JETS FROM TIDAL DISRUPTIONS OF STARS BY BLACK HOLES

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ABSTRACT

Tidal disruption of main-sequence stars by black holes has generally been thought to lead to a signal dominated by UV emission. If, however, the black hole spins rapidly and the poloidal magnetic field intensity on the black hole horizon is comparable to the inner accretion disk pressure, a powerful jet may form whose luminosity can easily exceed the thermal UV luminosity. When the jet beam points at Earth, its non-thermal luminosity can dominate the emitted spectrum. The thermal and non-thermal components decay differently with time. In particular, the thermal emission should remain roughly constant for a significant time after the period of maximum accretion, beginning to diminish only after a delay, whereas after the peak accretion rate, the non-thermal jet emission decays, but then reaches a plateau. Both transitions are tied to a characteristic timescale $t_{\text{edd}}$ at which the accretion rate falls below Eddington. Making use of this timescale in a new parameter-inference formalism for tidal disruption events with significant emission from a jet, we analyze the recent flare source Swift J2058. It is consistent with an event in which a main-sequence solar-type star is disrupted by a black hole of mass $\sim 4 \times 10^7 M_\odot$. The beginning of the flat phase in the non-thermal emission from this source can possibly be seen in the late-time light curve. Optical photometry over the first $\sim 40$ days of this flare is also consistent with this picture, but is only weakly constraining because the bolometric correction is very uncertain. We suggest that future searches for main-sequence tidal disruptions use methods sensitive to jet radiation as well as to thermal UV radiation.

Key words: accretion, accretion disks – black hole physics – stars: general

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1. INTRODUCTION

Now that we know that supermassive black holes frequently can be found in the centers of galaxies (Kormendy & Richstone 1995; Gültekin et al. 2009), we can also expect that ordinary stars occasionally pass near enough to them to be tidally disrupted. Theoretical estimates of the rate at which this might happen range from $3 \times 10^{-5} M_{\text{BH},0.45}^{0.45}$ galaxy$^{-1}$ yr$^{-1}$ (Brockamp et al. 2011) to $6.5 \times 10^{-4} M_{\text{BH},0.25}^{0.25}$ galaxy$^{-1}$ yr$^{-1}$ (Wang & Merritt 2004). Most observational estimates are in the same ballpark (Gezari et al. 2008; Esquej et al. 2008; Maksym et al. 2010), although Donley et al. (2002) estimated a rate of X-ray flares as low as $\sim 10^{-5}$ galaxy$^{-1}$ yr$^{-1}$.

More than 20 years ago, Rees (1988) outlined the basic dynamical processes relevant to these events. The most likely trajectory for a tidal disruption is one in which the star approaches on a nearly parabolic orbit. When the star is disrupted, its different parts acquire binding energies that span a very large range. The material of the star begins returning to the vicinity of the black hole after a delay corresponding to the orbital period of the most tightly bound matter. If the disrupted star has a uniform distribution in orbital binding energy per unit mass across that large range, matter continues to fall in at a rate $dM/dt \propto t^{-2.5}$ (Phinney 1989). This matter is captured into an accretion disk. Within this disk, gravitational energy is dissipated and the heat is radiated in the usual quasi-thermal fashion. At the peak accretion rate, the luminosity would be super-Eddington (Ulmer 1999); if it is thermally radiated by a disk of area corresponding to a few gravitational radii of the black hole, the associated temperature would be in the EUV or perhaps the soft X-ray band. Consequently, most searches for such events hitherto have been carried out in the EUV region.

Many of these expectations were upended by the remarkable object Swift J1644+57 (Levan et al. 2011; Burrows et al. 2011). During its time of peak luminosity (the first $\sim 2$ days of the outburst), its light curve, with its sequence of flares lasting $\sim 1000$ s alternating with quiescent periods $\sim 5 \times 10^4$ s long, bore little resemblance to the prediction of a rise time of $\sim 10^5$ s followed by a smooth rollover to a $t^{-5/3}$ falloff. Moreover, whereas all expectations were for the emission to peak in the EUV, the greatest part of the luminosity was carried in photons with energies $\sim 100$ keV, indicating electron energies far greater than the expected $\sim 100$ eV, while the shape of that spectrum ($\propto F_\nu$, gradually rising from 2 keV to 100 keV) certainly did not suggest any variety of thermal emission. In fact, the most successful emission models invoke non-thermal synchro-Compton radiation from a relativistic jet (Bloom et al. 2011; Burrows et al. 2011).

If most of the radiation actually comes from a relativistic jet, then clearly some of the conventional assumptions must be revised. The observed radiation has nothing to do with quiescent radiative cooling of an optically thick accretion flow. Indeed, if the jet is powered by something resembling the Blandford–Znajek mechanism, its real source of energy is rotation of the black hole; the only function of the mass accretion is to support the magnetic field threading the black hole’s horizon.

We have recently shown (Krolik & Piran 2011) that a possible solution to all these conundrums is that this event was caused by tidal disruption not of a main-sequence star, but of a white dwarf, and the disruption takes place over several passes, not in a single go. Most importantly for the question examined in that paper, a white dwarf disruption results in depositing so much mass in such a small region that the accretion flow is immensely optically thick. As a result, the matter’s cooling time is far longer.
than its inflow time. In such a state of affairs, thermal radiation from the accretion flow is suppressed, and the disk maintains a high pressure, one capable of confining a strong magnetic field close to the black hole.

Our goal in this paper is to inquire whether a similar state of affairs could occur when the victim of tidal disruption is a main-sequence star. We estimate here both the thermal emission of the disk and the jet power. The first is estimated using standard accretion disk theory. The latter is estimated in terms of the Blandford–Znajek mechanism, using the disk pressure near the ISCO as a measure of the black hole’s magnetic field. Although a tidal disruption of a main-sequence star takes place on a much larger lengthscale than the disruption of a white dwarf, for most of the relevant parameter space a similar situation holds: the accretion is super-Eddington, thermal radiation is suppressed, and conditions for emergence of a strong jet are established.

Our work is complementary to two related efforts. Giannios & Metzger (2011) investigated possible radio emission from a jet receiving a fixed (small) fraction of the accretion energy released by accreting tidally disrupted matter. Lei & Zhang (2011) have suggested a similar picture, but approach it rather differently. In particular, they use thin disk approximations for both the sub- and super-Eddington regimes, their scaling with black hole mass does not include the relation between disk thickness and accretion rate in the radiation-dominated sub-Eddington phase, and they do not discuss the luminosity of the thermal disk. We begin in Section 2 with a brief discussion of tidal disruption physics. We then discuss accretion dynamics in this context and the jet and disk outputs in Section 3. In Section 4, we show how this approach can be used to constrain a number of otherwise-unknown parameters of tidal disruptions and apply this method to Swift J2058 (Cenko et al. 2011), a second example of a jet-dominated tidal disruption. We summarize our results in Section 5.

