RADIAL ANGULAR MOMENTUM TRANSFER AND MAGNETIC BARRIER FOR SHORT-TYPE GAMMA-RAY-BURST CENTRAL ENGINE ACTIVITY

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ABSTRACT

Soft extended emission (EE) following initial hard spikes up to 100 s was observed with Swift/BAT for about half of known short-type gamma-ray bursts (SGRBs). This challenges the conventional central engine models of SGRBs, i.e., compact star merger models. In the framework of black-hole–neutron-star merger models, we study the roles of radial angular momentum transfer in the disk and the magnetic barrier around the black hole in the activity of SGRB central engines. We show that radial angular momentum transfer may significantly prolong the lifetime of the accretion process, which may be divided into multiple episodes by the magnetic barrier. Our numerical calculations based on models of neutrino-dominated accretion flows suggest that disk mass is critical for producing the observed EE. In the case of the mass being \( \sim 0.8 M_\odot \), our model can reproduce the observed timescale and luminosity of both the main and the EE episodes in a reasonable parameter set. The predicted luminosity of the EE component is lower than the observed EE within about one order of magnitude and the timescale is shorter than 20 s if the disk mass is \( \sim 0.2 M_\odot \). Swift/BAT-like instruments may be not sensitive enough to detect the EE component in this case. We argue that the EE component could be a probe for the merger process and disk formation for compact star mergers.

Key words: accretion, accretion disks – black hole physics – gamma-ray burst: general

Online-only material: color figures

1. INTRODUCTION

Gamma-ray bursts (GRBs) are sorted into two classes (Kouveliotou et al. 1993), i.e., short-duration \( (T_{90} < 2 \text{ s}) \), SGRBs) and long-duration GRBs \( (T_{90} > 2 \text{ s}) \). Their progenitors are thought to be mergers of two compact stars (Eichler et al. 1989; Paczynski 1991; Narayan et al. 1992; Zhang et al. 2007; Nakar 2007) and collapses of massive stars (Woosley 1993; Paczynski 1998; Piran 2004; Zhang & Mészáros 2004; Woosley & Bloom 2006), respectively. However, the observed burst duration is instrumentally dependent (Donaghy et al. 2006; Qin et al. 2012). Swift observations reveal that the short–long GRB classification scheme does not always match the physical origin classification scheme, i.e., mergers of compact binaries (Type I) versus collapses (Type II) (Zhang 2006; Zhang et al. 2007, 2009; Lü et al. 2010; Xin et al. 2010). With the CGRO/BATSE data, Lazzati et al. (2001) found an excess emission peaking \( \sim 30 \text{ s} \) after the prompt one, which is detectable for \( \sim 100 \text{ s} \) for some SGRBs (see also Connaughton 2002; Norris et al. 2010). About half of the light curves of those GRBs that are recognized as Type I GRBs with Swift/BAT show initial hard spikes following by an extended emission (EE) component of soft gamma rays up to \( \sim 100 \text{ s} \) post the BAT trigger (Barthelmy et al. 2005; Lin et al. 2008; Perley et al. 2009; Zhang et al. 2009). The most prominent case is GRB 060614. Its light curve is composed of some initial hard spikes and a long, soft gamma-ray tail, which leads to its \( T_{90} \) being \( \sim 110 \text{ s} \) (Gehrels et al. 2006). The initial hard spikes could be recognized as an SGRB with CGRO/BATSE-like instruments, since the soft EE is out of the instrument bands (Zhang et al. 2007). No accompanied supernovae were detected for this nearby long GRB (the redshift \( z = 0.1254 \); Della Valle et al. 2006; Fynbo et al. 2006; Gal-Yam et al. 2006), disfavoring the collapse of a massive star as the progenitor of this GRB. On the other hand, it is shown that some intrinsically short GRBs are likely of Type II origin (Zhang et al. 2009; Belczynski et al. 2010; Levesque et al. 2010; Lin et al. 2010; Xin et al. 2010; Virgili et al. 2011). These observations indicate that \( T_{90} \) may not be a good parameter for distinguishing two types of GRBs. A detailed analysis of the instrumental selection effect and energy dependence of \( T_{90} \) with Fermi/GBM data by Qin et al. (2012) further supports this idea. Lü et al. (2010) proposed a new classification parameter, i.e., \( \varepsilon \equiv E_{\text{iso}}/E_{p,\gamma}^{1.7} \), for grouping an observed GRB into the physically motivated Type I/II classification scheme,\(^{10}\) where \( E_{\text{iso}} \) is the isotropic gamma-ray energy and \( E_{p,\gamma} \) is the peak energy of the \( v f_{\nu} \) spectrum in the rest frame. They showed that some SGRBs are sorted into the high-\( \varepsilon \) group as typical Type II GRBs.

