

# FLARES FROM THE TIDAL DISRUPTION OF STARS BY MASSIVE BLACK HOLES

ANDREW ULMER

Princeton University Observatory, Peyton Hall, Princeton, NJ 08544; andrew@astro.princeton.edu

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## ABSTRACT

Tidal disruption flares are differentiated into two classes: those that are sub-Eddington and those that radiate near the Eddington limit. Flares from black holes above  $\sim 2 \times 10^7 M_\odot$  will generally not radiate above the Eddington limit. For a Schwarzschild black hole, the maximum bolometric luminosity of a tidal disruption is  $\sim L_{\text{Edd}}(5 \times 10^7 M_\odot)$ , which is substantially below the Eddington luminosities of the most massive disrupting black holes ( $\sim 2 \times 10^8 M_\odot$ ). Bolometric corrections to the spectra of the brightest flares are found to be large,  $\sim 7.5$  mag. Nevertheless, the brightest flares are likely to have absolute magnitudes in excess of  $-19$  in  $V$  and  $-21$  in  $U$  (in the absence of reddening). Because the spectra are so blue,  $K$ -corrections may actually brighten the flares in optical bands. If such flares are as frequent as believed, they may soon be detected in low- or high-redshift supernovae searches. The He II ionizing radiation produced in the flares may dominate that which is produced by all other sources in the centers of quiescent galaxies, creating a steady state, highly ionized, fossil nebula with an extent of  $\sim 1$  kpc that may be observable in recombination lines.

*Subject headings:* accretion, accretion disks — black hole physics — galaxies: nuclei — Galaxy: center — quasars: general

## 1. INTRODUCTION

Many lines of indirect evidence suggest that massive black holes reside in the centers of a substantial fraction of galaxies (e.g., Rees 1997; Kormendy & Richstone 1995 and references therein). For example, remarkable stellar-dynamical studies of M32 (van der Marel et al. 1997) and maser studies of NGC 4258 (Watson & Wallin 1994; Miyoshi et al. 1995) show extreme central-mass concentrations. Such concentrations are difficult to understand without a black hole. The most popular alternative explanation of extremely dense clusters of stars or stellar remnants runs into problems (e.g., physical collisions, dynamical relaxation, and even gravitational radiation) at such high densities (e.g., Goodman & Lee 1989), although dynamical “mixtures” of a black hole with a small, dense cluster cannot be excluded (e.g., van der Marel et al. 1997). A nagging theoretical problem with black holes in quiescent galactic centers was that the black hole luminosities were too low for the expected levels of accretion; however, the low luminosities can now be understood in the context of low-efficiency, advection-dominated models (e.g., Narayan, Yi, & Mahadevan 1995).

If black holes do, in fact, reside in the centers of galaxies, one clinching sign would be the tidal disruption of stars by the black hole (e.g., Rees 1988). The detection of a single tidal disruption event could even allow for a crude determination of the black hole mass (e.g., Loeb & Ulmer 1997). Such estimates would yield the black hole mass apart from any contribution from a dense stellar cluster.

Tidal disruption events are rare ( $\sim 10^{-4}$  yr $^{-1}$  per  $L_\star$  galaxy), bright, short (months to years), and are seemingly unavoidable consequences of  $10^6$ – $10^8 M_\odot$  black holes in galactic centers. The general picture for tidal disruption of a star has been developed over the last 20 yr (e.g., Hills 1975; Young, Shields, & Wheeler 1977; Rees 1988). Initially, the rate was thought to be so high as to fuel quasars, but presently more modest rates are predicted, such that disruptions are rare transients that may slowly grow a lower mass,  $\sim 10^7 M_\odot$  black hole (Goodman & Lee 1989).

If a star passes within the tidal radius,  $R_p$ , of the black hole, the star will be disrupted, as has been illustrated by a variety of numerical simulations (e.g., Nolthenius & Katz 1982; Evans & Kochanek 1989; Laguna et al. 1993; Diener et al. 1997). For the Sun,

$$R_t = R_\odot (M_{bh}/M_\odot)^{1/3} \sim 25 R_s M_6^{-2/3}, \quad (1)$$

where  $R_s$  is the Schwarzschild radius, and  $M_6$  is the black hole mass in units of  $10^6 M_\odot$ .

As the star is ripped apart by the tidal forces of the black hole, the debris is thrown into high-eccentricity orbits with a large range of periods, with

$$\Delta E \sim \frac{GM_{bh}R_\star}{R_p^2}, \quad (2)$$

where  $\Delta E$  is the range of debris energy, and  $R_p$  is the pericenter of the star’s orbit (Lacy, Townes, & Hollenbach 1982). The distribution of mass as a function of energy is nearly constant, as is borne out by simulations (e.g., Evans & Kochanek 1989; Laguna et al. 1993).

