Swift J1644+57 gone MAD: the case for dynamically important magnetic flux threading the black hole in a jetted tidal disruption event

Alexander Tchekhovskoy,1,2★†‡§ Brian D. Metzger,3 Dimitrios Giannios4 and Luke Z. Kelley5

1Lawrence Berkeley National Laboratory, 1 Cyclotron Rd, Berkeley, CA 94720, USA
2Department of Astronomy, University of California, Berkeley, CA 94720-3411, USA
3Department of Physics, Columbia University, 538 West 120th Street, 704 Pupin Hall, New York, NY 10027, USA
4Department of Physics, Purdue University, 525 Northwestern Avenue, West Lafayette, IN 47907, USA
5Astronomy Department, Harvard University, 60 Garden Street MS 10, Cambridge, MA 02138, USA

Accepted 2013 October 24. Received 2013 October 24; in original form 2013 January 10

ABSTRACT
The unusual transient Swift J1644+57 likely resulted from a collimated relativistic jet, powered by the sudden onset of accretion on to a massive black hole (BH) following the tidal disruption (TD) of a star. However, several mysteries cloud the interpretation of this event, including (1) the extreme flaring and ‘plateau’ shape of the X-ray/γ-ray light curve during the first \( t - t_{\text{trig}} \sim 10 \) d after the γ-ray trigger; (2) unexpected rebrightening of the forward shock radio emission at \( t - t_{\text{trig}} \sim \) months; (3) lack of obvious evidence for jet precession, despite the misalignment typically expected between the angular momentum of the accretion disc and BH; (4) recent abrupt shut-off in the jet X-ray emission at \( t - t_{\text{trig}} \sim 1.5 \) yr. Here, we show that all of these seemingly disparate mysteries are naturally resolved by one assumption: the presence of strong magnetic flux \( \Phi_1 \) threading the BH. Just after the TD event, \( \Phi_1 \) is dynamically weak relative to the high rate of fall-back accretion \( \dot{M} \), such that the accretion disc (jet) freely precesses about the BH axis = our line of sight. As \( \dot{M} \) decreases, however, \( \Phi_1 \) becomes dynamically important, leading to a state of ‘magnetically arrested disk’ (MAD). MAD naturally aligns the jet with the BH spin, but only after an extended phase of violent rearrangement (jet wobbling), which in Swift J1644+57 starts a few days before the γ-ray trigger and explains the erratic early light curve. Indeed, the entire X-ray light curve can be fitted to the predicted power-law decay \( \dot{M} \propto t^{-\alpha} (\alpha \simeq 5/3 - 2.2) \) if the TD occurred a few weeks prior to the γ-ray trigger. Jet energy directed away from the line of sight, either prior to the trigger or during the jet alignment process, eventually manifests as the observed radio rebrightening, similar to an off-axis (orphan) γ-ray burst afterglow. As suggested recently, the late X-ray shut-off occurs when the disc transitions to a geometrically thin (jetless) state once \( \dot{M} \) drops below ~the Eddington rate. We predict that, in several years, a transition to a low/hard state will mark a revival of the jet and its associated X-ray emission. We use our model for Swift J1644+57 to constrain the properties of the BH and disrupted star, finding that a solar mass main-sequence star disrupted by a relatively low-mass \( M_\bullet \sim 10^5-10^6 M_\odot \) BH is consistent with the data, while a white dwarf disruption (though still possible) is disfavoured. The magnetic flux required to power Swift J1644+57 is much too large to be supplied by the star itself, but it could be collected from a quiescent ‘fossil’ accretion disc that was present in

★E-mail: atchekho@berkeley.edu
†NASA Einstein Fellow.
‡Formerly at the Center for Theoretical Science, Jadwin Hall, Princeton University, Princeton, NJ 08544, USA.
§Center for Theoretical Science Fellow.
the galactic nucleus prior to the TD. The presence (lack of) of such a fossil disc could be a deciding factor in what TD events are accompanied by powerful jets.

**Key words:** accretion, accretion discs – black hole physics – MHD – gamma-rays: galaxies – X-rays: galaxies.

1 INTRODUCTION

The unusual soft $\gamma$-ray/X-ray/radio transient Swift J164449.3+573451 (hereafter Sw J1644+57) has been broadly interpreted as resulting from a relativistic outflow, powered by accretion following thetidal disruption (TD) of a star by a massive black hole (BH; Bloom et al. 2011; Burrows et al. 2011; Levan et al. 2011; Zauderer et al. 2011). Evidence supporting this model includes the rapid onset of Sw J1644+57 and its location at the centre of a compact galaxy at redshift $z \approx 0.353$ (Levan et al. 2011). At least until recently, the spectral energy distribution showed two distinct components, which led Bloom et al. (2011) and Burrows et al. (2011) to suggest different sources for the X-ray and radio emission (see also Liu, Pe’er & Loeb 2012). The X-ray/$\gamma$-ray emission is highly variable, which indicates an origin from relatively small radii, likely from a location internal to the jet itself (although see Socrates 2012). Fig. 1 shows that the emission was particularly variable for the first $\sim 10$ d after the *Swift/Burst Alert Telescope* (BAT) trigger, $t_{\text{lag}}$ (though roughly constant in a time-averaged sense), after which point undergoing a power-law decline,

$$L_X \propto (t - t_{\text{lag}})^{-\alpha},$$  

(1)

with $\alpha \sim 5/3$, consistent with the rate of fall-back accretion in simple TD models (e.g. Rees 1988; Lodato, King & Pringle 2009; Guillochon & Ramirez-Ruiz 2013; Stone, Sari & Loeb 2013; although see Cannizzo, Troja & Lodato 2011). The X-ray flux has recently abruptly dropped by more than two orders of magnitude, indicating that the jet has apparently ‘shut off’ approximately 500 d after the initial trigger (Zauderer et al. 2013). The total isotropic X-ray energy radiated to date is $\sim 5 \times 10^{51}$ erg.

In contrast to the X-ray emission, brightness temperature constraints place the radio emission from Sw J1644+57 at much larger radii, suggesting that it instead results from synchrotron emission powered by the shock interaction between the relativistic jet and the surrounding circumnucleus medium (Gianios & Metzger 2011; Zauderer et al. 2011, 2013; Berger et al. 2012; Metzger, Gianios & Mimica 2012; Wiersema et al. 2012). By modelling the observed radio emission based on the first several weeks of data, Metzger et al. (2012, hereafter MGM12) derived values for the bulk Lorentz Factor $\Gamma_j \approx 10$, opening angle $\theta_j \sim 1/\Gamma_j \sim 0.1$, and beaming fraction $f_b \approx 3 \times 10^{-3}$ of the jet which are remarkably similar to those of active galactic nuclei (AGN) jets. Berger et al. (2012, hereafter B12) presented updated radio light curves of Sw J1644+57, which showed a distinct rebrightening starting at $t - t_{\text{lag}} \sim 1$ month and peaking on a time-scale of several months. This behaviour is surprising since the emission is significantly brighter than expected if the blast wave were evolving with a relatively constant energy, as would be expected if the instantaneous jet power tracked the X-ray light curve. B12 proposed that this large additional energy results from slower material catching up to the forward shock at late times. Regardless of its interpretation, however, the radio rebrightening clearly indicates that the jet structure (angular or temporal) is more complex than that commonly and successfully applied to normal

gamma-ray burst (GRB) afterglows (Panaitescu & Kumar 2002; Cao & Wang 2012; Liu et al. 2012).

The discovery of a jetted tidal disruption event (TDE) presents several theoretical mysteries. Relativistic jets from AGN are thought to result from magnetic, rather than hydrodynamic, collimation and acceleration (e.g. Rees et al. 1982).\(^1\) If the jet energy is derived from the Penrose–Blandford–Znajek process, then the total jet power is given by (Tchekhovskoy, Narayan & McKinney 2010)

$$P_j = \frac{2\pi}{16\pi G^2 c^4} \Phi^2 \omega_0 \eta f(\omega_0)$$

$$= 1.2 \times 10^{47} \Phi^2 \omega_0 \eta \text{erg s}^{-1},$$

(2)

$$= 0.5 \times 10^{47} \Phi^2 \omega_0 \eta \text{erg s}^{-1},$$

(3)

where $\kappa \approx 0.045$, $r_g = G M_*/c^2$ is BH gravitational radius; $\Phi = 10^{40} \Phi_{\ast,30}$ cgs is the magnetic flux threading the hole; $M_* = M_\ast 10^5 M_\odot$ is the BH mass; $\omega_0 = a/(1 + (1 - a^2)^{1/2})$ is dimensionless angular frequency of BH horizon (equals unity for a maximally spinning BH); and $f(\omega_0) = 1 + 0.35 \omega_0^{-3} - 0.58 \omega_0^{-5}$ is a high spin correction, while the normalization in the third line has been calculated for $a = 0.9$.

Sw J1644+57 radiated $\approx 2 \times 10^{53}$ erg over the first $\sim 10$ d after the trigger, corresponding to an average isotropic X-ray luminosity $L_X^{\text{iso}} \approx 2 \times 10^{47}$ erg s$^{-1}$. The total (true) jet power during this interval was thus $P_j^{\text{iso}} = 2 f_b \epsilon_b \omega_0 \eta L_X^{\text{iso}} \approx 10^{48} (f_b \epsilon_b \omega_0 \eta L_X^{\text{iso}}/0.03)$ erg s$^{-1}$, where $\epsilon_b \sim f_b$ is a bolometric efficiency; $\epsilon_b \ll 1$ is the jet radiative efficiency; and the factor of 2 accounts for the jet beamed away from Earth. Equation (3) shows that in order to explain $P_j^{\text{iso}}$, the required magnetic flux $\Phi_{\ast,30} \sim M_\ast$, for $a = 0.9$, is several orders of magnitude larger than that through a typical main-sequence (MS) star (Bloom et al. 2011). In Section 5.2 we discuss possible alternative sources of magnetic flux, such as could be supplied by a pre-existing quiescent accretion disc.

Regardless of its origin, the high luminosity of Sw J1644+57 requires a large magnetic flux. In fact, at least two independent arguments suggest that such magnetic flux was actually present.\(^1\) First, note that a significant fraction of the mass of the disrupted star, and hence of its magnetic flux or that of a quiescent disc, is accreted on the characteristic fall-back time $t_{\text{fb}}$ [equation (12)] of the most tightly bound tidal debris (e.g. Ulmer 1999; Strubbe & Quataert 2009). Equation (3) would thus naively imply that the average jet power should be constant, or rising, at times $t \gg t_{\text{fb}}$, in contradiction to the observed power-law decline in the X-ray luminosity. Using 3D general relativistic magnetohydrodynamic (GRMHD) simulations, Tchekhovskoy, Narayan & McKinney (2011) have shown that if the magnetic flux $\Phi_\ast$ is sufficiently high, then magnetic forces impede accretion on to the BH, causing the flow to enter a ‘magnetically arrested disk’ (MAD; e.g. Narayan, Igumenshchev & Abramowicz

---

\(^1\) One theoretical objection to a hydrodynamic jet is the drag due to Compton cooling (e.g. Phinney 1982), which could in principle, however, be substantially reduced in the case of a TDE due to the lack of a previously existing broad-line region.
A. Tchekhovskoy et al.

Figure 1. X-ray light curve of Sw J1644+57, as measured by the Swift X-ray Telescope (XRT) and Chandra. When plotted as a function of time since the γ-ray trigger $t - t_{\text{trig}}$, the average emission shows a ‘plateau-like’ phase lasting for $\sim 10\, \text{d}$, which naively appears inconsistent with the predicted power-law decay $M_{\text{fb}} \propto (t - t_{\text{disr}})^{-3}$ in the rate of fallback accretion following a TDE ($t = t_{\text{disr}}$). However, if the trigger time is delayed with respect to the disruption, then a ‘plateau-like’ shape in $F_{\text{X}}(t - t_{\text{disr}})$ is naturally produced (Section 2.2). The two solid curves show $M_{\text{fb}} \propto t^{-\alpha}$ (arbitrary normalization) for complete (red, $\alpha = 5/3$ for trigger time delay $t_{\text{trig}} - t_{\text{disr}} = 15\, \text{d}$) and partial stellar disruption (blue, $\alpha = 2.2$, $t_{\text{trig}} - t_{\text{disr}} = 30\, \text{d}$), while the dotted curves show the conventional versions of power-law fits that neglect the trigger time delay. The trigger time delay for the solid lines is chosen to match $M_{\text{fb}}$ with the average luminosity of the early plateau phase ($t - t_{\text{disr}} \lesssim 10\, \text{d}$). If we instead match $M_{\text{fb}}$ to the ‘envelope’ created in the light curve by the brightest flares, then adopting a shorter trigger time delay is also consistent with the data, e.g., $t_{\text{trig}} - t_{\text{disr}} = 5\, \text{d}$ for a complete disruption (red dashed line, $\alpha = 5/3$) or a partial disruption (blue dashed line, $\alpha = 2.2$). This time-scale is consistent with the first evidence for activity from Sw J1644+57 (Burrows et al. 2011; Krimm & Barthelmy 2011; Zauderer et al. 2011) 4 d prior to the first BAT trigger.