The observed EE component challenges not only the short–long GRB classification scheme, but also the conventional central engine models for SGRBs. Popular central engine models of Type I GRBs are related to the accretion onto a central compact object that is formed from the merger of a stellar compact binary, namely, neutrino-dominated accretion

\(^{10}\) Similarly, Goldstein et al. (2010) used the ratio of gamma-ray fluence to \( E_p \) to create a GRB classification.
flows (NDAFs, e.g., Popham et al. 1999; Narayan et al. 2001; Di Matteo et al. 2002; Kohri & Mineshige 2002; Kohri et al. 2005; Gu et al. 2006; Chen & Beloborodov 2007; Liu et al. 2007, 2008, 2010a, 2010b, 2012; Lei et al. 2009; Sun et al. 2012). Merger of a black hole (BH) neutron star (NS) binary is the most favored scenario. Such a system would result in a rotating BH with several solar masses surrounded by a neutrino-cooled disk. The detection of the EE component likely suggests that the central engine did not die out rapidly. Several lines of evidence from Swift/BAT observations also support this idea. It was also proposed that the early shallow decay X-ray emission and internal X-ray plateau may be due to the spin-down energy release of the proto-magnetar of a stellar compact binary merger (Dai & Lu 1998; Zhang & Mészáros 2002; Lyons et al. 2010). Late X-ray flares may signal the restart of the GRB central engine and evolution of the disk (Burrows et al. 2006; King et al. 2005; Fan & Wei 2005; Dai et al. 2006; Perna et al. 2006; Proga & Zhang 2006; Lazzati et al. 2008; Lee et al. 2009). Lazzati et al. (2008) investigated the temporal evolution of the disk in GRB central engines to explain the observed decline of X-ray flare luminosity. They argued that it is the dynamics of the disk or the jet-launching mechanism that generates an intrinsically unsteady outflow on timescales much longer than the dynamical timescale of the system for late X-ray flares. It was also suggested that propagation instabilities, rather than variability in engine luminosity, are responsible for some X-ray flares (Lazzati et al. 2011).

Different from the early shallow decay X-ray emission and internal X-ray plateau, the EE component is usually highly variable and is usually not clearly separated from the burst itself in its light curves. It may be produced by the same process as the prompt emission in different episodes. Metzger et al. (2008) presented time-dependent models of the remnant accretion disks created during compact object mergers. They calculated the dynamics near the outer edge of the disk to study the evolution of the accretion rate on a long timescale (100 s or longer). They showed that late-time accretion can in principle provide sufficient energy to power the late-time activity observed by Swift/BAT for some SGRBs. In their models, the majority of the disk mass is in the outer edge and the disk becomes advective at late time. In this paper, we focus on the radial angular momentum transfer in the disk and the magnetic barrier in vicinity of the BH, where $n$ is for the $n$th emission episode.

