What Can We Learn about GRB from the Variability Timescale Related Correlations?

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Abstract

Recently, two empirical correlations related to the minimum variability timescale (MTS) of the light curves are discovered in gamma-ray bursts (GRBs). One is the anti-correlation between MTS and Lorentz factor $\Gamma$, and the other is the anti-correlation between the MTS and gamma-ray luminosity $L_\gamma$. Both of the two correlations might be used to explore the activity of the central engine of GRBs. In this paper, we try to understand these empirical correlations by combining two popular black hole central engine models (namely, the Blandford & Znajek mechanism (BZ) and the neutrino-dominated accretion flow (NDAF)). By taking the MTS as the timescale of viscous instability of the NDAF, we find that these correlations favor the scenario in which the jet is driven by the BZ mechanism.

Key words: accretion, accretion disks – gamma-ray burst: general – magnetic fields – stars: black holes

1. Introduction

The mechanism for launching a relativistic jet from the gamma-ray burst (GRB) central engine is still unclear. The leading GRB central engine model involves a stellar mass black hole (BH) surrounded by a hyperaccreting disk. An alternative model involves a rapidly spinning, strongly magnetized neutron star (also known as “millisecond magnetar,” e.g., Usov 1992; Dai & Lu 1998; Wheeler et al. 2000; Lü & Zhang 2014). In the BH scenario, there are two main energy reservoirs that provide the jet power: the gravitational energy in the neutrino-dominated accretion flow (NDAF) that is carried by neutrinos and antineutrinos, which annihilate and power a bipolar outflow (Popham et al. 1999; Narayan et al. 2001; Di Matteo et al. 2002; Kohri & Mineshige 2002; Gu et al. 2006; Chen & Beloborodov 2007; Janiuk et al. 2007; Liu et al. 2007, 2015; Lei et al. 2008, 2009, 2013; Yi et al. 2017); and the spin energy of the BH, which can be tapped by a magnetic field that is connecting to a remote astrophysical load through the Blandford & Znajek mechanism (BZ, Blandford & Znajek 1977). It is hard to clearly distinguish between these different models because there isn’t radiation (except for gravitational waves and neutrinos) that reaches the observer directly from the central engine. However, the empirical correlations of some observational variables (e.g., the Lorentz-factor–isotropic-luminosity/energy correlations (Liang et al. 2010; Ghirlanda et al. 2012; Lü et al. 2012) and the Lorentz-factor–beaming-corrected-energy/luminosity correlations (Yi et al. 2017)) may place constraints on GRB central engine models.

The prompt emission of GRBs can vary extremely. There are two empirical correlations involved with the minimum variability timescale $\tau_{\text{MTS}}$. One is the anti-correlation between $\tau_{\text{MTS}}$ and the bulk Lorentz factor $\Gamma$. The other is the anti-correlation between $\tau_{\text{MTS}}$ and the isotropic gamma-ray luminosity $L_\gamma$. By adopting the MTS measurements determined by the structure-function method (Golkhou & Butler 2014) and the Lorentz factors estimated from the conventional method takes the peak of the early afterglow lightcurve as the deceleration time of the external forward shock (e.g., Lü et al. 2012; Liang et al. 2015). Wu et al. (2016b) confirmed these correlations and extended them to the blazars, with resulted in $\tau_{\text{MTS}} \propto \Gamma^{-4.8 \pm 1.5}$ and $\tau_{\text{MTS}} \propto L_\gamma^{-1.0 \pm 0.1}$. All these parameters ($\tau_{\text{MTS}}$, $\Gamma$, and $L_\gamma$) are closely related to the central engine. Therefore, the correlations here are expected to shed light on the physics of jet acceleration and accreting activity.

