PROMPT TeV NEUTRINOS FROM THE DISSIPATIVE PHOTOSPHERES OF GAMMA-RAY BURSTS

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ABSTRACT

Recently it was suggested that a photospheric component that results from internal dissipation occurring in the optically thick inner parts of relativistic outflows may be present in the prompt $\gamma$/X-ray emission of gamma-ray bursts or X-ray flashes. We explore high-energy neutrino emission in this dissipative photosphere model, assuming that the composition of the outflow is baryon dominated. We find that neutrino emission from a proton–proton collision process forms an interesting signature in the neutrino spectra. Under favorable conditions for the shock dissipation site, these low-energy neutrinos could be detected by km$^3$ detectors, such as Icecube. Higher-energy ($\gtrsim$ 10 TeV) neutrino emission from proton–proton collision and photopion production processes could be significantly suppressed for dissipation at relatively small radii due to efficient Bethe–Heitler cooling of protons and/or radiative cooling of the secondary mesons in the photosphere radiation. As the dissipation shocks continue further out, high-energy neutrinos from the photoproduction process become dominant.

Key words: gamma rays: bursts -- neutrinos

1. INTRODUCTION

Although it has generally been accepted that the prompt gamma-ray emission of gamma-ray bursts (GRBs) results from internal dissipation, likely internal shocks, of a relativistic outflow (e.g., Paczyński & Xu 1994; Rees & Mészáros 1994), the dissipation site and the radiation mechanism for the gamma-ray emission are still largely unknown. Synchrotron and/or inverse-Compton scattering emission by shock-accelerated electrons in the optically thin region has been proposed as an efficient mechanism for the gamma-ray emission. However, this model does not satisfactorily account for a few observational facts, such as the low-energy spectral slopes that are steeper than synchrotron lower energy spectral indices (Preece et al. 2000; Lloyd et al. 2000), the clustering of peak energies, or the correlation between the burst’s peak energy and luminosity (Amati et al. 2002). It becomes recognized that an additional thermal component may play a key role and could solve these problems (e.g., Pe‘er et al. 2006; Ryde et al. 2006). It has also been pointed out that a hybrid model with both a thermal and a non-thermal component can describe the spectrum equally well as the Band function model (Band et al. 1993), but the former has a more physical meaning (Ryde 2005). Recently it was suggested that a strong quasi-thermal component could result from internal dissipation occurring in the optically thick inner parts of relativistic outflows (Rees & Mészáros 2005; Pe‘er et al. 2006; Thompson et al. 2007). Sub-synchroscopic shock dissipation can increase the radiative efficiency of the outflow, significantly boosting the original thermal photospheric component so that it may well dominate the nonthermal component from optically thin shocks occurring outside the photosphere.

Neutrino emission from gamma-ray bursts has been predicted at different stages of the relativistic outflow, such as the precursor phase (e.g., Bahcall & Mészáros 2000; Mészáros & Waxman 2001; Razzaque et al. 2003a, 2003b; Razzaque et al. 2004; Ando & Beacom 2005; Horiuchi & Ando 2008; Koers & Wijers 2008), the prompt emission phase (e.g., Waxman & Bahcall 1997; Dermer & Atoyan 2003; Guetta et al. 2004; Murase & Nagataki 2006; Gupta & Zhang 2007; Murase et al. 2006) and the afterglow phase (e.g., Waxman & Bahcall 2001; Dai & Lu 2001; Dermer 2002; Li et al. 2002; Murase & Nagataki 2006; Murase 2007; Dermer 2007). Based on the broken power-law approximation for the spectrum of the prompt emission, presumably from optically thin internal shocks, a burst of PeV neutrinos, produced by photomeson production, was predicted to accompany the prompt gamma-ray emission if protons are present and also accelerated in the shocks (Waxman & Bahcall 1997). The neutrino emission from proton–proton ($pp$) collisions was generally thought to be negligible due to the lower collision opacity for optically thin internal shocks. However, as we show below, if some part of the prompt emission arises from internal shocks occurring in the optically thick inner part of the outflow, as indicated by the thermal emission, a lower energy ($\lesssim$ 10 TeV) neutrino component may appear as a result of $pp$ collisions.