2003; Igumenshchev 2008) state. MAD flows achieve a quasi-steady state as the result of 3D instabilities, which allow matter to slip past the field lines towards the BH. This process regulates the jet power [equation (3)] to be proportional to the BH feeding rate, $P_{\text{j}} \propto M_{\text{lb}}$ [equation (17)]. Hence, the fact that the late time jet power in Sw J1644+57 faithfully tracks the expected rate of fall-back accretion $M_{\text{fb}} \propto t^{-3}$ is only naturally understood if the flow is in a MAD state.

Additional evidence for a strong magnetic flux is related to the mystery raised by Stone & Loeb (2012), who noted that in general a TD jet should precess if it is pointed along the angular momentum axis of the accretion disc.2 Lack of clear evidence3 for large-scale jet precession in Sw J1644+57 thus requires either a set of highly unlikely circumstances, such as an unphysically low BH spin or near-perfect alignment between the angular momentum of the BH and the original orbit of the disrupted star (Stone & Loeb 2012).

2 A disc-aligned jet is expected if outflows from the disc are sufficiently powerful to direct the jet on large scales, as is likely if the accretion rate is highly super-Eddington.

2 Sexton et al. (2012) and Lei, Zhang & Gao (2012) argue that the ‘dipping’ behaviour and other quasi-periodic features observed in the X-ray light curve of Sw J1644+57 are due to the jet precession/nutation. However, the high duty cycle of observed emission would still require a relatively small (and hence fine-tuned) inclination between the BH spin and stellar orbit ($\lesssim 10^\circ$ to $20^\circ$).

or some mechanism for aligning the angular momentum of the disc with the BH spin. In fact, recent numerical simulations by McKinney, Tchekhovskoy & Blandford (2013, hereafter MTB13) show that such an alignment between the disc and BH spin axis can occur due to magnetohydrodynamically (MHD) forces, but only if the strength of the magnetic field threading the BH is similar to that required for MAD accretion.

In this paper, we present a new physical scenario for Sw J1644+57 which addresses the seemingly disparate mysteries raised above, including the shape of the X-ray/γ-ray light curve; lack of jet precession; and late radio rebrightening. We show that all of these features are naturally expected given a single assumption: the presence of a strong magnetic flux threading the BH. We begin in Section 2 with some basic phenomenological considerations. Then in Section 3 we overview the time line of our proposed scenario for Sw J1644+57. In Section 4, we use our model to constrain the properties of the BH and stellar progenitor. In Section 5, we discuss our results, including the nature of the disrupted star (Section 5.1); the origin of the magnetic flux (Section 5.2); the nature and duration of the flaring state (Section 5.3); the origin of the radio rebrightening (Section 5.4); and future predictions (Section 5.5). We present our conclusions in Section 6. Throughout the paper we use Gaussian cgs units and set the zero time to the point of disruption, $t_{\text{disr}} = 0$.

2 PHENOMENOLOGICAL CONSIDERATIONS

We begin by discussing the origin of several features in the X-ray and radio light curves of Sw J1644+57 from a phenomenological perspective. Then in Section 3 we present a more systematic overview of our model.

2.1 Relativistic jet as the origin of X-ray emission

Fig. 1 shows the soft X-ray (0.3–10 keV) light curve of Sw J1644+57. Upper limits on the mass of the host galaxy and the observed variability time-scale place a rough upper limit of $M_{\text{r}} \lesssim 10^7\, M_{\odot}$ on the mass of the central BH (e.g. Bloom et al. 2011). The jet luminosity was thus highly super-Eddington over at least the first several days of activity, even after correcting for jet beaming (MGM12). At $t \gtrsim 10\, \text{d}$, the time-averaged emission follows a power-law decline with temporal index $\alpha \gtrsim 5/3$ (equation 1), similar to the rate of mass fall-back in standard TD scenarios (Rees 1988). This decay rate has been used as evidence that Sw J1644+57 was in fact a TDE, but as we now discuss, below, it is not clear a priori why the jet power should so faithfully track the mass accretion rate.

2.2 Early time light curve ‘plateau’

The first $\sim 10\, \text{d}$ of Sw J1644+57 were characterized by particularly intense flaring (Fig. 1). The luminosity during this period, though variable by several orders of magnitude, was approximately constant on average. Such a ‘plateau’ has no obvious explanation in TD scenarios, but, as we now discuss, it naturally results if the γ-ray trigger was delayed with respect to the time of disruption.

Fig. 1 shows that a plateau consistent with the data is reproduced simply by shifting the zero-point of time, even for a purely power-law decay in the assumed flux. In order to reproduce the duration of the plateau and match the predicted accretion rate to the average X-ray flux, we find that a trigger delay $t_{\text{trig}} - t_{\text{disr}} \sim \text{weeks/month}$.
2.5 Jet shutoff

The X-ray emission from Sw J1644+57 recently declined abruptly at \( t \approx 500 \text{ d} \), after which the XRT is able to place only upper limits on the flux (Fig. 1). Sw J1644+57 was subsequently detected by *Chandra* at a flux nearly two orders of magnitude lower than that just prior to the decline (Sbarufatti et al. 2012; Zauderer et al. 2013); however, this residual flux is consistent with it being the high-energy extension of the same forward shock synchrotron emission observed at radio frequencies (Zauderer et al. 2013). Such a jet ‘shutoff’ may occur once the accretion rate drops below a fraction of the Eddington accretion rate (Zauderer et al. 2013), since after this time the disc becomes geometrically thin and enters a thermally dominant accretion state, which is not observed to produce powerful jets (Russell et al. 2011). We note that in addition to the jets physically turning off, the X-ray emission drop could also be consistent with jets becoming misaligned with our line of sight. We discuss both possibilities in Section 3.5.

3 TIME LINE OF PROPOSED SCENARIO

We now describe the time line for our proposed scenario for Sw J1644+57, the stages of which are illustrated in Fig. 2.

3.1 Stage 0: tidal disruption and flux accumulation

A star of mass \( M_* \), mass \( M_\odot \) and radius \( R_* \), radius \( R_\odot \), is tidally disrupted by a BH of mass \( M_\bullet = 10^3 M_* M_\odot \) if its pericentre radius \( R_p \) lies within the tidal radius \( R_t \approx R_\odot (M_\bullet / M_\star)^{2/3} \), after disruption, approximately half of the star is immediately unbound, while the other half remains marginally bound and is placed on highly elliptic orbits, with the most tightly bound material falling back on a characteristic time-scale (Stone et al. 2013)

\[
t_{fb} \approx 17.3 \text{ d} M_*^{1/2} m_*^{-1/2} r_*^{3/2}.
\]

The fall-back accretion rate \( M_{fb} \) peaks at \( t_{peak} \approx t_{fb} \); Ulmer (1999) at a characteristic value

\[
M_{fb,peak} \approx \frac{\alpha - 1}{2/3} M_* \approx 2 \times 10^{36} \text{ g s}^{-1} \frac{\alpha - 1}{2/3} M_*^{1/2} m_*^{3/2} r_*^{3/2},
\]

or in terms of the Eddington accretion rate \( \dot{M}_{Edd} \equiv L_{Edd}/0.1c^2 \)

\[
\frac{M_{fb,peak}}{\dot{M}_{Edd}} \approx 4 \times 10^3 \frac{\alpha - 1}{2/3} M_*^{3/2} m_*^{3/2} r_*^{3/2},
\]

and subsequently decays as a power law

\[
M_{fb} \approx M_{fb,peak} \left( \frac{t}{t_{fb}} \right)^{\alpha},
\]

where \( \alpha = 5/3 \) for a complete disruption and \( \alpha = 2.2 \) for a partial disruption (Guillochon & Ramirez-Ruiz 2013; Figs 3 and 4). Until recently (e.g. Ulmer 1999; Strubbe & Quataert 2009), TD models predicted that \( t_{fb} \) and \( M_{fb,peak} \) should depend also on the impact parameter \( \beta \equiv R_p/R_\odot \) in addition to the stellar and BH properties; however, recent numerical (Guillochon & Ramirez-Ruiz 2013) and analytic (Stone et al. 2013) work has shown that these estimates were made on the faulty assumption that the energy distribution of the disrupted stellar material is frozen-in at pericentre instead of the tidal radius.

Matter falls back and circularizes (see, e.g., Ramirez-Ruiz & Rosswog 2009 for a hydrodynamic simulation of circularization process) to form an accretion disc at the radius

\[
R_{\text{circ}} \approx 2 R_g = 2 R_g \beta^{-1} \approx 430 r_* m_*^{-1/3} M_*^{2/3} \beta^{-1} R_g
\]
accreting soon thereafter. Equation (8) can be understood by considering the conservation of angular momentum during the process of circularization. The angular momentum of a general elliptical orbit is $J^2 = GMa(1 - e^2)$, where $a$ is semimajor axis, $e$ is ellipticity and $M$ is the mass of the system. The pericentre radius is given by $R_p = a(1 - e)$. So we can write this as $J^2 = GMa(1 - e)(1 + e) = GMR_p(1 + e)$. Thus, if we have an initially parabolic (zero energy) orbit with $e = 1$, then the pericentric ($=radius$) of the circular orbit ($e = 0$) with the same angular momentum $J$ is just twice as great. In other words, at constant $J$, the quantity $R_p(1 + e)$ is conserved.

When the accretion rate is super-Eddington at early times, the disc may be prone to outflows driven by radiation pressure (e.g. Ohsuga et al. 2005; Strubbe & Quataert 2009), in which case the accretion rate reaching the BH is less than $M_{fb}$. If the magnetic flux responsible for powering the jet in Sw J1644+57 originates from the star itself, then a substantial fraction of the total flux is accumulated on the relatively short fall-back time $t \lesssim t_{fb}$ (equation 4). However, flux accumulation can last significantly longer if the field is swept up from a quiescent fossil disc by the infalling debris (in which case $\Phi_\star \propto (t/t_{fb})^{3/2}$; Section 5.2). On time-scales $t \gtrsim t_{fb}$ the jet power $P_j \propto \Phi_\star^2$ [equation (3)] thus either saturates to a constant value (stellar flux), or increases as $P_j \propto (t/t_{fb})^{3/2}$ (fossil disc). In what follows, we express the flux in a dimensionless form,

$$\phi_\star \equiv \Phi_\star / (M_{fb}c)^{7/2} \approx 30 \Phi_{\star,30} \left( \frac{M}{M_{fb,peak}} \right)^{-1/2} M_{\star,5}^{3/4} m_{\star,4}^{3/4},$$

that quantifies its dynamical importance in relation to the accreting gas, where $r_g \equiv GM_\star/c^2$. Figs 3 and 4 show the time evolution of $\Phi_\star$ and $\phi_\star$ in our model for Sw J1644+57, based on two scenarios for the disrupted star (Section 4.1).

3.2 Stage 1: precessing disc-aligned jet (Fig. 2a)

In general the angular momentum of the initial stellar orbit and fall-back accretion disc will not be aligned with the spin of the BH. Such tilted, geometrically thick accretion discs undergo precession due to Lense–Thirring torques, and their jets are also expected to precess (Fig. 2a). Though expected, evidence for precession is not obviously present in the light curve of Sw J1644+57 (Stone & Loeb 2012; although see Lei et al. 2012; Saxton et al. 2012). This mystery is resolved if we postulate that our line of sight is aligned with the BH spin axis (Section 2), such that the emission from the tilted jet is beamed away from us and hence is not initially detectable.

Because the magnetic field is dynamically weakest relative to the accretion flow ($\phi_\star$ is smallest) when the accretion rate is highest $M \sim M_{fb,peak}$ (equation 9), the magnetic flux does not appreciably influence the disc inclination at these early times. This allows for a phase of jet precession as described above.