The origin of the MTS is debatable. One scenario proposes that the variation timescale hints the dissipation process related to the turbulence or magnetic reconnection in the jet itself (Narayan & Kumar 2009). The other possibility is that the variation timescale reflects the time intervals of the internal shocks modulated by the central engine activity (Kobayashi et al. 1997; Ramirez-Ruiz & Fenimore 2000). A possible process causing the temporal variation is the instability occurred in the disk. The thermal instability, for example, appearing in the inner region of the geometrically thin and optically thick accretion disk (Narayan & Piran 1997; Grupe et al. 1997; Ramirez-Ruiz & Fenimore 2000). A possible process causing the temporal variation is the instability occurred in the disk. The thermal instability, for example, appearing in the inner region of the geometrically thin and optically thick accretion disk, has been used to explain the “heartbeat”-like oscillations observed in some XRBs and AGNs (Grupe et al. 2015; Wu et al. 2016a). The NDAF has long been thought to be stable under parameters of interest (Di Matteo et al. 2002). However, based on the detailed treatment of the chemical equilibrium in the gas species, Janiuk et al. (2007) argued that the NDAF can be viscously and thermally unstable at an extremely high accretion rate. Moreover, Lei et al. (2009) proposed that the NDAF could be viscously unstable at a more moderate accretion rate when considering certain magnetic mechanism such as the magnetic coupling between the plumping region and the disk. More recently, Xie et al. (2016) argued that a none-zero inner boundary torque should be considered. This revised NDAF was found to be viscously unstable.

In this paper, we attribute the temporal variability MTS of the GRBs to the viscous instability of the NDAF. Meanwhile, on the basis of the investigation of the Lorentz factor of the GRB jet separately driven by BZ and NDAF models, we can infer that which jet production process is possible by checking whether it can reproduce these two MTS-related correlations.

This paper is organized as follows. In Section 2, we derive the MTS from the viscous timescale of the unstable NDAF, and then calculate the Lorentz factor of the jet driven by NDAF and BZ mechanisms, respectively. The theoretical predictions for the MTS–$\Gamma$ as well as MTS–$L_\gamma$ relation are compared with the empirical results. In Section 3, we summarize our conclusion.
and discuss the implications of the empirical MTS-related correlations in the GRB central engine.

2. The BH Central Engine Model and MTS–Γ Correlations

A hyperaccretion stellar BH system is the prevailing GRB central engine model. We summarize the main features of this model and focus on the timescale of the viscous instability in the NDAF, as well as the baryon loading and power of the jet. We equate MTS to the timescale of the instability. The initial Lorentz factors and the total power of the jets in NDAF and BZ mechanisms are calculated in the same way as they are in Lei et al. (2013).

2.1. MTS and the Viscous Instability of NDAF

According to previous works (Lei et al. 2009; Xie et al. 2016), the viscous instability can occur in the inner region of the NDAF when considering the plausible extra magnetic torque. Furthermore, the whole disk can be divided into five regions successively from the outer side to the inner side: region I (radiation-pressure dominated ADAF), region II (gas-pressure dominated, transparent NDAF), region III (gas-pressure dominated, opaque NDAF), region IV (radiation-pressure dominated, opaque NDAF), and region V (radiation-pressure dominated ADAF).

First, we’d like to estimate the location of the unstable region from the fact that the unstable region (region IV) is distinguished from its upper stream (region III) when the radiation pressure is important. The fraction of the radiation pressure in region III is expressed as

\[ P_{\text{rad}}/P = 6.6 \times 10^2 A^{-1} B^{-3/4} C^{-1/4} \tilde{\alpha}^{-1/2} m^{-5/2} \tilde{m}^2 R^{-15/4}, \]

(1)

where \( P \) is the total pressure, \( m \equiv M/\dot{M}_0 \) is the normalized BH mass, and \( \dot{m} \equiv \dot{M}/\dot{M}_0 \) s\(^{-1} \) is the disk accretion rate in unit of \( \dot{M}_0 \) s\(^{-1} \). \( A, B, C, \) and \( D \) are the relativistic correction factors for a disk around a Kerr BH (Riffert & Herold 1995), i.e.,

\[ A = 1 - 2(r/r_g)^{-1} + a^2(r/r_g)^{-2}, \quad (2a) \]

\[ B = 1 - 3(r/r_g)^{-1} + 2a(r/r_g)^{-3/2}, \quad (2b) \]

\[ C = 1 - 4a(r/r_g)^{-3/2} + 3a^2(r/r_g)^{-2}, \quad (2c) \]

\[ D = \int_{\ln \dot{m}/\dot{m}_0}^{\infty} \frac{x^2 - 6x + 8a_x x^{1/2} - 3a_x^2}{2\sqrt{R^x(x^2 - 3x + 2a_x x^{1/2})}} \, dx, \quad (2d) \]

