

## BLACK HOLE SPIN IN Sw J1644+57 and Sw J2058+05

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### ABSTRACT

Recently, a hard X-ray transient event, Sw J1644+57, was discovered by the *Swift* satellite. It likely marks the onset of a relativistic jet from a supermassive black hole (BH), possibly triggered by a tidal disruption event (TDE). Another candidate in the same category, Sw J2058+05, was also reported. The low event rate suggests that only a small fraction of TDEs launch relativistic jets. A common speculation is that these rare events are related to rapidly spinning BHs. We attribute jet launching to the Blandford–Znajek mechanism and use the available data to constrain the BH spin parameter for the two events. It is found that the two BHs indeed carry a moderate to high spin, suggesting that BH spin is likely the crucial factor behind the Sw J1644+57-like events.

*Key words:* accretion, accretion disks – black hole physics – magnetic fields

### 1. INTRODUCTION

A hard X-ray transient event, Swift J16449.3+573451 (Sw J1644+57 hereafter), was discovered by the *Swift* satellite (Gehrels et al. 2004) on 2011 March 28, initially reported as GRB 110328A (Cumplings et al. 2011). It was soon realized that it is not a regular gamma-ray burst (GRB). Long-term follow-up observations with *Swift* (Burrows et al. 2011) revealed that it has extended emission in the X-ray band without significant decay. A much longer variability timescale  $\delta t \sim 100$  s than normal for GRBs (Burrows et al. 2011) as well as its location near the center of a  $z = 0.354$  host galaxy (Levan et al. 2011) link the source to a supermassive black hole (BH) with  $M_{\bullet} \sim 10^7 M_{\odot}$  (Burrows et al. 2011; Bloom et al. 2011). The super-Eddington X-ray luminosity (Burrows et al. 2011), bright radio afterglow (Zauderer et al. 2011), and a historical stringent X-ray flux upper limit suggest that the event marks the onset of a relativistic jet from a supermassive BH. The sharp onset and gradual fadeaway of X-ray flux suggest that the transient may be triggered by tidal disruption of a star by the BH (Bloom et al. 2011; Burrows et al. 2011).

A second candidate in the same category, Swift J2058.4+0516 (Sw J2058+05 hereafter), was reported by Cenko et al. (2011), who argued that its observational properties are rather similar to those of Sw J1644+57.

These events are rare (two detected in seven years of *Swift* operation). Most previous studies of tidal disruption events (TDEs) from supermassive BHs did not expect an associated relativistic jet (e.g., Kobayashi et al. 2004; Strubbe & Quataert 2009; Lodato & Rossi 2011, but see Lu et al. 2008; Giannios & Metzger 2011). The neutrino-annihilation mechanism as invoked in GRBs is inadequate given a much lower accretion rate as compared with GRBs (e.g., Shao et al. 2011). A plausible mechanism to launch a jet from such an event is to tap spin energy of the BH through a magnetic field, which connects the BH event horizon and a remote astrophysical load (Blandford & Znajek 1977, hereafter BZ). This mechanism has been widely invoked to interpret active galactic nucleus (AGN) jets (e.g., Vlahakis & Konigl 2004). Within this scenario, a dominant magnetic jet composition is envisaged. Indeed, modeling the emission of Sw J1644+57 suggests that the jet is highly “particle starved,” i.e., most energy is not carried by particles.

A natural inference is that the outflow is Poynting-flux-dominated (Burrows et al. 2011).

If one identifies the BZ mechanism as the jet launching mechanism, the BH spin parameter (a dimensionless angular momentum of the BH, which ranges from 0 for no spin to 1 for the maximum spin) can be constrained from the data. This is the subject in this Letter.