At even earlier times, however, when $\phi_\star \lesssim \phi_{\star,30} \approx 20$, the ram pressure of the quasi-spherical accretion flow (as expected for super-Eddington accretion) is so high that it can stifle jet formation altogether (Komissarov & Barkov 2009). Since $\phi_\star \propto t^{3/2}$ if $\Phi_\star = constant$ (equation 9; assuming $M_{fb} \propto t^{-\omega}$), this limits the duration of Stage 1 to $\sim 4^{-2/3} t_{MAD} \sim (1/\text{few}) t_{MAD}$, where $t_{MAD}$ is the onset of

Figure 2. Stages in the proposed model for Sw J1644+57. [Panel (a)]: Stage 1. Shortly after the stellar disruption, the magnetic flux threading the BH $\Phi_\star$ increases, as it is dragged inward by the accreting material. Stellar debris returning to the BH forms a *tilted* accretion disc, with a rotation axis that is misaligned with the BH spin. The latter points towards our line of sight (shown as vertical in the figure). At early times, when the mass accretion rate $M$ is highest, the magnetic flux $\Phi_\star$ is not dynamically important. Under these conditions, the disc undergoes precession due to Lense–Thirring torques (Fragile et al. 2007) and the jets point along the disc axis. Their high-energy emission is beamed away from Earth and is not detectable. [Panel (b)]: Stage 2. As $M$ decreases with time, the magnetic flux eventually becomes dynamically important, leading to a state of MAD flow (Narayan et al. 2003; Tchekhovskoy et al. 2011). In the MAD state, the magnetic field is sufficiently strong to offset gravitational forces acting on the inner disc and to align the axis of the disc (and hence the jets) with the BH spin (so-called ‘magnetoalignment effect’; MTB13). However, it takes time for the entire disc and jet to align with the BH spin. As $M$ decreases, the BH magnetic flux and jet power $P_j \propto M$ also decrease. As excess magnetic flux leaks out of the BH into the disc, MAD encompasses a larger and larger fraction of the disc. During this process, the jet pushes against the disc and wobbles erratically (see movies in MTB13). As the jet comes in and out of our line of sight, this produces large-amplitude variations (‘bales’) in light curve of Sw J1644+57 during the first $\sim 10^{6}$ s after the trigger (see Fig. 1). [Panel (c)]: Stage 3. Once MAD encompasses the entire disc, the disc/jet alignment with the BH spin completes. Since the jet direction points steadily towards Earth, its X-ray emission becomes less variable as it continues to track the rate of fall-back accretion. [Panel (d)]: Stage 4. Once $M$ decreases below $\sim 30$ percent of the Eddington accretion rate $M_{\text{Edd}}$, the disc transitions to a geometrically thin state and the jet shuts off, producing an abrupt decline in X-ray emission at $t \gtrsim 500$ d. [Panel (e)]: Stage 5. At very late times, once the accretion rate drops below a few percent of $M_{\text{Edd}}$, the disc will again enter a geometrically thick regime, and the jet may turn back on, again analogous to state transitions in X-ray binaries. For Sw J1644+57 this ‘jet revival’ is estimated to occur between the years $\sim 2016–2022$ (Section 5.5).
Various quantities as a function of time since disruption, assuming a partial disruption of an MS star and a delay of $t_{\text{disr}} - t_{\text{delay}} = 60$ d between the disruption and the $\gamma$-ray trigger (Section 2.2). Different stages in our model for Sw J1644+57 are indicated with colour coding and numbers (Section 3; see Fig. 2 for a detailed explanation). [Panel (a)]: mass accretion rate $\lambda = M/M_{\text{Edd}} = f_{\text{acc}} M_\bullet / M_{\text{Edd}}$ as a fraction of the Eddington accretion rate $M_{\text{Edd}}$ (blue solid line; left-hand axis), which peaks at $t \approx \text{days} \lesssim 20$ d and subsequently decays as $\lambda \propto t^{-2.2}$ (see Section 4.2.4 for definitions of mass fall-back rate, $M_\text{fb}$, and the fraction of fall-back mass that reaches the hole, $f_{\text{acc}}$). The same curve (right-hand axis) gives the maximum value of the magnetic flux threading the BH, $\Phi_{\lambda,30} \propto \lambda^{1/2}$, in units of $10^{30}$ G cm$^2$ (red dashed line). As stellar debris returns to the vicinity of the BH, it drags in magnetic flux from a pre-existing, ‘fossil’ accretion disc (Section 5.2). This accumulated flux due to the swept-up fossil field increases with time as $\Phi_{\lambda,30} \propto t^{2/3}$ (green dotted line; equation 36). At early times, most of this flux ends up on the BH, so $\Phi_{\lambda,30} = \Phi_{\lambda,30}$ (red dashed line). For our choice of parameters (see below), BH receives a substantial amount of magnetic flux early on, enough to overcome the ram pressure of the infalling gas and produce the jets, $\Phi_{\lambda,30} > \Phi_{\lambda,30}$ (red dashed line lies above blue dotted line). Eventually, the central magnetic flux becomes dynamically important, i.e. $\Phi_{\lambda,30} \simeq \Phi_{\Lambda,30}$ (red dashed line crosses blue solid line). Since the inner disc can only hold $\Phi_{\Lambda,30}$ worth of flux on the hole (Stage 2), the rest leaks out and instead contributes to the disc flux (cyan long dashed line), i.e. $\Phi_{\Lambda,30} = \Phi_{\lambda,30} - \Phi_{\Lambda,30}$. However, the disc flux eventually also saturates (depending on its radial extent, equation 43), after which it also tracks the BH magnetic flux (Stage 3). Once the Eddington ratio, $\lambda$, falls below a critical value ($\lambda_{\text{crit}} = 0.2$ in this figure) the accretion disc becomes geometrically thin; the central BH and disc lose their magnetic flux; and the jets shutoff (Stage 4). At a much later time, when $\lambda \lesssim 0.02$, the accretion disc becomes geometrically thick again and produces powerful jets (Stage 5). [Panel (b)]: dimensionless values (equation 9) of the magnetic field are shown in panel (a). The magnetic flux accumulated on the BH is strong enough to launch the jets when $\Phi_{\text{c}} > \Phi_{\Lambda,30}$ (see Section 3.2). In this figure, we have assumed the following parameters: BH spin $a = 0.7$; BH mass $M_\bullet = 3 \times 10^6 M_\odot$; stellar mass $M_\star = 0.44 M_\odot$; fraction of star accreted by BH $f_{\text{acc}} = 0.75$; Eddington fraction of the fossil disc $\Phi_{\Lambda,30} = 1.8 \times 10^{30}$; and $\eta_{\text{jet}}/\eta_{\text{acc}} = 0.05$. For computing flux accumulation time-scale at early time (equation 37), we assume disc thickness $h/r = 1$ as expected for a highly super-Eddington flow. Since at later times $h/r$ decreases, for computing jet power we take a single representative value, $h/r = 0.3$.  

Figure 3. Various quantities as a function of time since disruption, assuming a partial disruption of an MS star and a delay of $t_{\text{disr}} - t_{\text{delay}} = 60$ d between the disruption and the $\gamma$-ray trigger (Section 2.2). Different stages in our model for Sw J1644+57 are indicated with colour coding and numbers (Section 3; see Fig. 2 for a detailed explanation). [Panel (a)]: mass accretion rate $\lambda = M/M_{\text{Edd}} = f_{\text{acc}} M_\bullet / M_{\text{Edd}}$ as a fraction of the Eddington accretion rate $M_{\text{Edd}}$ (blue solid line; left-hand axis), which peaks at $t \approx \text{days} \lesssim 20$ d and subsequently decays as $\lambda \propto t^{-2.2}$ (see Section 4.2.4 for definitions of mass fall-back rate, $M_\text{fb}$, and the fraction of fall-back mass that reaches the hole, $f_{\text{acc}}$). The same curve (right-hand axis) gives the maximum value of the magnetic flux threading the BH, $\Phi_{\lambda,30} \propto \lambda^{1/2}$, in units of $10^{30}$ G cm$^2$ (red dashed line). As stellar debris returns to the vicinity of the BH, it drags in magnetic flux from a pre-existing, ‘fossil’ accretion disc (Section 5.2). This accumulated flux due to the swept-up fossil field increases with time as $\Phi_{\lambda,30} \propto t^{2/3}$ (green dotted line; equation 36). At early times, most of this flux ends up on the BH, so $\Phi_{\lambda,30} = \Phi_{\lambda,30}$ (red dashed line). For our choice of parameters (see below), BH receives a substantial amount of magnetic flux early on, enough to overcome the ram pressure of the infalling gas and produce the jets, $\Phi_{\lambda,30} > \Phi_{\lambda,30}$ (red dashed line lies above blue dotted line). Eventually, the central magnetic flux becomes dynamically important, i.e. $\Phi_{\lambda,30} \simeq \Phi_{\Lambda,30}$ (red dashed line crosses blue solid line). Since the inner disc can only hold $\Phi_{\Lambda,30}$ worth of flux on the hole (Stage 2), the rest leaks out and instead contributes to the disc flux (cyan long dashed line), i.e. $\Phi_{\Lambda,30} = \Phi_{\lambda,30} - \Phi_{\Lambda,30}$. However, the disc flux eventually also saturates (depending on its radial extent, equation 43), after which it also tracks the BH magnetic flux (Stage 3). Once the Eddington ratio, $\lambda$, falls below a critical value ($\lambda_{\text{crit}} = 0.2$ in this figure) the accretion disc becomes geometrically thin; the central BH and disc lose their magnetic flux; and the jets shutoff (Stage 4). At a much later time, when $\lambda \lesssim 0.02$, the accretion disc becomes geometrically thick again and produces powerful jets (Stage 5). [Panel (b)]: dimensionless values (equation 9) of the magnetic field are shown in panel (a). The magnetic flux accumulated on the BH is strong enough to launch the jets when $\Phi_{\text{c}} > \Phi_{\Lambda,30}$ (see Section 3.2). In this figure, we have assumed the following parameters: BH spin $a = 0.7$; BH mass $M_\bullet = 3 \times 10^6 M_\odot$; stellar mass $M_\star = 0.44 M_\odot$; fraction of star accreted by BH $f_{\text{acc}} = 0.75$; Eddington fraction of the fossil disc $\Phi_{\Lambda,30} = 1.8 \times 10^{30}$; and $\eta_{\text{jet}}/\eta_{\text{acc}} = 0.05$. For computing flux accumulation time-scale at early time (equation 37), we assume disc thickness $h/r = 1$ as expected for a highly super-Eddington flow. Since at later times $h/r$ decreases, for computing jet power we take a single representative value, $h/r = 0.3$.  

MAD accretion (start of Stage 2), which occurs when $\Phi_{\Lambda,30} \sim 50$. The possibly substantial duration of this early jet smothering phase is illustrated by Fig. 4, which shows the time evolution of $\Phi_{\text{c}}$ in the case of a tidally disrupted white dwarf (WD). In this case, the disruption itself happens on a time-scale of $\sim$ tens of minutes, but Stage 1 begins only at $t \sim 10$ d. Also note that depending on how fast magnetic flux is accumulated, Stage 1 could be even briefer. If $\Phi_{\text{c}} \gtrsim \Phi_{\Lambda,30}$ at peak accretion, then the jet will enter the ‘wobbbling’ stage described next essentially from its onset.

3.3 Stage 2: MAD onset, erratic wobbling jet (Fig. 2b)

As $M$ decreases from its peak value, $\Phi_{\text{c}} \propto M^{-1/2}$ (equation 9) increases and the magnetic field becomes increasingly important dynamically. Once $\Phi_{\text{c}}$ exceeds a critical value $\Phi_{\text{c}} \sim 50$ (depending weakly on BH spin; equation 13), the field is sufficiently strong to give feedback on the accreting gas, leading to a state of ‘magnetically arrested disc’ accretion, or MAD (Narayan et al. 2003; Tchekhovskoy et al. 2011; McKinney, Tchekhovskoy & Blandford 2012). The strong magnetic flux also acts to orient the rotational axis of the inner accretion disc/jet with the BH spin axis (MTB13). This realignment does not, however, happen instantaneously, nor is it clean. The jets undergo a period of vigorous rearrangement during which they ram against the accretion disc, partially obliterate it, and intermittently punch holes through the disc as they work to reorient the disc’s angular momentum along BH spin axis (see Fig. 2b). Recent numerical simulations show that during this stage the jets...
wobble intensely between the orientation of the outer disc axis and the orientation of the BH axis (MTB13), so that jet emission transiently comes in and out of our line of sight. This jet wobbling state may explain the epoch of intense flaring comprising the first ~10 d of the Sw J1644+57 light curve (Section 2.4).

Alignment between the disc/jet and BH spin completes once sufficient magnetic flux is able to leak out of the immediate vicinity of the BH and new flux is brought in by the infalling debris, such that the entire disc (at least out to the circularization radius $R_{\text{circ}}$) becomes MAD. Depending on $R_{\text{circ}}$, this process requires $\dot{\phi}_* \approx 10^{-10}$ to decrease by a factor of a few, and hence $M$ to decrease by a factor of several. In Section 5.3, we show that this time-scale for the entire disc to ‘go MAD’ is broadly consistent with the observed duration of the early flaring state in Sw J1644+57. In Figs 3 and 4 this process is shown by the rising visibility of the disc flux $\Phi_0$ during Stage 2, until it eventually reaches a constant fraction of the BH flux $\Phi_*$ at the onset of Stage 3.

