold, which explains the none detection of such flashes in the EGRET era. We therefore expect that many sub-GeV photon flashes can be detected in 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 $\Delta$-resonance. In the RS region, DL01 have found that the typical neutrino energy is $3 \times 10^{15} - 3 \times 10^{17}$ eV.

We now estimate the neutrino emission from the FS region. As mentioned in [3] (see the paragraph below equation [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 $\nu_m^{\text{fs}}$ should be lowered by a factor $1/(1 + Y_{\text{SSC}})$. The typical synchrotron radiation frequency ($\nu_{m}^{\text{fs}}$) of the FS reads

$$\nu_{m}^{\text{fs}} \approx 3.1 \times 10^{15} \text{Hz} \ e_c^{1/2} \frac{1}{\text{L}_{52.6} A_{-1}^{1/2} R_{15}},$$

The protons that interact with the $h\nu_{m}^{\text{fs}}$ photons at the $\Delta$-resonance have an observed energy of

$$e_{p, \text{obs}}(1) \approx 0.3 \text{GeV} \frac{T^2}{(1+z)\nu_{m}^{\text{fs}}},$$

$$\approx 1.8 \times 10^{17} \text{eV} \ (1+z)^{-1} e_{c,-1}^{1/2} \frac{1}{\text{L}_{52.6} A_{1/2}^{1/2} R_{15}}.$$ (15)

The photomeson interactions include (1) production of $\pi$ mesons, $p\gamma \rightarrow p + \pi^0$ and $p\gamma \rightarrow n + \pi^+$, and (2) decay of $\pi$ 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 typical energy $h\nu_{m}^{\text{fs}}$, these processes produce neutrinos with energy $e_{p, \text{obs}}(1) \approx 0.05 e_{p, \text{obs}}(1)$ (e.g. Waxman & Bahcall 1997), i.e.

$$e_{p, \text{obs}}(1) \approx 9 \times 10^{15} \text{eV} \ (1+z)^{-1} e_{c,-1}^{1/2} \frac{1}{\text{L}_{52.6} A_{1/2}^{1/2} R_{15}}.$$ (17)

In principle, $\gamma$ -rays of similar energies are produced by $\pi^0$ decay. But as shown in [3] the MeV photon flow is optically thick for these high energy photons, and an $e^\pm$ pair cascade develops. The bulk $\gamma$-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 $e_{p, \text{obs}}(1)$, the fraction of energy loss of protons through photomeson interaction with the $h\nu_{m}^{\text{fs}}$ photons, $f_\pi(1)$, can be parameterized as (in the calculation, we have taken $Y_{\text{SSC}} \approx 3.2 e_c^{1/2} e_{c,-1}^{1/2}$)

$$f_\pi(1) \sim \sigma_\Delta n_\pi(1) \Delta R_{\pi} \ < p_{\pi,\text{c}} >$$

$$\approx \frac{e_c E}{(1+Y_{\text{SSC}}) h\nu_{m}^{\text{fs}} A_{1/2}^{1/2} R_{15}} \sigma_\Delta < p_{\pi,\text{c}} >$$

$$\approx 0.66 e_c^{3/2} \frac{1}{\text{L}_{52.6} A_{1/4}^{1/2}}.$$ (18)

where $\sigma_\Delta \approx 5 \times 10^{-28}$ cm$^2$ is the cross section of the $\Delta$-resonance, $< p_{\pi,\text{c}} > > 0.2$ is the average fraction of the energy transferred from the initial proton to the produced pion, $\Delta R_{\pi}$ is the width of the shock, and $n_\pi(1) \approx (R/R_s)(e_c E/(2h\nu_{m}^{\text{fs}}))/[4\pi (1+Y_{\text{SSC}}) R^2 \Delta R_{\pi}]$ (for $R \leq R_s$) is the photon number density, $E \sim 4 \times 10^{15}$ eV 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 on $t$ since $\nu_{m}^{\text{fs}} \sim R^{-1}$.