### 2. METHOD

The bipolar BZ jet power from a BH with mass  $M_{\bullet}$  and angular momentum  $J$  is (Lee et al. 2000; Li 2000; Wang et al. 2002; McKinney 2005)

$$L_{\text{BZ}} = 1.7 \times 10^{44} a_{\bullet}^2 M_{\bullet,6}^2 B_{\bullet,6}^2 F(a_{\bullet}) \text{ erg s}^{-1}, \quad (1)$$

where  $a_{\bullet} = Jc/(GM_{\bullet}^2)$  is the BH spin parameter,  $M_{\bullet,6} = M_{\bullet}/10^6 M_{\odot}$ ,  $B_{\bullet,6} = B_{\bullet}/10^6 \text{ G}$ , and

$$F(a_{\bullet}) = [(1 + q^2)/q^2][(q + 1/q) \arctan q - 1]; \quad (2)$$

here  $q = a_{\bullet}/(1 + \sqrt{1 - a_{\bullet}^2})$  and  $2/3 \leq F(a_{\bullet}) \leq \pi - 2$  for  $0 \leq a_{\bullet} \leq 1$ . It apparently depends on  $M_{\bullet}$ ,  $B_{\bullet}$ , and  $a_{\bullet}$ . However, since an isolated BH does not carry a magnetic field,  $B_{\bullet}$  is closely related to the accretion rate  $\dot{M}$  and the radius of the BH, which depends on  $M_{\bullet}$ . Combining these dependences, one finds that  $L_{\text{BZ}}$  is essentially independent of  $M_{\bullet}$ , but is rather a function of  $\dot{M}$  and  $a_{\bullet}$ . This may be proven with the following rough scalings. For a Newtonian disk, angular momentum equation states  $\dot{M} r^2 (GM_{\bullet}/r^3)^{1/2} \simeq -4\pi r^2 \tau_{r\phi} h$ , where  $\tau_{r\phi}$  is the viscous shear and  $h \propto r$  is the half disk thickness, with  $h/r \ll 1$  for a thin disk and  $h/r \sim 1$  for a thick disk. Adopting the  $\alpha$ -prescription for viscosity, the viscous shear can be expressed as  $\tau_{r\phi} = -\alpha P$ , where  $P$  is the total pressure in the disk. One therefore derives  $P \simeq \dot{M} r^2 (GM_{\bullet}/r^3)^{1/2} / (4\pi r^2 \alpha h) \propto (M_{\bullet}/r^5)^{1/2} \propto M_{\bullet}^{-2}$ , where we have applied the scaling  $r \propto r_s \propto M_{\bullet}$  ( $r_s = 2GM_{\bullet}/c^2$  is the Schwarzschild radius). To make an efficient BZ jet, the accretion inflow should carry a large magnetic flux (e.g., Tchekhovskoy et al. 2011). It is reasonable to assume that magnetic fields in the disk are in close equilibrium with the total pressure, so that  $B_{\bullet}^2 \propto P \propto M_{\bullet}^{-2}$ . Inserting this dependence to Equation (1), one finds that  $L_{\text{BZ}}$  is essentially independent of  $M_{\bullet}$ .

More precisely, we adopt the following prescription to treat the problem. The total pressure in the disk can be expressed as  $P = P_{\text{rad}} + P_{\text{gas}} + P_{\text{B}}$ , where  $P_{\text{rad}} = aT^4/3$ ,  $P_{\text{gas}} = nkT$ , and  $P_{\text{B}} = B_{\text{disk}}^2/8\pi$  are radiation, gas, and magnetic pressure, respectively. Here  $T$  is the temperature of the disk,  $n$  is the gas particle number density,  $a$  is the radiation density constant, and  $k$  is the Boltzmann constant. We denote  $P_{\text{B}} = \beta P$  and take  $\beta \sim 0.5$  in this work. This corresponds to a maximized magnetic flux. A smaller  $\beta$  would demand an even larger  $a_{\bullet}$  to reach the same BZ power. So taking  $\beta \sim 0.5$  gives an estimate of a conservative lower limit of  $a_{\bullet}$ . For the two sources (Sw J1644+57 and Sw J2058+05) we are interested in, the accretion rate is estimated close to or even larger than the Eddington accretion rate (Burrows et al. 2011; Cenko et al. 2011). In this regime, a thin disk model may be adequate to describe the disk, and we apply it to estimate disk pressure for simplicity. The disk pressure peaks in the inner region where radiation pressure may dominate. As a result, for  $\beta \sim 0.5$ , one has

