Swift J164449.3+573451: A Plunging Event with a Poynting-Flux-Dominated Outflow

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ABSTRACT

Swift J164449+573451 is a peculiar outburst which is most likely powered by the tidal disruption of a star by a massive black hole. Within the tidal disruption scenario, we show that the periastron distance is considerably smaller than the disruption radius and the outflow should be launched mainly via magnetic activities (e.g., the Blandford–Znajek process), otherwise the observed long-lasting X-ray afterglow emission satisfying the relation $L_X \propto M$ cannot be reproduced, where $L_X$ is the X-ray luminosity and $M$ is the accretion rate. We also suggest that $L_X \propto M$ may hold in the quick decline phase of gamma-ray bursts.

Key words: accretion, accretion disks – black hole physics – gamma-ray burst: general – radiation mechanisms: non-thermal

1. INTRODUCTION

Swift J164449.3+573451 (Sw J1644+57) triggered the Swift Burst Alert Telescope on 2011 March 28 (Cummings et al. 2011; Burrows et al. 2011). As revealed by late optical observations, this transient lay at the center of a galaxy at a redshift $z = 0.3534$ (Levan et al. 2011a, 2011b). In the first few weeks, the average isotropic luminosity in the 0.3–10 keV band was about $10^{37} – 10^{38}$ erg s$^{-1}$. Several months later, it was still about a few $\times 10^{45}$ ergs$^{-1}$, well above the Eddington limit. The X-ray emission declined as $L_X \propto t^{-5/3}$ during the time interval from $10^3$ s to $10^5$ s after the trigger, while the subsequent decline can be approximated as $L_X \propto t^{-4/3}$ (Levan et al. 2011a, 2011b; Bloom et al. 2011; Cannizzo et al. 2011).

The super-long duration (>8 months) of the X-ray activities essentially rules out gamma-ray burst (GRB) models (Shao et al. 2011). Instead it strongly favors the model of tidal disruption of a (giant) star by a massive black hole (BH; Bloom et al. 2011; Burrows et al. 2011; Levan et al. 2011b; Cannizzo et al. 2011; Shao et al. 2011; Wang & Cheng 2012). As for the central BH, it is impossible to measure the mass directly. Indirect constraints indicate that the mass should be in the range $10^5 M_\odot – 10^7 M_\odot$ (detailed analysis can be seen in Cannizzo et al. 2011). In the tidal disruption scenario, if a star passes within the disruption radius $R_T \approx R_s(M_{BH}/M_*)^{1/3}$, the BH’s tidal gravity exceeds the star’s self-gravity and consequently the star is disrupted, where $R_s$ ($M_*$) is the radius (mass) of the disrupted star and $M_{BH}$ is the mass of the BH. After about hours to weeks, part of the remnants remain on bound and will return to the pericenter of the orbit where the material starts to be accreted inward, releasing a flare of energy (Rees 1988, 1990; Phinney 1989; Strubbe & Quataert 2009).

Theoretical calculations of the tidal disruption events suggest that the immediately accreted unbinding gas falls back to the pericenter at a rate of $M_{fb} \propto t^{-5/3}$ (Rees 1988; Phinney 1989), and the subsequent disk accretion follows a rate of $M_{fb} \propto t^{-4/3}$ if the disk is thick (Cannizzo & Gehrels 2009). So for $t > 10^5$ s, the X-ray emission luminosity of Sw J1644+57 is proportional to the accretion rate, i.e., $L_X \propto M_{fb}$. In this work, we pay special attention on the physical implications of such a relation. The physical parameters of the disrupted star are also investigated.

2. THE PHYSICAL PARAMETERS OF THE DISRUPTED STAR

There is an express limit derived in the Newtonian for the timescale of the return of the most bound stellar material to the pericenter (Rees 1988; Phinney 1989):

$$t_{fb} = 2\pi G M_{BH}(2\Delta E)^{-3/2}$$

$$= 0.048 \text{yr} \left(\frac{M_{BH}}{5 \times 10^6 M_\odot}\right)^{1/2} \left(\frac{M_*}{M_\odot}\right)^{-1} \left(\frac{R_s}{R_\odot}\right)^{3/2} \mu^3,$$

(1)

where $\Delta E = k(GM_s/R_s)(M_{BH}/M_*)^{1/3}$, the dimensionless coefficient $k$ depends on the spin-up state of the star. If the star is spun up to the break-up spin angular velocity, we have $k \approx 3$. If the spin-up effect is negligible then we have $k \approx 1$ (Rees 1988; Ayal et al. 2000). The dimensionless coefficient $\mu \equiv R_p/R_T$ is taken to be a free parameter in the following discussion, where $R_p$ is the periastron distance of the star.