Assuming that the progenitor of SGRBs is a BH–NS binary, the merger of this system would result in a rotating BH surrounded by a neutrino-cooled disk. We focus on the roles of radial angular momentum transfer in the disk and the magnetic barrier in prolonging the lifetime of the SGRB central engine (e.g., Proga & Zhang 2006). A fraction of the disk matter may carry part of the angular momentum of the accretion matter to form a radial outflow. The competition of the radial angular momentum transfer to the gravity of the central BH would increase the timescale of the accretion process. On the other hand, the magnetic field accumulated near the horizon of the BH may be strong enough to prevent the gravity and accretion process. The magnetic barrier would dissipate quickly as the accretion rate drops. These processes likely lead to multiple well-connected accretion episodes (e.g., Narayan et al. 2003; Cao 2011). We illustrate the processes of our model and corresponding schematic light curve in Figure 1. Our model is elaborated below.

2. MODEL

2.1. Outward Angular Momentum Transfer

Without considering the mass (energy) and angular momentum lost in jet production, the conservation of the mass (energy) and angular momentum read (e.g., Bardeen 1970; Thorne 1974; Wang et al. 2002):

$$M_{n+1} - M_n = (M^n_n - M^n_{n+1})e_{in,n} = M_n T_n e_{in,n}.$$  \hspace{1cm} (1)

$$J_{n+1} = J_n - J^n_{n+1} = M_n T_n l_{in,n},$$  \hspace{1cm} (2)

where $M_n$ ($J_n$) and $M^n_n$ ($J^n_n$) are masses (angular momentums) of the BH and the disk, respectively, $M_n$ is the mass accretion rate,
$T_n$ is the accretion timescale, and $\varepsilon_{m,n}$ and $l_{m,n}$ are the specific energy and the angular momentum at the inner boundary orbit in the $n$th episode ($n = 1, 2, 3, \ldots$). The mass of the disk can be written as

$$M_n^* = 2\pi \int_{r_{m,n}}^{r_{o,n}} \Sigma_n r dr,$$

where $r_{i,n}$ and $r_{o,n}$ are the inner and outer boundaries of the disk in the $n$th episode, respectively. The angular momentums of the BH and disk can be calculated with

$$J_n = \frac{a_n GM_n^2}{c},$$

$$J_n^* = 2\pi \int_{r_{i,n}}^{r_{o,n}} \Sigma_n l_n r dr,$$

where $l_n$ is the specific angular momentum per unit mass. It is known that $a_{ep}$ cannot exceed unity (e.g., Janiuk et al. 2008). Since $M_n^*$ supplies all of the accretion process, we have

$$M_n^* = \sum_{n} M_n T_n,$$

where $T_n$ is the timescale of the $n$th epoch, which depends on the competition between the magnetic flux pressure and the gravity of the accreting mass.

2.2. Switch of the Magnetic Field

Since the predecessor of the accretion disk is a highly magnetic NS, the conservation of magnetic flux requires an inherited magnetic field in the disk. The successive magnetic flux from the remaining NS can be given by

$$\Phi_{NS} = 2\pi \int_{r_{i,1}}^{r_{o,1}} B_1 r dr,$$

where $B_1$ is the magnetic induction strength of the disk in the initial condition. The instabilities become effective when the magnetic pressure in the radial direction can counteract the gravity of the BH (e.g., Spruit & Uzdensky 2005; Proga & Zhang 2006). The accretion process may be closed if the magnetic induction strength satisfies a critical value $B_{\text{crit},n}$.

$$\frac{B_{\text{crit},n}^2}{4\pi} \sim \frac{GM_n \Sigma_n}{r_{m,n}},$$

where $M_n$ and $\Sigma_n$ are the mass of the BH and the surface density, respectively, in the $n$th episode. The critical radius $r_{\text{crit},n}(\leq r_{o,n})$ can be estimated with

$$2 \int_{r_{i,n}}^{r_{o,n}} B_n r dr = B_{\text{crit},n}^2 r_{m,n},$$

where $B_n$ is calculated by the conservation of magnetic flux. Therefore, the timescale is obtained by

$$T_n \sim \frac{r_n}{|\vec{v}_n|},$$

where $|\vec{v}_n|$ is the absolute value of the average radial velocity from $r_{\text{crit},n}$ to $r_{m,n}$.