where \( r_g \equiv GM/c^2 \) denotes the gravitational radius, and \( r_{ms} \) is the radius of the innermost stable circular orbit (ISCO). The symbol \( \mathcal{D} \) in Equation (1) is defined as \( \mathcal{D} \equiv D + \eta \chi_{ms} \ln \dot{m}/r_g^2 \Omega_x \), in which \( \eta \) is introduced to parameterize the magnitude of the extra magnetic torque. \( L_{ms} = 2GM(3\chi_{ms} - 2a_x)/\sqrt{3}c\chi_{ms} \) is the specific angular momentum of a particle in the disk, and \( \chi_{ms} = \sqrt{c^2r_{ms}/GM} \). According to Xie et al. (2016), the NDAF becomes unstable at the inner region when a significant nonzero boundary torque is applied. In the unstable region, the coefficients in Equation (1) can be approximated as

\[ A \approx 1, \quad B \approx 1, \quad C \approx \tilde{\alpha}^{-1}, \quad D \approx \tilde{m}^2 R^{-5.9}. \]

(3)

For the unstable region, the \( \tilde{m} - \Sigma \) profile presents an S-shaped curve (where \( \Sigma \) is the surface density of the disk), and the flow oscillates repeatedly between the “high state” (high \( \tilde{m}_{cr,b} < \tilde{m} < \tilde{m}_{high} \), radiation pressure, and advective cooling dominate) and the “low state” (low \( \tilde{m}_{low} < \tilde{m} < \tilde{m}_{cr,l} \), gas pressure, and neutrino cooling dominate), which corresponds to the upper branches and the lower branches in Figure 1 through a series of limit cycles one after another. Note that the “high state,” “unstable state,” and “low state” have the solutions of region V, region IV, and region III, respectively (refer to Xie et al. (2016) for further details). The expressions of the surface density of the disk for these branches are listed as follows.

\[ \Sigma_{HS} = 1.3 \times 10^{57} A^{-1} B^{1/2} C^{1/2} \tilde{\alpha}^{-2} \tilde{m}^{-1} \times \tilde{m}^{5/2} \tilde{R}_{ms}^{5/2} \text{ g cm}^{-2}. \]

(5a)

\[ \Sigma_{US} = 1.0 \times 10^{55} A^{-1} B^{1/2} C^{1/2} \tilde{\alpha}^{-1} \tilde{m}^{5/3} \times \tilde{m}^{1/3} \tilde{R}_{ms}^{5/2} \text{ g cm}^{-2}. \]

(5b)

\[ \Sigma_{LS} = 2.4 \times 10^{57} A^{-4/3} B^{2/3} C^{1/3} \tilde{\alpha}^{-2/3} \tilde{m}^{-1/3} \times \tilde{m}^{5/2} \text{ g cm}^{-2}. \]

(5c)

In each cycle, the accretion rate will decrease from the high accretion rate \( \tilde{m}_{high} \) to the higher critical value \( \tilde{m}_{cr,b} \), and then suddenly drops to the low accretion rate \( \tilde{m}_{low} \). The flow then enters the “low state” and its accretion rate will gradually increase due to the mass feeding from the upper stream (i.e., region III), until it reaches the lower critical value \( \tilde{m}_{cr,l} \). It jumps back directly to “high state” with accretion rate \( \tilde{m}_{high} \).

Modulated by a series of limit cycles, the luminosity of the jet varies, which gives the temporal behavior of GRB prompt emission. From Equation (5), we get the four characteristic values of accretion rate in the S-shaped curve, i.e.,

\[ \tilde{m}_{high} = 4.7 \times 10^{-14} B^{2/3}/C^{1/3} \tilde{\alpha}^{-4} \tilde{m}_{low}^{11/12} \tilde{R}_{ms}^{9/8} \]

\[ \sim 5.8 \times 10^{-14} \tilde{m}_{low}^{11/12} \tilde{R}_{ms}^{0.009}. \]

(6a)
\[ m_{\text{cr},h} = 8.8 \times 10^{-2} A_{12}^2 B^{-1/6} \mathcal{G}^{-1/2} \alpha_1^{1/3} m_4^{4/3} R_{\text{ms}}^{3/2} \]
\[ \sim 7.6 \times 10^{-2} \eta^{-0.5} \alpha_1^{1/3} m_4^{4/3} R_{\text{ms}}^{1.8} \] (6b)
\[ m_{\text{cr},l} = 1.6 \times 10^{-2} A_{12}^2 B^{3/8} C_1^{1/2} \mathcal{G}^{-1/2} \alpha_1^{1/3} m_4^{4/3} R_{\text{ms}}^{15/8} \]
\[ \sim 1.1 \times 10^{-2} \eta^{-1/2} \alpha_1^{1/3} m_4^{5/4} R_{\text{ms}}^{2.6} \] (6c)
\[ m_{\text{low}} = 1.1 \times 10^{-4} B^2 C_1^{1/2} \mathcal{G}^{-5/2} m_4 R_{\text{ms}}^3 \]
\[ \sim 5.5 \times 10^{-5} \eta^{-2.5} m_4 R_{\text{ms}}^{4.6}. \] (6d)