2. TIDAL DISRUPTION

In order for the tidal gravity of a black hole to strongly affect a star, the pericenter of the star’s orbit cannot be much larger than

\[ R_T \approx 50(k/f)^{1/6} M_\odot^{2/3-\xi} M_{BH,6}^{-2/3} R_\odot, \]  

where \( k \) is the apsidal motion constant (determined by the star’s radial density profile) and \( f \) is its binding energy in units of \( GM_*^2/R_* \) (Phinney 1989). Here \( M_* \) is the mass of the star in solar units, \( M_{BH,6} \) is the black hole’s mass in units of \( 10^6 M_\odot \), and \( R_\odot \) is the black hole’s gravitational lengthscale \( GM_{BH}^2/c^2 \). We have approximated the main-sequence mass–radius relation by \( R_\odot \approx R_\odot M_*^{1-\xi}/\xi \approx 0.2 \) for \( 0.1 < M_* < 1 \); increases to \( \approx 0.4 \) for \( 1 < M_* < 10 \) (Kippenhahn & Weigert 1994). The analogous expression for non-main-sequence stars can be found by substituting the appropriate mass–radius relation; for example, for white dwarfs with masses well below the Chandrasekhar mass, \( \xi = 4/3 \), and the radius of a \( 1 M_\odot \) star is \( \approx 0.011 R_\odot \). The ratio \( k/f \) \( \approx 0.02 \) for radiative stars, but is \( \approx 0.3 \) for convective stars (Phinney 1989). Although the ratio \( k/f \) appears in the expression for \( R_T \) raised to only the 1/6 power, and would therefore appear to be an innocuous correction factor of order unity, it can be quantitatively significant because the timescale for the flare is \( \propto R_T^2 \) and the contrast between the values of \( k/f \) for radiative and convective stars is \( \approx 15 \). In numerical estimates, we scale \( k/f \) to the value for fully radiative stars because this is a reasonable approximation for main-sequence stars with \( 0.4 M_\odot < M < 10 M_\odot \) (Kippenhahn & Weigert 1994).

The work done by the tides stretching and rotating the star is removed from the orbit, so on average, the star’s material becomes bound to the black hole with a specific binding energy of order the star’s original specific binding energy \( \sim GM_*^2/R_* \). Because the semi-major axis \( a \) of an orbit with binding energy \( E_B \) is \( GM_{BH}/(2E_B) \) when \( M_* \ll M_{BH} \), the semi-major axis of material with the mean energy is much larger than \( R_T \),

\[ a_{\text{mean}} \approx (1/2)(k/f)^{-1/6}(M_{BH}/M_*)^{2/3} R_T. \]  

However, when the star is disrupted, its different parts acquire binding energies that span a very large range. Suppose that the actual pericenter is \( R_p = \beta R_T \), where \( \beta \lesssim 1 \) is called the “penetration factor.” The most bound matter may have a binding energy as great as \( \sim GM_{BH}R_\epsilon/R_p^2 \), leading to a semi-major axis only,

\[ a_{\min} \approx (1/2)\beta^2 (k/f)^{1/6} M_{BH,6}^{2/3} M_*^{1/3-\xi} R_T, \]

much smaller than \( a_{\text{mean}} \). Tidal torques during the initial stage of the encounter could also lead to a somewhat larger binding energy for the most bound material (Rees 1988); in terms of the discussion in this paper, the effects are indistinguishable from those due to a decrease in \( \beta \).

On the assumption that the disrupted star has a uniform distribution in orbital binding energy per unit mass, matter returns to the region \( \sim R_p \) from the black hole at a rate \( dM/dt \propto (t/t_0)^{-5/3} \) (Phinney 1989). The characteristic timescale \( t_0 \) for initiation of this power-law accretion rate is the orbital period for the most bound matter:

\[ P_{\text{orb}}(a_{\min}) \approx 5 \times 10^5 M_\odot^{1-3\xi/2} M_{BH,6}^{1/2} (k/f)^{1/2} \beta^{1/3} \text{ s}, \]  

the period for an orbit with semi-major axis \( a_{\min} \). More detailed calculations (e.g., Lodato et al. 2009) indicate that, almost independent of the star’s internal structure, the mass return rate does follow this behavior at late times, but at early times the rate rises rapidly after a delay \( \sim P_{\text{orb}}(a_{\min}) \) and then gradually rolls over to the \( t^{-5/3} \) proportionality, transitioning more slowly when the star’s pre-disruption structure is more centrally concentrated.

In this picture, once the matter returns to the vicinity of \( R_p \), it is captured into an accretion disk whose inflow time \( t_\text{in} \ll P_{\text{orb}}(a_{\min}) \). As the matter moves inward through this disk, there is local dissipation of the conventional accretion disk variety, and the heat is radiated in the usual quasi-thermal fashion. At the peak accretion rate estimated by Lodato et al. (2009) the luminosity would be

\[ L_{\text{peak}} \sim 1 \times 10^{47} (\eta/0.1) M_\odot^{(1+3\xi)/2} M_{BH,6}^{-1/2} (k/f)^{-1/2} \beta^{-3} \text{ erg s}^{-1}, \]  

for radiative efficiency \( \eta \). In Eddington units, this luminosity is \( \approx 800 R_\epsilon^{-3}(\eta/0.1) M_*^{(1+3\xi)/2} M_{BH,6}^{-3/2} \) for radiative stars and \( \approx 4 \) times smaller for convective stars (see Ulmer 1999; Strubbe & Quataert 2009); as we will discuss in Section 3, super-Eddington accretion may result in only a fraction of this inflow rate reaching the black hole. If such a luminosity were thermally radiated by a disk of area \( \sim O(10 \pi R_\epsilon^2) \), the associated temperature would be

\[ T_{\text{char}} \sim 6 \times 10^6 M_\odot^{(1+3\xi)/8} M_{BH,6}^{-5/8} (k/f)^{-1/8} \beta^{-3/4} \text{ K}. \]
implying that most of the radiation would emerge in the soft X-ray band.