2. THE DISSIPATIVE PHOTOSPHERE MODEL

Photosphere models have been widely discussed in relation to the prompt emission of GRBs (e.g., Thompson 1994; Ghisellini & Celotti 1999; Rees & Mészáros 2005; Thompson et al. 2007; Ioka et al. 2007). The potential advantage of photosphere models is that the peak energy can be stabilized, which is identified as the thermal or Comptonization thermal peak (see Ioka et al. 2007 and references therein). The photosphere radiation may also produce a large number of electron–positron pairs, which may lead to a pair photosphere beyond the baryon-related photosphere (e.g., Rees & Mészáros 2005), and may also enhance the radiative efficiency (Ioka et al. 2007). On the other hand, it is also suggested that the number of pairs produced does not exceed the baryon-related electrons by a factor larger than a few (Pe‘er et al. 2006). For simplicity, we here only consider the dissipation below the baryon-related photosphere, which is more favorable for $pp$ neutrino production.

Following Rees & Mészáros (2005) and Pe‘er et al. (2006), we assume that during the early stage of prompt emission, internal shocks of the outflow occur at radii below the baryonic photosphere. Initially, the internal energy is released at the base of the outflow, $r_0 \sim a r_g = 2G M / c^2$, where $a \gtrsim 1$ and $r_g$ is the Schwarzschild radius of a central object of mass $M$. The internal energy is then converted to the kinetic energy of the flow, whose bulk Lorentz factor grows as $\gamma \sim r$ up to a saturation radius at
outflow rates. Above the saturation radius, the observer-frame Thomson scattering by the baryon-related electrons is \( \sim \) for protons with energies above the like spectrum, the number density of photons in the sub-photosphere emission at the dominant energy is converted into magnetic fields, we have a magnetic field \( \mathbf{B}' = 2.5 \times 10^4 \mathbf{B}_\odot \), for which the optical depth to Thomson scattering by the baryon-related electrons is \( \tau_T = \sigma_T n_e (\mathbf{B}/\mathbf{B}_\odot)^2 \), where \( \sigma_T \) is the Thomson scattering cross section and \( n_e \) is the electron density. The photosphere is further out, at radius \( R_{ph} = 1.2 \times 10^{13} \mathbf{L}_{k,2} \mathbf{G}^{-2} \) cm. A detailed calculation taking into account electron/positron cooling and the Comptonization effect leads to a quasithermal emission which peaks at energy \( \sim 300-500 \) keV for dissipation at a Thomson optical depth of \( \tau_T \sim 10-100 \) (Pe`r et al. 2006). This temperature is consistent with the observed peak energies of prompt gamma-ray emission of a majority of GRBs.

Assuming that a fraction of \( \mathcal{E}_B \simeq 0.1 \) of the shock internal energy is converted into magnetic fields, we have a magnetic field \( \mathbf{B}' = 2.5 \times 10^4 \mathbf{B}_\odot \mathbf{L}_{k,2} \mathbf{G}^{-1} \mathbf{R}_{1,2} \mathbf{G} \). Protons accelerated by internal shocks are assumed to have a spectrum \( dn/d\epsilon_p \sim \epsilon_p^{-\gamma} \) with \( \gamma \simeq 2 \), as often assumed for non-relativistic or mildly relativistic shock acceleration. The maximum proton energy is set by comparing the acceleration timescale \( t_{\text{acc}} \approx (\mathcal{E}_p/m_p c^2) \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \) with the energy-loss timescales. The synchrotron loss time is \( t_{\text{syn}} \approx (\mathcal{E}_p/m_p c^2) \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_p^{1/3} c^2 \right)^{-1} \) s. Assuming that the sub-photosphere emission at the dissipation site peaks at \( \epsilon_{\gamma} = 300 \) keV with a thermal-like spectrum, the number density of photons in the comoving frame is \( n_\gamma' = \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right)^{-1} \) s. The proton energy is approximately \( \epsilon_p' = \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right)^{-1} \), where \( \mathcal{L}_{\gamma,51} \simeq 0.2 \) is the inelasticity and \( \mathcal{E}_p \simeq 5 \times 10^{-28} \) cm\(^2\) is the peak cross section at the \( \Delta \) resonance. By comparison with the synchrotron loss time, it is found that the most effective cooling mechanism for protons is the \( \mathcal{P} \gamma \) process for protons with energies above the \( \gamma \) threshold, but below \( \sim 10^8 \) GeV. Equating \( t_{\text{acc}} = t_{\text{syn}}' \), we obtain the maximum proton energy in the shock comoving frame