We estimate the magnetospheric radius from

$$r_m \approx 6 \times 10^5 (v/v_{ff})^{3/2} (M/M_\odot)\Sigma^{-2/3} (M/3 M_\odot)^{-4/3} r_g,$$

where $v$, $v_{ff}$, $M$, and $r_g$ are the radial velocity of the accretion disk, the free fall velocity, the mass accretion rate, and the Schwarzschild radius, respectively (e.g., Narayan et al. 2003; Proga & Zhang 2006). Thus, the timescale of magnetic field dissipation can be estimated as $\sim r_m/v_{ff}$. We simplify the magnetic field as a uniform field with a magnetic flux of $10^{29}$ G cm$^2$ in our calculation (e.g., Proga & Zhang 2006). For $M_1 = 3 M_\odot$, $a_{e1} = 0.9$, $M_1 = 0.05 M_\odot$, $\alpha_1 = 0.01$, and $v/v_{ff} \sim 10^{-2}$ to $10^{-3}$, the timescale of the first episode is $T_1 \sim 2$ s from Equation (10), and the timescale of magnetic field dissipation is $\sim 0.1$–1 s. If the remaining magnetic flux is reduced to $5 \times 10^{28}$ G cm$^2$ and $M_2 \sim M_1$, $a_2 \sim \alpha_1$, and $a_{e2} \sim a_{e1}$, the timescale of the second episode is $T_2 \sim 27$ s in the case of $M_2 = 0.01 M_\odot$ s$^{-1}$. Therefore, our model can potentially explain the EE of SGRBs.

2.3. NDAF Model

We adopt the method presented by Riffert & Herold (1995). This method is dedicated to numerically investigating the NDAF in the vicinity of a rotating BH. The method defines general relativistic correction factors as quoted below,

$$A = 1 - \frac{2GM}{c^2 r} + \left(\frac{a_n GM}{c^2 r}\right)^2,$$

$$B = 1 - \frac{3GM}{c^2 r} + 2 a_n \left(\frac{GM}{c^2 r}\right)^\frac{1}{2},$$

$$C = 1 - 4 a_n \left(\frac{GM}{c^2 r}\right)^\frac{1}{2} + 3 \left(\frac{a_n GM}{c^2 r}\right)^2,$$

$$D = \int_{r_m}^{r} \frac{\sqrt{\frac{\Sigma}{4\pi}}} {2GM} \left[ \frac{\Sigma^2 (\Sigma_0^2 M c^2)}{G} + 2 \left(\frac{x \alpha M c^2}{G}\right)^\frac{1}{2} \right] dx,$$

where $M$, $a_n$, and $r_m$ are the mass, the dimensionless spin parameter, and the radius of the marginally stable orbit of the BH (e.g., Kato et al. 2008), respectively. The continuity equation remains valid,

$$\dot{M} = -2 \pi r \Sigma v.$$ 

The hydrostatic equilibrium in the vertical direction leads to a corrected expression for the half-thickness of the disk (Riffert & Herold 1995; Lei et al. 2009; Liu et al. 2010b),

$$H \approx c_s \left(\frac{P}{GM}\right)^\frac{1}{2} \left(\frac{B}{C}\right)^\frac{1}{2},$$

where $c_s$ is $(p/\rho)^{1/2}$ is the isothermal sound speed, and $p$ and $\rho$ are the total pressure and the density of the disk, respectively. The viscous shear $T_{\phi \theta}$ is also corrected as (Liu et al. 2010b)

$$T_{\phi \theta} = -\alpha \rho A \frac{BC}{(BC)^2},$$

where $\alpha$ is a dimensionless constant that absorbs all of the detailed microphysics of the viscous processes. The angular momentum equation can be simplified as