3.4 Stage 3: steady spin-aligned jet (Fig. 2c).

Once the entire disc is MAD, the system enters a scale-free MAD solution that depends on just one parameter: the mass accretion rate (Fig. 2c). The strong centrally accumulated magnetic field not only aligns the jets along the direction of BH spin axis (Section 3.3), it also explains why the jet luminosity faithfully tracks mass fall-back rate. When $\phi_*$ is small at early times (Stage 1), the BH flux $\Phi_*$ and jet power $P_j$ (equation 3) are approximately constant or increase in time (Stage 1), which is inconsistent with the power law decrease in the X-ray light curve of Sw J1644+57 between ~10 and ~500 d since the trigger (Fig. 1). In contrast, when the central magnetic field is sufficiently strong ($\phi_* \gtrsim \phi_*^{\text{MAD}} \sim 50$; equation 13) to be dynamically important (MAD state), then the magnetic flux $\Phi_*$ threading the BH is not determined by initial conditions, but instead by the ram pressure of the accretion flow. Its value regulates such that the jet power $P_j \sim \dot{\eta}_j M c^2$ is a constant fraction $\dot{\eta}_j \approx 1.3\sigma_2$ of the accretion power (equation 17). Thus, as the fall-back rate decreases as a power law in time, so do the jet power and luminosity, as is observed.

Another expected feature of MAD is a stable quasi-periodic oscillation (QPO) at a frequency that is 1/4 of BH horizon angular frequency (McKinney et al. 2012). Reis et al. (2012) detected a potential QPO with period $\tau_{\text{QPO}} \sim 200$ s in the power-law decay portion of the Sw J1644+57 light curve, which could also be evidence for MAD. We discuss the constraints implied by this period (Section 4.2.6) as a part of our analysis in Section 4.

3.5 Stage 4: no jet (Fig. 2d).

Once the accretion rate drops below a few tens of percent of $M_{\text{Edd}}$, the flow transitions to a geometrically thin disc state, or thermal state, which is not observed (nor theoretically expected) to power a jet (e.g. Fender, Belloni & Gallo 2004; Russell et al. 2011). The abrupt decrease in the X-ray flux at $t \sim 500$ d (Fig. 1; by more than two orders of magnitude) can thus plausibly be associated with the point at which $M \sim 0.3 M_{\text{Edd}}$ (Steiner et al. 2009, 2010; Abramowicz et al. 2010). The jet luminosity and time-scale of this transition also constrain the properties of the disrupted object and central BH (Section 4.2.3).

However, jet emission can shut off if the jets become misaligned with our line of sight before the accretion flow transitions to a geometrically thin jetless state. Such a misalignment can occur because as mass accretion rate decreases, the disc dimensionless thickness, or the ratio of height to radius, $h/r$ decreases as well. Since $h/r$ directly determines the dimensionless strength of BH’s magnetic flux $\phi_*$ (see equation 13), which is responsible for the alignment of the jets along the direction of BH spin (MTB13), if $\phi_*$ falls below a certain value, jet alignment might be no longer possible. MTB13 observed efficient jet alignment for two values of disc thickness: $h/r \approx 0.3$ and 0.6. While by continuity such an alignment is clearly possible for somewhat thinner discs, the process of magnetic alignment of such discs is not understood. It is possible that the critical value of $h/r$ at which such alignment no longer occurs is similar to that of the state transition to a geometrically thin disc. For the purposes of this work, we will therefore not make a distinction between the two jet shutoff scenarios described above.

3.6 Stage 5: jet revival (Fig. 2e).

As $M_0$ continues to decrease as a power law in time, eventually it will reach a few per cent of $M_{\text{Edd}}$. After this point the disc will again transition to a geometrically thick disc, analogous to the ‘low/hard’ state observed in X-ray binaries. Since this state is conducive to jet formation (Narayan & Yi 1995; Fender et al. 2004), the jet in Sw J1644+57 and its associated X-ray emission may suddenly turn back on (this is estimated to occur sometime around 2016–2022; Section 5.5).

4 CONSTRAINTS ON Sw J1644+57

4.1 Stellar progenitor scenarios

We begin by overviewing the possible classes of stellar progenitors which could be responsible for Sw J1644+57. TD scenarios usually consider the disruption of a lower MS star. However, in principle the star could have been a giant (MacLeod, Guillochon & Ramirez-Ruiz 2012), WD (e.g. Krolik & Piran 2011; Haas et al. 2012; Shcherbakov et al. 2013), or even a planet. Low-mass planets seem to be ruled out because the (beaming-corrected) energy budget of Sw J1644+57 of $\gtrsim 10^{31}$ erg already exceeds the rest mass energy of Jupiter. A giant star is also unlikely, because the fall-back time of the stellar debris would greatly exceed the observed trigger time (Section 4.2.1), incompatible with the observed X-ray light curve. However, a WD companion cannot be ruled out a priori, especially considering their potential for harbouring a large magnetic flux.

In what follows, we thus consider constraints based on two progenitor scenarios: a lower MS star and a WD. We employ approximate mass–radius relations for each given by (Nauenberg 1972; Shapiro & Teukolsky 1986)

$$r_s = m_s, \quad (\text{MS})$$

$$r_s = 0.013(m_s/0.6)^{-1/3}, \quad (\text{WD}),$$

(10)

where $m_s$ is in solar masses and $r_s$ is in solar radii. Since the stellar radius must exceed the tidal radius, only low-mass BHs with $M_\star \lesssim 10^7$–$10^9 M_\odot$ are capable of disrupting a WD, depending on WD mass and BH spin (Kesden 2012). For a centrally concentrated
MS star, the disruption process can be either ‘full’ or ‘partial’ (Guillochon & Ramirez-Ruiz 2013), resulting in different predictions for the rate of fall-back accretion (Section 4.2.1).

In what follows, we thus consider three scenarios: (i) a complete or (ii) partial TD of a lower mass MS star by a supermassive BH; or (iii) a complete disruption of a WD by an intermediate-mass BH.

### 4.2 Observational constraints

For each progenitor scenario, our goal is to determine the following four unknowns: BH mass, $M_\star$, BH spin, $\alpha$, stellar mass, $m_\star$, and magnetic flux threading the hole, $\Phi_\ast$. We now discuss what observational data constrain these properties within the framework of our model for Sw J1644+57.

#### 4.2.1 Shape of the X-ray light curve

First consider what constraints can be placed on Sw J1644+57 based on the shape of the X-ray light curve. Recall that although the time-averaged light curve does not follow a single power-law decay if plotted against the time since the $\gamma$-ray trigger $t - t_{\text{trig}}$, a much better fit is achieved by moving the TD ‘zero-point’ prior to the trigger (Fig. 1). We thus first consider what range of trigger delay time $t_{\text{trig}} - t_{\text{dist}}$ produces a light curve shape consistent with the predicted power-law decline in the X-ray luminosity $L_X \propto \dot{M}_h \propto (t - t_{\text{dist}})^{-\alpha}$ given the expected range in $\alpha$.

Although the canonical value of $\alpha = 5/3$ is often quoted, $\alpha$ actually depends on the fraction of mass lost by the star during TD. $\dot{M}_h / \dot{M}_\star = 0.4 \alpha$, which differs between complete and partial disruptions (Guillochon & Ramirez-Ruiz 2013). If $\dot{M}_h / \dot{M}_\star \leq 50$ per cent then $\alpha = 2.2$, but if $\dot{M}_h / \dot{M}_\star \geq 50$ per cent then the value of $\alpha$ approaches that for a complete disruption, $\alpha = 5/3$ (Guillochon & Ramirez-Ruiz 2013). By considering the two limiting cases of complete disruptions with $\alpha = 5/3$ and partial disruptions with $\alpha = 2.2$, we bracket the allowed range of possibilities.

For complete disruptions ($\alpha = 5/3$) we find an allowed range of $t_{\text{trig}} - t_{\text{dist}} \sim 15_{-10}^{+15}$ d when fitting to the shape of the time-averaged X-ray light curve. Fig. 5 shows the trigger delay-corrected light curve for a fiducial value $t_{\text{trig}} - t_{\text{dist}} = 15$ d. For partial disruptions ($\alpha = 2.2$), we find required delay times that are somewhat longer during the ‘wobbling jet’ stage (Stage 2) as opposed to 15 d when fitting to the shape of the time-averaged light curve (see also Fig. 5). Note that just this simple shift in the zero-point in time causes the early time ‘plateau’ in the light curve (see Fig. 1) to disappear and the entire light curve to become consistent with a single power law dependence in time, $\propto t^{-3/5}$. Various stages of the disruption process are indicated with colour coding and text. These stages are explained in Fig. 2 and its caption.

![Figure 5](https://example.com/fig5.png)

**Figure 5.** Sw J1644+57’s soft X-ray light curve accounting for the best-fitting time of the disruption, $t_{\text{trig}} - t_{\text{dist}} = 15$ d before the trigger, for a complete disruption (see also Fig. 1). Note that just this simple shift in the zero-point in time causes the early time ‘plateau’ in the light curve (see Fig. 1) to disappear and the entire light curve to become consistent with a single power law dependence in time, $\propto t^{-3/5}$. Various stages of the disruption process are indicated with colour coding and text. These stages are explained in Fig. 2 and its caption.

For partial disruptions ($\alpha = 2.2$) we find an allowed range of $t_{\text{trig}} - t_{\text{dist}} \sim 30_{-15}^{+30}$ d when fitting to the shape of the time-averaged X-ray light curve. Fig. 6 shows the corrected light curve in this case for a fiducial delay $t_{\text{trig}} - t_{\text{dist}} = 30$ d.

Our fits quoted above (Figs 5 and 6) were made by matching the fall-back accretion rate $\alpha$ jet power to the time-averaged X-ray flux. However, if the early phase of high variability indeed results from a ‘wobbling’ jet (Section 3.3), then bright flares may arise from transient episodes when the jets points towards our line of sight. In this case, if we assume that the beaming factor remains constant throughout the event (and is therefore the same in the initial ‘flaring state’ and in the late time ‘steady spin-aligned jet’ state), it is more physical to match the fall-back rate $\alpha$ jet power to the ‘envelope’ in the X-ray light curve comprised by the flare peaks, since these better characterize the true jet power. Since the observed duty cycle of the flaring state is $\sim 10$ per cent, this implies that the jet points towards us only about 1/10 of the time if the wobbling scenario is correct. The total jet power (accounting for both on- and off-axis jet emission) could thus exceed the observed (on-axis) emission by an order of magnitude (see also Section 5.4). Figs 1 and 7 show that such fits would allow for a trigger delay $t_{\text{trig}} - t_{\text{dist}}$ as short as a few days, consistent with the first evidence for a jet in Sw J1644+57 $\sim 4$ d prior to the $\gamma$-ray trigger (Krimm & Barthelmy 2011).

However, there is no reason to expect that the beaming angle remains the same during the ‘wobbling jet’ stage (Stage 2) and the ‘steady spin-aligned jet’ stage (Stage 3). This is because in the ‘wobbling jet’ stage the jet has to drill fresh holes in the disorganized flow of material. This causes the jet in this stage to be much more strongly collimated than in the ‘spin-aligned’ stage in which the jet freely escapes via a pre-drilled channel along the BH axis (as was demonstrated using first-principles GRMHD simulations, see MTB13). In particular, the opening angle of the jet at $r = 40\, r_\text{g}$ is $5\degree$ during the ‘wobbling jet’ stage (MTB13) as opposed to $15\degree$
Figure 7. Soft X-ray light curve of Sw J1644+57 for a complete disruption in fitting the peaks in the light curve as opposed to the time-averaged X-ray flux (as was done in Figs 5 and 6). Such fits allow for a trigger delay time as short as a few days (fit for $t_{\text{lag}} - t_{\text{sh}} = 5$ d is shown). We consider this scenario to be unlikely as the jet is more tightly collimated during the ‘wobbling’ Stage 2 than in the late time light curve see MTB13. This is because at early times the jet, which is forced to drill holes through the misaligned accretion flow, will be more tightly collimated than at later times, when it is spin aligned and can escape along the poles without obstruction. This means that the effective beaming correction $f_0$ is higher in the ‘wobbling’ stage. We implicitly take this into account by fitting the average light curve, as shown in Figs 1, 5 and 6.

when the jet is steady (Tchekhovskoy & McKinney, 2012). The factor of 3 tighter collimated jet during the wobbling stage gives a factor of $\sim 10$ higher beaming factor and justifies fitting the average light curve which effectively gives a fit at the $\sim 10$ per cent level of the peaks. Also, while a trigger time delay of the order of several days appears consistent with the X-ray data shown in Fig. 1, we will see that it requires the presence of an intermediate-mass BH, a less likely possibility than a supermassive BH. For these reasons, we consider scenarios with longer trigger time delays to be more likely.