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

(1) The comoving shock acceleration time $t_s' \sim e_p/\Gamma e B c$ should be smaller than the comoving wind expansion time $t_d' \sim R/\Gamma c$, which yields $e_p^{M}(1) \sim e B R$. The numerical value reads

$$e_p^{M}(1) \approx 5 \times 10^{20} \text{eV} \ t_d'^{1/2} L_{52.6} A_{-1}^{1/4}.$$ (19)

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

(2) The comoving proton synchrotron cooling timescale

$$t_{\text{cool}} = [6 \pi m_p^4 e^3/\sigma_T m_e^3 \Gamma e_c B^-2] \text{B}^2$$

should be longer than the comoving acceleration timescale $t_{\text{a}}'(e.g., L i e t al. 2002)$, which results in

$$e_p^{M}(2) \approx 4 \times 10^{20} \text{eV} \ t_d'^{-1/4} L_{52.6} A_{-1}^{1/8} R_{15}^{-1/2}.$$ (20)

(3) The comoving proton cooling timescale due to photomeson interaction $t_s'$ should also be longer than the comoving acceleration timescale $t_{\text{a}}'$ of $f_\pi(1)$ (e.g. Waxman & Bahcall 1997), i.e.,

$$t_{\text{a}}' \approx (R/\Gamma c) f_\pi^{-1},$$

where $f_\pi = f_\pi(1)(e_p^{M}/e_p^{M}(1))^{2/3}$ for $e_p < e_p^{M}(1)$. Requiring $t_{\text{a}}' = t_{\text{a}}'$ yields $e_p/\Gamma e B c = R/[f_\pi(1)\Gamma c]$, which can be simplified as $e_p^{M}(3) \approx f_\pi^{-1} e_p^{M}(1) \approx f_\pi(1) e_p^{M}(3)/e_p^{M}(1)^{1/2}$. We then have

$$e_p^{M}(3) \approx f_\pi^{-1/2} e_p^{M}(1)^{2/3},$$

$$\approx 4 \times 10^{19} \text{eV} \ e_c^{1/3} e_{c,-1}^{1/6} L_{52.6} A_{-1}^{1/3} R_{15}^{-1/3}.$$ (21)

Since $e_{p, \text{obs}} < e_p^{M} = e_p^{M}(3) \sim 4 \times 10^{19}$ eV, the predicted flux of neutrino emission is insensitive to the actual value of $e_p^{M}$.

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 $\tau_{\pi} \approx 2.6 \times 10^{-8} e_c^{1/2} (m_{\pi} c^2)$ is comparable to the synchrotron loss time $t_{\pi, \text{syn}}' = 6 \pi m_e^4 e^3/(\sigma_T m_e^3 e_p^{M} B^2)$. We can then define a critical pion energy $e_\pi c$ by requiring $t_{\pi} = t_{\pi, \text{syn}}'$, which gives

$$e_\pi c \approx 1.1 \times 10^{17} \text{eV} \ e_c^{1/2} L_{52.6} A_{-1}^{1/4} R_{15},$$ (22)
above which the neutrino emission is suppressed significantly. Correspondingly, we have a critical observed energy for $\nu_\mu$ (since $E^\mu_{\nu_\mu} \approx 4E^\nu_{\nu_\mu}$),
\[ \epsilon^\nu_{\nu_\mu, \text{obs}} \approx \frac{\Gamma \epsilon^\nu_{\nu_\mu}}{4(1+z)} \approx 1.2 \times 10^{18} \text{eV} \left( \frac{1+z}{2} \right)^{-1} \epsilon^\nu_{B, -2} R_{15}, \]
above which the slope of the $\nu_\mu$ spectrum steepens by 2 since $\epsilon^\nu_{\nu_\mu} / \epsilon^\nu_{\nu_\mu, \text{obs}} \propto \epsilon^\nu_{\nu_\mu}^{-2}$.