\[
\epsilon_{p,\text{max}}' = 10^7 a_{1,6}^{-1/2} L_{k,52}^{-1/2} L_{\gamma,51}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right) \text{GeV}. \tag{2}
\]

3. PROTON AND MESON COOLING

The shock-accelerated protons produce mesons via \( \mathcal{P} \gamma \) and \( \mathcal{P} \gamma \) interactions. Since the meson multiplicity in \( \mathcal{P} \gamma \) interactions is about 1 for pions and 0.1 for kaons, neutrinos contributed by pion decay are dominant when the cooling effect of pions is not important, which is applicable to the low-energy \( \mathcal{P} \gamma \) neutrinos. Therefore we here consider only pion production in \( \mathcal{P} \gamma \) interactions. Pion production by \( \mathcal{P} \gamma \) interaction in the subphotosphere dissipation is efficient since the cooling time in the shock comoving frame,

\[
t_{\mathcal{P} \gamma} = 1/(\mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right)) \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right), \tag{3}
\]

can be shorter than the shock dynamic time, \( t_{\text{dyn}} = R/\Gamma c = 0.03 \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{-1} \right) \), where \( \mathcal{L}_{\gamma,51} \simeq 4 \times 10^{-26} \) cm\(^2\) is the cross section for \( \mathcal{P} \gamma \) interactions, \( n_\gamma = \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{-1} \right) \) is the proton number density, and \( \mathcal{L}_{\gamma,51} \simeq 0.5 \) is the inelasticity. Protons also cool through Bethe–Heitler interactions (\( \gamma \gamma \rightarrow p \mathcal{E}^2 \)) and \( \gamma \gamma \) interactions when the target photon energy seen by the protons is above the threshold energy for each interaction. Denoting by \( n(\epsilon_{\gamma})\mathcal{E}_{\gamma} \) the number density of photons in the energy range \( \epsilon_{\gamma} \) to \( \epsilon_{\gamma} + d\epsilon_{\gamma} \), the cooling time in the shock comoving frame for \( \mathcal{P} \gamma \) and Bethe–Heitler cooling processes is given by

\[
t_{\gamma,\gamma,\text{BH}} = \frac{c}{2\mathcal{E}_{\gamma}^2} \int_{\epsilon_{\gamma}}^{\infty} \mathcal{E}_{\gamma} d\epsilon_{\gamma} \int_{\epsilon_{\gamma}}^{\infty} dx x^2 n(x), \tag{4}
\]

where \( \Gamma_{\gamma} = \epsilon_{\gamma}/m_p c^2 \), \( \sigma \) and \( \mathcal{K} \) are, respectively, the cross section and the inelasticity for the \( \mathcal{P} \gamma \) (or Bethe–Heitler) process. As a rough estimate, the Bethe–Heitler cooling time is

\[
t_{\gamma,\gamma,\text{BH}} \approx 0.5 \left( \frac{28/9}{\mathcal{E}_{\gamma}} \right) \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right) \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right) \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right). \tag{5}
\]