$$\frac{B_{\text{disk}}^2}{8\pi} \sim \frac{1}{3} aT^4, \quad (3)$$

where the temperature of the disk ( $a_{\bullet}$ - and  $r$ -dependent) can be written as

$$T(a_{\bullet}, r) = \left( \frac{3GM\dot{M}}{8\pi r^3 \sigma} f \right)^{1/4}, \quad (4)$$

where  $\sigma$  is the Stephan–Boltzmann constant and the general relativistic correction factor (Page & Thorne 1974)

$$f(a_{\bullet}, r) = \frac{\chi^2}{(\chi^3 - 3\chi + 2a_{\bullet})} \left[ \chi - \chi_{\text{in}} - \frac{3}{2} a_{\bullet} \ln \left( \frac{\chi}{\chi_{\text{in}}} \right) - \frac{3(\chi_1 - a_{\bullet})^2}{\chi_1(\chi_1 - \chi_2)(\chi_1 - \chi_3)} \ln \left( \frac{\chi - \chi_1}{\chi_{\text{in}} - \chi_1} \right) - \frac{3(\chi_2 - a_{\bullet})^2}{\chi_2(\chi_2 - \chi_1)(\chi_2 - \chi_3)} \ln \left( \frac{\chi - \chi_2}{\chi_{\text{in}} - \chi_2} \right) - \frac{3(\chi_3 - a_{\bullet})^2}{\chi_3(\chi_3 - \chi_1)(\chi_3 - \chi_2)} \ln \left( \frac{\chi - \chi_3}{\chi_{\text{in}} - \chi_3} \right) \right], \quad (5)$$

where  $\chi = (r/r_g)^{1/2}$ ,  $\chi_{\text{in}} = (r_{\text{in}}/r_g)^{1/2}$ , and  $r_g = GM_{\bullet}/c^2$ . For a Kerr BH, the disk inner edge  $r_{\text{in}}$  is expressed as (Bardeen et al. 1972)

$$r_{\text{in}}/r_g = 3 + Z_2 - [(3 - Z_1)(3 + Z_1 + 2Z_2)]^{1/2}, \quad (6)$$

for  $0 \leq a_{\bullet} \leq 1$ , where  $Z_1 \equiv 1 + (1 - a_{\bullet}^2)^{1/3}[(1 + a_{\bullet})^{1/3} + (1 - a_{\bullet})^{1/3}]$  and  $Z_2 \equiv (3a_{\bullet}^2 + Z_1^2)^{1/2}$ . In Equation (5),  $\chi_1$ ,  $\chi_2$ , and  $\chi_3$  are the three roots of  $\chi^3 - 3\chi + 2a_{\bullet} = 0$ , i.e.,  $\chi_1 = 2 \cos(\frac{1}{3} \cos^{-1} a_{\bullet} - \pi/3)$ ,  $\chi_2 = 2 \cos(\frac{1}{3} \cos^{-1} a_{\bullet} + \pi/3)$ , and  $\chi_3 = -2 \cos(\frac{1}{3} \cos^{-1} a_{\bullet})$ . It is easy to check that  $f(r = r_{\text{in}}) = 0$  and  $f(r \gg r_{\text{in}}) \simeq 1$ . For a Newtonian disk,  $f$  can be simply written as  $f = 1 - \sqrt{r_{\text{in}}/r}$ .

We assume that the strength of the magnetic field threading the BH is comparable to the largest field strength in the disk, i.e.,  $B_{\bullet} \sim B_{\text{disk}}^{\text{max}}$ . For given  $\dot{M}$  and  $a_{\bullet}$ , we solve numerically the temperature profile of the thin disk and identify the radius  $r_{\text{peak}}$ , where  $T$  reaches the maximum. It is found that  $r_{\text{peak}}$  is very close to  $r_{\text{in}}$ , the innermost radius of the accretion disk. We then calculate  $B_{\text{disk}}^{\text{max}}$  from Equation (3), which is assigned to  $B_{\bullet}$ . Applying Equations (1) and (2), one then obtains  $L_{\text{BZ}}$  once  $\dot{M}$  and  $a_{\bullet}$  are specified.