Muons have a lifetime $\sim 100$ times longer than that of pions, so that the energy cutoff of $\bar{\nu}_\mu$ and $\nu_e$ will therefore be 10 times smaller than $\epsilon^\nu_{\nu, \text{obs}}$ (e.g. Guetta et al. 2004). With equation (24), we have
\[ \epsilon^\nu_{\nu, \nu_\mu, \text{obs}} \approx \frac{\epsilon^\nu_{\nu, \nu_\mu, \text{obs}}}{10} \approx 1.2 \times 10^{17} \text{eV} \left( \frac{1+z}{2} \right)^{-1} \epsilon^\nu_{B, -2} R_{15}. \]

Equation (18) is valid only for protons at the break energy $\epsilon^b_{p, \text{obs}} (1)$. Considering emission from all the protons, the resulting neutrino spectrum would then be
\[ f_\nu \approx \left\{ \begin{align*}
&f_\nu (1) \frac{\epsilon^0_{\nu, \nu_\mu, \text{obs}}}{\epsilon^0_{\nu, \nu_\mu, \text{obs}} (1)} \left[ \frac{1}{1+z} \right]^{1.1} - 0.5, \quad \epsilon_{\nu, \text{obs}} < \epsilon^b_{p, \text{obs}} (1), \\
&f_\nu (1) \frac{\epsilon^0_{\nu, \nu_\mu, \text{obs}}}{\epsilon^0_{\nu, \nu_\mu, \text{obs}} (1)} - 1, \quad \epsilon^b_{p, \text{obs}} (1) < \epsilon_{\nu, \text{obs}} < \epsilon^0_{\nu, \nu_\mu, \text{obs}}, \\
&1, \quad \epsilon^0_{\nu, \nu_\mu, \text{obs}} \leq \epsilon_{\nu, \text{obs}} < \epsilon^M_{\nu, \text{obs}},
\end{align*}\right. \]
where $\epsilon^0_{\nu, \nu_\mu, \text{obs}}$ is defined as $(\epsilon^0_{\nu, \nu_\mu, \text{obs}} / \epsilon^0_{\nu, \nu_\mu, \text{obs}} (1))^{0.5} = 1 / f_\nu (1)$; $\epsilon^M_{\nu, \text{obs}} = \epsilon^M_p (1+z)$. The neutrino spectrum finally reads
\[ \frac{\epsilon^0_{\nu, \nu_\mu, \text{obs}} d^2 N_{\nu}}{d\epsilon_{\nu, \nu_\mu, \text{obs}} d\Omega} \sim \frac{1}{8} \frac{(1+z) L}{4 \pi D_L^2} \left\{ \begin{align*}
&f_\nu (1) \frac{\epsilon^0_{\nu, \nu_\mu, \text{obs}}}{\epsilon^0_{\nu, \nu_\mu, \text{obs}} (1)} \left[ \frac{1}{1+z} \right]^{1.1} - 0.5, \quad \epsilon_{\nu, \text{obs}} < \epsilon^b_{p, \text{obs}} (1), \\
&f_\nu (1) \frac{\epsilon^0_{\nu, \nu_\mu, \text{obs}}}{\epsilon^0_{\nu, \nu_\mu, \text{obs}} (1)} - 1, \quad \epsilon^b_{p, \text{obs}} (1) < \epsilon_{\nu, \text{obs}} < \epsilon^0_{\nu, \nu_\mu, \text{obs}}, \\
&1, \quad \epsilon^0_{\nu, \nu_\mu, \text{obs}} \leq \epsilon_{\nu, \text{obs}} < \epsilon^M_{\nu, \text{obs}},
\end{align*}\right. \]
where the factor $1/8$ takes into account the fact that charged and neutral pions are produced with roughly equal probabilities, and that each neutrino carries $1/4$ of the pion energy, i.e. $\epsilon^\nu_{\nu_\mu, \text{obs}} \approx 0.05 \epsilon^p_{\text{obs}}$. In equation (20), $\epsilon^c_{\nu, \nu_\mu, \text{obs}} < \epsilon^M_{\nu, \nu_\mu, \text{obs}}$ is assumed. In case of $\epsilon^c_{\nu, \nu_\mu, \text{obs}} > \epsilon^M_{\nu, \nu_\mu, \text{obs}}$, the neutrino flux is 0 for $\epsilon_{\nu, \text{obs}} > \epsilon^M_{\nu, \nu_\mu, \text{obs}}$.