when \( K = \epsilon_{\gamma}/m_p c^2 \) is a large value (a good approximation when \( K > 10 \)). Chodorowski et al. (1992), where \( \epsilon_{\gamma}' \) is the thermal peak energy of photons in the comoving frame. So the proton energy is larger than \( \epsilon_{\gamma}' \) over the \( \mathcal{P} \gamma \) cooling dominates over the \( \mathcal{P} \gamma \) cooling. At even higher energies near the threshold for \( \mathcal{P} \gamma \) interactions at \( \epsilon_{\gamma}' \) is a large value (a good approximation when \( K > 10 \)). Chodorowski et al. (1992), where \( \epsilon_{\gamma}' \) is the thermal peak energy of photons in the comoving frame. So the proton energy is larger than \( \epsilon_{\gamma}' \) over the \( \mathcal{P} \gamma \) process. As a rough estimate, the Bethe–Heitler cooling time is

\[
t_{\gamma,\gamma,\text{BH}} = 1/\left( \mathcal{L}_{\gamma,51}^{-1} \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right) \right) \mathbf{G} \mathbf{R}_{1,2}^{-1} \left( \epsilon_{\gamma}^{1/3} c^2 \right). \tag{6}
\]

We compare these three cooling timescales for protons in Figure 1 for representative parameters, using more accurate cross sections for the \( \gamma \gamma \) and Bethe–Heitler processes in Equation (4). For the photopion production cross section, we take the Lorentzian form for the resonance peak (Mücke et al. 2000) plus a component contributed by multi-pion production at higher energies, while for the Bethe–Heitler process we use the cross section given by Chodorowski et al. (1992). The number density of photons used in the calculation has been assumed to have a blackbody distribution. The numerical result confirms that \( \mathcal{P} \gamma \) cooling and \( \mathcal{P} \gamma \) cooling are dominant, respectively, at lowest and highest energies, while at intermediate energies, Bethe–Heitler cooling is dominant.
adiabatic expansion (which is equal to the dynamic timescale \( t \) and \( p \gamma \) line, and the dashed line are for proton–proton collision, Bethe–Heitler cooling, internal shock as functions of proton energy. The straight solid line, the dotted with the lifetime of pions\( \tau_{\pi} \) is the energy lost by the meson per collision. 1 The mean pion energy is roughly equal probability and muon neutrinos carry roughly one-fourth of the pion energy in pion decay. 1 The mean pion energy is about 20% of the energy of the proton producing the pion, so the mean energy of neutrinos is \( \epsilon_{\nu} \approx 0.05 \epsilon_{p} \). Assuming that protons are efficiently accelerated in shocks with an energy density of \( U_{p} = 100 U'_{p} \), the number of TeV neutrinos from one GRB is about \( N_{\nu} \approx 0.1(\Phi_{\nu}/10^{-4} \text{erg cm}^{-2} \text{sr}^{-1}) \) for \( \eta_{p} \approx 1 \), according to Equation (8). So only from very strong bursts with gamma-ray fluence \( \Phi_{\nu} \approx 10^{-3} \text{erg cm}^{-2} \), which are very rare events, can neutrinos from a single GRB be detected.

The aggregated muon neutrino flux from all GRBs is approximately given by

\[
\epsilon_{\nu}^{2} J_{\nu(pp,py)}(\epsilon_{\nu}) \approx \frac{c}{4 \pi H_{0}} \epsilon_{\nu}^{2} J_{\nu(pp,py)}(\epsilon_{\nu}) R_{GRB(0)} f_{z} = 1.5 \times 10^{-9} \epsilon_{p,53} \left( \frac{R_{GRB(0)}}{1 \text{Gpc}^{-3} \text{yr}^{-1}} \right) \left( \frac{f_{z}}{3} \right) \times \eta_{p} \delta_{(pp,py)} \zeta_{\pi} \text{GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1},
\]

where \( f_{z} \) is the correction factor for the contribution from high redshift sources and \( R_{GRB(0)} \) is the overall GRB rate at redshift \( z = 0 \). Assuming that the GRB rate traces the star-formation rate in the universe, the calculation gives \( f_{z} \approx 3 \) (Waxman & Bahcall 1999). It is not clear how efficiently the protons are accelerated in GRB shocks. Assuming an optimistic case that protons are efficiently accelerated in shocks and that half of the maximum and minimum energies of acceleration protons.

