EPISODIC JETS AS THE CENTRAL ENGINE OF GAMMA-RAY BURSTS

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ABSTRACT

Most gamma-ray bursts (GRBs) have erratic light curves, which demand that the GRB central engine launches an episodic outflow. Recent Fermi observations of some GRBs indicate a lack of the thermal photosphere component as predicted by the baryonic fireball model, which suggests a magnetic origin of GRBs. Given that powerful episodic jets have been observed along with continuous jets in other astrophysical black hole systems, here we propose an intrinsically episodic, magnetically dominated jet model for the GRB central engine. Accumulation and eruption of free magnetic energy in the corona of a differentially rotating, turbulent accretion flow around a hyperaccreting black hole lead to ejections of episodic, magnetically dominated plasma blobs. These blobs are accelerated magnetically, collide with each other at large radii, trigger rapid magnetic reconnection and turbulence, efficient particle acceleration, and radiation, and power the observed episodic prompt gamma-ray emission from GRBs.

Key words: accretion, accretion disks – gamma-ray burst: general – ISM: jets and outflows – magnetic reconnection

1. INTRODUCTION

Observations of gamma-ray bursts (GRBs) suggest that the GRB central engine is able to launch an ultraluminous, highly relativistic jet. Most GRBs have erratic, rapidly varying light curves (Fishman & Meegan 1995), typically lasting tens to hundreds of seconds for long-duration GRBs and less than 2 s for short-duration GRBs. Recent observations of GRBs pose some important constraints on models of the GRB central engine. First, recent Fermi Large Area Telescope observations suggest that most GRBs have featureless, smoothly joined broken power-law spectra (i.e., Band-function spectra (Band et al. 1993)) over a wide energy band (Abdo et al. 2009; Zhang et al. 2011). The standard fireball model predicts a strong thermal emission component from the fireball photosphere (Paczynski 1986; Mészáros & Rees 2000; Pe’er et al. 2006). The non-detection likely suggests that the GRB outflow is magnetically dominated (Zhang & Pe’er 2009; Fan 2010). Second, data analysis suggests that the GRB light curves not only can be decomposed into many pulses (Norris et al. 1996; Hakilla et al. 2003), but most of them can also be decomposed into the superposition of a fast, spiky component and a slow, smooth component (Gao et al. 2012; Vetere et al. 2006). A radiation model that invokes magnetic turbulent reconnection triggered by collisions of magnetically dominated blobs has been proposed (Zhang & Yan 2011); this can interpret the new observational data. This radiation model requires a central engine that can eject an episodic, magnetically dominated jet.

The leading model of the GRB central engine invokes a hyperaccreting black hole (Narayan et al. 1992, 2001; Mészáros et al. 1999). In most previous GRB central engine models, the rapid variability in the erratic light curves is attributed either to the intermittency of the accretion flow, i.e., a time-dependent accretion rate $M$ (MacFadyen & Woosley 1999), or to instability during the propagation of the jet in the stellar envelope (Zhang et al. 2003; Morsony et al. 2010). For a magnetically dominated central engine, the leading model is the Blandford–Znajek (BZ) mechanism (Blandford & Znajek 1977), which requires a large-scale open magnetic field connecting the black hole and an external astrophysical load, the origin of which is an open question. This model also tends to generate a continuous jet unless the accretion rate is highly variable. Screw or kink instabilities (Li 2000; Mizuno et al. 2009) are invoked to disrupt a continuous jet and produce discrete blobs. Since the BZ mechanism is powered by accretion, the immediate advantage of involving black hole spin rather than the accretion disk as the source of jet power is not evident.

On the other hand, in addition to the continuous jets, episodic jets, which are intermittent and in the form of discrete moving blobs, have been observed in other black hole systems. A magnetohydrodynamical (MHD) model has been proposed to explain the formation of these episodic jets (Yuan et al. 2009). In this paper, we propose a central engine model for GRBs based on this idea. A review of observations and the model of episodic jets are presented in Section 2. Our model is delineated in detail in Section 3. The salient features of the model are summarized in Section 4.