The star could be spun up via tidal interaction. In linear perturbation theory, the spin-up angular velocity is given by (Press & Teukolsky 1977; Alexander & Kumar 2001; Alexander & Livio 2001; Li et al. 2002)

$$\frac{\omega_s}{\omega_p} \approx \frac{T_2(\mu^{3/2})}{2\mu^3},$$

(2)

where $\omega_s \equiv v_p/r_p$ is the orbit angular velocity of the star at the pericenter, $I$ is the stellar momentum of inertia in units of $M_* R_s^2$; and $T_2$ is the second tidal coupling coefficient and depends on the structure of the star and the eccentricity of the orbit. For an $n = 1.5$ polytrope star of mass $0.76 M_\odot$ and radius $0.75 R_\odot$, Alexander & Kumar (2001) found $I \approx 0.21$ and $T_2(1) \approx 0.36$, corresponding to $\omega_s/\omega_p \approx 0.86$ for $\mu = 1$. Furthermore, they showed that the numerical simulations including nonlinear effects led to a larger energy transfer from the orbit to the star and a larger spin-up than that predicted by linear theory. Therefore, we take $k = 3$ as the fiducial value in our analysis. For completeness we also present the results in the non-spinning case (i.e., $k = 1$).
The bound material returns to pericenter at the rate (Phinney 1989)

\[ M_{fb} = \frac{2\Delta M}{\dot{M}_{fb}} \left( \frac{t - t_f}{t_{fb}} \right)^{-5/3} = 5.6 \times 10^{23} \text{ g s}^{-1} \left( \frac{f}{0.1} \right) \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{1/3} \left( \frac{M_*}{M_\odot} \right)^{1/3} \left( \frac{R_*}{R_\odot} \right) \left( \frac{t - t_f}{1 \text{ yr}} \right)^{-5/3} \mu^2, \]  

(3)

where \( t_f \) is the time of initial tidal disruption, \( \Delta M \) is the mass that falls back to pericenter, and the dimensionless factor \( f \) is defined as \( f = \Delta M / M_* \).

When the accretion rate of the fall-back material is highly super-Eddington, only a fraction \((1 - f_{out})\) of such material forms a disk and can be accreted all the way down to the central BH, i.e., \( M = (1 - f_{out}) M_{fb} \). The remaining part will instead leave the system, undergoing a strong radiation pressure. Strubbe & Quataert (2009) took a constant \( f_{out} = 0.1 \). However, numerical simulation indicates that the parameter \( f_{out} \) is a growing function of \( M_{fb}/M_{Edd} \), reaching \( f_{out} \approx 0.7 \) for \( M_{fb}/M_{Edd} = 20 \) (Dotan & Shaviv 2010). Considering the observed X-ray luminosity \( L_X \sim 10^{37} - 10^{38} \text{ erg s}^{-1} \) during the first \( 10^5 \) s after the trigger (Burrows et al. 2011), even for a radiation efficiency as high as 0.1, we need an accretion rate \( M = 10L_X/c^2 \sim 5 \times 10^{-7} - 2 \times 10^{-5} M_\odot \). The Eddington luminosity can be scaled as \( L_{Edd} \approx 6.25 \times 10^{38} M_6 \text{ erg s}^{-1} \), so the Eddington rate \( M_{Edd} \equiv L_{Edd}/c^2 = 3.5 \times 10^{-7} M_6 \text{ s}^{-1} \). We then have \( M/M_{Edd} \sim 1 - 60 \). We define a free dimensionless parameter \( \xi = f(1 - f_{out}) \), the fraction of the material that is actually accreted onto the central BH.