Half of the debris is unbound, with velocities around  $5 \times 10^3$  km s $^{-1}$ . This unbound debris may possibly produce some long-term observable effects in the surrounding medium. Nucleosynthesis due to the extreme compression of a star may occur in a small fraction of tidal disruptions (Luminet & Pichon 1989; Luminet & Barbury 1990) and lead to observable enrichment within the Galactic center. In active galactic nuclei and Seyferts, a highly enhanced rate of tidal disruptions may produce enough unbound debris to form broad-line and narrow-line quasar-like regions (Roos 1992). Khokhlov & Melia (1996) show that interactions between the high-velocity unbound debris and the interstellar medium (ISM) may result in a supernova-like remnant with properties similar to those of Sagittarius A East. In extremely low angular momentum encounters, the material could deposit over  $10^{52}$  ergs into the ISM, creating a very large supernova remnant. More typically, the energy deposited would be  $\sim 10^{50}$  ergs.

The bound half may create a bright flare as it accretes onto the black hole. The most bound material first returns to pericenter after a time

$$t_{\min} = \frac{2\pi R_p^3}{(GM_{bh})^{1/2}(2R_\star)^{3/2}} \approx 0.11 \left(\frac{R_p}{R_t}\right)^3 \left(\frac{R_\star}{R_\odot}\right)^{3/2} \left(\frac{M_\star}{M_\odot}\right)^{-1} M_6^{1/2} \text{ yr} . \quad (3)$$

The bound material returns to pericenter at a rate (Rees 1988; Phinney 1989) given by

$$\dot{M} \sim \frac{1}{3} \frac{M_\star}{t_{\min}} \left(\frac{t}{t_{\min}}\right)^{-5/3} . \quad (4)$$

The peak return rate, according to simulations by Evans & Kochanek (1989), is

$$\text{Peak Rate} \sim 1.4 \left(\frac{R_p}{R_t}\right)^{-3} \left(\frac{R_\star}{R_\odot}\right)^{-3/2} \left(\frac{M_\star}{M_\odot}\right)^2 M_6^{-1/2} M_\odot \text{ yr}^{-1} , \quad (5)$$

which occurs at  $t \sim 1.5 t_{\min}$ .

The scalings given above are for  $R_p \gg R_s$ . When  $R_p \sim R_s$ , the tidal interactions are stronger (e.g., Frolov et al. 1994). These effects of the Schwarzschild metric can be approximated in the pseudo-Newtonian potential of Paczyński & Wiita (1980),  $\Psi(R) = -GM/(R - R_s)$ . For instance, close to a black hole, the tidal radius,  $R_{t,s} = R_\odot(M_{bh}/M_\odot)^{1/3} + R_s$ , differs from that of the Newtonian potential. The deepest, noncaptured orbit has  $R_p = 2R_s$ , so, as is well known, the maximum mass black hole that can disrupt a solar-type star is  $\sim 1 \times 10^8 M_\odot$ . As is discussed in § 4, encounters more distant than the tidal radius may still strip significant matter from the star, allowing disruptions for black holes up to  $\sim 2 \times 10^8 M_\odot$  in the Schwarzschild case. Also altered is the range of energy of the debris (eq. [2]),

$$\Delta E_s \sim \frac{GM_{bh} R_\star}{(R_p - R_s)^2} . \quad (6)$$

The increase in binding requires that equations (3)–(5) be modified [for the pseudo-Newtonian potential,  $R_p$  is replaced by  $(R_p - R_s)$ ] and has the effect of reducing the minimum return time and of increasing the peak debris return rate.

Another relativistic phenomenon, orbital precession, may lead to the crossing of debris orbits and strong shocks (Rees 1988). The explicit evolution of the debris is complex and is dependent on the black hole's properties as well as those of the star and the stellar orbit. As was discussed by Kochanek (1994), at intersections between the debris streams, energy and angular momentum may be transferred. The computational complexities of the problem prohibit a direct numerical solution, but as an approximation, a large fraction of the debris can be said to circularize in a time

$$t_{\text{cir}} = n_{\text{orb}} t_{\min} , \quad (7)$$

where  $n_{\text{orb}}$  is a number greater than 1 and probably between 2 and 10.

After, and perhaps while, the debris circularizes, it may accrete onto the black hole. The dynamical time at the tidal radius is short relative to the return time, so it is likely that

the disk accretes quickly, resulting in a luminous flare. Cannizzo, Goodman, & Lee (1990) investigated the accretion of a tidally disrupted star in the thin disk case.