Now, we evaluate the variability timescale due to each limit cycle, which is taken as the minimum temporal variability (MTS). The timescale of the “unstable state” can be ignored, and state transitions are assumed to take place immediately once the accretion rate reaches the critical values. The duration of each limit cycle depends primarily on the evolution of the mass accretion rate, namely the viscous timescale in the “low state” and the “high state.”

\[ t_{\text{vis}}^{\text{HS}} = 8.9 \times 10^{-6} C \mathcal{G}^{-1} \alpha_1^{1/3} m_4 R_{\text{vis}}^{3/2} \text{ s}, \] (7a)
\[ t_{\text{vis}}^{\text{LS}} = 1.7 \times 10^{-5} A_{12}^2 B^{-5/6} C_5^{5/6} \mathcal{G}^{-2/3} \alpha_1^{-2/3} m_4^2 \times \dot{m}^{-2/3} R_{\text{vis}}^2 \text{ s}, \] (7b)

where the viscous timescale is estimated by \( t_{\text{vis}} \sim r^2 / \nu = (\alpha R_\mathcal{K})^{-1}(h/r)^{-2} \). The Keplerian angular velocity is \( \Omega_k \equiv \sqrt{GM / r^3} \).

Generally we have \( t_{\text{vis}}^{\text{HS}} > t_{\text{vis}}^{\text{LS}} \), i.e., MTS is dominated by the timescale in “low state” of the limit cycle MTS \( t_{\text{vis}} \propto m_{\text{low}}^{-2/3} \).

Substituting Equation (4) into Equation (6d), we have the relation between the low state accretion rate \( \dot{m}_{\text{low}} \) of the unstable region and the disk accretion rate \( \dot{m} \) as
\[ \dot{m}_{\text{low}} \sim 4.5 \times 10^{-2} \eta^{-0.94} \alpha_1^{-0.37} m^{-0.93} m_4^{1.56}. \] (8)

Combining Equations (7b) and (8), we approximately have
\[ \text{MTS} \sim 1.3 \times 10^{-4} A_{12}^2 B^{-5/6} C_5^{5/6} \mathcal{G}^{-2/3} \eta^{0.63} \alpha_1^{-0.42} \times m_4^{2.62} \dot{m}^{-1.04} R_{\text{vis}}^2 \text{ s}. \] (9)

\( t_{\text{vis}} \) under the typical parameters is about 100 ms, which is in the same order of magnitude as the observations (e.g., MacLachlan et al. 2013; Golkhou & Butler 2014).

2.2. The Jet Driven by Neutrino Annihilation

The neutrino annihilation (\( \nu \bar{\nu} \rightarrow e^+ e^- \)) process above an NDAF can launch a relativistic jet reaching the GRB luminosity. An approximate expression for the neutrino annihilation power \( \dot{E}_{\nu} \) is given by Zalamea & Beloborodov (2011) as,
\[ \dot{E}_{\nu} \approx 6.2 \times 10^{49} \left( \frac{R_{\text{ms}}}{2} \right)^{-4.8} \left( \frac{m}{3} \right)^{-3/2} \dot{m}^{-1/4} m_4 \text{ erg s}^{-1}. \] (10)