3. ACCRETION DYNAMICS AND THE JET AND DISK POWERS

We begin by estimating conditions at the time of peak accretion rate. Lodato et al. (2009) estimate that 1/3 of the star’s mass should arrive at $R_p$ after a delay $\sim P_{\text{orb}}(a_{\text{min}})$. Because $M$ falls $\propto t^{-5/3}$ thereafter, $dM/d\ln t$ is greatest during this period, so this is both the period of greatest luminosity and the period in which most of the energy of the entire event is radiated. As the matter streams in, its orbital velocity is roughly the free-fall velocity at $R_p$. Collision with any mass in a more circular orbit involves velocity differences of the order of the orbital speed, so the post-shock temperature is comparable to the virial temperature. Following conventional accretion disk theory in the supposition that the inflow rate is controlled by angular momentum transport and writing the vertically integrated stress as some number or times the similarly vertically integrated pressure, we estimate the inflow time at $R_p$ as

$$t_{\text{in}} \approx \alpha^{-1}(R_p/H)^3(R_T/R_G)^{3/2}(R_g/c)\left(k/f\right)^{1/4}(0.02)^{1/2} \beta^{3/2}$$  (7a)

$$\sim 2 \times 10^3(\alpha/0.1)^{-1}(k/f)^{1/4}(R_p/H)^2M_*^{1-3/2}\beta^{3/2} \text{ s.}$$  (7b)

The quantity $H$ is the disk thickness, which cooling may reduce to be less than $R_p$, while $\alpha$ is the usual ratio of stress to pressure. Note that $t_{\text{in}}$ depends only on $M_*$ and not at all on $M_{\text{BH}}$, because it is fundamentally a dynamical time, and at the tidal radius that is determined by the dynamical time of the star $\sim (R_p^3/GM_*)^{1/2}$. Because $t_{\text{in}} \ll P_{\text{orb}}(a_{\text{min}})$, the mass held at $R_p$ is determined by a balance between the slowly changing rate at which returning tidal streams deliver mass and the more rapid inflow rate. The characteristic Thomson optical depth from the middle of the material to the outside at $R_p$ is then

$$\tau_T \sim 3 \times 10^3(\alpha/0.1)^{-1}(R_p/H)^2M_*^{1+12}\beta^{-1/12}$$  \times M_{\text{BH}}^{1/6}

$$\left(k/f\right)^{-1/2}(0.02)^{-1/2} \beta^{-7/2},$$  (8)

and the cooling time is

$$t_{\text{cool}} \sim \tau_T H/c \sim 4 \times 10^6(\alpha/0.1)^{-1}(R_p/H)M_*^{1+6/2}\beta^{-1/12}$$  \times M_{\text{BH}}^{-5/6}

$$\left(k/f\right)^{1/2}(0.02)^{1/2} \beta^{-5/2} \text{ s.}$$  (9)

The effectiveness of cooling can be measured by comparing $t_{\text{cool}}$ to the inflow time $t_{\text{in}}$ at that radius:

$$t_{\text{cool}}/t_{\text{in}} \sim 500H/R_pM_*^{1+6/2}\beta^{-1/2}M_{\text{BH}}^{3/6}M_*^{1/6}(k/f)^{-1/2}(0.02)^{-1/2} \beta^{-4}.$$  (10)

If the initial temperature of the matter arriving in the disk is close to the virial temperature, it cools efficiently only if the black hole mass is $\sim 10^8M_\odot$ or the star’s mass is $\sim 0.1M_\odot$. Although it is possible for stars on specific trajectories to be tidally disrupted by spinning black holes with masses up to $\sim 10^9M_\odot$, the probability of tidal disruption (as opposed to direct capture) declines rapidly when $M_{\text{BH}} \sim 100M_*^{(-3/2)\beta}$ or more (Kesden 2011). For this reason, unless the matter arrives at $R_p$ substantially cooler, we expect that at the time of peak accretion rate the inflow will never be able to cool efficiently.

Another way of viewing this result is to note that photon trapping is expected whenever the accretion rate is super-Eddington (Begelman 1979; Abramowicz et al. 1988). As we have already seen, the characteristic peak luminosity in these events exceeds the Eddington luminosity by a factor $\sim 800(\eta/0.1)M_*^{1+3}\beta M_{\text{BH}}^{3/2}(k/f)/0.02)^{-1/2} \beta^{-3}$. Thus, to order-of-magnitude accuracy, the dividing line between super-Eddington accretion in main-sequence tidal disruption events falls at roughly the same place as the dividing line between the radiatively efficient and photon-trapping regimes, and both are near the maximum black hole mass at which there is a significant probability of tidal disruption (Ulmer 1999; Struble & Quataert 2009).

3.1. Jet Power

Rotating black holes whose horizons are threaded by large-scale poloidal field can drive relativistic jets (Blandford & Znajek 1977; McKinney & Gammie 2004; Hawley & Krolik 2006). At the order-of-magnitude level, the Poynting luminosity of such a jet can be estimated through a simple dimensional argument: it must be $\sim B^2$, where $B^2$ is the poloidal field intensity on the horizon. The actual magnitude of the luminosity can then be written as $L_{\text{jet}} = f(\alpha/M)c(B^2/8\pi)^2$, where $f(\alpha/M)$ is a dimensionless function that should increase with $|\alpha/M|$. To estimate $B^2$, we follow Beckwith et al. (2009), who demonstrated that the magnetic pressure near the horizon is generally bounded above by the midplane total pressure near the ISCO and bounded below by the magnetic pressure at that location. Our first task is therefore to estimate the disk pressure near the ISCO.

When the black hole mass is very large (i.e., roughly the maximum black hole mass capable of disrupting the star before it enters the black hole), the inflow can be both sub-Eddington and cool efficiently even at the time of peak accretion rate. If $t_{\text{cool}}/t_{\text{in}} < 1$ at $R_p$ (Equation (10)), the matter can cool to the temperature at $R_p$ associated with steady accretion at the prevailing rate. At smaller radii, the temperature profile of the accretion flow, as well as its other properties, can be described by standard stationary disk solutions (e.g., as in Krolik 1999). Although the accretion rate is sub-Eddington, it is nonetheless high enough that the flow will certainly be radiation dominated.