$$T_{\phi \theta} = \frac{\dot{M}}{4\pi H} \left(\frac{GM}{r^3}\right)^\frac{1}{2} \left(\frac{D}{A}\right)^\frac{1}{2}.$$
The total pressure includes the gas pressure from nucleons \( p_{\text{gas}} \), the radiation pressure of photons \( p_{\text{rad}} \), the degeneracy pressure of electrons \( p_{e} \), and the radiation pressure of neutrinos \( p_{\nu} \) (see, e.g., Liu et al. 2007):

\[
p = p_{\text{gas}} + p_{\text{rad}} + p_{e} + p_{\nu}.
\]  

(20)

The energy equation is given by

\[
Q_{\text{vis}} = Q_{\text{adv}} + Q_{\text{photo}} + Q_{\nu},
\]  

(21)

where \( Q_{\text{vis}} \), \( Q_{\text{adv}} \), \( Q_{\text{photo}} \), and \( Q_{\nu} \) are the viscous heating rate, the advective cooling rate, the cooling rate due to photodisintegration of \( \alpha \)-particles, and the cooling due to the neutrino radiation, respectively (see, e.g., Liu et al. 2007). The heating rate \( Q_{\text{vis}} \) is expressed as

\[
Q_{\text{vis}} = \frac{3GMM\alpha D}{8\pi r^3}.
\]  

(22)

The radiation luminosity of the neutrinos released from the disk is obtained with the neutrino cooling rate \( Q_{\nu} \), i.e.,

\[
L_{\nu} = 4\pi \int_{r_{\text{in}}}^{r_{\text{out}}} Q_{\nu} rd\tau.
\]  

(23)

We follow the approach by Ruffert et al. (1997), Popham et al. (1999), and Rosswog et al. (2003) to calculate the neutrino annihilation luminosity. The disk is modeled as a grid of cells in the equatorial plane. A cell \( k \) has its mean neutrino energy \( \epsilon_{\nu}^k \), neutrino radiation luminosity \( L_{\nu}^k \), and distance to a point in space above (or below) the disk \( d_k \). The angle at which neutrinos from cell \( k \) encounter antineutrinos from another cell \( k' \) at that point is denoted as \( \theta_{kk'} \). Then the neutrino annihilation luminosity at that point is given by the summation over all pairs of cells,

\[
L_{\nu\nu} = A_{1} \sum_{k} \frac{L_{\nu}^k}{d_k} \sum_{k'} \frac{L_{\nu}^{k'}}{d_{k'}} (\epsilon_{\nu}^k + \epsilon_{\nu}^{k'}) (1 - \cos \theta_{kk'})^2 + A_{2} \sum_{k} \frac{L_{\nu}^k}{d_k} \sum_{k'} \frac{L_{\nu}^{k'}}{d_{k'}} \epsilon_{\nu}^k \epsilon_{\nu}^{k'} (1 - \cos \theta_{kk'}),
\]  

(24)

where \( A_{1} \approx 1.7 \times 10^{-44} \text{ cm erg}^{-2} \text{s}^{-1} \) and \( A_{2} \approx 1.6 \times 10^{-56} \text{ cm erg}^{-2} \text{s}^{-1} \) (e.g., Popham et al. 1999). The total neutrino annihilation luminosity is integrated over the whole space outside the BH and the disk,

\[
L_{\nu\nu} = 4\pi \sum_{i} \int_{r_{\text{in}}}^{r_{\text{out}}} \int_{H} L_{\nu\nu} rd\tau dz.
\]  

(25)