2. EPISODIC JETS IN BLACK HOLE SYSTEMS AND THEIR FORMATION MECHANISM

Episodic jets are most evidently observed in black hole X-ray binaries, e.g., in GRS 1915+105 and GRO J1655−40 (Mirabel & Rodríguez 1994; Hjellming & Rupen 1995; Fender & Belloni 2004), and also in active galactic nuclei (AGNs), as manifested as knots or blobs, e.g., in 3C 120 (Marscher et al. 2002; Chatterjee et al. 2009) and NGC 4258 (Doi et al. 2011). In the case of 3C 120, knots have been directly resolved as due to episodic ejections, and their ejection times have been determined by extensive radio monitoring observations. In the case of the supermassive black hole in our Galactic center, we have strong evidence for episodic jets, but no continuous jets have ever been detected (Yuan 2011). The differences in the properties of the two types of jets are summarized in Fender & Belloni (2004). Amongst other things, episodic jets are much faster and more powerful. The collision between blobs is often invoked to explain the flares detected in AGN jets.
Yuan et al. (2009) proposed an MHD model for the formation of episodic jets by analogy with coronal mass ejection (CME) in the Sun. The basic scenario is as follows: closed magnetic field lines continuously emerge out of the accretion flow into the corona. Because of shear and turbulent motion of the accretion flow, the field line is twisted and deformed, resulting in the formation of a flux rope in the corona. The flux rope is initially in force balance between magnetic tension and magnetic compression forces. Energy and helicity are accumulated and stored until a threshold is reached. The system then loses its equilibrium, and the flux rope is thrust outward by the magnetic compression force in a catastrophic way, which causes an episodic jet. After a magnetic blob is ejected, the magnetic tension is temporarily relaxed. Later, magnetic field emergence and distortion restart, and the process described above repeats. Within this scenario, no large-scale open magnetic field lines are needed. The basic picture has gained support from some numeric simulations (e.g., Romanova et al. 1998; Kudoh et al. 2002; Machida et al. 2004); see also Ouyed et al. (1997) and Ouyed & Pudritz (1997) for a related model of episodic jets. Within the GRB context, several magnetic central engine models along similar lines have been proposed. In the context of their neutron star central engine model, Blackman et al. (1996) pointed out that magnetically dominated blobs rather than a continuous jet have the advantage in reproducing the time structure of GRBs. Kluzniak & Ruderman (1998) and Dai & Lu (1998) also emphasized the importance of accumulation of magnetic fields and subsequent emergence to account for an episodic jet. A similar idea was proposed to interpret X-ray flares (Dai et al. 2006) in some short GRBs. Uzdensky & MacFadyen (2006, 2007) proposed a magnetic tower model for GRBs that invokes extraction of the rotational energy of an accretion disk by amplification of the toroidal magnetic field (Lynden-Bell 1996). The model is essentially time independent, different from the model discussed in Yuan et al. (2009) and below.

3. EPISODIC MAGNETIC JETS IN GRBs

We now extend the Yuan et al. (2009) scenario to GRBs. A schematic picture of our model is shown in Figure 1. We consider a stellar-size black hole surrounded by a hyperaccretion flow with a typical accretion rate \( \dot{M} \sim 0.1 M_\odot \, \text{s}^{-1} \). Neutrino cooling is important at \( \lesssim 200 \, \text{G} \mu \text{m}^2 / \text{c}^2 \) (neutrino-dominated accretion flow, NDAF); outside of this radius the accretion flow is advection dominated (ADAF; Chen & Beloborodov 2007). The solutions for NDAF and ADAF can be described by (Narayan et al. 2001; Beloborodov 2003)

\[
\rho = 6 \times 10^{13} \alpha_{-2}^{-1/3} M_{32}^{-1/3} r^{-2.55} \, \text{g cm}^{-3},
\]

\[
T_c = 3 \times 10^{10} \alpha_{-2}^{0.2} M_{3}^{-0.2} r^{-0.3} \, \text{K},
\]

\[
v_k = 2 \times 10^{10} r^{-0.5} \, \text{cm s}^{-1},
\]

\[
v_r = 2 \times 10^{9} \alpha_{-2}^{1/2} M_{3}^{-0.2} r^{0.2} \, \text{cm s}^{-1},
\]

and

\[
\rho = 3 \times 10^{11} \alpha_{-2}^{-1/4} M_{32}^{-1/4} r^{-1.5} \, \text{g cm}^{-3},
\]

\[
T_c = 3 \times 10^{11} \alpha_{-2}^{-1/4} M_{32}^{1/4} M_{3}^{-0.5} r^{-5/8} \, \text{K},
\]

respectively. Here, \( \rho \) is the density, \( T_c \) is the temperature at the equatorial plane, \( v_k = \sqrt{GM/R} \) and \( v_r \) are the Keplerian and radial velocities, \( \alpha = 0.01 \alpha_{-2} \) is the viscous parameter, \( r \) is radius \( R \) in units of \( 2GM/c^2 \), \( M_3 \) is the mass accretion rate in units of \( 10^{32} \, \text{g} \, \text{s}^{-1} \), and \( M_3 \) is the black hole mass in units of \( 3 M_\odot \). The viscous timescales of the inner NDAF and outer ADAF are then