Given the allowed range in $t_{\text{lag}}$, we require that the fall-back accretion rate must be past its peak at the trigger:

$$t_{\text{peak}} \approx t_b < t_{\text{lag}}/(1+z),$$

where the fall-back time (equation 4) specialized to WD and MS scenarios is given by

$$t_b = 0.02 \, M_{\odot}^{1/2} \, t_{\text{acc}}^{-3/2}, \quad \text{(WD)}$$

$$t_b = 17.3 \, M_{\odot}^{1/2} \, t_{\text{acc}}^{-1/2}, \quad \text{(MS)},$$

(12a)

(12b)

since otherwise the light curve would still be rising and hence would not match the $L_X \propto M_b \propto (t/t_b)^{-1}$ decay predicted for $t \gg t_b$.

4.2.2 Magnetic flux

Although it does not represent an independent constraint, the magnetic flux $\Phi_1$ is determined once the spin of the BH is known, if one assumes that the jet is in a MAD state after the trigger, up until the point when the jet shuts off. During MAD accretion, the BH magnetic flux is regulated by the mass accretion rate to a dimensionless value (Tchekhovskoy, McKinney & Narayan, 2012; Tchekhovskoy & McKinney, in preparation), which is well approximated as in a MAD:

$$\Phi_{\text{MAD}}^{\text{MAD}} \equiv \frac{\Phi_{\text{MAD}}}{(M^2 c^3)^1/2} \approx \frac{70(1 - 0.38 \bar{\omega}_0) \bar{\eta}_0^{1/2}}{h_0^{1/3}},$$

(13)

corresponding to absolute flux

in a MAD:

$$\Phi_{\text{MAD}}^{\text{MAD}} = 0.067 \, M_{\odot}^{1/2} \, h_0^{1/2} \, (1 - 0.38 \bar{\omega}_0) \bar{\eta}_0^{1/2},$$

(14)

where $\lambda = M_{/\text{Edd}}$ is the Eddington ratio and $h/r = 0.3 h_0$ is the thickness of the accretion flow. Thus, if the dimensionless BH spin and the jet luminosity at a given Eddington ratio are known (such as at the point of jet shut-off; see equation 15 below), then $\Phi_1$ can be determined.

4.2.3 Jet shut-off power

Another constraint is that the time at which the jet was observed to shut-off $t_{\text{off}} - t_{\text{lag}} \approx 500$ d (see Fig. 1) happens simultaneously with the expected state transition occurring at a fraction of the Eddington accretion rate (Section 2.5), i.e.

$$\dot{M}_{\text{cr}} = 0.3 \frac{\lambda_{\text{cr}}}{0.3} \dot{M}_{\text{Edd}},$$

(15)

where theoretical and observational uncertainties place the threshold value in the range $0.1 \lesssim \dot{\lambda} \lesssim 0.5$.

Just prior to when the jet shut off, its X-ray luminosity $\propto \dot{M}$ power $P_{\text{off}}$ was $\sim 200$ times smaller than its value at the end of the $\sim 10$ d plateau phase after the trigger,

$$P_{\text{off}} \approx (1/200) P_{\text{lag}} \approx 5 \times 10^{41} \frac{f_0 \bar{\eta}_0 \bar{\omega}_0^{-1}}{0.03} \mbox{ erg s}^{-1}.$$  

(16)

Now, by combining equations (3) and (13), the jet power can be written as

in a MAD:

$$P_j \approx F_{(0.9)} h_0 \bar{M} c^2 \approx 1.3 h_0 \bar{M} c^2 \frac{0.3}{0.3} \approx 1.6 \times 10^{44} h_0 \bar{M} c^2 \frac{0.3}{0.3},$$

(17)

where $\lambda = M_{/\text{Edd}}$ and we used the fact that the spin-dependent factor entering jet power, $F_{(0.9)} = 4.4 \bar{\omega}_0^2 (1 - 0.38 \bar{\omega}_0)^2$, can be approximated as $F \approx 1.3 a^2$ to 10 per cent accuracy for $0.3 \leq a \leq 1$ (Tchekhovskoy & McKinney, in preparation).

Matching the jet power with the observed power (equation 16) at $M = M_{\text{cr}}$ (equation 15) thus requires

$$a^2 M_{\odot} \approx \frac{f_0 \bar{\eta}_0 \bar{\omega}_0^{-1}}{0.03} \frac{0.3}{0.3},$$

(18)

This constraint is independent of the nature of the disrupted object.

4.2.4 Jet shut-off accretion rate

Another constraint is that the BH mass accretion rate must equal the critical accretion rate (equation 15) at the observed jet shutoff time, $t_{\text{off}} - t_{\text{lag}} \approx 500$ d, viz

$$M(t_{\text{off}}) = \dot{M}_{\text{cr}},$$

(19)

where $M = f_{\text{acc}} M_b$, and the predicted fall-back accretion rate (equations 4–6) specialized to the MS and WD scenarios:

$$\dot{M}_{\text{off}} = 4.2 \times 10^{28} M_{\odot}^{1/3} M_{\odot}^{4/3} t_0 - 3/3 \, \mbox{ g s}^{-1}, \quad \text{(MS, complete)}$$

$$\dot{M}_{\text{off}} = 1.2 \times 10^{29} f_0 M_{\odot}^{1/3} M_{\odot}^{2/3} t_0 - 2/2 \, \mbox{ g s}^{-1}, \quad \text{(MS, partial)}$$

$$\dot{M}_{\text{off}} = 4.6 \times 10^{29} M_{\odot}^{1/3} M_{\odot}^{2/3} t_0 - 2/2 \, \mbox{ g s}^{-1}, \quad \text{(WD, complete)}$$

(20)

(21)

(22)

where $t_0 \approx t \, d^{-1}$. 


The factor $f_{\text{acc}} < 1$ accounts for the possibility that only a fraction of the fallback material actually reaches the BH horizon, with the rest either expelled in the form of the accretion disc winds and or lost during the circularization of the tidal streams. Super-Eddington accretion is susceptible to outflows driven by radiation pressure (e.g. Ohsuga et al. 2005, Strubbe & Quataert 2009), but the magnitude of this effect is uncertain. Fortunately, we find that the lower limit on $f_{\text{acc}}$ is not constraining (Section 4.3), so without the loss of generality we allow the full range $0 < f_{\text{acc}} \leq 1$ in our calculations.

4.2.5 Spin sufficient for alignment

We require that the spin of the BH is sufficiently high,

$$a \gtrsim 0.5,$$

such that the BH is able to magnetically align the disc and the jets (MTB13). We note that the alignment in the simulations was observed over a limited range in radius (MTB13): essentially the entire MAD part of the flow, $r \lesssim 30r_g$, became aligned with the direction of BH spin. It is unknown if the size of the aligned region would increase, if the simulations were run for a longer period of time and the MAD region grew in size. In this work, we assume that provided equation (23) is fulfilled, the entire MAD region of the disc and the jets become aligned with the direction of BH spin (however, see Section 3.5). Thus, once the entire disc ‘goes MAD’, the jets in a time-averaged sense point straight along the direction of BH spin.

4.2.6 MAD QPO

A final possible constraint is to match the QPO period measured in the power-law X-ray light curve of Sw J1644+57 (Reis et al. 2012),

$$\tau_{\text{QPO}} = 210 \pm 30 \text{ s},$$

with the predicted MAD QPO, which occurs at a frequency that is 1/4 of BH horizon angular frequency (McKINNEY et al. 2012), $\Omega_{\text{BH}} = 2\pi/(2\nu_{\text{QPO}}) = \omega_{\text{BH}}/(2r_g)$, resulting in a predicted period

$$\tau_{\text{MAD}} = 2\pi / 0.25\Omega_{\text{BH}} = \frac{16\pi r_g}{5\omega_{\text{QPO}} c} = 24.8 M_{\text{BH}} \omega_{\text{QPO}}^{-1} \text{ s}.$$ (25)

Equating (25), multiplied by $(1+z) \approx 1.353$, with (24) gives a final constraint,

$$M_{\star} \lesssim (6.3 \pm 0.9) M_{\odot}.$$ (26)

We show below that this requirement is constraining only in the WD scenario.

4.3 Results

Equations (11), (18), (19), (23) and (26) provide four or five constraints on the unknown parameter space ($M_\star$, $M_\bullet$, $a$), depending on whether one adopts the (speculative) QPO constraint described in Section 4.2.6. Once $M_\bullet$ and $a$ are determined, the magnetic flux $\Phi_\ast$ follows from equation (14). We now apply these constraints individually to each of our proposed scenarios (Section 4.1): Sections 4.3.1, 4.3.2 and 4.3.3.

4.3.1 Complete disruption of a main-sequence star

Fig. 8 summarizes the constraints on the BH mass and spin for the complete TD of a low-mass MS star. The first constraint, based on the shape of the X-ray light curve (equation 11), can be written as

$$M_{\star} < 3 m_\ast^{-1},$$ (27)

where we have adopted the highest value for $t_{\text{short}} = 30 \text{ d}$ that allows us to reproduce the observed shape of the X-ray light curve (Figs 1 and 5). This constraint is shown in Fig. 8 as the dark (light) red shaded regions for $m_\ast > 1$ ($m_\ast > 0.1$), respectively. The QPO period constraint [equation (26)] gives the green curve. The common region seems to favour a lower mass star, $m_\ast \sim 0.1$, and a light BH, $M_{\bullet} \sim 0.5$, with $a \gtrsim 0.5$.

![Sw J1644+57 gone MAD](https://example.com/swj1644_57_gone_mad.png)

Figure 8. Constraints on BH mass, $M_\bullet$, and spin, $a$, in the scenario of a complete disruption of a lower mass MS star. The constraint on jet power at jet shut off [equation (18)] gives the blue shaded regions, with the darker shaded region more likely. The constraint on disruption time-scale [equation (11)] gives the dark (light) shaded red regions for $m_\ast > 1$ ($m_\ast > 0.1$), respectively. The QPO period constraint [equation (26)] gives the green curve. The common region seems to favour a lower mass star, $m_\ast \sim 0.1$, and a light BH, $M_{\bullet} \sim 0.5$, with $a \gtrsim 0.5$. The second constraint (equation 18) cuts out a stripe in the $M_\ast-a$ plane, shown as a blue coloured region in Fig. 8. The width of the dark (light) blue region reflects an optimistic (conservative) factor of 20 (1000) uncertainty in the value of the right-hand side of equation (18). The chief effect of this constraint is to place a lower limit on the BH mass and spin, the latter consistent with the fourth constraint (Section 4.2.5). The third constraint (equation 19) does not place an interesting limit on BH mass or spin due to the uncertainty in the fraction of the stellar material accreted $f_{\text{acc}}$. However, it does pick out a preferred range in the value of $f_{\text{acc}}$ given the other parameters

$$f_{\text{acc}} = 0.02 \frac{\lambda_{\ast}}{0.3} M_{\ast}^{2/3} m_\ast^{-4/3} < 0.024 \frac{\lambda_{\ast}}{0.3} m_\ast^{-2},$$ (28)

where in the inequality we have used equation (27). Equation (28) shows that if the tidally disrupted star was solar-like ($m_\ast \sim 1$), then large mass-loss is required, as could be the result of outflows from the disc or at the impact point of tidal streams. Alternatively, a higher...
value $f_{\text{acc}} \sim 1$ is allowed if the progenitor is instead a low-mass M star $m_0 \sim 0.1-0.2 M_{\odot}$ near the hydrogen-burning limit.

The fifth constraint, on the QPO frequency, produces an allowed region shown with green in Fig. 8, which is consistent with (and does not appreciably alter) our conclusions above. If taken seriously, this constraint places an upper limit on the BH mass $M_\bullet < 6 \times 10^5 M_{\odot}$, but otherwise a wide range of BH spin, $a \gtrsim 0.5$, is allowed.

In the above, we assumed a power-law fit through the time-averaged X-ray light curve. An alternative but less likely possibility, as we discuss in Section 4.2.1, is to obtain a fit through the peaks of the light curve, as shown with the dashed red line in Fig. 1. For $t_{\text{triq}} = 5$ d. This is a factor of $\sim 6$ shorter fall-back time than in the scenario of a complete disruption of an MS star (5 d instead of 30 d, see Section 4.3.1). By equation (11), this would require a factor $6^2 \sim 40$ times smaller BH than given by equation (27), $M_\bullet < 8.3 \times 10^5 M_{\odot} m_0^{-1}$, i.e. an intermediate-mass BH. This upper limit is marginally consistent with the lower limit on the hole mass due to equation (18).