Assuming that the jet efficiency is \( \epsilon \) during the stellar debris fall-back accretion, the intrinsic jet luminosity then can be given by

\[ L_j = \epsilon \dot{M} c^2 \approx 2.5 \times 10^{42} \text{ erg s}^{-1} \left( \frac{\epsilon}{0.01} \right) \left( \frac{\xi}{0.05} \right) \times \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{1/3} \left( \frac{M_*}{M_\odot} \right)^{1/3} \left( \frac{R_*}{R_\odot} \right) \left( \frac{t - t_f}{1 \text{ yr}} \right)^{-5/3} \mu^2, \]  

(4)

where the efficiency is normalized to \( \epsilon \sim 0.01 \) (in Section 3.3, we will show that the jet should be launched mainly via magnetic activities (e.g., the Blandford–Znajek (BZ) effect) and the efficiency is about 0.01). As shown below, the conclusion drawn in this section is independent of the value of \( \epsilon \).

When most of bound debris falls back to the pericenter, the jet luminosity peaks at \( t - t_f = t_{fb} \) and can be estimated as

\[ L_{j, peak} \approx 3.9 \times 10^{44} \text{ erg s}^{-1} \left( \frac{\epsilon}{0.01} \right) \left( \frac{\xi}{0.05} \right) \times \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{-1/2} \left( \frac{M_*}{M_\odot} \right)^{2} \left( \frac{R_*}{R_\odot} \right)^{-3/2} \mu^{-3}. \]  

(5)

The observed maximal X-ray luminosity is \( \sim 4 \times 10^{48} \text{ erg s}^{-1} \). For a collimated emitting region with a half-opening angle \( \theta_* \), we have the constraint \( L_{X, peak} \geq 2 \times 10^{48} (\theta_*/0.1)^2 \text{ erg s}^{-1} \). Assuming that most of the radiated energy is in the X-ray band during the time interval, we have

\[ \left( \frac{\epsilon}{0.01} \right) \left( \frac{\theta_*}{0.1} \right)^{-2} \left( \frac{\xi}{0.05} \right) \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{-1/2} \times \left( \frac{M_*}{M_\odot} \right)^{2} \left( \frac{R_*}{R_\odot} \right)^{-3/2} \mu^{-3} \geq 51. \]  

(6)

The observed X-ray fluence \( S_X (0 < t < 10^7 \text{ s}) \) suggests a total energy \( \Delta E_X = \int_0^t L_X (t) dt / 2 \sim 1 \times 10^{51} (\theta_*/0.1)^2 \text{ erg} \). The total mass of the accreted material is thus

\[ M_* = \frac{\Delta E_X}{\xi \epsilon c^2} \approx 1.1 M_\odot \left( \frac{\epsilon}{0.01} \right)^{-1} \left( \frac{\xi}{0.05} \right)^{-1} \left( \frac{\theta_*}{0.1} \right)^2. \]  

(7)

The approximated mass–radius relationship can be scaled as \( (R_* / R_\odot) = (M_* / M_\odot)^{3/4} \). For the main-sequence stars, we have \( \eta \approx 0.8 \) for \( 0.1 M_\odot < M_* < 1 M_\odot \) and \( \eta \approx 0.6 \) for \( 1 M_\odot < M_* < 10 M_\odot \) (Kippenhahn & Weigert 1994). With Equations (6) and (7), we obtain

\[ \left( \frac{M_*}{M_\odot} \right)^{1 - 3\eta/2} \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{-1/2} \geq 4.64, \]  

(8)

which then yields

\[ \mu \leq 0.36 \left( \frac{M_{BH}}{5 \times 10^6 M_\odot} \right)^{-1/6} \left( \frac{M_*}{M_\odot} \right)^{1/3 - \eta/2}. \]  

(9)

Interestingly, the parameter \( \mu \) is independent of \( \epsilon, \xi, \) and \( \theta_* \). Its dependence on both the stellar mass and the BH mass is also very weak. The above analysis is under the condition that the star is spun-up to the break-up spin angular velocity, i.e., \( k = 3 \). For the non-spinning case \( k = 1 \), we can obtain a more stringent result \( \mu \leq 0.16 ((M_{BH})/(5 \times 10^6 M_\odot))^{-1/6} (M_* / M_\odot)^{1/3 - \eta/2} \). Therefore, we conclude that the periastron distance is likely well within the tidal disruption radius (i.e., it is a plunging event), in agreement with Cannizzo et al. (2011).