\[
t_{\text{vis,NDAF}} = R/v_r = 17 \alpha_{-2}^{-1/2} M_{3}^{1/2} (r/100)^{0.8} \, \text{s}
\]

and

\[
t_{\text{vis,ADAF}} = R/v_r = 30 \alpha_{-2}^{-1} M_3 (r/200)^{1.5} \, \text{s}
\]

respectively. Hereafter, we normalize the NDAF solutions at \( r = 100 \) where the disk solution transitions from NDAF to ADAF. All our ADAF solutions are normalized to \( r = 200 \), which corresponds to a viscous timescale of \( \sim 30 \, \text{s} \), the typical duration of a long GRB. We will show below that powerful episodic jets are launched in the ADAF regions, i.e., at radii larger than \( r \sim 100 \), since the ejection power of magnetic blobs is relatively suppressed in the NDAF region (see discussions below Equations (20) and (21)).

Now we estimate the energetics and time intervals of the magnetic blobs. The region in which a magnetic blob occurs, i.e., the flux rope region in the disk corona, is special because the available free magnetic energy is large due to the topological structure of the magnetic field. By analogy with the CME theory.
of solar physics, the total available free magnetic energy of one blob is (Lin et al. 1998)

\[ E_{\text{free}} \approx 0.5 \times \left( \frac{1}{12} B_0^2 v \right) \]  

(11)

Here, \( B_0 \) is the magnetic strength in the accretion disk and \( V \sim \pi R_1^3 \) is the volume of the system. Defining \( \beta = P_{\text{mag}} / P_{\text{tot}} \), where \( P_{\text{tot}} = (P_{\text{gas}} + P_{\text{rad}} + P_{\text{mag}} + P_{\text{c}}) \sim P_{\text{gas}} \) in NDAF and \( P_{\text{tot}} = (P_{\text{gas}} + P_{\text{rad}} + P_{\text{mag}}) \sim P_{\text{rad}} \) in ADAF (where \( P_{\text{mag}} = B^2 / (8\pi) \), \( P_{\text{gas}} \), \( P_{\text{rad}} \), \( P_{\text{c}} \), and \( P_{\text{tot}} \) are magnetic, gas, radiation, neutrino, and total pressure in the accretion disk, respectively), one can estimate the strength of the magnetic field \( B_0 \) for a given \( \beta \). Numerical simulations show that if the viscosity \( \alpha \) is intrinsically the magnetic stress associated with the MHD turbulence driven by the magnetorotational instability (Balbus & Hawley 1991, 1998), as widely accepted, then the values of \( \alpha \) and \( \beta \) are not independent (Blackman et al. 2008).

This is confirmed by recent numeric simulations (Hawley et al. 2011; Sorathia et al. 2012). Noticing the different definitions between our work and Blackman et al. (2008), one should have roughly \( \alpha / \beta \sim 0.1 \) according to Blackman et al. (2008). Therefore, \( \alpha = 0.01 \) implies that the value of \( \beta \) should be \( \beta = 0.1 \beta_{-1} = 0.1 \). In the following, we adopt \( \alpha = 0.01 \) and \( \beta = 0.1 \) as typical values. We note, however, that a much higher \( \beta \gg 1 \) can be achieved if the accretion material is already moderately magnetized at the beginning of accretion (Shibata et al. 1990; Johansen & Levin 2008).

For the inner gas-pressure-dominated NDAF and outer radiation-pressure-dominated ADAF, one has

\[ \frac{B^2}{8\pi} \approx 0.1 \beta_{-1} P_{\text{gas}} \]  

(12)

and

\[ \frac{B^2}{8\pi} \approx 0.1 \beta_{-1} P_{\text{rad}} = 0.1 \beta_{-1} \frac{4\alpha}{3\pi} T_{c}^4, \]  

(13)

respectively, so that the available energy in the NDAF and ADAF regions are

\[ E_{\text{free},\text{NDAF}} = 7.4 \times 10^{48} \alpha_{-2}^{-1.5} \beta_{-1}^{1.5} M_3^{1.5} M_{\odot} (r/100)^{0.15} \text{ erg} \]  

(14)

and

\[ E_{\text{free},\text{ADAF}} = 8.2 \times 10^{49} \alpha_{-2}^{-1.5} \beta_{-1}^{1.5} M_3 M_{\odot} (r/200)^{0.5} \text{ erg}, \]  

(15)

respectively.