### 4.3.2 Partial disruption of a main-sequence star

Fig. 9 summarizes the constraints on the BH mass and spin based on our second scenario, the partial TD of a lower mass MS star. As in equation (27), the first constraint can be written as

$$M_{\bullet,5} < 12 m_0^{-1} \quad (29)$$

where we have taken $t_{\text{triq}} = 60$ d (again as the maximum allowed by fitting the observed shape of the light curve; Section 4.3.1). Constraint (29) is a factor of several less restrictive than that for a complete disruption (equation 27). For example, the data now allow a relatively massive BH with $M_\bullet \gtrsim \text{few } 10^5 M_{\odot}$, even for a Sun-like star, $m_0 \sim 1$.

The region allowed by the second constraint (equation 18; again shown in blue) is exactly the same as for the full disruption (Fig. 8).

The QPO condition is again only moderately constraining, placing an upper limit $M_\bullet,5 < 6 \times 10^5 M_{\odot}$ on the BH mass.

Similar to the full disruption, the third constraint does not place any interesting limits on the BH or stellar parameters, but it does tell us the fraction

$$f_{\text{acc}} = 0.2 f_{\text{bd}} \frac{\lambda_\gamma}{0.3} M_\bullet^{2/5} m_0^{-8/5} \lesssim 0.39 f_{\text{bd}} \frac{\lambda_\gamma}{0.3} m_0^{-2} \quad (30)$$

of the fallback material reaching the BH, where equation (29) is used in the last inequality. For low-mass stars the required value of $f_{\text{acc}}$ may even exceed unity, potentially ruling out such progenitors depending on the value of $\lambda_\gamma \sim 0.1-0.5$. This behaviour is opposite to that in the complete MS disruption scenario, where we were forced to conclude that $f_{\text{acc}} \lesssim 1$ (equation [28]). As we discussed in Section 4.2.1, the two extreme possibilities – of a complete and partial disruption – bracket a continuous family of allowed solutions that, by continuity, are consistent with the data.

Just like in a scenario of complete stellar disruption (Section 4.3.1), a (less likely, see Section 4.2.1) fit through the peaks of the early time light curve instead of the average light curve provides a reasonably good description of the light curve (dashed blue line in Fig. 1). In this scenario, the trigger time delay is 5 d. This is much shorter than the delay time of 60 d discussed in the fiducial scenario in the rest of this subsection. Consequently, equation (29) requires an intermediate-mass BH, $M_\bullet < 8.3 \times 10^5 M_{\odot} m_0^{-1}$, i.e. the same as for the complete disruption scenario.

### 4.3.3 Complete disruption of a white dwarf

Fig. 10 summarizes the constraints on the BH properties based on our third scenario, the disruption of a WD. Since a WD is much denser than an MS star, the fall-back time in the WD scenario is much shorter than in the MS scenario. In fact, $t_0 \ll t_{\text{triq}}$, where $t_{\text{triq}} \gtrsim 5$ d is required to explain the shape of the X-ray light curve to within a factor of few (see Section 4.2.1). For these reasons, the first constraint, (11), does not place interesting limits on the system parameters.

The most important constraint is that on the accretion rate, equation (19), which gives

$$f_{\text{acc}} = 2 \frac{\lambda_\gamma}{0.3} M_\bullet^{2/5}. \quad (31)$$

The physical requirement that $f_{\text{acc}} < 1$ places an upper limit on the BH mass, $M_\bullet,5 \lesssim 2$ (for $\lambda_\gamma > 0.1$), thus requiring an intermediate-mass BH.

Since the mass fall-back rate is independent of the WD mass (equation 22), this makes the second constraint, equation (18), on jet power particularly constraining on the BH spin,

$$a = 1.6 \left( \frac{f_{\text{bd}} m_0^{1/3}}{0.03} \right)^{1/2} \left( \frac{\lambda_\gamma}{0.3} \right)^{1/4} f_{\text{acc}}^{3/4} m_0^{-1/2}. \quad (32)$$

The constraints resulting from equations (31) and (32) are illustrated in Fig. 10 with the dark (light) blue area, whose width reflects optimistic (conservative) modelling and observational uncertainties of a factor of 20 (100) in the right-hand side of equation (18). The most likely (dark blue region) favours a high BH spin, $a \gtrsim 0.9$.

The addition of the less certain QPO constraint (equation 26) in combination with the spin constraint (equation 23) favours a BH mass in the range $M_\bullet,5 \sim 0.6-2$ with an intermediate spin $a \simeq 0.5$. However, we caution against overinterpreting the QPO constraint, as the observed QPO could be caused by other processes than those resulting from MAD accretion (McKinney et al. 2012).
observed jets, which thus greatly restricts the allowed parameter space (Section 4.3.3).

Although the WD scenario is formally allowed, there are several reasons to favour the MS scenario. In addition to the narrow allowed parameter space in the WD case, the probability of WD disruption is much lower than in the MS since its smaller tidal radius limits the allowed range of impact parameters capable of resulting in disruption. Furthermore, the intermediate-mass BH required in the WD case is an exotic, and probably rare, object (Greene 2012); it is also unclear how frequently such an object would be expected to reside near the nucleus of a galaxy. Given the convergence of several rare events in the WD scenario, we conclude that an MS disruption is more likely. The existence of plausible and constrained solutions provides a consistency check on our model for Sw J1644+57.

In addition to jetted emission, TDEs are accompanied by thermal optical/UV/soft X-ray emission from the accretion disc (e.g. Ulmer 1999) or super-Eddington outflows (Strubbe & Quataert 2009, 2011). No such optical/UV emission was detected from Sw J1644+57, possibly due to substantial dust extinction (Bloom et al. 2011). However, if an otherwise similar event with less extinction were to occur in the future, a measurement of the disc flux just after the jet shuts off would directly determine $\lambda_{\text{ed}}$, the critical Eddington ratio at which the disc transitions from being thick to thin. Such a measurement would also better determine the properties of the BH and disrupted star.

5.2 Origin of the magnetic flux

Regardless of the nature of the disrupted star, the magnetic flux threading the BH must be sufficient to power jet responsible for Sw J1644+57. Equation (2) can be written as

$$\Phi_{\text{mag}}^{\text{BH}} \approx 0.4 M_{*} \left( \frac{f_{\text{j}}}{10^{48} \text{erg s}^{-1}} \right)^{1/2} \left( \frac{\omega_{\text{H}} f^{1/2}(\omega_{\text{H}})}{0.64} \right)^{-1},$$

where $\omega_{\text{H}} f^{1/2}(\omega_{\text{H}})$ is normalized to its value for BH spin $a = 0.9$. If the onset of MAD indeed occurred near the time of the γ-ray trigger (at Eddington ratio $\lambda_{\text{ed}} \approx 200 \lambda_{\text{cr}} \approx 60 \lambda_{\text{cr}}/0.3$; see Section 4.2.3), then we require $\Phi_{\text{mag}}^{\text{BH}} \approx 0.1-10$ for BH masses $M_{*} \sim 0.1-0.1$ consistent with our modelling of Sw J1644+57 (Figs 8–10). This flux is $\sim 4-6$ orders of magnitude greater than that through a solar-type star $\Phi_{\gamma} \sim \pi R_{\odot} L_{\odot} = 10^{37} \text{erg s}^{-1} (B_{z}/k\text{G})(R_{\odot}/R_{\odot})^{2}$ G cm$^{-2}$, even for an optimistically large $B_{z} \sim$ kG average stellar magnetic field. Likewise, a WD $(R_{\odot} \sim 0.1 R_{\odot})$ would require a field $B_{z} \gtrsim 10^{11}$ G for $M_{*} \sim 1$ (Fig. 10) which exceeds the largest measured (surface) fields by two orders of magnitude (e.g. Kepler et al. 2013).

If the star itself cannot explain the flux, what could be its source? It has been suggested (Krolik & Piran 2011, 2012) that the requisite flux is generated by a turbulent dynamo in the accretion flow. However, since the net vertical magnetic flux is a conserved quantity in ideal MHD, then the vertical component of the field must undergo a random walk about zero (if not, then what determines a preferred direction?), such that the BH periodically receives patches of random polarity. The characteristic time-scale between such flips in the mean field would presumably be set by the accretion time-scale near the outer radius of the disc $R_{\text{acc}}$, which is $t_{\text{acc}} \sim 10^{5}$ s and $\sim 10^{6}$ s in the case of WD and MS stars, respectively (see Section 5.3). Without a large-scale magnetic flux to produce a sustained jet, such polarity flips would cause the jet power to transiently switch off each time the flux changes sign (Beckwith, Hawley & Krolik 2008; McKinney & Blandford 2009; McKinney et al. 2012) at
characteristic intervals $\sim t_{\text{acc}}$. Such large-scale variability is not obviously seen in Sw J1644+57 light curve after the first $\sim 10\,\text{d}$, once it has settled into a MAD state. Instead, the fact that Sw J1644+57 was continuously active over nearly $1.5\,\text{yr}$ may suggest that the BH must be threaded by large-scale magnetic flux of the same sign. Although it is difficult to rule out a dynamo process completely, since the physics of large-scale magnetic field generation in accretion discs is at best poorly understood, we instead focus on the possibility that a reservoir of large-scale magnetic flux is required.

A possible source of large-scale flux is that contained in a pre-existing (‘fossil’) accretion disc, which was present at the time of TD but was not detectable in pre-imaging of the host of Sw J1644+57 since its accretion rate was relatively low (Eddington ratio $\lambda_{\text{fossil}} \approx 10^{-2}$). During the disruption process, stellar debris is flung outwards on to a series of highly eccentric orbits with an apocentre radius (e.g. Strubbe & Quataert 2009)

$$r_a \approx \frac{\text{ct}}{2\pi r_s} \approx 1.3 \times 10^5 M_{\ast,5}^{-1/3} m_s^{-2/3} r_e \left( \frac{t}{t_{\text{fb}}} \right)^{2/3}$$

that increases for material with fall-back times longer than that of the most bound stellar debris $t_{\text{fb}}$, where the prefactor in equation (34) is calculated for a solar mass star. Thus, as debris returns to the BH, it sweeps up a significant fraction$^5$ of the magnetic flux threading the fossil disc at radii $r < r_a(t)$.

We now estimate how luminous the jet from such a fossil disc would have in order to supply the required flux, assuming that the fossil disc is itself in a MAD state. In a MAD with a vertical thickness $h/r \approx 0.3 - 0.6$ (as characterizes both super-Eddington and highly sub-Eddington accretion, every $\Delta r \sim 30 h_0 r_g$ of the accretion disc contains as much magnetic flux as the BH itself (Tchekhovskoy & McKinney 2012; MTB13), where $h/r = 0.3 h_0$. Thus, the magnetic flux contained by a MAD fossil disc out to a distance $r$ is given by (McKinney et al. 2012; Tchekhovskoy & McKinney 2012),

$$\Phi_{\text{fossil}}(r) \approx \frac{r}{30 r_s} \left( \frac{\lambda_{\text{fossil}}}{\lambda_{\text{MAD}}} \right)^{1/2} h_0^{-1/2} \Phi_{\ast,5}^{\text{MAD}}$$

where we used equation (14) to relate the flux threading the BH by the quiescent disc to that in the MAD state of Sw J1644+57, viz. $\Phi_{\ast,5}^{\text{MAD}} = (\lambda_{\text{fossil}}/\lambda_{\text{MAD}})^{1/2} \Phi_{\ast,5}^{\text{MAD}}$. Since the apocentre distance of the infalling debris increases as $r_e \propto t^{2/3}$ (equation 34), the cumulative amount of ‘fall-back’ magnetic flux, which is brought to the BH by the infalling tidal streams, is given by

$$\Phi_{\text{fossil}}^0(t) \approx 0.43 \left( \frac{\lambda_{\text{fossil}}/\lambda_{\text{MAD}}}{10^{-6}} \right)^{1/2} \Phi_{\ast,5}^{\text{MAD}} M_{\ast,5}^{-1/3} m_s^{-2/3} r_e^{-1/3} \left( \frac{t}{t_{\text{fb}}} \right)^{2/3}$$

$$\approx 0.38 \left( \frac{\lambda_{\text{fossil}}/\lambda_{\text{MAD}}}{10^{-6}} \right)^{1/2} M_{\ast,5}^{2/3} m_s^{-2/3} r_e^{-1/3} \times \left( \frac{P_j}{10^{46} \text{ erg s}^{-1}} \right)^{1/2} h_0^{-1/2} \left( \frac{t}{t_{\text{fb}}} \right)^{2/3}$$

where in the last line we have used equation (33) and have assumed $\alpha = 0.9$. Figs 3 and 4 show the time evolution of $\Phi_{\ast,5}$ in fiducial MS partial disruption and WD scenarios, respectively.