However, as shown by Popham et al. (1999) and Liu et al. (2007), the neutrino annihilation would inject a highly beaming outflow around the inner part of the disk. For \( M = 3 M_{\odot} \), \( M = 0.01-1 M_{\odot} \text{s}^{-1} \), and \( \alpha = 0.01-0.1 \), the opening angle of the ejection \( \theta \) is about \( 10^\circ-20^\circ \). We conservatively assume that the opening angle of ejection is \( 10^\circ \) and the efficiency of the fireball is \( \eta = 0.1 \) in our calculations. The observed isotropic luminosity of the \( n \)th step \( L_{\text{iso},n} \) can be estimated with

\[
L_{\text{iso},n} = \eta L_{\nu\nu, n}/(1 - \cos \theta).
\]  

(26)

3. NUMERICAL RESULTS

Observationally, half of the Type I GRBs have EE component detection. We show the light curves of some SGRBs with EE components in Figure 2. The light curves are visually recognized as different emission episodes. We apply our model to these SGRBs. For simplicity, we consider only two emission episodes, i.e., the initial hard spikes and the EE component. We describe our numerical method for the two emission episodes as follows.

There are seven unknown variables in our model, i.e., \( \alpha_{1}, \alpha_{2}, M_{1}, M_{2}, M_{1}', M_{2}', \) and \( M_{\ast}^{\alpha} \). Therefore, a group of seven equations is required for our calculations. Three of them are from the conservation of mass (energy), i.e.,

\[
M_{2} = M_{1} + M_{1} \delta_{\text{in}},
\]  

(27)

\[
M_{2}' = M_{1}' - M_{1},
\]  

(28)

and

\[
M_{\ast}^{2} = M_{2}' T_{2} + \delta M^{\ast},
\]  

(29)

where \( \delta M^{\ast} \) is the mass of the residual disk after the second episode. Two other equations are from the conservation of angular momentum, i.e.,

\[
J_{2} = J_{1} + M_{1} T_{1} l_{\text{in}},
\]  

(30)

and

\[
J_{2}' = J_{1}' - M_{1} T_{1} l_{\text{in}},
\]  

(31)

The values of \( J_{1} \) and \( J_{2} \) are calculated with Equation (4). \( J_{1}' (J_{2}') \) is determined by \( M_{1}(M_{2}) \), \( a_{1}(a_{2}) \), \( M_{1}'(M_{2}') \), and \( \alpha \), which is calculated by Equation (5). The other two equations are related to \( L_{\nu\nu} \), which is a function of \( M, a, \) and \( \alpha \) in the NDAF model,

\[
L_{\nu\nu,1} = f(M, a_{1}, M_{1}, \alpha),
\]  

(32)

\[
L_{\nu\nu,2} = f(M_{2}, a_{2}, M_{2}, \alpha),
\]  

(33)

where the function \( f \) can be obtained from Equations (12)–(25). The average neutrino annihilation luminosity in each episode is calculated assuming that the accretion rate and the spin parameter of the BH are constant. The accretion timescale in each step depends on the initial mass of the disk and the switch in polarity of the magnetic field. The procedure of our calculations is as follows.

First, we assign the initial parameters of the BH and the disk. We fix the initial mass and the spin parameter of the BH as \( 3 M_{\odot} \) and 0.9, respectively. The viscous parameter is assumed to be \( \alpha = 0.01 \). The initial mass of the disk is adjustable in our calculation.

Second, we take the observed average luminosity and the timescale of the initial hard spikes to be \( L_{\nu\nu,1} \) and \( T_{1} \), and then calculate \( M_{1} \) with Equation (32). The values of \( M_{2}, M_{2}', a_{2}, \) and \( M_{2} \) are derived from Equations (27), (28), (30), and (33), respectively.

Third, we calculate \( T_{2} \) and \( L_{\nu\nu,2} \) with Equations (6) and (33) by taking a small value of \( \delta M^{\ast} \), i.e., about \( 0.1 M_{\odot} \).