The time interval between two consecutive ejections can be estimated as the timescale to accumulate and release \( E_{\text{free}} \) in the flux rope system in the corona of the accretion flow. The first timescale is the energy accumulation time. The magnetic energy of the flux rope system is converted from the rotational energy of the accretion flow by Alfvén waves propagating along the magnetic field lines (Yuan et al. 2009). The energy density of the rotational energy is \( \sim \rho c^2 \) and the energy conversion speed is the Alfvén speed \( v_A \equiv B_0 / \sqrt{4\pi \rho} \). The energy transfer rate by the magnetic field is then \( E_{\text{mag}} = \rho c^2 v_A R^2 \). The corresponding energy transfer timescales in the NDAF and ADAF regimes are

\[ t_{\text{tran},\text{NDAF}} = 4.2 \times 10^{-3} \alpha_{-2}^{0.1} \beta_{-1}^{0.5} M_3^{0.9} (r/100)^{0.85} \text{ s} \]  

(16)

and

\[ t_{\text{tran},\text{ADAF}} = 4.8 \times 10^{-1} \beta_{-1}^{0.5} M_3 (r/200)^{1.5} \text{ s}, \]  

(17)

respectively.

The second relevant timescale is the emergence timescale of the magnetic field line from the disk to the corona due to Parker instability (Horiuchi et al. 1998):

\[ t_{\text{parker, NDAF}} \approx \frac{5H}{v_A} \approx 4\alpha_{-2}^{-0.1} \beta_{-1}^{-0.5} M_3 (r/100)^{1.5} s \gg t_{\text{tran},\text{NDAF}}, \]  

(18)

and

\[ t_{\text{parker, ADAF}} \approx \frac{5H}{v_A} \approx 3.4 \beta_{-1}^{-0.5} M_3 (r/200)^{1.5} s \gg t_{\text{tran},\text{ADAF}}, \]  

(19)

respectively. The time interval between two consecutive ejections is thus \( t_{\text{int}} = t_{\text{tran}} + t_{\text{parker}} \approx t_{\text{parker}} \) for both NDAF and ADAF. The power output from these two regions is then

\[ P_{\text{NDAF}} = 1.8 \times 10^{48} \alpha_{-2}^{-1} \beta_{-1}^{1.5} M_3^2 (r/100)^{-1} \text{ erg s}^{-1} \]  

(20)

and

\[ P_{\text{ADAF}} = 2.4 \times 10^{49} \alpha_{-2}^{-1} \beta_{-1}^{1.5} M_3^2 (r/200)^{-1} \text{ erg s}^{-1}, \]  

(21)

respectively. One can see that the power output increases significantly beyond the "transition region" (\( r \sim 100–200 \)) from NDAF to ADAF. So the ADAF region is the main region for powerful magnetic blob injection. 3 From now on, we focus on the ADAF only and neglect the subscript "ADAF" in the equations.

For long GRBs, the jet-corrected total energy is of the order of \( 10^{51} \) erg (Frail et al. 2001; Liang et al. 2008; Racusin et al. 2009). Since a GRB may be powered by multiple collisions among blobs, each blob would carry an energy of \( \sim 10^{50} \) erg. For \( r \sim 100–200 \), our estimate of Equation (15) matches this observational fact well. Comparing Equations (19) and (10), we find that \( t_{\text{int}} < t_{\text{vis}} \). This suggests that the system has enough time to store magnetic energy and to eject multiple magnetic blobs before being accreted into the black hole.

Following Uzdensky & MacFadyen (2006, 2007), we assume that a magnetic jet can penetrate the star and remain intact. For an episodic jet, this is possible as long as the time interval between two consecutive blobs is shorter than the time for the funnel to close, which is \( \sim 10 \) s (Wang & Mészáros 2007). This condition is satisfied for our model. After escaping the star, the magnetic blob undergoes acceleration under its own magnetic pressure gradient (Tchekhovskoy et al. 2010). An impulsive magnetic blob can reach \( \Gamma \sim \sigma_0 1/3 \) quickly and then gradually accelerate as \( \Gamma \propto R^{1/3} \) until reaching \( \Gamma \sim \sigma_0 \) (Granot et al. 2011), where \( \sigma_0 \equiv B_c^2 / 8\pi \rho_c c^2 \) is the initial magnetization parameter at the base of the flow, \( \rho_c \) is the density of the flux rope in the corona, and \( B_c \) is the magnetic field in the corona near the flux rope region. Unfortunately, both \( B_c \) and especially \( \rho_c \) are poorly constrained by current analytical theory and numeric simulations of accretion flows. MHD numeric simulations of both optically thin ADAFs and standard thin disks show that the density in the corona is strongly inhomogeneous and stratified. A density contrast as large as \( \sim 5 \) orders of magnitude has been achieved in numeric simulations (Hirose et al. 2009). This can