The pressure in the midplane of the inner disk is then

$$p_{\text{mid}} \approx \frac{2\epsilon c^2}{\kappa k} (R/R_g)^{-3} X(R/R_g),$$  (11)

where $\kappa$ and $\kappa_T$ are the actual and the Thomson opacity, respectively, and the function $X(R/R_g)$ is a place holder to account for how these scalings are altered as the ISCO is approached. Note that $p_{\text{mid}}$ in the sub-Eddington radiation-dominated regime is inversely proportional to $M_{\text{BH}}$. Although the outgoing thermal radiation flux is proportional to the accretion rate, the optical depth in this regime is inversely proportional; as a result, the midplane radiation pressure $p_{\text{mid}}$ is independent of the accretion rate (Moderski & Sikora 1996).

Conversely, in the low-mass limit (which applies to the majority of relevant black holes), the inflow is both super-Eddington and radiatively inefficient because photon trapping prevents effective radiative cooling. In this regime, the effects of strong radiation forces and possible radiation-driven outflows make the theory less clear (consider, e.g., the range of views presented...
in Abramowicz et al. 1988; King & Pounds 2003; Strubbe & Quataert 2009; Dotan & Shaviv 2011; Begelman 2011). For our purposes, it is sufficient to estimate the pressure using dimensional analysis and simple scaling arguments. Precisely because the gas cannot cool itself, the ratio of total (i.e., radiation plus gas) pressure to density remains at a level close to the virial temperature as the flow moves inward, so that, unlike the sub-Eddington radiation pressure-dominated regime, the pressure when the flow is super-Eddington is proportional to the accretion rate and the flow configuration is nearly round. The midplane pressure (again, dominated by radiation) is then

\[ p_{\text{mid}} \sim \frac{q m c^2}{\kappa T R E} \frac{c_s^2}{v_{\text{orb}}^2} (R/R_p)^{-5/2}, \]  

(12)

where \( c_s \) is the effective sound speed (including radiation pressure) and \( v_{\text{orb}} \) is the speed of a circular orbit at that radius. Here we have substituted \( \alpha \) for the ratio of inflow speed \( v_i \) to orbital speed because \( v_i/v_{\text{orb}} \sim \alpha \) when the disk is geometrically thick. We define \( \dot{m} \) to be the accretion rate delivered to the outer edge of the disk normalized to Eddington units with unit efficiency, i.e., \( \dot{m} = M c^2/L_E \) because the uncertain effect of photon trapping makes the radiative efficiency difficult to estimate. When the distinction is important, we will use \( \dot{m}_0 \) for the value of this normalized accretion rate at its peak value. The parameter \( q \) gives the fraction of \( \dot{m} \) arriving at the black hole; when \( \dot{m} > 1 \), it is possible for \( q \leq 1 \), but when \( \dot{m} < 1 \), in general we expect \( q = 1 \). Modulo corrections of order unity, when electron scattering dominates the opacity and the relevant radii are \( \sim R_g \), the only difference between this expression and the previous one is that this one is also \( \propto q \dot{m} \).

Because our estimates for the total pressure in both the sub- and super-Eddington regimes are so similar, it is convenient to write an estimate for the jet luminosity in a single form:

\[ L_{\text{jet}} = \frac{c^3 R_p}{\alpha \beta_h \kappa_T} f (a/M) \left\{ \begin{array}{ll} 1 & q \dot{m} \ll 1 \\ \frac{m}{q m} & q m \gg 1 \end{array} \right. \]  

(13)

where \( \alpha \) should be interpreted as the usual stress parameter in the sub-Eddington regime, but only as \( v_i/v_{\text{orb}} \) in the super-Eddington case. When \( M_{\text{BH}} f (a/M)/(\alpha \beta_h) = 1 \) (in the super-Eddington regime \( q \dot{m} = 1 \)), the magnitude of the jet power is \( 1 \times 10^{43} \) erg s\(^{-1}\). The quantity \( \beta_h \) is the ratio of the midplane total pressure near the ISCO to the magnetic pressure in the black hole’s stretched horizon. In very rough terms, we might expect \( \beta_h \alpha \sim 0.1-1 \). Thus, for fixed black hole spin, the power in the jet is a constant fraction of the Eddington luminosity, independent of accretion rate, in the sub-Eddington regime, but rises in proportion to the accretion rate in the super-Eddington regime. The reason is that at lower accretion rates, the radiative losses curb any increase in pressure with increasing accretion rate, whereas once the accretion rate becomes super-Eddington, the photons are trapped in the accretion flow.

However, there are several subtleties hidden in this estimate. The first is that \( f (a/M) \) is not well determined. Blandford & Znajek (1977) worked out an expansion in small \( a/M \) for time-steady split-monopolar field distributions. Tchekhovskoy et al. (2010) developed a similar expansion in terms of the black hole rotation rate. Both required the field rotation rate to be exactly half the black hole rotation rate. Tchekhovskoy et al. (2010) found that their expression compared well to axisymmetric force-free MHD simulations with an imposed magnetic flux function. Unfortunately, neither axisymmetry nor an imposed flux function may be a good approximation to the behavior of time-dependent three-dimensional relativistic jets whose field configuration is created self-consistently by accretion dynamics, which, in turn, may be influenced by black hole spin (Hawley & Krolik 2006). It is therefore difficult to be confident about much more than that jet power should increase with faster black hole rotation.

The second subtlety is that the relevant field intensity on the horizon is the time average of the poloidal component. When the sign of the vertical field through the equatorial plane changes frequently, this intensity can be strongly suppressed (Beckwith et al. 2008). Tidal disruptions, however, may be a particularly favorable case for jet launching because the extreme elongation of tidal streams might create very large-scale structures in the field with exactly the sort of coherent direction required to support powerful and long-lived jets.

Third, up to this point we have been careful to discuss only the two extreme limits: high-mass, sub-Eddington, and radiatively efficient; and low-mass, super-Eddington, and radiatively inefficient. That is because the transition between them is not well understood. For estimating both the jet luminosity and (later) the emergent thermal luminosity, the key parameter is the ratio \( r_{\text{cool}}/r_{\text{in}} \). Several problems must be solved before its transition between the two extreme limits can be defined. There may be parameter regimes in which this ratio is \( > 1 \) at \( R_p \) if the gas is deposited there with \( T \sim T_{\text{in}} \) even though the temperature in a steady-state flow at the specified accretion rate is sufficiently cooler that the gas could cool in less than an inflow time. On the other hand, there are accretion rates for which the flow in steady state is radiatively efficient at \( R_p \), but not at smaller radii. The quantitative character of the transition is also sensitive to exactly how the dissipation profile varies with radius near the ISCO. In view of these uncertainties, we will adopt here the simplest solution, understanding that it is only provisional: we will define the transition point by the intersection between the two expressions for \( L_{\text{jet}} \), we will adopt a similar prescription for the thermal luminosity.