As in many publications, we assume that the protons in the FS distribute as $\frac{dn}{dt_p} \propto \gamma_p^{-2}$ for $\gamma_p < \gamma_p^m$ and $\gamma_p^m < \gamma_p < \gamma_p^M$, where $\gamma_p = (1 - \epsilon)(1 - \epsilon) / (\ln(\gamma_p, M / \gamma_p^m))$ and $\gamma_p^m = \epsilon^p / M \Gamma / m_c^2$. For a neutrino detector with a surface area of $S_{\text{det}} \sim \text{km}^2$, the detectable events for one burst can be estimated by (taking $R = 0.5 D_c / \text{wind}$)
\[ N_{\text{events}} (1) \approx \frac{1}{8} \frac{(1 - \epsilon)}{\epsilon^p_{\text{obs}} (1)} \left[ \frac{1}{1+z} \right]^{1.1} S_{\text{det}} P_{\nu - \mu} \frac{\ln(\epsilon^0_{\nu, \nu_\mu, \text{obs}} / \epsilon^0_{\nu, \nu_\mu, \text{obs}} (1)) f_{\nu} (1) + \ln(\epsilon^c_{\nu, \nu_\mu, \text{obs}} / \epsilon^0_{\nu, \nu_\mu, \text{obs}})}{\ln(\gamma_p, M / \gamma_p^m)} \sim 0.002 E_{52.6, D_L^2}^{-2} L_{28.34}^{-2}. \]

where $P_{\nu - \mu} \sim 6 \times 10^{-4} (\epsilon_{\nu, \text{obs}} / 3 \times 10^{15} \text{eV})^{0.5}$ is the probability that a neutrino produces a detectable high energy muon for $\epsilon_{\nu, \text{obs}} > 10^3 \text{TeV}$ (Gaiser et al. 1995). The factor $1/8$ takes into account the fact that there are three neutrinos produced with each carrying about 1/8 of the energy loss of the proton, and that due to neutrino flavor oscillation, only 1/3 of the neutrinos are the detectable muon neutrinos. For a nearby GRB with $z < 0.1 \left( D_L \sim 1.4 \times 10^{27} \text{cm} \right)$ and $E = 4 \times 10^{53} \text{ergs}$, the resulting detectable events is $\sim 0.4$. For such a nearby energetic GRB, the prompt $\gamma$-ray energy fluence would be $F_{\gamma} = (1+z)E_{\gamma} / (4\pi D_L^2) \sim 10^{-3} \text{ergs cm}^{-2} E_{\gamma, 53} D_L^{-2}$. This means that only for those very bright GRBs can the predicted neutrinos be detected. Similar conclusions have also been 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 none 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; Vieth et al. 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 ultrahigh energy cosmic rays are mostly produced by GRBs, which might not be the case.

4.2. Neutrinos from photomeson interaction with the MeV photon flow: $\sim 10^{14} \text{eV}$

Because of the significant overlapping of the MeV photon flow and the shocked regions, high energy protons with energy larger than $2 \times 10^{15} \text{eV}$ 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, $\gamma$-ray photons with the observed typical energy $\epsilon^b_{\gamma, \text{obs}} \approx 300 \text{keV}$ have an energy $\epsilon^b_{\gamma, \text{obs}} \sim (1+z) \epsilon^b_{\gamma, \text{obs}} / \Gamma$. The protons that interact with the photons with energy $\epsilon^b_{\gamma}$ at the $\Delta$-resonance have an observed energy of
\[ \epsilon^b_{\nu, \text{obs}} (2) \sim 0.3 \text{GeV}^2 T^2 / (1+z)^2 \epsilon^b_{\gamma, \text{obs}} \]
\[ = 2 \times 10^{15} \text{eV} L_{52.6}^{1/2} A_{-1/2}^{-1/2} (1+z)^{-2} (\epsilon^b_{\gamma, \text{obs}} / 300 \text{keV})^{-1}, \]