We adjust the initial mass of the disk and find that in the case of \( M_{1}' = 0.8 M_{\odot} \), our results are roughly consistent with the observations for some typical SGRBs with EE, i.e., 050724, 060614, 061006, 061210, 070714B, and 071227. Our results are shown in Figure 3. The first episode of these GRBs is in the top-left circle of the figure. We correspondingly derive a region of the second episode in a circle with the universal parameters mentioned above. We note that besides GRB 050724 and 071227, the second episode of the other four GRBs is in the region. Note that the EE component of the two GRBs is
likely a late flare, which may have a different physical origin as mentioned in Section 1.

The mass of the disk is a very important factor for the EE. The decrease of $M^*_1$ would result in the significant decrease of the timescale and luminosity in the second episodes. For the case of $M^*_1 = 0.2 \, M_\odot$, we find that the timescale of the second episode is shorter than 20 s\textsuperscript{11} and the luminosity is roughly one order of magnitude lower than that in the case of $M^*_1 = 0.8 \, M_\odot$, as shown in Figure 3. Our results indicate that a bright EE component may only be detected for SGRBs with a massive disk.

4. CONCLUSIONS AND DISCUSSION

We have proposed that both the outward angular momentum transfer and the switch in polarity of the magnetic barrier of a rotating BH–neutrino-cooled disk system may result in a long-lasting impulse engine, producing several radiation episodes as observed in Type I GRBs. Based on the NDAF model, we have presented detailed calculations for some typical GRBs

\textsuperscript{11} A low viscosity would increase the accretion time (e.g., Metzger et al. 2008). In our calculations, we take $\alpha_1 = \alpha_2 = 0.01$. 

Figure 2. BAT light curves of six SGRBs with detection of EE components. (A color version of this figure is available in the online journal.)
The detection or non-detection of the EE component might be due to an instrumental selection effect. Our model suggests that a disk with \( M^* < 0.2 \, M_\odot \) would not produce an EE component that can be detectable with \textit{Swift}/BAT-like instruments. Assuming a constant radiation efficiency, the observed luminosity is proportional to the neutrino annihilation luminosity. Since the peak energy of the \( v_\nu \) spectrum is tightly correlated with the radiation luminosity (Yonetoku et al. 2004; Liang et al. 2004), the EE component of some SGRBs may be out of the instrument bands and only the initial hard–short episode can be detectable as usually seen in typical SGRBs. With the results for \( M^* = 0.8 \, M_\odot \) and \( M^* = 0.2 \, M_\odot \), one can observe a luminosity–duration correlation for the EE component. This seems to be at odds with observations that show the opposite trend: longer events have a lower luminosity. Note that the observations for the EE component greatly suffer from the selection effect of instrumental sensitivity. We check the observed luminosity–duration correlation for both the initial hard spikes and the EE components for the GRBs with EE detection reported in Zhang et al. (2009), but no statistically accepted correlation is found. The relation between the luminosity and the duration of our model is parameter dependent. This calls for a detailed analysis based on Monte Carlo simulations in order to elaborate this relation and its observational biases.

Some caveats for our model should be discussed. The neutrino annihilation luminosity is correlated positively with the accretion rate. In order to ensure that the mass of an NS is less than \( 2 \, M_\odot \), a moderate accretion rate is required in our model. The derived accretion rates for both the initial and late episodes from our model are less than \( 0.05 \, M_\odot \, \text{s}^{-1} \), which are lower than the rates for the typical ones, i.e., \( 0.1–1 \, M_\odot \, \text{s}^{-1} \). Note that the typical values are for a spatially uniform neutrino annihilation luminosity. In our calculations, the spatial distribution of the neutrino annihilation luminosity is collimated (e.g., Popham et al. 1999; Liu et al. 2007). The jet opening angle is conservatively assumed to be \( 10^\circ \) in our calculations.\(^{12}\) Therefore, the isotropic-equivalent accretion rate would be larger than our results by one to two orders of magnitude, which is consistent with the typical values used by previous authors.