Fourth, there is the question of what fraction of the returning mass actually reaches the vicinity of the black hole, i.e., the correct value of \( q \). This is important in the super-Eddington regime where \( L_{\text{jet}} \propto q \dot{m} \). If a significant fraction of \( \dot{m} \) is ejected as a wind (as could well happen during the super-Eddington period), \( q \) could be rather small and the jet power would be reduced accordingly.

Last, jets in Galactic black hole binaries turn off when \( \dot{m} \) rises toward \( \sim 1 \) (Remillard & McClintock 2006). On the other hand, the existence of radio-loud active galactic nuclei (AGNs) with strong optical/UV continua and emission lines suggests that larger black holes can somehow support jets even when the accretion rate is near Eddington. Luminous quasars are thought to accrete at rates such that \( 0.01 \lesssim L/L_E \lesssim 1 \) (the lower end of this range favored by, e.g., Kelly et al. (2010), who estimate black hole masses from the quasar luminosity and broad-line width; the higher end favored by, e.g., Liu et al. (2009), whose black hole masses come from the bulge dispersion in obscured quasars). Although these statistics are dominated by the radio-quiet variety, they likely apply to radio-loud quasars as well because there are only slight differences between the optical/UV continua of radio-loud and radio-quiet quasars (Richards et al. 2011), a fact suggesting that the inner disks of these two categories of quasars are very similar. A more direct indication comes from narrow-line Seyfert 1 galaxies (NLS1s). These
related to the accretion flow. The pressure in the accretion flow confines the magnetic field, but the magnetic field taps the rotational kinetic energy (i.e., the difference between the total mass and the irreducible mass) of the black hole. For the same reason, the jet efficiency measured in rest-mass units, which is \( f(a/M)/(4\pi \beta_0 \alpha) \), is likely to be \( \sim O(0.1) \), but can in principle be greater than unity.

It is also important to emphasize that the jet power does not translate directly into the photon luminosity we see. On the one hand, the photon luminosity of the jet is generally a small fraction \( \eta_J \) of the initial Poynting luminosity. In the conditions relevant to blazar jets, for example, \( \eta_J \sim 10^{-2} \) to \( 10^{-1} \) (Celotti & Ghisellini 2008). On the other hand, whatever photon luminosity emerges will in general be strongly beamed due to jet collimation and relativistic kinematics: \( B \sim \min(4\pi/\Delta\Omega, 2\Gamma^2) \), where \( \Delta\Omega \) is the solid angle occupied by the jet and \( \Gamma \) is its Lorentz factor. The beaming factor \( B \) can easily be \( \sim 10^3 \).

### 3.2. Thermal Luminosity

The thermal photon luminosity behaves differently as a function of black hole mass. It can be written as \( L_{\text{therm}} = \eta q \dot{m} L_E \), where we expect that \( \eta \) depends on spin for \( \dot{m} < 1 \), but declines once \( \dot{m} \) is large enough that the portion of the accretion flow where the majority of the luminosity would otherwise emerge enters the photon-trapping regime. In this limit, the photon luminosity is capped at \( \sim L_E \). As we have already discussed in the context of the jet power, the critical \( \dot{m} \) above which this occurs is not well determined. Thus, we estimate

\[
\dot{L}_{\text{therm}} \simeq 1.5 \times 10^{44} \frac{f(a/M)}{\beta_0 \alpha} \text{ erg s}^{-1}
\]

In other words, the jet luminosity has a minimum as a function of \( M_{\text{BH}} \). When the black hole mass is smaller, \( L_{\text{jet}} \propto M_{\text{BH}}^{-1/2} \), because that is the photon-trapping regime, making \( p \propto q \dot{m} \), and tidal disruption mechanics make \( \dot{m}_0 \propto M_{\text{BH}}^{3/4} \); when the black hole mass is larger, the accretion flow is in the conventional sub-Eddington regime, and \( L_{\text{jet}} \propto M_{\text{BH}} \) because \( p \) is independent of \( M \) but \( \propto M_{\text{BH}}^{-1} \), while the black hole area \( \propto M_{\text{BH}}^3 \). The transition mass between the two regimes, \( M_{\text{BH, \text{jet}}} \), decreases with decreasing accretion rate. Thus, a system that is initially in the photon-trapping regime will eventually become sub-Eddington later. On the other hand, a large black hole system that is sub-Eddington initially will remain so throughout the event. Even if only a small fraction of the fallback accretion rate reaches the black hole (i.e., \( q \ll 1 \)), the jet luminosity would still increase toward smaller black hole masses provided \( q \) scales less rapidly with black hole mass than \( \propto M_{\text{BH, \text{jet}}}^{-1/2} \).

It is important to recall that at the time of peak accretion, the black hole mass dividing the low- and high-mass regimes is almost as large as the very largest black hole mass permitting tidal disruption of main-sequence stars (Kesden 2011). This conclusion would be weakened somewhat for stars with large internal convection zones (stars at either the low or high end of the mass distribution), but \( M_{\text{BH, \text{jet}}} \) decreases by only a factor \( \approx 2.5 \). At the peak of the flare, therefore, we expect essentially all events to be effectively in the low-mass, super-Eddington regime. Only at later times, as the accretion rate falls, does the high-mass regime become relevant.

Note also that the jet luminosity is not directly limited by the Eddington luminosity because its energy is not directly
However, as the figure shows, if the black hole rotates rapidly enough to make $f(a/M)$ not too small, for $M_{\text{BH}} \ll 10^5 M_\odot$, a very large proportion of the returning mass would have to be ejected in order for $L_{\text{jet}}$ to fall to a level at which it is merely comparable with $L_{\text{therm}}$. For rapidly spinning black holes, $L_{\text{therm}} \sim L_{\text{jet}}$ only when the mass is the largest permitting any tidal disruptions at all.

Thus, the peak output is almost always dominated by the jet unless the black hole rotates rather slowly and the black hole mass is $\sim 10^8 M_\odot$. a mass so large that the tidal disruption probability is considerably less than unity. However, whether non-thermal radiation from the jet or thermal radiation from the disk dominates the observed spectrum depends on the interplay of the jet radiative efficiency and beaming with the intrinsic jet/disk ratio determined by $a/M$ and $M_{\text{BH}}$. The jet is favored by viewing angles along its axis, low black hole mass, high spin, and large $\eta_{\text{jet}}$: the disk is favored by viewing angles outside the favored cone, high black hole mass, slow spin, and small $\eta_{\text{jet}}$.