Based on the observation data, the peak accretion rate can be derived with Equations (5) and (7),

\[ M_{peak} = \frac{L_{j, peak}}{\epsilon c^2} = \frac{\xi M_* L_{X, peak}}{\Delta E_X} \approx 4.4 \times 10^{-5} M_\odot s^{-1} \left( \frac{\xi}{0.05} \right) \left( \frac{M_*}{1.1 M_\odot} \right). \]  

(10)

The accretion rate \( \dot{M} \) can be scaled as \( \dot{M} = M_{peak} (t - t_f) / 1 \text{ yr} / 5^{3/5} \) during the fall-back accretion process.

In this plunging event, the disrupted star’s orbit is likely to misalign with the equatorial plane of the spinning central BH. A tilted accretion disk should be formed and the jet aligned with the disk normal vector is expected to precess (Stone & Loeb 2012; Lei et al. 2012). Saxton et al. (2012) have analyzed the X-ray timing and spectral variability of Sw J1644+57 and found periodic modulation, possibly due to jet precession.

3. THE RADIATION MECHANISM IN THE FALL-BACK PHASE

After the time \( t \sim 10^5 \) s, the observed X-ray luminosity followed the fall-back accretion rate (Levan et al. 2011a, 2011b; Bloom et al. 2011), i.e., \( L \propto \dot{M} \). Such a relationship has shed some light on the underlying physics.
3.1. Thermal X-Ray Radiation from the Disk?

While the fall-back accretion rate is super-Eddington, the stellar material returning to pericenter is so dense that it cannot radiate and cool. The gas is most likely to form an advective-dominated accretion flow accompanied by powerful outflow, which dominates the emission. Most of the radiation will be emitted from the outflow’s photosphere. When the photosphere lies inside the outflow, the photosphere’s radius and temperature can be written as (Strubbe & Quataert 2009; Rossi & Begelman 2009)

\[ R_{\text{ph}} \sim 4 f_{\text{out}} f_v^{-1} \left( \frac{M_{\text{fb}}}{M_{\text{edd}}} \right) R_{p,3}^{-1/2} R_s \]

and

\[ T_{\text{ph}} \sim 1 \times 10^5 K f_{\text{out}}^{-1/3} f_v^{-1/3} \left( \frac{M_{\text{fb}}}{M_{\text{edd}}} \right)^{-5/12} M_{s,7}^{-1/4} R_{p,3}^{-7/24}, \tag{12} \]

where \( R_s \equiv (2GM_{\text{BH}}/c^2) \) is the Schwarzchild radius and \( f_v \) is the ratio of the terminal velocity of gas with the escape velocity at a radius of \( \sim 2R_p \). However, the photons escape from the photosphere are mainly in the UV optical band and have a blackbody spectrum. The UV optical emission luminosity can be scaled as \( v L_v \sim 4\pi R_{\text{ph}}^2 v B_v (T_{\text{ph}}) \propto R_{\text{ph}}^2 T_{\text{ph}} \propto M_{\text{fb}}^{19/12} \).

These optical photons could be Compton-scattered by the relativistic electrons in the outflow. The energy of the photons getting scattered is \( h\nu_{\text{C}} \approx D^2 \gamma^2 h\nu \), where \( D = 1/\Gamma(1 - \beta \cos \theta) \) is the Doppler factor, \( \Gamma \) is the Lorentz factor of the outflow, \( \theta \) is the angle between the outflow axis and the observer’s line of sight, and \( \gamma \) is the Lorentz factor of the relativistic electrons. The energy of inverse-Compton-scattered photons can peak in the X-ray band if the parameter \( D \sim 1 \) and \( \gamma \sim 10 \). However, even in this case, the X-ray luminosity does not satisfy the relation \( L_v \propto M_{\text{fb}} \), which is inconsistent with the observational data.

3.2. Neutrino-annihilation-launched Outflow?

When the mass accretion rate is high enough, the accretion proceeds via neutrino cooling and neutrinos can carry away a significant amount of energy from the inner regions of the disk. This mechanism is used to explain the launch of at least some significant amount of energy from the inner regions of the disk. The emission from the BH may be possible through the BZ mechanism (Blandford & Znajek 1977). Such a process is based on the expectation that the differential rotation of the disk will amplify pre-existing magnetic fields until they approach equipartition with the gas kinetic energy. For a BH of mass \( M_{\text{BH}} \) and angular momentum \( J \), with a magnetic field \( B_1 \) normal to the horizon at \( R_h \), the power arising from the BZ mechanism is given by (e.g., Thorne et al. 1986)