We have considered the switch in polarity of the magnetic field for the accretion processes, but the effects of confining the jet opening angle by the BZ mechanism (Blandford & Znajek 1977) and of enhancing the neutrino annihilation are ignored. These effects may also greatly impact the jet luminosity. Lei et al. (2009) investigated the magnetic coupling between the BH and the disk. They found that the luminosity of neutrino annihilation is much larger than the luminosity of NDAF without magnetic field. Recently, Barkov & Pozanenko (2011) proposed a two-jet model that describes both main and EE components using different off-axis observer positions. Their model involves a short-duration jet powered by heating due to neutrino annihilation and a long-lived BZ jet with a significantly narrow opening angle. The BZ mechanism can replace the neutrino annihilation to produce the main emission and EE in our model. Because of the lack of information regarding the intensity and distribution of the magnetic field, we restrict our discussion to neutrino annihilation as the main power source in our model.

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\(^{12}\) The jet opening angle of SGRBs may even be much smaller than \( 1^\circ \), such as \( \sim 0.3 \) for GRB 090510 (Corsi et al. 2010; He et al. 2011; Fan & Wei 2011).

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\[ L_{\nu} = L_{\nu}^\text{obs} \times (1+z)^{\delta} \]

\[ T_{\text{obs}} = T_{\text{acc}} \]

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\[ M^* = 0.8 \, M_\odot \]

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\[ M^* = 0.2 \, M_\odot \]

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\[ M^* \text{ in the top left circle of the figure. We correspondingly derive a region of the second episode with the universal parameters for } M^* = 0.8 \, M_\odot \text{. Besides GRB 050724 and 071227, the second episode of the other four GRBs are in the region. The EE component of the two GRBs is likely a late flare, which may have a different physical origin as mentioned in Section 1. The corresponding region for } M^* = 0.2 \, M_\odot \text{ is also shown. It indicates that the luminosity of the second episode is roughly one order of magnitude lower than that for the case of } M^* = 0.8 \, M_\odot \text{ and the timescale is shorter than } 20 \, \text{s. The switch in polarity of the timescale of the magnetic barrier is } 0.1–1 \, \text{s. This timescale is much smaller than the observed timescale of the EE component. We thus ignore it here.} \]

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\[ T_{\text{obs}} = T_{\text{acc}} \]

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\[ M^* \text{ in the top left circle of the figure. We correspondingly derive a region of the second episode with the universal parameters for } M^* = 0.8 \, M_\odot \text{. Besides GRB 050724 and 071227, the second episode of the other four GRBs are in the region. The EE component of the two GRBs is likely a late flare, which may have a different physical origin as mentioned in Section 1. The corresponding region for } M^* = 0.2 \, M_\odot \text{ is also shown. It indicates that the luminosity of the second episode is roughly one order of magnitude lower than that for the case of } M^* = 0.8 \, M_\odot \text{ and the timescale is shorter than } 20 \, \text{s. The switch in polarity of the timescale of the magnetic barrier is } 0.1–1 \, \text{s. This timescale is much smaller than the observed timescale of the EE component. We thus ignore it here.} \]

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\[ M^* \text{ in the top left circle of the figure. We correspondingly derive a region of the second episode with the universal parameters for } M^* = 0.8 \, M_\odot \text{. Besides GRB 050724 and 071227, the second episode of the other four GRBs are in the region. The EE component of the two GRBs is likely a late flare, which may have a different physical origin as mentioned in Section 1. The corresponding region for } M^* = 0.2 \, M_\odot \text{ is also shown. It indicates that the luminosity of the second episode is roughly one order of magnitude lower than that for the case of } M^* = 0.8 \, M_\odot \text{ and the timescale is shorter than } 20 \, \text{s. The switch in polarity of the timescale of the magnetic barrier is } 0.1–1 \, \text{s. This timescale is much smaller than the observed timescale of the EE component. We thus ignore it here.} \]
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