Ulmer (1997) has investigated accretion in a thick-disk case. The emergent spectra may have a strong dependence on the structure and orientation of the disk as well as on the possible existence of winds. A starting point for models of the emergent radiation is the temperature associated with an Eddington luminosity emanating from the tidal radius,

$$T_{\text{eff}} \approx \left(\frac{L_E}{4\pi R_t^2 \sigma}\right)^{1/4} \approx 2.5 \times 10^5 M_6^{1/12} \left(\frac{R_\star}{R_\odot}\right)^{-1/2} \left(\frac{M_\star}{M_\odot}\right)^{-1/6} \text{ K} , \quad (8)$$

or from an order of unity factor times the Schwarzschild radius,

$$T_{\text{eff}} \approx \left[\frac{L_E}{4\pi(5R_s)^2 \sigma}\right]^{1/4} \approx 5 \times 10^5 M_6^{-1/4} \text{ K} . \quad (9)$$

In either case, the temperature scaling with black hole mass is slight. Therefore, in the simplest models all disruption flares would have comparable characteristic temperatures.

A tidal disruption event has yet to be definitively detected, but there are a number of suggestive observations. First, a UV flare from the center of an elliptical was observed on the timescale of 1 yr by Renzini et al. (1995), but this flare is  $10^{-4}$  as luminous as might be expected. Furthermore, the spectrum flattens in the near-UV ( $\sim 2000$  Å), which would not be expected except perhaps in the case of a very strong wind (see § 5). The possibility remains, as they suggest, that they have observed the disruption of only the outer layer of a star. Second, Peterson & Ferland (1986) observed the optical brightening of an optically variable Seyfert along with the appearance of a strong, abnormally broad He II recombination line. The event is intriguing, but without a light curve or other definitive evidence, the event is not conclusive. As was pointed out by Rees (1997), the most convincing disruption would be one coming from a quiescent galaxy. Third, observations of variable double-peaked H $\alpha$  lines in the spectrum of Seyfert NGC 1097 may be modeled as emission lines from the appearance of an elliptical disk (possibly formed through tidal disruption) close to a black hole, but other explanations exist (Eracleous et al. 1995; Storchi-Bergmann et al. 1995). Last, Brandt, Pounds, & Fink (1995) observed the X-ray flux from an active galaxy to drop a factor of  $\sim 70$  between observations separated by 1 yr. This event, too, could be related to tidal disruptions.

The outline of the paper is as follows. In § 2 a crude but useful phenomenology of accretion disks is discussed. In § 3 timescales in the tidal disruption event are presented. In § 4 these timescales are used to delineate Eddington-level tidal-disruption flares from the generally less-luminous sub-Eddington flares. In § 5 characteristics of the emergent spectra from these flares are investigated. Long-term emission as a result of tidal disruption events is discussed in § 6. Conclusions are given in § 7.

## 2. ACCRETION DISK PHENOMENOLOGY

In this section, the known relationships between luminosity and mass accretion rate are examined in a variety of accretion regimes that are important in determining the

luminosity of tidal disruption flares. The mass accretion rate is given in terms of the Eddington accretion rate,

$$\dot{M}_E \equiv L_{\text{Edd}}/\epsilon c^2, \quad (10)$$

where  $\epsilon$ , taken to be 10%, is the efficiency of conversion of gravitational binding energy to radiation in the accretion process.

Figure 1 shows a schematic of luminosity versus mass accretion rate. The brightest known accretion disks have high mass accretion rates and are optically and geometrically thick, but thick disks are never much brighter than the Eddington luminosity because at very high accretion rates, pressure gradients push the inner edge of the disk closer to the marginally bound orbit, so less binding energy is released during accretion and more is carried into the black hole (e.g., Paczyński & Wiita 1980). The luminosity increases only logarithmically with the mass accretion rate for thick disks (Paczynski 1980). Thin disks are less luminous, with luminosity scaling linearly with the mass accretion rate (e.g., Shakura & Sunyaev 1973). Slim accretion disks (Abramowicz et al. 1988; Szuszkiewicz, Malkan, & Abramowicz 1995) may form a transition between the thin and thick accretion regime. At very low accretion rates, optically thin, advection-dominated models may exist and are even less luminous (Narayan & Yi 1995). They exist at a maximum accretion rate of  $\sim \alpha^2 \dot{M}_E$  and have  $L \sim \dot{M}^2$ , so the luminosity falls off faster with a decreasing accretion rate ( $\alpha$  is the standard viscosity parameter; e.g., Shakura & Sunyaev 1973). A possibility exists that at accretion rates very much higher than the Eddington rate (and probably beyond those produced by stellar disruptions), the accretion