3.3. Time Dependence

Now consider what happens at later times, as matter of progressively smaller binding energy returns to the disk. If the mass per unit energy of tidally disrupted matter is independent of binding energy, the accretion rate into the disk at radii $\sim R_p$ falls $\propto (t/t_0)^{-5/3}$, where $t_0 \sim P_{\text{orb}}(a_{\text{min}})$ (Rees 1988; Phinney 1989). Lodato et al. (2009) found that, depending on the degree of density concentration in the star, the mass accretion rate might fall more shallowly than this when $t \gtrsim t_0$, but gradually tends toward $t^{-5/3}$ when $t \gg t_0$. If nearly all the matter returning to the disk makes it all the way to the black hole, the time dependence of the accretion relevant to the jet and jet powers should match the time dependence of the accretion rate reaching the outer disk; if, however, super-Eddington conditions lead to a significant fraction being expelled, the accretion rate reaching the event horizon falls more slowly because the expelled fraction can also be expected to diminish.

When $\dot{m}$ first begins to decrease, the thermal radiation hardly changes because photon trapping continues to limit $L_{\text{therm}}$ to $L_E$. Consequently, in the period near and shortly after the peak luminosity, the thermal light curve should be almost flat and only slowly roll over toward $(t/t_0)^{-5/3}$.

On the other hand, $L_{\text{jet}}$ in those sources is $\propto q\dot{m}$, so the jet power falls, provided that the accretion rate reaching the black hole is truly super-Eddington. The rate of decrease may be slower than $t^{-5/3}$ if $q$ increases as $\dot{m}$ falls. With declining $\dot{m}$, the black hole mass at which the flow switches from non-radiative to radiative also decreases. Consequently, even though virtually every source begins in a non-radiative state, those in which the black hole mass is relatively large may ultimately become efficient radiators at later times. When that changeover occurs, the jet power remains constant, while the thermal luminosity begins to fall. This timescale, when the accretion rate and therefore the thermal luminosity fall below Eddington, we call $t_{\text{Edd}}$. It should be emphasized that all of these remarks pertain to the expectation value for the jet luminosity; relativistic jets are generically unsteady, so fluctuations at the orderunity level should be expected around all these trends.

Figure 2 gives a schematic view, beginning at the time of peak accretion rate, of what might be expected in terms of the light curves for the jet power (before allowance for beaming and radiative efficiency) and the thermal luminosity. For the parameter values chosen ($M_{\text{BH}}=10^5 M_\odot$, all other scaling parameters unity), $L_{\text{jet}}$ falls to the level of $L_{\text{therm}}$ at almost the same time, $t \sim 7 t_0$, as $L_{\text{therm}}$ enters the sub-Eddington regime and also begins to decline. From that time to $t \approx 30 t_0$, both fall together, maintaining similar power levels. Finally, after $t \approx 30 t_0$ (i.e., a time larger by $\eta^{-3/5} \approx 4$ than the time at which the thermal luminosity begins to decline), the jet luminosity stabilizes, while $L_{\text{therm}}$ continues to fall.

The ratio $L_{\text{jet}}/L_{\text{therm}}$ as a function of time for fixed black hole mass can be described much more simply. In rough terms, it is just

$$L_{\text{jet}} / L_{\text{therm}} \propto \frac{\max(1, q\dot{m}) f(a/M)}{\min[1, \eta(a/M) \dot{m}]}$$

(18)

where $\eta(a/M)$ is the intrinsic (i.e., without photon trapping) radiative efficiency for a given spin parameter. At high accretion rate, the ratio is $\propto q\dot{m}$; in the transition range where photon trapping is marginal, the ratio may change relatively slowly; at low accretion rate, the ratio is $\propto \dot{m}^{-1}$. In any particular event, $\dot{m}$ rises rapidly (on a timescale $\sim t_0$) at first, but then declines slowly. The ratio of observed jet luminosity to thermal luminosity follows the same trend, modulated by any evolution of $\eta B_{\text{jet}}$.

Eventually, the accretion rate falls so low that the disk is no longer radiation-dominated, and our estimate for the pressure in the inner disk is no longer valid. Assuming that this occurs when the gas pressure becomes comparable to
the radiation pressure at \( r \approx 10r_g \), the critical accretion rate is \( \dot{m} \simeq 0.05(\eta/0.1)^{-1}(\alpha/0.1)^{-1/8}M_{\mathrm{BH},6}^{-1/8} \). If the accretion rate declines as \( (t/t_0)^{-5/3} \), it happens at a time

\[
t_0 \simeq 1300t_0\beta^{-1/5}(\eta/0.1)^{1/5}(\alpha/0.1)^{3/40}M_{\mathrm{BH},6}^{-33/40}M_\ast^{(1+3\beta)/10}.
\]

After this time, although the thermal luminosity continues to decline in proportion to the accretion rate, the midplane pressure (and therefore jet luminosity) declines \( \propto \dot{m}^{-5/3} \). This break time is quite late in the development of the flare for low-mass events, but when the black hole mass is relatively large, it could take place as early as \( \simeq 30t_0 \).

4. CONSTRAINING TIDAL DISRUPTION PARAMETERS: THE CASE OF J2058.4+0516

Although the theory of jet and disk emission from tidal disruptions as we have presented it contains a large number of free parameters, in any given event there are also potentially a sizable number of observable properties that may be used to constrain these parameters. In nearly every event, it is possible to measure the characteristic time \( t_0 \). Optical and/or ultraviolet observations often give an indicator of the peak thermal disk luminosity, \( L_{\mathrm{200}} \); we describe this as only an “indicator” because, as we discuss below, there is likely to be a sizable—and uncertain—bolometric correction. When there is hard X-ray emission, the peak jet luminosity \( L_{\mathrm{200}} \) can also be obtained; if an event can be followed long enough, it may also be possible to measure two times related to the transition from super- to sub-Eddington accretion: the time \( t_{\mathrm{m}} \) at which the jet luminosity flattens, and the time \( t_{\mathrm{d}} \), at which the disk luminosity begins to diminish. In this section, we will lay out a general formalism for using these observables as parameter constraints and then apply that method to a specific example, Swift J2058.4+0516.