Observationally, a time-dependent isotropic X-ray luminosity  $L_{\text{X,iso}}$  was measured for the two sources. This luminosity can be connected to the BZ power through

$$\eta L_{\text{BZ}} = f_b L_{\text{X,iso}}, \quad (7)$$

where  $\eta$  is the efficiency of converting BZ power to X-ray radiation,

$$f_b = \frac{\Delta\Omega}{4\pi} = \max \left( \frac{1}{2\Gamma^2}, \frac{\theta_j^2}{2} \right) < 1 \quad (8)$$

is the beaming factor of the jet (Burrows et al. 2011),  $\Delta\Omega$  is the solid angle of the bipolar jet, and  $\Gamma$  and  $\theta_j$  are the Lorentz factor and opening angle of the jet.

In this work, we adopt  $\eta \sim 0.5$ , motivated by the potentially high radiation efficiency of a magnetically dominated jet (e.g., Drenkhahn & Spruit 2002; Zhang & Yan 2011). Again this gives a conservative lower limit of the inferred  $a_{\bullet}$ , since a less efficient jet would demand an even higher spin rate in order to interpret the same observed luminosity.

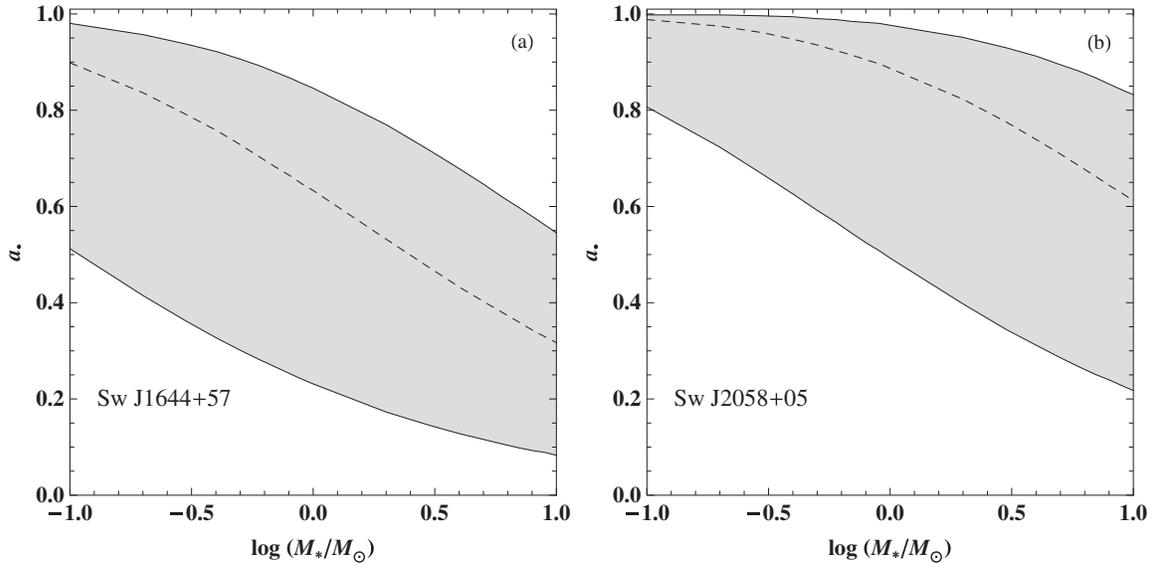
The parameter  $f_b$  has a large uncertainty. Based on the current data, one cannot precisely measure  $\Gamma$  and  $\theta_j$ . Here, we apply a statistical method to estimate the range of  $f_b$  within the TDE framework (see also Burrows et al. 2011). First, the facts that two such events (Sw J1644+57 and Sw J2058+05) were detected by *Swift* in  $\sim 7$  yr and that the field view of the *Swift* Burst Alert Telescope (BAT; Barthelmy et al. 2005) is  $\sim 4\pi/7$  sr suggest that the all-sky rate of such events is  $R_{\text{obs}} \sim 2 \text{ yr}^{-1}$ , with a 90% confidence interval of  $(0.44\text{--}5.48) \text{ yr}^{-1}$  (Kraft et al. 1991). Next, the TDE rate is estimated as  $\sim 10^{-5}$  to  $10^{-4} \text{ yr}^{-1} \text{ galaxy}^{-1}$  based on observational (Donley et al. 2002; Gezari et al. 2009) and theoretical (Wang & Merritt 2004) studies. The galaxy number density is  $n_{\text{gal}} \sim 10^{-3}$  to  $10^{-2} \text{ Mpc}^{-3}$  (Tundo et al. 2007). Sw J2058+05 was marginally detected at  $z = 1.1853$  (Cenko et al. 2011). We then obtain the total TDE event rate within the volume ( $z \leq 1.1853$ )  $R_{\text{tot}} \sim 10^4\text{--}10^5 \text{ yr}^{-1}$ . Considering that only  $\sim 10\%$  of the population can launch a jet (the ‘‘radio-loud’’ AGN fraction; Kellerman et al. 1989; Cirasuolo et al. 2003), the beaming factor can be estimated as