\( \text{3 The power output from the innermost NDAF region with } 10 > r > 3 \text{ can become comparable to that from the outer ADAF. However, the blobs emitted from these regions are closely packed and would likely collide below the photosphere and would not give rise to significant variability. In addition, this power can be consumed to penetrate the star and open a funnel for blobs ejected later from the ADAF region.} \)
be regarded as a lower limit due to the “density floor” imposed in the simulations by hand to stabilize the simulations. Given the uncertainties, we estimate the density of the corona by analogy with the case of the Sun. The density at the bottom of the solar corona is about 7–12 orders of magnitude lower than the density of the turbulence layer of the Sun (the counterpart of the corona is about 7–12 orders of magnitude lower than the density with the case of the Sun. The density at the bottom of the solar uncertainties, we estimate the density of the corona by analogy be regarded as a lower limit due to the “density floor” imposed

\[ \sigma_0 \approx 3500 \left( \frac{\rho/\rho_c}{10^9} \right)^{\beta-1} \left( \frac{r}{200} \right)^{-1}. \]  

For a higher \( \beta \) value as discussed above, a smaller density contrast can be incorporated to achieve the same \( \sigma_0 \) value. The Lorentz factor of the outflow in the emission region is \( \Gamma \lesssim \sigma_0 \) depending on the radius of the emission region, \( R_{\text{GRB}} \), from the central engine. Because of the inhomogeneity of the density of the corona, the value of \( \sigma_0 \) is expected to be variable for different blobs. This implies that their collisions will be efficient.

The initial size of each blob is of the order of the size of the disk, i.e., \( \Delta_0 \sim R_{\text{disk}} \sim 2 \times 10^8 \) cm \( M_3 \left( r/200 \right) \). The blob is initially at rest and is accelerated by the magnetic pressure gradient. First, it undergoes a non-relativistic phase (Yuan et al. 2009), during which the bubble continues to expand adiabatically. This phase lasts for a timescale of reconnection near the base of the flux rope, i.e., \( t_{\text{rec}} \sim R_{\text{disk}}/v_{\text{rec}} \sim 0.7 \) s \( r/(200) \). The expansion speed is essentially \( \sim c \) for a high-\( \sigma \) bubble. The size of the bubble grows to \( \Delta \sim c t_{\text{rec}} \sim 2 \times 10^{10} \) cm before entering the relativistic phase, during which the laboratory-frame width essentially stops growing.

During the relativistic phase, the Lorentz factor at the distance \( R_0 \sim \Delta \) is \( \Gamma(R_0) \sim \sigma_0^{1/3} \sim 15 \left( 10^3 \rho/\rho_c \right)^{1/3} \rho_3^{1/3} \left( r/200 \right)^{-1/3} \). With a slow increase, \( \Gamma \propto R^{1/3} \) (Granot et al. 2011), one would reach the full Lorentz factor \( \Gamma_{\text{full}} = 3500 \) at a radius \( R_{\text{full}} \sim 3.8 \times 10^{17} \) cm. GRB prompt emission likely occurs at much smaller radii where \( \Gamma(R_{\text{GRB}}) \ll \sigma_0 \). A plausible scenario would be the Internal Collision-induced MAgnetic Reconnection and Turbulence (ICMART) scenario conjectured by Zhang & Yan (2011). Within this scenario, most magnetic energy is discharged in the ICMART region, so that after dissipation \( \sigma \) drops to around or below unity, and the final bulk Lorentz factor \( \Gamma \sim \Gamma(R_{\text{GRB}}) \ll \sigma_0 \). For example, for \( R_{\text{GRB}} \sim 10^{15} \) cm, \( \Gamma_{\text{GRB}} \sim 550 \) for \( \sigma_0 = 3500 \) and \( \Gamma_{\text{GRB}} \sim 260 \) for \( \sigma_0 = 350 \).