This energy is well below the observed maximum proton energy $\epsilon^M_p$ (equations 19-21). The generated neutrinos have a typical energy
\[ \epsilon^b_{\nu, \text{obs}} (2) \sim 10^{14} \text{eV} L_{52.6}^{1/2} A_{-3/2}^{-1/2} (1+z)^{-2} (\epsilon^b_{\gamma, \text{obs}} / 300 \text{keV})^{-1}. \]
Similar to Eq. 11, at \( \epsilon_{\nu,\text{obs}}^b \), the fraction of the energy loss of protons through photoemission interaction with the MeV photons can be parameterized as

\[
f_\pi(2) \approx 0.18L_{52}^{-1/2}L_{52}^{-1/2} A_{1/4}^{-1} \left[ \frac{1 + z}{600 \text{keV}} \epsilon_{\gamma,\text{obs}}^b \right]^{-1} R_{15}^{-1} \tag{30}
\]

The MeV flux photons has a broken power law spectrum, i.e., \( n(\epsilon_{\nu}) \propto \epsilon_{\gamma,\text{obs}}^{-1} \) below \( \epsilon_{\gamma,\text{obs}} \) and \( n(\epsilon_{\nu}) \propto \epsilon_{\gamma,\text{obs}}^{-2} \) above \( \epsilon_{\gamma,\text{obs}} \). So the general form of the fraction of the energy loss of protons through photoemission interaction with the MeV photons is

\[
f_\pi \approx f_\pi(2) \left\{ \begin{array}{ll}
1, & \epsilon_{\text{p,obs}} > \epsilon_{\nu,\text{obs}}^b(2), \\
\epsilon_{\text{p,obs}} \leq \epsilon_{\nu,\text{obs}}^b(2), & \epsilon_{\text{p,obs}} \leq \epsilon_{\nu,\text{obs}}^b(2).
\end{array} \right. \tag{31}
\]

As a result, the differential neutrino spectrum is also a broken power law, i.e. \( n(\epsilon_{\nu}) \propto \epsilon_{\gamma,\text{obs}}^{-1} \) below \( \epsilon_{\gamma,\text{obs}} \) and \( n(\epsilon_{\nu}) \propto \epsilon_{\gamma,\text{obs}}^{-2} \) above \( \epsilon_{\gamma,\text{obs}} \) (see also Waxman & Bahcall 1997). However, for \( \epsilon_{\nu,\text{obs}} > \epsilon_{\nu,\text{obs}}^b \), 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 detail). So the neutrino spectrum cannot extend to \( \epsilon_{\nu,\text{obs}} > \epsilon_{\nu,\text{obs}}^b \).

Similar to equation 24, we now estimate the detectability of such 10^14eV neutrinos in the early afterglow phase. For one typical GRB, taking \( R = 0.5R_\odot \) (wind), one has

\[
N_{\text{events}}(2) \approx \frac{1}{8} \left[ \frac{1 - \epsilon_{\nu,\text{obs}}}{4\pi D_L^2 S_{\text{det}} P_{\nu-\mu}} \right] \ln \left( \frac{\epsilon_{\nu,\text{obs}}}{\epsilon_{\nu,\text{obs}}^b(2)} \right) \ln \left( \frac{D_{28,34}}{E_{53.6D_{28,34}}^2} \right) \tag{32}
\]