Neutrino heating via neutrino absorption on baryons (\( p + \nu_e \rightarrow n + e^+ \) and \( n + \nu_e \rightarrow p + e^- \)) in the atmosphere of an NDAF can drive a baryonic wind (e.g., Metzger et al. 2008). Since the majority of the mass lies in large radii, the main part of the wind originates from the region that is dominated by gas pressure and URCA cooling (Region II). According to Lei et al. (2013), the neutrino-heating driven baryon loading rate of the jet can be estimated as
\[ \dot{M}_{\nu} \approx 7.0 \times 10^{-7} A_{12}^{1.35} \mathcal{C}^{0.22} \dot{m}_1^{0.105} \mathcal{C}_1^{1.17} \times \left( \frac{R_{\text{ms}}}{2} \right)^{0.32} \dot{m}_1^{1.7} \left( \frac{m}{3} \right)^{-0.9} \left( \frac{\xi}{2} \right)^{0.32} M_{\odot} \text{ s}^{-1}, \] (11)

where \( \xi \equiv r / r_{\text{ms}} \) is the disk radius in terms of \( r_{\text{ms}} \), \( \varepsilon \approx (1 - E_{\text{ms}}) \) denotes the neutrino emission efficiency, \( E_{\text{ms}} = (4 \sqrt{R_{\text{ms}} - 3 a_0}) / \sqrt{3} R_{\text{ms}} \) is the specific energy at ISCO, and \( \theta_i \) is the jet opening angle.

If most neutrino annihilation energy is converted to the kinetic energy of baryons after acceleration, the ultimate Lorentz factor of the jet is determined by the dimensionless “entropy” parameter, \( \eta_0 \), i.e.,
\[ \Gamma_{\text{max}} \approx \eta_0 \equiv \dot{E}_{\nu} / \dot{\mathcal{M}}_\nu \sqrt{c^2} \]
\[ = 50 A_{12}^{-1.13} B^{-3.5} \mathcal{C}_1^{-0.22} \dot{m}_1^{0.57} \mathcal{C}_1^{1.7} \left( \frac{\xi}{2} \right)^{-0.32} \times \left( \frac{R_{\text{ms}}}{2} \right)^{-5.12} \left( \frac{m}{3} \right)^{-0.6} m_4^{0.58}. \] (12)

Inspecting Equations (9), (10) and (12), one finds MTS, \( \dot{E}_{\nu} \) and \( \eta_0 \) are functions of \( m \) and \( a \). Usually the BH mass in GRBs \( m \) varies in a narrow range of (2.5, 10) (e.g., Popham et al. 1999). Considering the hyperaccreting process during the prompt emission phase, the BH is quickly spun up so that the \( a_0 \)-dependence is not significant (Lei et al. 2013). Therefore, the \( m \)-dependence may be the key to define the MTS–\( \eta_0 \) correlation (MTS \( \propto \eta_0^{1.85} \)) and MTS–\( \dot{E}_{\nu} \) correlation (MTS \( \propto \dot{E}_{\nu}^{-0.5} \)). Comparing with the data, the predicted indices 1.8 (0.5) in MTS–\( \eta_0 \) correlation (MTS–\( \dot{E}_{\nu} \) correlation) are significantly smaller than the observational value 4.8 ± 1.5 (1.0 ± 0.1).

However, the \( m \)-dependence will add to the scatter in the correlations. In Figure 2, we plot MTS versus \( \Gamma \) (left), and MTS versus \( L \) (right) with dotted lines for \( m = 2.5 \) (bottom) and 10 (top). For each \( m \) case, we change \( m \) in a wide range of values and keep other parameters to fixed values, see Figure 2. If we allow \( m \) to randomly vary in the range of (2.5, 10), the simulated GRBs should be scattered in the light shaded region between the dotted lines. From Figure 2, we find that the predictions with the NDAF model is inconsistent with the data.

2.3. The Jet Driven by BZ Mechanism

We now consider the BZ scenario for the jet production. The magnetic field in the BZ mechanism is established and supported by the magnetized NDAF around the BH. The baryons in the jet are also loaded from the neutrino-driven wind. Unlike the NDAF scenario, the existence of the magnetic barrier will significantly suppress the baryon loading from the disk (e.g., Li 2000). Therefore, a baryon-poor jet will be built in the BZ model.

The BZ power from a BH with mass \( M \), and spin \( a \) is (e.g., Lee et al. 2000; Lei et al. 2013; Wu et al. 2013)
\[ \dot{E}_{\text{BZ}} = 9.3 \times 10^{52} a_0^2 a_1 (X(a)) \text{ erg s}^{-1}, \] (13)

where \( X(a) = F(a)/[1 + (1 - a_0^2)^2] \), \( F(a) = [1 + a_0^2]/a_0^2 \)
\[ \times [(q + 1)/q] \arctan(q - 1), \text{ and } q = a_0/\sqrt{1 - a_0^2}. \]
As for the baryon loading process, while considering the blocking effect of the magnetic barrier on the protons, Lei et al.
we separately take slow acceleration phase in a hybrid out and Turbulence through the Internal-Collision induced MAgnetic Reconnection magnetic dissipation, such as the magnetic energy relaxion with the speci…