As shown by Lodato et al. (2009), the accretion rate peaks at a time \( P_{\mathrm{orb}}(a_{\mathrm{min}}) \) past pericenter passage (see Equation (3)) and diminishes thereafter as a power law in time. If we identify the time of peak flare luminosity \( t_0 \) with that orbital period, we obtain the constraint

\[
t_{0,d} \simeq 5.8M_\ast^{(1-3\beta)/2}M_{\mathrm{BH},6}^{1/2}f/kf/0.02 \beta^{-1/6}.
\]

where \( t_0 \) is measured in days.

Because the peak accretion rate is likely to be super-Eddington for almost the entire range of possible black hole masses, we expect the peak disk luminosity to be close to the Eddington luminosity. Unfortunately, however, the characteristic temperature of thermal disk radiation at the Eddington luminosity is \( \sim 1 \times 10^6 \) \( M_{\mathrm{BH},6}^{-1/4} \) K, indicating that the bulk of the light may emerge in the EUV, where direct measurements are very difficult. Consequently, a bolometric correction that could well be \( \sim O(10) \) or greater must be applied to any observable measure of the disk luminosity. Unfortunately, given our current limited understanding of disk spectra even when the accretion rate is sub-Eddington, not to mention potential dust extinction in the host galaxy or possible reprocessing in a wind (Strubbe & Quataert 2011), a sizable uncertainty must be attached to any such correction. For this reason, we give higher priority to the use of other observables.

In our description of the time dependence of the jet and disk luminosities, the transition from super- to sub-Eddington behavior occurs earlier for the disk than for the jet. The ratio of the accretion rates at these two timescales is \( \sim \eta \), the radiative efficiency of the disk in the trans-Eddington regime. Consequently, the ratio of these two timescales primarily constrains the black hole spin, but not any of the other parameters. For the other parameters, no additional information is gained by using both timescales; either one will do.

Suppose, then, that we choose to use the timescale at which the jet luminosity flattens as a function of time, \( t_{\mathrm{jet}} \). This timescale is typically a few times \( t_{\mathrm{fud}} \). For generality, let the accretion rate scale with time as \( (t/t_0)^{-n} \), where \( n \simeq 5/3 \) is expected. Then we find

\[
t_{\mathrm{jet,d}} \simeq 5.8(8 \times 10^3)^{1/n}M_\ast^{(1+11n)/2+33(1/n-1)/2}M_{\mathrm{BH},6}^{-33/8}M_\ast^{(1+3\beta)/10}.\]

where \( t_{\mathrm{jet}} \) is likewise scaled in days. Combining this timescale constraint with the one based on the characteristic flare timescale, we may solve for the black hole mass and the penetration factor:

\[
M_{\mathrm{BH},6} \simeq 4.6 \times 10^4M_\ast t_{0,d}^{-2/9}t_{\mathrm{jet,d}}^{n-1}t_{\mathrm{d}}^{-n}.
\]

and

\[
\beta \simeq 0.093M_\ast t_{0,d}^{-2/9}t_{\mathrm{jet,d}}^{n-1/6}t_{\mathrm{d}}^{n/6}.
\]

Last, using both of these estimates, the (numerous) parameters governing the jet luminosity may be constrained:

\[
B_{\eta \mathrm{jet}}=f(a/M)/(\beta_\ast a) \simeq 0.022L_{\eta \mathrm{jet},46}M_\ast t_{0,d}.\]

where \( L_{\eta \mathrm{jet},46} \) is the peak luminosity of the jet in units of \( 10^{46} \) erg s\(^{-1} \).

The recently discovered flare source J2058.4+0516 (Cenko et al. 2011) may be an example of exactly the sort of tidal disruption event to which this formalism is applicable. Its light curve is shown in Figure 3. For \( \simeq 10 \) days, its flux stayed nearly constant; for the next three months, the flux declined roughly as a power law in time. Around \( t \sim 100 \) days, the decline became shallower; unfortunately, the flux at that point was already near Swift’s detection limit, so the light curve beyond \( \simeq 120 \) days is a mix of detections and upper bounds whose 1\( \sigma \) limits are not very different from the level of some of the last detections. To describe this light curve parametrically, we fit the data for \( t > 7 \) days to a model of the form \( F(t)=At^{-n} \), where \( A \) and \( n \) are the best-fit model is shown as the solid line in Figure 3. That \( \chi^2 \) is reduced by a value of \( t_0 \) within the span of the data supports our prediction that the jet luminosity should become nearly constant after an initial period of decline.

To use the procedure just outlined, we take \( t_0=10 \) days, as that appears to be the point at which power-law decline begins. Following the results of our fit to the light curve, we set \( t_{\mathrm{fud}}=120 \) days, uncertain as that identification may be, and \( n=1.9 \) rather than \( 5/3 \). The difference in the parameter inferences due to the latter choice is probably smaller than the intrinsic error in the method. Using these numbers, we find

\[
M_{\mathrm{BH},6} \simeq 40M_\ast
\]

and

\[
\beta \simeq 0.6M_\ast t_{0,d}^{-2/9}t_{\mathrm{jet,d}}^{n-1/6}t_{\mathrm{d}}^{n/6}.
\]
That is, we estimate that the black hole is in the upper range of masses at which tidal disruptions of main-sequence stars can occur, and the penetration factor is quite modest. The expected mass might be a bit smaller if the stellar mass is less than solar; in that case, $\xi \approx 0.2$, and $\beta$ would increase.

The luminosity in this flare appears to be dominated by hard X-rays (photon energies $\gtrsim 10$ keV) and peaked (in nominal isotropic terms) at $\gtrsim 3 \times 10^{47}$ erg s$^{-1}$ (or perhaps a factor of a few more when contributions from still harder X-rays are considered). The constraint on the jet parameters (Equation (24)) then becomes

$$qB_{\text{jet}} f(a/M)/(\beta_\text{jet} a) \simeq 6 M_\odot \text{ s}^{-1}.$$  

(27)

For self-consistency, the jet must be reasonably beamed and efficient, and the black hole must be spinning rapidly enough to make $f(a/M)$ not too small. If $M_\odot < 1$, the beaming, jet efficiency, spin, etc., must be somewhat greater.

Although the available optical/UV observations are only very rough guides for further inference, they are consistent with these values. Optical photometry suggests a steeply rising spectrum in the observed frame between the $g'$ band and the $u$ band (Cenko et al. 2011). The $u$-band flux on its own translates to a luminosity $\nu L_\nu \simeq 1 \times 10^{35}$ erg s$^{-1}$ given the object’s redshift ($z = 1.18$) and assuming a flat cosmology with $H_0 = 70$ km s$^{-1}$ Mpc$^{-1}$. If $L_{\text{therm}}$ is Eddington-limited and isotropic, this luminosity alone requires a black hole mass $\gtrsim 7 \times 10^6 M_\odot$. However, the mass estimated in this fashion is proportional to the bolometric correction, and, as we have already noted, this may well be at least an order of magnitude. If so, the mass required would be in the range just estimated on the basis of $t_0$ and $t_{\text{jet}}$.