where \( P_{\nu-\mu} \approx 1.7 \times 10^{-6} (\epsilon_{\nu,\text{obs}}/1\text{TeV})^{0.8} \) (e.g., Halzen & Hooper 2002) for 1TeV < \( \epsilon_{\text{e,obs}} < 10^{15}\text{eV} \). The resulting detectability is comparable with that of 5 \times 10^{15}\text{eV} (see equation 24). These neutrinos can be only marginally detectable above the atmospheric neutrino background. For a typical GRB with \( T_{90} \approx 20s, E \approx 4 \times 10^{53}\text{ergs} \) and \( z = 1 \), the 10^{14}\text{eV neutrino background event number is} \( \sim 9 \times 10^{-5} \) in the time span. It is \( < 0.001 \) (Eq. 24), but not significantly. For a detected \( \sim 10^{14}\text{eV neutrino} \) 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 be also enhanced significantly when fluctuations in distance, energy are considered (e.g., Haltzen & 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 km^3 high energy neutrino detectors such as ICECUBE, ANTARES, and NESTOR.

One caveat is that the \( \sim 10^{14}\text{eV} \) neutrinos would also overlap with the \( \sim 5 \times 10^{14}\text{eV} \) neutrinos generated from the internal shocks (e.g., Waxman & Bahcall 1997), so that it might not be easy to distinguish both species. We notice, 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 progenitor of long, soft GRBs. For an ultra-relativistic 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 modification of the early RS and FS emission significantly, since the dominant cooling process for electrons is likely the 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 very early R-band lightcurve).

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 \( \epsilon^\pm \) pairs. These \( \epsilon^\pm \) pairs re-scatter the soft X-rays in the FS region and the prompt \( \gamma \)-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 the GLAST. In the wind model, the total thermal energy contained in the FS region is larger 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 RS is suppressed). The predicted fluence of these sub-GeV flashes is high enough to be detected by GLAST.

Equally interestingly, 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 \( \sim 5 \times 10^{15}\text{eV} \) component powered by the photomeson interaction in the FS region, and a \( \sim 10^{14}\text{eV} \) component powered by photomeson interaction with the MeV photon flow. The detectability of the two components is comparable. The spectrum of 10^{14}eV neutrinos 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 \( \sim (T/T_0)^{2} \). The spectrum of the 5 \times 10^{15} eV neutrinos takes a more complicated form (see equation 24). A caveat is that these neutrinos overlap with the \( \sim 5 \times 10^{14}\text{eV} \) internal shock neutrino-component (Waxman & Bahcall 1997), and hence, is 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 \( \epsilon^\pm \) pair loading (The pairs are created by the interaction of the prompt \( \gamma \)-rays with the back-scattered \( \gamma \)-rays by the wind medium), the medium will be accelerated and the ejecta moves into a cavity until it reaches a radius \( R_{\text{gap}} \sim 3.3 \times 10^{15}E_{53,32}^{1/2} \) (e.g., Beloborodov 2002). As a
result, the shock crossing radius (see equation (2)) should be $\sim R_{\text{gap}} + R_{\times}\text{(wind)} \sim 2R_{\times}\text{(wind)}$. The overlapping of the MeV photon flow and the shocked regions is still significant since the ejecta essentially does not decelerate at $R < R_{\text{gap}}$. However, $f_\pi(1)$ and $f_\pi(2)$ involved in the above calculation should be smaller by a factor of about 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 re-scatter the initial MeV $\gamma-$rays to the sub-GeV energy range. Notice 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 the pairs is $k$ times of the electrons associated with protons, $\gamma_{\text{m}}^{i}$ (see eq.[3]) should be reduced to $\gamma_{\text{m}}^{i}/(1 + 2k)$. 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 sub-GeV energy range. Therefore sub-GeV flashes are still expected when the radiation front effect is taken into account.

In this work, only the wind model is discussed. In the ISM model, if the GRB is long enough, the overlapping of the MeV photon flow and the shocked regions could be also tight. In this case, $R_{\times}\text{(ISM)} \sim 10^{17}$cm, and the optical depth of 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 for 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 the one 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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