\[ L_{\text{BZ}} = \frac{\pi}{8} \frac{\omega_F^2}{4\pi} \left( \frac{B_1^2}{2\pi} \right) R_h^2 c \left( \frac{J}{J_{\text{max}}} \right)^2, \tag{15} \]

where \( J_{\text{max}} = GM^2/c \) is the maximal angular momentum of the BH. The factor \( \omega_F^2 = \Omega_F (\Omega_F - \Omega_0)^2 \) depends on the angular velocity of field lines \( \Omega_F \) relative to that of the BH, \( \Omega_0 \). Usually we adopt \( \omega_F = 1/2 \), which maximizes the power output (Macdonald & Thorne 1982; Thorne et al. 1986). We follow the common assumption that the magnetic field in the disk will rise to some fraction of its equipartition value \( P_{\text{mag}} = (B^2/8\pi) \sim \alpha P \) in the inner disk. The pressure \( P = \rho c^2 \) is given by (Armitage & Natarajan 1999)

\[ P = \frac{\sqrt{2} M}{12\pi \alpha} (5 + 2\varepsilon)^{1/2} (GM)^{1/2} R^{-5/2}, \tag{16} \]

where \( \varepsilon \) is the parameter governing the property of the disk. For a thick disk we have \( \varepsilon < 1 \) otherwise \( \varepsilon > 1 \). In the inner region of the disk, we assume \( B_1 \approx B, R \approx R_h = GM/c^2 \). The BZ power for the case of a maximally rotating BH (i.e., \( J = J_{\text{max}} \)) can be estimated as

\[ L_{\text{BZ}} \approx 7 \times 10^{-3} (5 + 2\varepsilon)^{1/2} \frac{M_{\text{BH}}}{M_{\odot}} c^2, \tag{17} \]

corresponding to an efficiency \( \varepsilon_{\text{BZ}} = L_{\text{BZ}}/M_{\text{BH}} c^2 \sim 10^{-2} \) for the thick-disk model. In the thin-disk scenario the radiative energy can be as high as \( \sim 0.1 \). For Sw J1644+57, at the time \( t - t_\bullet = t_\bullet + 10^8 \) s, the mass accretion rate is about \( 3.3 \times 10^{-8} M_{\odot} \) s\(^{-1}\), and the observed luminosity is \( 2L_{\text{BZ}} \theta_j^{-2} \sim 5 \times 10^{37} \) erg s\(^{-1}\) (\( \theta_j / 0.1 \) )\(^{-2}\), consistent with the observation.

Our conclusion that the outflow powering the super-long X-ray emission should be launched via magnetic activities (e.g., BZ mechanism) is consistent with that of Shao et al. (2011). Lei & Zhang (2011) have also analyzed the jet launched by the BZ mechanism and then constrained the physical parameter of the central BH. One interesting finding is that the central BH should have a moderate to high spin.

In the Poynting-flux-dominated outflow, the X-ray emission could be due to the dissipation of the magnetic field (Usov 1994; Thompson 1994). There are several magnetic field dissipation models that could produce the observed emission, such as the global MHD condition breakdown model (Usov 1994), the gradual magnetic reconnection model (Drenkhahn & Spruit 2002), the magnetized internal shock model (Fan et al. 2004), and the collision-induced reconnection model (Zhang & Yan 2011). For illustration, here we take the global MHD condition
breakdown model to calculate the emission. By comparing with the pair density ($\propto r^{-2}$, $r$ is the radial distance from the central source) and the density required for co-rotation ($\propto r^{-1}$ beyond the light cylinder of the compact object), one can estimate the radius at which the MHD condition breaks down, which reads (Usov 1994; Zhang & Mészáros 2002)

$$r_{\text{MHD}} \sim 5 \times 10^{20} \text{ cm} \left(\frac{L}{10^{37}}\right)^{1/2} \left(\frac{\sigma}{10}\right)^{-1} \left(\frac{t_{\text{v,m}}}{10^2}\right) \left(\frac{\Gamma}{10}\right)^{-1},$$

(18)

where $\sigma$ is the ratio of the magnetic energy flux to the particle energy flux, $\Gamma$ is the bulk Lorentz factor of the outflow, and $t_{\text{v,m}}$ is the minimum variability timescale of the central engine. Beyond this radius, intense electromagnetic waves are generated and outflowing particles are accelerated (e.g., Usov 1994). Such a significant magnetic dissipation process converts the electromagnetic energy into radiation.