Last, we consider what information may be gleaned from the time dependence of the optical/UV flux. Cenko et al. (2011) report $g'$-band optical photometry taken within 10 days of the flare’s initiation and $\approx 10$ days, $\approx 30$ days, and $\approx 40$ days later. They also report the $u$-band magnitudes at $\approx 10$ days and $\approx 40$ days. In $g'$, there was no detectable dimming until $\approx 30$ days after the flare, but at $\approx 40$ days, the latest time reported, the flux had diminished by $\approx 0.5$ mag, $\approx 40\%$. The $u$-band flux drops about 0.5 mag from $\approx 10$ days to $\approx 40$ days. Thus, in rough terms, the optical/UV luminosity appears to have varied very little over the first month or so, even though the X-ray luminosity fell by a factor $\sim 8$; such behavior is in good agreement with our prediction that the disk output should remain steady while the jet power falls during the super-Eddington phase. If the drop in the 30 days and 40 days makes the disk transition to sub-Eddington behavior, the factor $\sim 3$–4 between that timescale and the possible timescale of the jet transition is also consistent with our suggestion that the accretion rate at the time of disk transition is a factor $\sim \eta^{-1}$ greater than at the time of jet transition.

5. CONCLUSIONS

We have demonstrated that tidal disruption of main-sequence stars by rotating black holes can lead to the launching of a powerful relativistic jet. Because the peak accretion rate in most of these events is likely to be super-Eddington, a regime in which photon trapping is significant, the thermal luminosity may be held to no more than roughly the Eddington luminosity; that is, the thermal luminosity is suppressed well below what the accretion rate would predict on the basis of the usual relativistic radiation efficiency. As a result, the thermal light curve (in bolometric terms) should have a peak that is rather flatter than the curve describing the accretion rate. On the other hand, the jet power suffers no such suppression, so it can exceed the thermal luminosity of the accretion flow.

As the accretion rate declines, the jet power diminishes until the sub-Eddington regime is reached; after that point, we expect that its power changes little until the accretion rate is so low that the disk is no longer radiation-dominated. By contrast, the thermal luminosity can be expected to remain roughly constant at roughly the Eddington luminosity as the accretion rate falls, diminishing with the classic $r^{-5/3}$ scaling only after the flow becomes sub-Eddington and radiatively efficient. The time at which the thermal luminosity begins to fall we have named $t_{\text{f}}$, the time at which the jet light curve levels out is a few times $t_{\text{f}}$.

It follows that in many main-sequence tidal disruptions, the jet luminosity may be comparable to or greater than the thermal luminosity. Particularly when the black hole rotates rapidly, the non-thermal jet photon luminosity may substantially exceed the thermal disk luminosity even after allowing for comparatively inefficient radiation of the jet’s kinetic power. The degree to which we see this non-thermal luminosity depends, of course, on our viewing angle. When it is favorable, the non-thermal emission might be significantly larger than the thermal.

Particularly when the reach of the survey extends only to low redshifts, beaming also enhances the accessible population.

![Figure 3](http://www.swift-psu.edu/monitoring). $F_x$ is Swift counts per second. Triangles show 1$\sigma$ upper bounds. The solid line is the result of a $\chi^2$ fit to the light curve for $t > 7$ days. (A color version of this figure is available in the online journal.)
Although the fraction of sources we see is reduced by $B$, the luminosity distance to which we can see them (for negligible $k$-correction) is increased by $B^{1/2}$, implying an increase in accessible comoving volume of $B^{3/2}/(1 + \zmax)^3$, for $\zmax$ the greatest redshift at which the luminosity of interest can be detected (provided $\zmax \lesssim 1$, so that cosmological corrections are small). Thus, if $B \sim 100$ (as from relativistic beaming with a Lorentz factor $\Gamma \sim 10$), beaming results in an increase in the number of detectable sources by a factor $\sim O(10)$, provided the population has redshifts smaller than $\lesssim 1$.

For these reasons, it may therefore be desirable to augment traditional UV-based methods of searching for stellar tidal disruptions (e.g., Gezari et al. 2008) with other methods more sensitive to jet radiation, perhaps based on the Swift system of all-sky monitoring in hard X-rays coupled to 2–10 keV X-ray follow-up (e.g., using possible future missions such as the Space-based multi-band astronomical Variable Objects Monitor (SVOM; Paul et al. 2011) or the Advanced X-ray Timing Array (AXTAR; Ray et al. 2010)). Because tidal disruption flares last much longer than the $\gamma$-ray bursts for which Swift was designed, the follow-up (for events involving main-sequence stars) can be delayed by as much as $\sim 1$–10 days.

Applying this analysis to the recently discovered example of Swift J2058, we find that this event is consistent with a tidal disruption of a main-sequence star of roughly solar mass by a black hole of mass $\sim 4 \times 10^5 \, M_{\odot}$. Its hard spectrum suggests a jet of the sort we have described, while its characteristic timescale and X-ray luminosity are likewise consistent with this picture. Our model predicts that the jet luminosity should fall when the accretion rate begins to fall, but become constant later in the flare when the accretion rate becomes sub-Eddington; the observed light curve suggests this sort of behavior. Published optical photometry is also consistent with our prediction that the optical/UV thermal luminosity should initially be almost independent of time, although this consistency is weakened by the large, and plausibly time-dependent, bolometric correction that must be applied. Observations using more sensitive detectors might be able to test these predictions more definitively.

Finally, we remark that, although the timescales of Galactic black hole transients are such that we can observe a rich phenomenology of spectral states related to the magnitude of the accretion rate, the corresponding timescales for supermassive black holes in AGNs are far too long to permit human observation. Main-sequence tidal disruptions, on the other hand, present us with a remarkable laboratory for studying how the properties of accretion onto supermassive black holes depend on accretion rate. Beginning well above Eddington, the accretion rate in these systems declines on timescales from weeks to years, and in a fashion we (at least partially) understand. We can hope that further study of these events may shed light on the relation between the magnitude of the accretion rate and thermal disk emission, jets, and coronal X-rays.

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