$$f_b \sim \frac{R_{\text{obs}}}{10\% R_{\text{tot}}} \in (4.4 \times 10^{-5}, 5.5 \times 10^{-3}). \quad (9)$$

Finally, in order to infer  $a_{\bullet}$  from  $L_{\text{X,iso}}$ ,  $f_b$  (given  $\beta = 0.5$  and  $\eta = 0.5$ ), one needs to know the accretion rate  $\dot{M}$ . This is an even more loosely constrained parameter. However, if one assumes that the luminosity history of the light curve delineates the accretion history of the BH (noticing  $L_{\text{BZ}} \propto \dot{M}$ ) well, one can normalize the accretion rate using the total accreted mass based on the observed flux and fluence of the event. For Sw J1644+57, since the source has entered a decaying phase, during which the residual fluence no longer contributes significantly to the total fluence, one can take the current total X-ray fluence as a good proxy of the total mass of the tidally disrupted star. Taking the peak flux as an example, the peak accretion rate can be estimated as

$$\dot{M}_{\text{peak}} = \frac{F_{\text{X}}^{\text{peak}}(1+z)}{S_{\text{X}}} M_{\star} = \frac{L_{\text{X,iso}}^{\text{peak}}}{E_{\text{X,iso}}} M_{\star}, \quad (10)$$

where  $F_{\text{X}}^{\text{peak}}$  is the peak X-ray flux,  $S_{\text{X}}$  is the total X-ray fluence,  $L_{\text{X,iso}}^{\text{peak}}$  is the peak isotropic X-ray luminosity,  $E_{\text{X,iso}}$  is the isotropic X-ray energy, and  $M_{\star}$  is the total mass of the



**Figure 1.** Parameter space of  $M_*$ - $a_*$  for Sw J1644+57 (left panel) and Sw J2058+05 (right panel) with  $\beta = 0.5$  and  $\eta = 0.5$ . The shaded region indicates the range for BH spin bracketed by  $f_b = 4.4 \times 10^{-5}$  (lower) and  $f_b = 5.5 \times 10^{-3}$  (upper). The dashed line corresponds to the most probable value  $f_b = 10^{-3}$ . For Sw J2058+05,  $\zeta = 0.1$  is adopted. It is shown that the most probable values are  $a_*(\text{J1644}) = 0.63, 0.90$  and  $a_*(\text{J2058}) = 0.89, 0.99$  for  $M_* = 1, 0.1 M_\odot$ , respectively.

star. The factor  $(1+z)$  was applied to convert the observed time to the time in the source rest frame. The accretion rate at other epochs can be estimated similarly. The range of  $M_*$  may be between  $0.1 M_\odot$  and  $10 M_\odot$ . One can then derive a mass-dependent constraint on  $a_*$ .

### 3. Sw J1644+57 AND Sw J2058+05

Now we apply the above method to the two sources.