we take \( \alpha_1 = 0.8, \alpha_2 = 0.1, \eta_1 = 1, \) and \( R = 2.5 \) denoting region IV, and 20 denoting region II, \( \eta_2 = 0.03, \eta_3 = 0.01, \) here \( \eta_2 \equiv L_c (1 - \cos \theta_i) / E_{\text{BZ}}. \) The top (bottom) dotted lines correspond to the cases with \( m = 10 \) (2.5). For BZ mechanism, we take \( \alpha_1 = 0.2, \alpha_2 = 0.2, \eta_1 = 1, \) and \( R = 6 \) denoting region IV, and 20 denoting region II, \( r_c = 5 \times 10^{11} \text{cm}, f_p = 0.01, \theta_1 = 0.4, \theta_B = 0.011, \eta_1 = 0.1, \) here \( \eta_1 \equiv L_c (1 - \cos \theta_i) / E_{\text{BZ}}. \) The top (bottom) dotted lines correspond to the cases with \( m = 10 \) (2.5).

(2013) suggested that the neutron drift rate into the jet as

\[
M_{\text{J,BZ}} \simeq 3.5 \times 10^{-7} A^{23/30} B^{-33/40} C^{7/120} \dot{R}_{p-1}^{-1/2} \dot{\theta}_{B-1}^{-1} \theta_{B-2}^{-1} \times \alpha^{23/60} \epsilon^{5/6} \dot{m}^{3/2} \left( \frac{R_{\text{ms}}}{2} \right)^{1/120} \left( \frac{m}{3} \right)^{11/20} \dot{r}_{c,11}^{1/2} \left( \frac{\xi}{2} \right)^{1/120} M_\odot \text{s}^{-1}.
\]

where \( f_p \) denotes the fraction of the protons in the wind, \( r_c \) is the distance from the BH in the jet direction, and \( \dot{\theta}_B \) is introduced here to reflect the fact that only the protons with small ejected angles (\( \leq \theta_B \)), with respect to the field lines, can come into the disk atmosphere.

For such a magnetized central engine, the maximum available energy per baryon in the jet can be evaluated as

\[
\dot{\mu}_0 \simeq \frac{E_{\text{BZ}}}{M_{\text{J,BZ}} c^2} = 1.5 \times 10^5 A^{-23/30} B^{33/40} C^{-7/120} \dot{R}_{p-1}^{1/2} \times \dot{\theta}_{B-1}^{-1} \theta_{B-2}^{-1} \alpha^{23/60} \epsilon^{5/6} \dot{m}^{3/2} \left( \frac{R_{\text{ms}}}{2} \right)^{-1/120} \left( \frac{m}{3} \right)^{-11/20} \dot{r}_{c,11}^{-1/2} \left( \frac{\xi}{2} \right)^{-1/120} \left( \frac{\dot{\mu}}{\dot{\mu}_0} \right)^{1/2} \dot{\mu}_0^{1/6}.
\]

(15)

Since the acceleration process of the jet suffer from uncertainties, the terminating Lorentz factor of the jet generally satisfies

\[
\Gamma_{\min} < \Gamma < \Gamma_{\max}.
\]

with the specific value depending on the detailed process of magnetic dissipation, such as the magnetic energy relaxion through the Internal-Collision induced MAgnetic Reconnection and Turbulence (ICMART, Zhang & Yan 2011) or the shearing interaction at the inner/outer layer interface in the possible two-componet jet (e.g., Wang et al. 2014). Following Lei et al. (2013), we separately take \( \Gamma_{\min} = \max \left( \frac{\dot{\mu}_0}{\dot{\mu}_0} / N_0 \right) \) \( (n_0 = E_{\nu,0} / (M_{\text{J,BZ}} c^2)) \) and \( \Gamma_{\max} = \dot{\mu}_0 \), which corresponds to the start and the end of the slow acceleration phase in a hybrid outflow (a detailed discussion of the acceleration dynamics of an arbitrarily magnetized relativistic or hybrid jet is referred to Gao & Zhang 2015).