By demanding in equation (36) that at $t = t_{\text{peak}} \approx t_{\text{fb}}$ the accreted flux $\Phi_{\ast,5}$ equals $\Phi_{\ast,5}^{\text{MAD}}$, i.e. sufficient to magnetically arrest the disc at $t = t_{\text{MAD}}$, this places a lower limit on the Eddington ratio of the fossil disc (using $\lambda_{\text{MAD}} \approx 60$):

$$\lambda_{\text{fossil}} > 3.5 \times 10^{-4} M_{\ast,5}^{3/5} m_s^{-1/3} r_e^{-2/3} (t_{\text{peak}}/t_{\text{MAD}})^{2/3}$$

corresponding to a jet power of the fossil disc given by (equation 17)

$$P_j^{\text{fossil}} \gtrsim 5 \times 10^{40} M_{\ast,5}^{3/5} m_s^{-1/3} r_e^{-2/3} (t_{\text{peak}}/t_{\text{MAD}})^{2/3} \text{ erg s}^{-1}$$

If we adopt bolometric and beaming corrections similar to that applied in Sw J1644+57, then the resulting X-ray luminosity $L_X \approx 1.5 \times 10^{42} M_{\ast,5}^{3/5} m_s^{-1/3} r_e^{-2} \text{ erg s}^{-1}$ to $t_{\text{peak}}$ is comfortably consistent with ROSAT upper limits $L_X \lesssim 2 \times 10^{44} \text{ erg s}^{-1}$ on prior activity from the host galaxy of Sw J1644+57 (Bloom et al. 2011) for $1 \lesssim M_\ast \lesssim 10^6$. Thus, a fossil disc sufficiently dim to go undetected prior to Sw J1644+57 nevertheless could supply sufficient magnetic flux to power the observed jet.

What could be the origin of such a fossil disc? One possibility is that the host galaxy of Sw J1644+57 contained a subluminous AGN that previously escaped detection; this is possible since discs with Eddington ratios $\lambda \lesssim 10^{-3}$ are actually quite common (Ho 2009). If Sw J1644+57 was produced by the partial disruption of a star, then it is also possible that the star was partially disrupted on at least one previous orbit as well. In this case, since the orbital period would be $\sim 10^3 \text{ yr}$, the current expected accretion rate from this relic disruption event would also be in the range $\lambda \sim 10^{-3} - 10^{-3}$ necessary to produce the required magnetic flux. The much larger size of the viscously spreading disc by the time of Sw J1644+57 could aid in the production of the required large-scale field.

5 Although the tidal debris traverses only a small fraction of the azimuthal extent of the disc, the time it spends at pericentre is comparable to the local orbital time. Thus, a large fraction of the disc will have sufficient time to rotate into, and be collected by, the debris before it falls back.

6 That said, it is not at all clear whether the jet from the fossil disc would indeed be pointed along our line of sight (= spin axis of the BH) since although the disc is in a MAD state at small radii, the jet direction could be set by the plane of the disc on larger scales, where it is not MAD and in general is misaligned with our line of sight.
flows due to periodic accumulation and expulsion of flux by the BH on semiregular intervals of $(0.5-2) \times 10^3 M_\odot$ s (Tchekhovskoy et al. 2011; Tchekhovskoy & McKinney 2012). To understand this result, note that the inner $\sim 30 h_0 r_g$ of a MAD accretion disc contains as much flux as the BH itself (McKinney et al. 2012; Tchekhovskoy & McKinney 2012). Therefore, when the BH expels an order unity fraction of its flux, this flux is replenished on a characteristic time-scale set by the accretion time at $r \simeq 15 h_0 r_g$.

$$\Delta t_{\text{flare}} \sim t_{\text{acc}} = \alpha_{\text{vis}}^{-1} \left( \frac{h}{r} \right)^{-2} \Omega_k^{-1}$$

$$\simeq 3 \times 10^3 \text{s} \left( \frac{\alpha_{\text{vis}}}{0.1} \right)^{-1} h_{0.3}^{-1/2} M_\odot^{-1/2} r_g^{3/2},$$  \hspace{1cm} (MS)

$$\simeq 10^3 \text{s} \left( \frac{\alpha_{\text{vis}}}{0.1} \right)^{-1} h_{0.3}^{-1/2} M_\odot^{-1} r_g^{-3/2},$$  \hspace{1cm} (WD)

where $\alpha_{\text{vis}}$ parametrizes the disc viscosity and $\Omega_k = (GM_\odot/r^3)^{1/2}$ is the Keplerian orbital velocity. This time-scale roughly agrees with the flaring time-scale seen in the simulations (MTB13). Equation (40) shows that $\Delta t_{\text{flare}}$ is consistent with the observed interval few times $10^3$ s between the large-amplitude flares in Sw J1644+57 if the BH is moderately massive, then $M_\odot \gtrsim 10^3$ (depending on $\alpha_{\text{vis}}$). Our model for variability contrasts with that of De Colle et al. (2012), who suggest that variability may arise from instabilities at the nozzle point, when the tidally disrupted star returns to pericentre.

In contrast to the interval between flares, the total duration of the flaring state must be at least as long as the accretion time-scale of the entire disc near its outer radius $\sim$ the circularization radius. Substituting $R_{\text{disc}}$ (equation (8)) into equation (40), we find a flaring duration $\tau_{\text{acc}} \equiv t_{\text{acc}}(r = R_*)$.

$$\simeq 5 \times 10^3 \text{s} \left( \frac{\alpha_{\text{vis}}}{0.1} \right)^{-1} h_{0.3}^{-1/2} M_\odot^{-1/2} r_g^{3},$$  \hspace{1cm} (MS)

$$\simeq 10^3 \text{s} \left( \frac{\alpha_{\text{vis}}}{0.1} \right)^{-1} h_{0.3}^{-1/2} M_\odot^{-1} r_g^{-3/2},$$  \hspace{1cm} (WD)

Again, $\tau_{\text{acc}}$ is consistent with the duration of the flaring state in Sw J1644+57 ($\sim 10^3$ s) for a solar-type star with impact parameter $\beta \sim 1$, independent of the BH mass. However, for a WD $\tau_{\text{acc}}$ appears to be too short.

Although $\tau_{\text{acc}}$ [equation (42)] sets the minimum duration of the flaring state, the state can last longer if the BH requires more time to fully align the accretion disc with the BH spin. Alignment completes only once the entire disc is MAD. At the onset of a MAD, most of the (large-scale) magnetic flux in the system is concentrated near the BH. As time goes on, two processes take place. First, the mass accretion rate drops, causing the centrally concentrated magnetic flux to be redistributed to the outer regions of the accretion disc. Secondly, new flux is added at the outer edge of the disc by the infalling stellar debris that has picked up magnetic flux from the fossil accretion flow (Section 5.2). Since in a MAD every $30 h_0 r_g$ in radius holds roughly the same amount of magnetic flux as the BH itself (Tchekhovskoy & McKinney 2012), the whole disc goes MAD once its flux reaches a value

$$\Phi^\text{max}_D \approx \frac{R_{\text{circ}}}{30 r_g} h_{0.3}^{-1} M_\odot^{1/3} r_g^{-1/2} \beta^{-1/2} \Phi_\odot,$$  \hspace{1cm} (43)

where $\Phi_D$ is the total flux through the mid-plane of the disc and the last equality makes use of equation (8).

Now, the flux through the BH evolves as (see equation 14)

$$\Phi_\odot = \Phi^\text{MAD}(t/t_{\text{MAD}})^{\nu/2},$$  \hspace{1cm} (44)

where $\Phi_{\text{MAD}}$ is BH magnetic flux and $t_{\text{MAD}}$ is the time at the onset of MAD near the hole. The fall-back magnetic flux, brought in by the infalling debris, evolves as (see equation 37)

$$\Phi^\text{fb} = \Phi^*_{\text{MAD}} \left( \frac{t}{t_{\text{MAD}}} \right)^{2/3}.$$  \hspace{1cm} (45)

The flux through the disc evolves as

$$\Phi_D = \Phi^\text{MAD}_{\text{tot}} + \Phi^\text{fb}_{\text{MAD}},$$  \hspace{1cm} (46)

where the disc flux increases as the result of both flux leaving the BH (first term in parentheses) and new flux brought in by the infalling debris (second term in parentheses).

Condition (43) can now be written as

$$\tau_{\text{align}} \approx \left( 1 + \frac{\Phi^\text{max}_D}{\Phi^*_{\text{MAD}}} \right)^{\gamma} \left[ 1 + \frac{1}{1 + 15 r_g M_\odot^{1/3} M_*^{2/3} \beta^{-1} h_{0.3}^{-1}} \right]^{2/3},$$  \hspace{1cm} (47)

$$\approx \left[ 1 + 0.8 \frac{M_*}{0.5} \gamma \left( \frac{M_\odot}{3} \right)^{2/3} \left( \frac{\beta}{2} \right)^{-1} \left( \frac{h_{0.3}}{3} \right)^{-1} \right]^{0.57},$$  \hspace{1cm} (MS, complete)

$$\approx \left[ 1 + 1.9 \frac{M_*}{0.5} \gamma \left( \frac{M_\odot}{3} \right)^{2/3} \left( \frac{\beta}{0.8} \right)^{-1} \left( \frac{h_{0.3}}{3} \right)^{-1} \right]^{0.57},$$  \hspace{1cm} (MS, partial)

$$\approx \left[ 1 + 0.4 \frac{M_*}{0.6} \gamma \left( \frac{M_\odot}{0.1} \right)^{2/3} \beta^{-1} \left( \frac{h_{0.3}}{3} \right)^{-1} \right]^{0.57},$$  \hspace{1cm} (WD, complete)

where $\gamma = 6/(4 + 3\beta)$, and the last three lines are evaluated for our three progenitor scenarios (Section 4.1). Equation (47) shows that the jet alignment (flaring) phase can last for a time-scale comparable (or somewhat longer) than that required for the MAD to form. For an MS star, both $\tau_{\text{acc}}$ [equation (42)] and $\tau_{\text{align}}$ [equation (47)] are sufficiently long to account for the observed duration of the flaring state in Sw J1644+57. However, for a WD, although $\tau_{\text{acc}}$ is very short, the duration of the flaring state is set by $\tau_{\text{align}} \sim t_{\text{MAD}}$ and hence (since $t_{\text{MAD}} \sim t_{\text{align}} \sim \text{days}$) is also consistent with observations.

Many of the points above are illustrated explicitly in Figs 3 and 4, which show the time evolution of $\Phi_D$ and $\Phi_\odot$ in fiducial MS partial disruption and WD scenarios, respectively. In both cases, $\Phi_\odot$ initially rises with the accumulated flux $\Phi_{\text{fb}}$ until the inner disc near the BH becomes MAD. After this point, flux leaks out of the BH into the surrounding disc and new flux is added by fall-back material at the outer edge of the disc ($\Phi_D$ rises). However, eventually $M$ drops sufficiently for the entire disc to become MAD (magnetic field marginally dynamically important everywhere). At this point, even the disc itself cannot hold the accumulated flux, which begins to leak out, causing $\Phi_D$ to fall. The jet alignment/flaring phase described above (Stage 2) occurs during the interval when $\Phi_D$ is still rising.

### 5.4 Origin of radio rebrightening

Our fits to the X-ray light curve of Sw J1644+57 (Figs 1, 5 and 6) show that the jet could have been active weeks–month prior to the first detection. However, since the jet was pointed away from – and possibly precessing about – our line of sight (i.e., spin axis of the BH),
its emission was not initially observable due to geometric or relativistic beaming. Nevertheless, since the jet still injects relativistic kinetic energy into the surrounding interstellar medium during this phase, this gives rise to delayed radio afterglow emission on a timescale of months–year for a typical misalignment angle (Section 2.3), consistent with observed radio rebrightening several months after the trigger (B12).