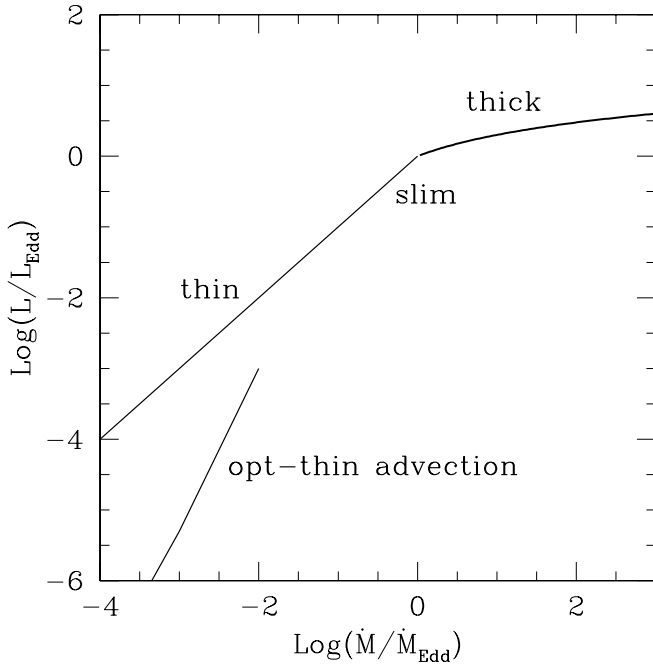


FIG. 1.—Schematic of luminosity vs. mass accretion rate (where  $\dot{M}_{\text{Edd}} \equiv L_{\text{Edd}}/\epsilon c^2$ ), where  $\epsilon = 0.1$  is the efficiency of the accretion process. The brightest accretion disks have high mass accretion rates and are thick. Thick disks are not much brighter than the Eddington luminosity, because at very high accretion rates pressure gradients push the inner edge of the disk closer to the marginally bound orbit, and energy is advected (because of geometry; e.g., Paczyński & Wiita 1980). Thin disks (e.g., Shakura & Sunyaev 1973) and optically thin advection dominated models (e.g., Narayan & Yi 1995) are less luminous.

may become more spherically symmetric and highly optically thick, so the radiative diffusion time could exceed the infall time. In this case, much of the binding energy could be carried into the black hole (e.g., Shvartsman 1971; Ipser & Price 1977; Flammang 1984), and the luminosity could actually fall below the Eddington limit. An additional complication is the possibility of beaming in the thick disk regime, but for simplicity and because of the absence of a preferred beaming model, beaming is set aside for the time being.

The simplified view of accretion disks that I adopt in this paper is that the luminosity increases with increasing mass accretion rate until the Eddington mass accretion rate (eq. [10]), at which point the luminosity reaches the Eddington limit and remains at that level for all higher mass accretion rates. Accretion events with sub-Eddington luminosities occur in thin disks, and accretion proceeds in thick disks for all super-Eddington accretion rates. For sub-Eddington accretion rates, I assume that luminosity falls off at least as fast as linearly with the mass accretion rate (it falls off even faster for advection dominated disks).

### 3. DERIVATION OF TIMESCALES

Although many of the results derived in this section are already known, I briefly rederive them here for completeness and to scale all relevant parameters. The timescales considered are (1) the radiation time, the time to radiate the energy of the bound debris at the Eddington limit; (2) the circularization time, the time for a large fraction of the debris to circularize; and (3) the accretion time, the time to accrete a large fraction of the disk. Throughout the paper,  $m = M_\star/M_\odot$ ,  $r = R_\star/R_\odot$ ,  $M_6$  as the black hole mass in units of  $10^6 M_\odot$ ,  $R_s$  as the Schwarzschild radius,  $2GM_{bh}/c^2$ , and  $R_p$  as the pericenter of the orbit in Schwarzschild coordinates.

The main timescales of the problem can now be written in terms of the system parameters.

The time to accrete all of the bound debris (which is half of the stellar mass) if the black hole radiates at the Eddington luminosity is

$$t_{\text{rad}} = \frac{0.5m M_\odot c^2 \epsilon}{L_E} \approx 21m M_6^{-1} \epsilon_{0.1} \text{ yr} \quad (11)$$

for a hydrogen mass fraction,  $X = 0.74$ , and efficiency,  $\epsilon = 0.1$ .