At $r_{\text{MHD}}$, the corresponding synchrotron radiation frequency can be estimated as (Fan et al. 2005; Gao & Fan 2006)

$$v_{\nu,\text{MHD}} \sim 1.5 \times 10^{18} \text{ Hz} \left(\frac{1 + z}{1.35}\right)^{-1} \left(\frac{\xi}{0.1}\right) \times C_p^2 \left(\frac{\sigma}{10}\right)^{3} \left(\frac{t_{\text{v,m}}}{10^2}\right),$$

(19)

where $C_p \equiv (\epsilon_{e}/0.1)[13(p - 2)/3(p - 1)]$, $\epsilon_e$ is the fraction of the dissipated comoving magnetic field energy converted to the comoving kinetic energy of the electrons, the accelerated electrons distribute as a single power-law $dn/d\gamma \propto \gamma^{-p}$, and $\xi < 1$ reflects the efficiency of magnetic energy dissipation. Most of the energy is radiated in the X-ray band.

4. A CLUE TO THE X-RAY STEEP DECLINE FOLLOWING THE PROMPT EMISSION IN GAMMA-RAY BURSTS

The observations of Sw 1644+57 suggest that the long-lasting fall-back accretion onto a BH can produce energetic X-ray emission and the radiation luminosity traces the accretion rate (i.e., $L_X \propto M$). One interesting question is whether or not a similar process takes place in GRBs. The answer may be yes. Here, we only discuss the X-ray steep decline (quicker than $t^{-3}$; see Figure 1 of Zhang et al. 2006 for illustration) following the prompt emission in GRBs. In the collapsar model, a fraction of the gas in the core of the collapsing star does not have sufficient centrifugal support and directly forms a central BH. The rest of the material will have sufficient angular momentum to go into orbit around the BH. The fall-back accretion rate is tightly related to the pre-collapse stellar density profile, which is of the form $\rho \propto t^{-1}$. A numerical simulation has revealed that when the outermost 0.5 $M_\odot$ layer of the star (where $t > 5$) is accreted, the fall-back accretion rate can reach a value of $M_{\dot{m}} \propto t^{-3}$ or steeper (MacFadyen et al. 2001; Kumar et al. 2008). If the relation $L_X \propto M$ still holds, one has an X-ray emission decline steeper than $t^{-3}$, in agreement with observational data.\footnote{An alternative scenario is that the central engine turns off abruptly and the quick decline is dominated by the high latitude emission of the prompt emission pulses (see Zhang et al. 2006 and references therein). With future X-ray polarimetry data, we may be able to distinguish between these two kinds of models if our line of sight is not along the center of the ejecta (Fan et al. 2008).}

5. CONCLUSION AND DISCUSSION

Swift J164449+573451 is a peculiar outburst which is most-likely powered by the tidal disruption of a star by a massive BH. In this work, we find that the ratio of the periapsis distance to the disruption radius $\mu < 1$, implying that SW J1644+77 is a plunging event (see Section 2). The mass of the plunging star, however, cannot be tightly constrained due to its strong dependence on the poorly understood parameters $\epsilon$ (the radiation efficiency), $\xi$ (the fraction mass of the star that is eventually accreted into the central), and $\theta_j$ (the half opening angle of the collimated outflow).

As a tidal disruption event, the accretion rate $\dot{M}$ at late times (say, $t > 10$ day) is relatively well understood and is widely believed to be $\propto t^{-4/3}$. The detected X-ray emission $L_X$ has a rather similar decline behavior. Since the forward shock origin of the long lasting and highly variable X-ray emission has already been convincingly ruled out (Shao et al. 2011), the X-ray emission has to be from an outflow launched by the accreting BH. These two facts strongly suggest that $L_X \propto M$, which can shed valuable light on the underlying physics, in particular, the energy extraction process. Three kinds of possible mechanisms have been examined and only the Poynting-flux-dominated outflow model is found to be able to account for the data (see Section 3 for details). Therefore, the magnetic activity at the central engine (e.g., Blandford–Znajek process or Blandford–Payne process) plays the main role in extracting the rotation energy of the BH and then launching the outflow. We suggest that $L_X \propto M$ may also hold in the quick decline phase of GRBs.

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