Based on Equations (9), (13), and (15), we find the MTS–\( \mu_0 \) correlation (MTS \( \propto \mu_0^{-6} \)) and MTS–\( E_{\text{BZ}} \) correlation (MTS \( \propto E_{\text{BZ}}^{-1.04} \)). With the consideration to the spin-up process due to accretion and the spin-down process due to the BH process, the BH spin parameter always evolves to an equilibrium value (Lei et al. 2005) so that the \( \alpha_1 \)-dependence essentially does not enter the problem. As shown in Figure 2, the \( m \)-dependence will add to the scatter in the correlations. The top (bottom) dashed lines represent the MTS–\( \Gamma \) (left) and MTS–\( L_c \) correlations (right) for \( m = 10 \) (2.5). GRBs with BH mass in the range of (2.5, 10) should be scattered in the empirical region between the dashed lines. We find that the predictions from the BZ mechanism are more consistent with the empirical correlations than that of the \( \nu/\Delta \)—annihilation mechanism.

3. Conclusions and Discussions

In this paper, we compare the theoretical predictions from neutrino annihilation and BZ processes with the empirical MTS–\( \Gamma \) and MTS–\( L_c \) correlations. We find that both empirical correlations favor the BZ scenario. These correlations may provide us with a clue about the GRB central engine (i.e., that a good fraction of GRBs may be driven by a hyperaccretion system consisting of a stellar BH and a surrounding NDAF). Furthermore, the jet power of GRBs may be supplied by the BH rotating energy through the BZ mechanism. This result is consistent with the implication from the \( \Gamma - E_{\nu,0} \) correlation proposed by Liang et al. (2015), which suggested that the GRB jet might be Poynting-flux dominated. In addition, our result is also in agreement with Yi et al. (2017), who investigated the BH central engines from the correlation between \( \Gamma \) and the beaming-corrected luminosity. Even in the BH model, the disk is still an NDAF. The latter plays two important roles in the BZ scenario. First, the hyperaccretion is necessary to maintain the strong magnetic field for the BZ process. Second, the neutrino-heating wind in the surface of the NDAF acts as a significant contribution to the jet baryon loading.
In this work, we adopt $m$ as the primary correlating variable. As mentioned in Section 2, the dependence of the BH spin and mass may also enter the correlations. However, the BH spin parameter will quickly evolve to the maximum value (for the NDAF model) or equilibrium value (for the BZ scenario) so that the $a$ dependence is not important. On the other hand, the empirical correlations were obtained from the average $\Gamma$ and $L_\gamma$. In order to compare with the data, one needs to calculate the average $\Gamma$, $E_{\text{ne}}$, and $E_{\text{BZ}}$, which smears the $a$ dependence (Lei et al. 2013). The BH mass $m$ is generally believed to vary from 2.5 to 10 in GRBs (e.g., Popham et al. 1999). Therefore, with $m$-dependence only, the theoretical models can hardly reproduce the wide range of observed $L_\gamma$ and MTS. As shown in Figure 2, it adds to the scatter in the correlations, but does not change our main conclusion. A detailed study of the effects due to the dependences on the model parameters (e.g., $m$, $\alpha$, $\theta_p$, $\eta_p$) will be addressed in future work.

It is interesting that the BZ mechanism and the NDAF process have mutual influence on each other. On one hand, the magnetic field in the BZ mechanism is supported (or even established, e.g., Cao et al. 2014) by the NDAF. On the other hand, such a magnetic field has strong effects in suppressing the baron-loading from the neutrino-driven wind. A further discussion for such a co-existence relation is beyond the scope of this work.

In our model, we define MTS as the viscous timescale due to viscous instability. There are other possible disk origins for the temporal variation. Cao et al. (2014) investigated the capability of the NDAF to drag the large-scaled magnetic field inward to the BH, and the result shows that the competition between the inherent diffusion and the accretion driven advection of the magnetic field leads to a oscillating accretion as well as an episodic jet, the oscillation timescale can be about one second which is comparable with the observed variation in the soft extended emission of short GRBs. An alternative mechanism involving the temporal variation is referred to the disk’s inertial-acoustic oscillation which has been extensively discussed (e.g., Kato 1978; Wagoner et al. 2001; Yu & Lai 2015). For aperoic variation, some authors thought it reflects the possible propagating fluctuations in the disk (e.g., Lyubarskii 1997; Lin et al. 2016). It is worthwhile to check the feasibility of these different models in interpreting the MTS related correlations in the future.

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