The absence of pion cooling loss, the neutrinos produced by pion decay carry one-eighth of the energy lost by protons to pion production, since charged and neutral pions are produced with roughly equal probability and muon neutrinos carry roughly one-fourth of the pion energy in pion decay. 1 The mean pion energy is about 20% of the energy of the proton producing the pion, so the mean energy of neutrinos is \( \epsilon_{\nu} \approx 0.05 \epsilon_{p} \). Assuming that protons are efficiently accelerated in shocks with an energy density of \( U_{p} = 100 U'_{p} \), the number of TeV neutrinos from one GRB is about \( N_{\nu} \approx 0.1(\Phi_{\nu}/10^{-4} \text{erg cm}^{-2} \text{sr}^{-1}) \) for \( \eta_{p} \approx 1 \), according to Equation (8). So only from very strong bursts with gamma-ray fluence \( \Phi_{\nu} \approx 10^{-3} \text{erg cm}^{-2} \), which are very rare events, can neutrinos from a single GRB be detected.

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\]

1 As an approximate estimate, we have neglected the effect of multi-pion production, the muon contribution decay to the neutrino flux, and the neutrino oscillation effect, which may affect the estimate of the \( pp \) neutrino flux within a factor of 2 (the factor, however, could be larger for \( py \) neutrino flux).
the kinetic energy dissipation occurs below the photosphere.\footnote{This is based on the analysis by Ryde et al. (2006) and also in a very recent paper by Ryde & Pe'er (2008), who find that the thermal photons carry a fraction of 30\% to more than 50\% of the prompt emission energy.} we take a mean value $E_p = 1.5 \times 10^{53}$ erg for the isotropic equivalent energy in accelerated protons in one GRB during the dissipative photosphere phase, based on a typically used value $L_k = 10^{52}$ erg s$^{-1}$ for the isotropic kinetic energy luminosity and a typical long GRB duration of $\Delta T = 30$ s. The GRB rate\footnote{There is large uncertainty in the estimate of the local GRB rate. Some suggest a lower GRB rate based on the analysis of Swift bursts with $R_{\text{GRB}}(0) = 0.05 \div 0.27$ Gpc$^{-2}$ yr$^{-1}$ (e.g., Guetta & Piran 2007; Le & Dermer 2007), while others get a higher rate comparable to earlier estimate before Swift (e.g., Liang et al. 2007).} at redshift $z = 0$ is taken to be $R_{\text{GRB}}(0) = 1$ Gpc$^{-2}$ yr$^{-1}$ (Guetta et al. 2005; Liang et al. 2007). The isotropic luminosity is taken to be $L_\nu = 10^{53}$ erg\footnote{Some observations have indicated rather high radiative efficiency, and the importance of pp neutrinos may be reduced if $L_\nu$ is smaller than 10$L_{\nu,\gamma}$.}.\footnote{The energy-dependent neutrino fluxes contributed by pp and $p\gamma$ interactions are plotted in Figure 2 for a set of representative parameters of the dissipative photosphere model and three different dissipation radii. If the kinetic energy is dissipated at a radius of $R = 10^{11}$ cm, the calculation (the red solid curves) shows that at energies below tens of TeV, the neutrino flux is dominated by a pp component. Taking the detection probability of $P_{\nu\mu} = 10^{-4}(\epsilon_\nu/1$ TeV) for TeV neutrinos (Gaisser et al. 1995), the expected flux of upward moving muons contributed by this pp component is about 8 to 10 events each year for a km$^3$ neutrino detector, such as Icecube. We can also estimate the atmospheric neutrino background expectation in coincidence with these GRB sources, noting that the search for neutrinos accompanying GRBs requires that the neutrinos be coincident in both direction and time with gamma rays. Taking an average GRB duration of $\approx 30$ s, an angular resolution of Icecube of $\approx 1^\circ$ and an atmospheric neutrino background flux of $\approx 10^{-4}$ GeV cm$^{-2}$ s$^{-1}$ sr$^{-1}$ at 1 TeV (Ahrens et al. 2004), the atmospheric neutrino background expectation is $\approx 5 \times 10^{-3}$ events from 500 GRBs (in one year). Such a low background in coincidence with GRBs allows the claim of detection of TeV neutrinos from GRB sources. At energies from a few TeV to tens of TeV, the Bethe–Heitler cooling suppresses the $pp$ cooling, resulting in a steepening at several TeV in the neutrino spectrum. At energies above the threshold for $p\gamma$ interactions, the neutrino from $p\gamma$ process is heavily suppressed due to the strong radiative cooling of secondary pions. For a larger dissipation radius at $R = 10^{12}$ cm (the blue dashed curves), the neutrino emission flux from $p\gamma$ process is no longer suppressed and in this case both pp and $p\gamma$ neutrino components can be detected by km$^3$ detectors. We also calculated the neutrino flux, shown by the green dotted lines in Figure 2, for shock at the photosphere radius $R_{\text{ph}}$, assuming that the radiation spectrum is a broken power-law (due to Comptonization) peaking at $\epsilon_\nu = 100$ keV, with lower-energy and higher-energy photon indexes given by $-1$ and $-2$, respectively. In this case, the $pp$ neutrino flux becomes small (may be marginally detectable), while the $p\gamma$ neutrino flux spectrum is similar to the analytic result obtained by Waxman & Bahcall (1997), as expected for a broken power-law photon spectrum. Note that in one burst the shock dissipation could be continuous from small to large radii, as indicated by the larger variability timescales seen in GRBs. By comparing the three cases of different dissipation radii in Figure 2, one can see that as the shock radius increases, the neutrino emission from the $pp$ component decreases, while the $p\gamma$ component increases until it reaches the saturation level. The total $pp$ neutrino flux from such continuous dissipation is thus contributed predominantly by the deepest internal shocks below the photosphere. In the whole neutrino spectrum, a “valley” is seen between the $pp$ and $p\gamma$ components of the spectrum, which may be a potential distinguishing feature of the sub-photosphere dissipation effect.}