The trigger of ICMART events is through internal collisions. It is expected that \( \sigma_0 \) of the blobs may vary from case to case due to the fluctuation of \( \rho_c \) and \( B_c \). The faster blobs ejected later would inevitably catch up with the preceding slower ones. Such collisions of magnetic blobs have been frequently observed in the Sun (Gopalswamy et al. 2002) and are often invoked to explain the bright knots or flares commonly observed in other astrophysical systems such as AGN jets and the Crab nebula. The separation between the blobs is \( d \sim c t_{\text{int}} \sim 10^{11} \) cm. The typical collision radius is

\[ R_{\text{GRB}} \sim \Gamma^2 c t_{\text{int}} \sim 10^{15} \left( \frac{\Gamma}{100} \right)^2 \beta_{-1}^{-0.5} M_3 \left( \frac{r}{200} \right)^{1.5} \text{cm}. \]  

For a typical GRB \( \Gamma \gtrsim 100 \) (Lithwick & Sari 2001; Liang et al. 2010), this radius is consistent with various observational constraints that suggest a relatively large \( R_{\text{GRB}} \) (Kumar et al. 2006; Nishikori et al. 2006; Zhang & Pe’er 2009; Shen & Zhang 2009; Fan 2010).

The polarities of the magnetic field lines trapping the two colliding blobs are in general opposite or have some large angles. This greatly eases the trigger of ICMART events. The initial fast reconnection would induce turbulence and a cascade of turbulent reconnection (Zhang & Yan 2011; Lazarian & Vishniac 1999). A large fraction of the magnetic energy stored in the blobs would be efficiently converted into lepton energy and then into photon energy, giving rise to radiatively efficient GRB emission (Zhang & Yan 2011; Uzdensky & McKinney 2011). The model can also account for the two variability components observed in GRBs (Gao et al. 2012; Vetere et al. 2006): the angular spreading time (Piran 1999)

\[ t_{\text{ang}} \sim \frac{R_{\text{GRB}}}{c \Gamma^2} \sim t_{\text{int}} \sim 3.4 \beta_{-1}^{-0.5} M_3 \left( \frac{r}{200} \right)^{1.5} \text{s} \]

corresponds to the timescale of the slow component, while the fast variability is related to relativistic turbulence (Narayan & Kumar 2009).

4. SUMMARY AND DISCUSSION

Observations of relatively well-studied black hole systems such as AGNs and black hole X-ray binaries show the existence of two types of jets: continuous and episodic ones. GRB observational data require that the central engine launches a magnetically dominated episodic jet. The traditional BZ jet model requires that the accretion rate is highly variable or that a continuous jet is disrupted by instabilities during propagation. In this paper, we propose an intrinsically episodic central engine model for GRBs by invoking ejections of episodic magnetic blobs from a hyperaccretion flow around a black hole. Our basic calculations suggest that the predicted energetics and timescales are consistent with the observations. More detailed numeric simulations are required to validate the scenario.

This model has several appealing features. First, the episodic jets invoked in our model have obtained strong observational support in other black hole systems (e.g., Yuan et al. 2009, and references therein). Observations show that they are more powerful than continuous jets. Second, this model naturally satisfies the observational requirement for the lack of a bright photosphere component in the GRB spectra (Zhang & Pe’er 2009), since the engine naturally launches a magnetized outflow. Third, episodic jets intrinsically consist of individual blobs whose collisions are naturally expected. These collisions are an important ingredient in interpreting GRB variability within the magnetically dominated model (Zhang & Yan 2011). Fourth, the directions of the magnetic fields surrounding the two adjacent blobs in general have some angles. This greatly eases triggering fast reconnection and the subsequent turbulence, which can efficiently convert magnetic energy into radiation (Zhang & Yan 2011; Lazarian & Vishniac 1999; McKinney & Uzdensky 2011).

We add two notes here. First, we did not discuss collimation of episodic jets in the progenitor stellar envelope. We point out that the same physics invoked in the magnetic tower model (Uzdensky & MacFadyen 2006) should apply equally to our model. Second, the scenario we propose should also work for a neutron star, not necessarily a black hole, as has been discussed.
by some previous authors (e.g., Blackman et al. 1996; Kluzniak & Ruderman 1998; Dai et al. 2006; Metzger et al. 2011).

It may be difficult to differentiate this model from other GRB central engine models using GRB observational data. Due to their large distances and small scales, it is essentially impossible to witness ejection of magnetic blobs from GRBs. Nonetheless, observing episodic blobs may be possible for nearby AGNs and X-ray binaries with future high spatial and temporal resolution observations. These observations may be used to verify this generic episodic central engine model.

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