According to Burrows et al. (2011), during the first 50 days after the first BAT trigger the total X-ray energy corrected for live-time fraction for Sw J1644+57 is  $E_{X,\text{iso}}(\text{J1644}) \sim 2 \times 10^{53}$  erg in the 1.35–13.5 keV rest-frame energy band (corresponding the energy band of the *Swift* X-Ray Telescope, Burrows et al. 2005). The peak luminosity in the same energy band is  $L_{X,\text{iso}}^{\text{peak}}(\text{J1644}) \sim 2.9 \times 10^{48}$  erg s $^{-1}$ . The accretion rate (Equation (10)) can be estimated as

$$\dot{M}_{\text{peak}}(\text{J1644}) \simeq 1.45 \times 10^{-5} M_* \text{s}^{-1}. \quad (11)$$

Figure 1(a) presents the constraint on  $a_*$  for Sw J1644+57 as a function of  $M_*$ , with  $\beta = 0.5$  and  $\eta = 0.5$ . The two boundary lines correspond to two ends of the range of  $f_b$ , with the lower and upper boundaries corresponding to  $f_b = 4.4 \times 10^{-5}$  and  $f_b = 5.5 \times 10^{-3}$ , respectively. The middle dashed line corresponds to the most probable value  $f_b \sim 10^{-3}$ . One can see that the supermassive BH must have a moderate to high spin rate. Given the standard stellar initial mass function, the number of low-mass stars is much more abundant than the number of high-mass stars. If one takes  $M_* = 1 M_\odot$ , the required range of  $a_*$  is (0.23, 0.85), with the most probable value  $a_* = 0.63$ . For  $M_* = 0.1 M_\odot$  (more probable), the range of  $a_*$  is (0.51, 0.98) with the most probable value  $a_* = 0.90$ .

Sw J2058+05 was discovered by *Swift*/BAT through a four-day (2011 May 17–20) integration. A subsequent target-of-opportunity (ToO) observation eight days after the end of the four-day integration (2011 May 28) revealed an X-ray source that behaves very similarly to Sw J1644+57 (Cenko et al. 2011), suggesting that it is very likely another Sw J1644+57-like event. The 0.3–10 keV peak flux is  $F_{X,\text{iso}}^{\text{peak}} \simeq 7.9 \times$

$10^{-11}$  erg cm $^{-2}$  s $^{-1}$ , corresponding to a peak luminosity of  $L_{X,\text{iso}}^{\text{peak}} \simeq 3 \times 10^{47}$  erg s $^{-1}$ . The total X-ray fluence from the beginning of the ToO observation to 2011 July 20 is  $S_X \simeq 1.0 \times 10^{-4}$  erg cm $^{-2}$  (Cenko et al. 2011). Since we did not catch the X-ray emission during the first 11 days (May 17–27) when it is supposed to be much brighter, the registered X-ray fluence only corresponds to a small fraction of the total mass of the star. One can still apply Equation (10) to estimate the accretion rate, except that one should replace  $M_*$  by  $\zeta M_*$ , where  $\zeta < 1$  is the fraction of stellar mass that is accreted after May 28. Noticing  $z = 1.1853$  (Cenko et al. 2011), this gives a peak accretion rate

$$\dot{M}_{\text{peak}}(\text{J12058}) \simeq 1.72 \times 10^{-7} \zeta_{-1} M_* \text{s}^{-1}, \quad (12)$$

where  $\zeta = 0.1 \zeta_{-1}$  has been adopted.

The constraint on BH spin for Sw J2058+05 is shown in Figure 1(b). We find that the demand for BH spin is even more stringent for this source. For  $M_* = 1 M_\odot$ , the required range of  $a_*$  is (0.49, 0.98) with the most probable value  $a_* = 0.89$ . For  $M_* = 0.1 M_\odot$  (more probable), the range of  $a_*$  is (0.81, 0.998) with the most probable value  $a_* = 0.99$ .

### 4. CONCLUSION AND DISCUSSION

Sw J1644+57 and Sw J2058+05 are the first two prototype objects in this newly identified astrophysical phenomenon, namely, a relativistic jet associated with a TDE from a supermassive BH. A straightforward question is why only some TDEs launch jets. Based on observational properties, it has been argued that the jets must be Poynting-flux-dominated (Burrows et al. 2011; Shao et al. 2011). Invoking the BZ mechanism as the power of the jet, we show here that both events need to invoke a BH with a moderate to rapid spin in order to interpret the observations: the most probable values are  $a_*(\text{J1644}) = 0.63, 0.90$  and  $a_*(\text{J2058}) = 0.89, 0.99$  for  $M_* = 1, 0.1 M_\odot$ , respectively. We therefore suggest that BH spin is the key factor behind the Sw J1644+57-like events, although other factors may also play a role (e.g., Cannizzo et al. 2011).