5. DISCUSSIONS AND CONCLUSIONS

Waxman & Bahcall (1997) as well as later works have studied neutrino emission from the photomeson process during the prompt internal shocks of GRBs, assuming that the radiation in the shock region has a broken power-law nonthermal spectrum. It was found that the neutrino emission peaks at energies above 100 TeV. Toward lower energies, the neutrino emission intensity decreases as $\epsilon_\nu^2\Phi_\nu \sim \epsilon_\nu$ (Waxman & Bahcall 1997). However, if internal shocks, especially at the early stage of the prompt emission, occur below the photosphere, a quasi-thermal spectrum will arise. In this Letter, we have discussed the neutrino emission associated with the dissipative photosphere that produces such prompt thermal emission. We find that $pp$ interaction process becomes important for shock-accelerated protons and provides a new neutrino component, which dominates at energies below tens of TeV. The neutrino emission from photopion processes of protons interacting with sub-photosphere radiation could be significantly suppressed due to radiative cooling of secondary pions, when the dissipation radius is relatively small. Nevertheless, the total contribution by photopion processes will not be suppressed since the shock dissipation could be continuous and occur at large radii as well. Although TeV neutrinos may also be produced during the early precursor stage of a GRB, i.e., before the jet breaking out the progenitor star (e.g., Razzaque et al. 2004; Ando & Beacom 2005), we want to point out that the TeV neutrino component discussed here can be distinguished from these, because, in our case, the neutrino emission is associated in time with the prompt emission.

After this work was completed and later put onto the arXiv Web site (arXiv:0807.0290), we became aware that K. Murase was also working on the sub-photosphere neutrino independently (Murase 2008). X.Y.W. would like to thank P. Mészáros, S. Razzake, K. Murase, E. Waxman, Z. Li, and K. Ioka for useful comments or discussions. This work is supported by the National Natural Science Foundation of China under grants 10221001, 10403002, and 10873009, and the National Basic Research Program of China (973 program) under grants No. 2007CB815404 and 2009CB824800.

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