To produce radio rebrightening of the magnitude observed in Sw J1644+57, the additional energy released by the early off-axis jet must exceed that released later during the on-axis phase (i.e. after the initial γ-ray detection). Since the mass accretion rate decreases as $M \propto t^{-1}$ and the jet power $P_j \propto M$ [equation (17)] during MAD accretion, then the maximum’ energy released by the jet prior to time $t$ is given by

$$E(t) \propto \int_{t_{\text{peak}}}^{t} t^{-\alpha} \, dt = \frac{1}{\alpha - 1},$$

where $\alpha \approx 5/3 - 2.2$. If one demands that the jet energy released before the trigger exceeds the energy released after the trigger,

$$1 < \frac{E(t_{\text{trigger}})}{E(\infty)} = \left( \frac{t_{\text{peak}}}{t_{\text{-trigger}}} \right)^{\alpha - 1} - 1,$$

then this requires

$$t_{\text{trigger}} > 2^{\frac{1}{\alpha - 1}} t_{\text{peak}}.$$ 

Aside from a different pre-factor, this constraint is identical to that based on the shape of the X-ray light curve (equation (11)). Self-consistency of our model thus requires that the off-axis power be similar to that required to explain the observed radio rebrightening.

Although we focus above on energy injected off-axis prior to the γ-ray trigger, this is not the only means by which the jets could inject ‘invisible’ energy into the ambient medium. The onset of MAD accretion after the γ-ray trigger may cause the jet to wobble in and out of our line of sight, giving rise to high amplitude variability (Section 3; Stage 2 in Fig. 2). Since during this process the jets are pointed towards our line of sight only a fraction of the time, most of their energy is released during a misaligned state. Indeed, the $\sim 10$ percent duty cycle of the observed flaring (Fig. 1) suggests that the energy injected during misaligned phases could exceed that injected along our line of sight by an order of magnitude. Since $\sim 1/2$ of the total X-ray flux occurred during the first $t = t_{\text{trigger}} < 10$ d (wobbling jet phase), misaligned jets could produce off-axis relativistic ejecta with $\sim 5$ times more energy than that directly probed by the observed γ-ray/X-ray emission. This alone might be sufficient to power the observed rebrightening, without the need for significant energy injection prior to the trigger.

5.5 Predictions

5.5.1 X-ray Transients

We postulate that Sw J1644+57 resulted from a somewhat special geometric configuration, in which the BH spin axis was pointed along our line of sight. This favourable geometry resulted in several bright flares over the first $\sim 10$ d, which are produced as the jet settles down into its fully aligned configuration (Stage 2 in Fig. 2; Sections 4.2.1 and 5.3). If, however, we had instead been positioned along a more ‘typical’ line of sight not aligned with the BH, then we might only have observed a small number of flares, possibly just one in the majority of cases. Such a single flare would appear as an X-ray/soft γ-ray transient of duration $\sim 10^3$ s, yet unaccompanied by an extended luminous X-ray tail as characterized Sw J1644+57.

Could such a population of off-axis jetted TDE flares be contributing to the known population of high-energy transients? Given their long duration and low luminosity, such flares could be mistaken for a long-soft GRB or an X-ray flash. Perhaps the closest known analogue is the class of ‘low-luminosity gamma-ray bursts’ (LLGRBs; e.g. Cobb et al. 2006; Soderberg et al. 2006), although most of these are probably not standard TDEs due to their locations off the nucleus of their host galaxies and their observed association with core-collapse supernovae (SNe; e.g. Chornock et al. 2010; but see Shcherbakov et al. 2013). The volumetric rate of LLGRBs actually exceeds that of classical GRBs (Soderberg et al. 2006), but only a handful are known since they are more challenging to detect than luminous high-redshift GRBs (see Nakar & Sari 2012 for a recent compilation). Thus, even if the rate of single-flare off-axis TD jet is a factor of $\sim 10$ times higher than the rate of Sw J1644+57-like events, it is unclear how many should have been detected yet. If such an event is eventually detected, perhaps by the next generation of wide-field XRTs, it may be distinguished from normal LLGRBs by its (1) nuclear position; (2) associated optical/UV/soft X-ray emission produced by isotropic thermal emission from the accretion disc or ionized stellar ejecta (Strubbe & Quataert 2011; Clausen et al. 2012), rather than a core-collapse SNe and (3) delayed radio emission from the off-axis jet (Giannios & Metzger 2011; Section 5.5.3).

5.5.2 Jet Revival

Although the jet in Sw J1644+57 is presently off, our model predicts that the BH continues to accrete through a geometrically thin accretion disc in a jetless ‘thermal state’. However, eventually the mass accretion rate will decrease below $\approx 2$ per cent of $M_{\text{Edd}}$ (Maccarone 2003), after which point the flow may transition to a ‘hard’ state, as characterized by a radiatively inefficient geometrically thick accretion flow. Once this transition occurs, magnetic flux can once again accumulate near the BH on a short timescale ($\sim t_{\text{acc}}$; equation 42), resulting in a new MAD accretion phase (Stage 5 in Fig. 2). A jet aligned with the Earth and BH spin axis will thus again form, with a power that again tracks the accretion rate $M$. From Figs 5 and 6, we estimate that the X-ray flux will be $\sim 2 \times 10^{-14} - 10^{-13} \text{erg cm}^{-2} \text{s}^{-1}$ when the jet turns back on, well within the detection limits of current X-ray observatories.

In order to test this idea, we strongly encourage regular X-ray follow-up of Sw J1644+57 over the next decade. The time-scale and flux of the observed revival would inform (1) the rate at which the accretion rate is decreasing, and therefore help to distinguish between partial and complete disruption scenarios, or whether the disc has transitioned to a spreading evolution (e.g. Cannizzo et al. 2011) (2) the ratio of accretion rates which characterize the thick→thin disc and thin→thick disc transitions, respectively; and (3) whether the jet beaming corrections (related to the opening angle and bulk Lorentz factor) during the low–hard state are similar to those during the super-Eddington state.
5.5.3 Ubiquity of radio transients from TD events

Due to the enormous energy released in relativistic ejecta, Sw J1644+57 remains a bright radio source \((F_{\nu} \sim 1-10 \text{ mJy at } \nu \sim 1-50 \text{ GHz})\) even now, almost two years after the TDE. Since most of the jet activity occurred \(\sim \) months around the trigger time, by now the ejecta has slowed down appreciably due to interaction with the circumnuclear medium. The current expansion velocity of the ejecta is at most mildly relativistic \(\gamma \sim 2\) \((B12)\), in which case the radio emission should be relatively isotropic (i.e. flux varying by less than an order of magnitude from front to side).

The fact that Sw J1644+57 is a bright isotropic radio source has two implications: first, given the relatively close proximity \((z \sim 0.1)\) of most candidate events detected over the past several years, any jet from these events as remotely as powerful as that in Sw J1644+57 should be easily detectable by now, even if the jet remains always pointed away from our line of sight. Bower et al. \(2012\) and van Velzen et al. \(2013\) have conducted radio follow-up observations of previous thermal TD flares; since most of these observations produced only deep upper limits (with a few interesting exceptions), this already constrains the fraction of TDEs accompanied by powerful jets to be \(\lesssim 10\) per cent. In hindsight, it is perhaps unsurprising that Sw J1644+57 is unique, given the enormous magnetic flux required to power its jet, which could require special conditions not satisfied by most TDEs \(\text{Section 5.2; De Colle et al. 2012}\). A second consequence of the radio luminosity of Sw J1644+57 is that off-axis emission from other jetted TDEs (albeit rare) is isotropic and hence should be detectable out to much higher redshifts \((z \sim 1)\). Such events are one of the most promising sources for future wide-field radio surveys (Giannios & Metzger 2011; van Velzen, Körding & Falcke 2011; Frail et al. \(2012\); see Cenko et al. \(2012\) for a potential high-redshift analogue to Sw J1644+57), which will help constrain the rate and diversity of jetted TDEs.

6 CONCLUSIONS

We have presented a self-consistent model that explains many of the previously disparate puzzles associated with the jetted TDE Sw J1644+57 \(\text{Fig. 2; Section 3}\). This model relies on just one major assumption: the accumulation of a large, dynamically important magnetic flux near the central BH, such that the accretion flow from the returning stellar debris becomes MAD \(\text{Section 3.3}\) on a time-scale of \~{}month after the TDE \(\text{Figs 5 and 6}\). The onset of MAD accretion in Sw J1644+57 naturally accounts for (i) the period of intense flaring that lasted for the first \~{}10 d after the trigger; (ii) the approximate constancy (in a time-averaged sense) of the luminosity during this period; (iii) the subsequent transition of X-ray luminosity to a steady \(\text{(non-precessing)}\) jet with a power that tracks the predicted power-law decline in the accretion rate; (iv) the sudden shut off of the jet emission at \sim{}500 d after the trigger; and (v) the potential origin of the mysterious late time radio rebrightening that started about a month after the trigger. Our model also naturally predicts a QPO (at a frequency closely tied to that of a BH spin) that is consistent with that seen in the light curve of Sw J1644+57.

We emphasize that the strong magnetic flux required by our model is not an entirely independent assumption: a flux of at least this magnitude is necessary to explain the observed jet power in the first place \(\text{equation 3; Section 5.2}\); given plausible values of the BH mass and spin, while a much weaker flux would be unable to produce a jet at all \(\text{(the jet would be ‘stiffed’) against the powerful ram pressure of the misaligned disc (Komissarov & Barkov 2009)}\).

ACKNOWLEDGEMENTS

We thank Binbin Zhang, David Burrows, Michael Eracleous, Jonathan Granot, James Guillochon, Michael Kasden, Serguei Komissarov, Julian Kroklik, Pawan Kumar, Morgan MacLeod, Jonathan C. McKinney, Petar Mimica, Ramesh Narayan, Ryan O’Leary, Asaf Pe’er, Tsvi Piran, Enrico Ramirez-Ruiz, Eliot Quataert, Roman Shcherbakov, Steinn Sigurdsson, Nicholas Stone and Sjoert van Velzen for insightful discussions. We thank the anonymous referee for suggestions that helped improve the manuscript. AT was supported by a Princeton Center for Theoretical Science fellowship, by NASA through Einstein Postdoctoral Fellowship grant number PF3-140115 awarded by the Chandra X-ray Center, which is operated by the Smithsonian Astrophysical Observatory for NASA under contract NAS8-03060, and an XSEDE computational time allocation TG-AST100040 on NICS Kraken and Nautilus and TACC Ranch.

REFERENCES

Abramowicz M. A., Jaroszyński M., Kato S., Lasota J.-P., Różańska A., Sadowski A., 2010, A&A, 521, A15
Beckwith K., Hawley J. F., Krolik J. H., 2008, ApJ, 678, 1180
Berger E., Zauderer A., Pooley G. G., Soderberg A. M., Sari R., Brunthaler A., Bietenholz M. F., 2012, ApJ, 748, 36 (B12)
Bloom J. S. et al., 2011, Science, 333, 203
Bower G. C., Metzger B. D., Cenko S. B., Silverman J. M., Bloom J. S., 2012, ApJ, 763, 84
Burrows D. N. et al., 2011, Nature, 476, 421
Cannizzo J. K., Troja E., Lodato G., 2011, ApJ, 742, 32
Cao D., Wang X.-Y., 2012, ApJ, 761, 111
Cenko S. B. et al., 2012, ApJ, 753, 77
Chabrier G., 2003, ApJ, 586, L133
Chornock R. et al., 2010, preprint (arXiv:1004.2262)
Claussen D., Sigurdsson S., Eracleous M., Irwin J. A., 2012, MNRAS, 424, 1268
Cobb B. E., Bailyn C. D., van Dokkum P. G., Nararajan P., 2006, ApJ, 645, L113
De Colle F., Guillochon J., Naiman J., Ramirez-Ruiz E., 2012, ApJ, 760, 103
Fender R. P., Belloni T. M., Gallo E., 2004, MNRAS, 355, 1105
Frail D. A., Kulkarni S. R., Ofek E. O., Bower G. C., Nakar E., 2012, ApJ, 747, 70
Giannios D., Metzger B. D., 2011, MNRAS, 416, 2102
Greene J. E., 2012, Nat. Commun., 3, 1304
Guillochon J., Ramirez-Ruiz E., 2013, ApJ, 767, 25
Haas R., Shcherbakov R. V., Bode T., Laguna P., 2012, ApJ, 749, 117
Kepler S. O. et al., 2013, MNRAS, 429, 2934
Krolik J. H., Piran T., 2011, ApJ, 743, 134
Krolik J. H., Piran T., 2012, ApJ, 749, 92
Lei W.-H., Zhang B., Gao H., 2012, ApJ, 762, 98
Levan A. J. et al., 2011, Science, 333, 199
Li J., Pe’er A., Loeb A., 2012, preprint (arXiv:1211.5120)
Lodato G., King A. R., Pringle J. E., 2009, MNRAS, 392, 332
Maccarone T. J., 2003, A&A, 409, 697
MacLeod M., Guillochon J., Ramirez-Ruiz E., 2012, ApJ, 757, 134
McKinney J. C., Blandford R. D., 2009, MNRAS, 394, L126
McKinney J. C., Tchekhovskoy A., Blandford R. D., 2012, MNRAS, 423, 3085