An elegant feature of the method we employ is that the inferred BH spin parameter  $a_*$  essentially does not depend on

the BH mass. On the other hand, knowing  $a_\bullet$  leads to a better constraint on BH mass based on the variability argument. For example, the observed minimum variability timescale of J1644+57 is  $\delta t_{\text{obs,min}} \sim 100$  s. If one relates  $\delta t_{\text{min}} = \delta t_{\text{obs,min}}/(1+z)$  to the timescale defined by the innermost radius of the accretion disk,  $r_{\text{in}}/c$ , one can derive the BH mass of Sw J1644+57:

$$M_{\bullet,6} \simeq 15 \left( \frac{r_{\text{in}}}{r_g} \right)^{-1} \left( \frac{\delta t_{\text{obs,min}}}{100 \text{ s}} \right). \quad (13)$$

One has  $2.5 < M_6 < 15$  for  $0 \leq a_\bullet \leq 1$ , with the most probable value  $M_6 = 6.5$  for  $a_\bullet = 0.9$ . This is consistent with the constrained BH mass from the  $M-L_{\text{bulge}}$  relation, which gives an upper limit of  $2 \times 10^7 M_\odot$  (Burrows et al. 2011).

From Equation (1), one can also infer the strength of the magnetic field at the BH horizon:

$$B_{\bullet,6} \simeq 131 f_b^{1/2} \eta^{-1/2} F(a_\bullet)^{-1/2} a_\bullet^{-1} M_6^{-1}. \quad (14)$$

Taking BH mass  $M_\bullet = 6.5 \times 10^6 M_\odot$  and the most probable value for BH spin  $a_\bullet = 0.9$ , one finds that the magnetic field threading BH would be  $B_\bullet \sim 1.1 \times 10^6$  G, which is much higher than the average field strength of a typical main-sequence star ( $< 10^3$  G). The accumulation of magnetic flux by accretion flow and instability in the disk may account for such high magnetic field strength (e.g., Tchekhovskoy et al. 2011).

For Sw J2058+05, due to the low X-ray flux at the late epochs a much looser constraint on variability,  $\delta t_{\text{obs,min}} < 10^4$  s, was obtained (Cenko et al. 2011), so that the precise values of  $M_\bullet$  and  $B_\bullet$  cannot be derived.

An alternative scenario to interpret the Sw J1644+57-like event may be the onset of an AGN (Burrows et al. 2011). This scenario, which predicts that the two sources will be active at least in the following millennium, may be less favorable in view of the rapid onset of emission in Sw J1644+57 and the gradual decay in both events, but is not ruled out. Our method can be applied to this scenario as well, except that  $f_b$  becomes much larger (due to the intrinsic rarity of AGN onset events), and  $M_*$  is no longer limited to the range of  $0.1\text{--}10 M_\odot$ . Our analysis suggests that the BZ power is likely not adequate to interpret the data (because of the large emission power demanded by the small  $f_b$  factor) even for maximum spin ( $a_\bullet \sim 1$ ), unless the accretion rate is much higher, so that the total amount of fuel  $M_* \gg 1 M_\odot$ . This is not impossible since the fuel in the AGN onset scenario is from a gas cloud near the BH, whose mass is not specified.

Krolik & Piran (2011) proposed a model for Sw J1644+57 invoking a white dwarf being tidally disrupted by a smaller BH ( $M_\bullet \sim 10^4 M_\odot$ ). Regardless of how this model may interpret the  $\delta t_{\text{min}}$  and the apparent association of the source with the center of host galaxy, the derived  $a_\bullet$  range also applies to their

model for the  $M_*$  range of a white dwarf since our constraint is  $M_\bullet$ -independent.

Finally, in our calculations we did not consider the evolution of  $a_\bullet$  during the accretion phase. This is justified given the large  $M_\bullet/M_*$  ratio.

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