As was discussed in § 1, the debris are given a range of periods, and the circularization time is closely related to the orbital time of the most tightly bound debris,  $t_{\text{min}}$  (see eqs. [3] and [7]):

$$t_{\text{cir}} \approx n_{\text{orb}} t_{\text{min}} \approx 0.11 n_{\text{orb}} \left( \frac{R_p}{R_t} \right)^3 r^{3/2} m^{-1} M_6^{1/2} \text{ yr}, \quad (12)$$

where  $n_{\text{orb}}$  is the small number of orbits necessary for circularization to occur.

The time in which a large fraction of the matter in a disk can accrete can be written

$$t_{\text{acc}} \approx \frac{2R_p}{v_r} \quad (13)$$

$$\approx \frac{t_{\text{Kep}}(2R_p)}{\alpha \pi h^2} \quad (14)$$

$$\approx \frac{3 \times 10^{-4}}{\alpha h^2} \left( \frac{R_p}{R_t} \right)^{3/2} r^{3/2} m^{-1/2} \text{ yr}, \quad (15)$$

where  $h$  is the ratio of disk height to radius and is approximately one for thick disks, and  $t_{\text{Kep}}$  is the rotation period in a Keplerian potential. In § 2 the geometrical thickness of the disk is discussed, but unless noted otherwise, it is the thick disk case that is of interest.

#### 4. COMPARISON OF TIMESCALES

By considering the main timescales of tidal disruption events, the gross characteristics of the flares are investigated. Of particular interest is the luminosity and whether the flare is super-Eddington or sub-Eddington. The argument that will be made stems from the system's timescales.

Consider the three possibilities: (1) the accretion rate is so high that the material is drained from the disk nearly as fast as it circularizes, (2) the accretion rate is slower than circularization but faster than the radiation time so the material forms a reservoir that accretes in a thick disk, and (3) the material accretes at a rate slower than it circularizes and on a timescale longer than the radiation time, so the disk is not thick, and the debris accretes in the less-luminous form of a thin disk.

Let us first consider these scenarios for a solar-type star with a pericenter at the tidal radius that is disrupted by a  $10^6 M_\odot$  black hole. The parameter,  $n_{\text{orb}}$ , which is related to the efficiency of circularization, is taken as 2. The only unspecified parameter is  $\alpha$ . For a  $10^6 M_\odot$  black hole: (1) if  $\alpha \gtrsim 1.5 \times 10^{-3}$ , then  $t_{\text{acc}} \lesssim t_{\text{cir}}$ , and a thick disk will form and drain as the material circularizes; (2) if  $1.5 \times 10^{-5} \lesssim \alpha \lesssim 1.5 \times 10^{-3}$ , then  $t_{\text{cir}} \lesssim t_{\text{acc}} \lesssim t_{\text{rad}}$ , and a reservoir of debris will form and accrete in a thick disk; and (3) if  $\alpha \lesssim 1.5 \times 10^{-5}$ , then  $t_{\text{acc}} \gtrsim t_{\text{rad}}$ , and a reservoir of debris will form and accrete in a thin disk. Therefore, unless  $\alpha$  is extremely small, much smaller than ( $\sim 10^{-1}$ – $10^{-3}$ ) as is estimated from, for example, DN and FU outbursts (Cannizzo 1992; Lin & Papaloizou 1996 and references therein) and from numerical experiments on the Balbus-Hawley mechanism (Hawley, Gammie, & Balbus 1995; Stone et al. 1996), thick disks are the norm. Therefore, it appears likely that for a  $10^6 M_\odot$  black hole the tidal disruption of a star produces a thick disk that will radiate near the Eddington limit (although it is still uncertain in what energy band the energy will be released).

In Figure 2 these timescales are shown as a function of black hole mass, again for the case of a solar-type star with a pericenter equal to the tidal radius. Because of the different scalings with black hole mass, the circularization time becomes longer than the radiation time at  $M_{\text{bh}} > 2 \times 10^7 \epsilon_{0.1}^{2/3} M_\odot$ . This finding is independent of  $\alpha$  as long as  $\alpha$  is larger than  $\sim 1.5 \times 10^{-4}$ . When the peak accretion rate falls below the Eddington rate above this critical mass, the accretion should occur in a slim disk (Abramowicz et al. 1988) or a thin disk rather than a thick disk. Although little has been said yet about the spectra, it is clear that there may be observational differences between the two classes.

A similar conclusion may be reached by considering when the peak return rate of the stellar debris (eq. [5]) is super-Eddington. For an accretion efficiency of 0.1, the maximum mass is  $1.5 \times 10^7 M_\odot$ . This estimate is slightly smaller than that obtained by considering the timescales because the circularization time is an underestimate, since a

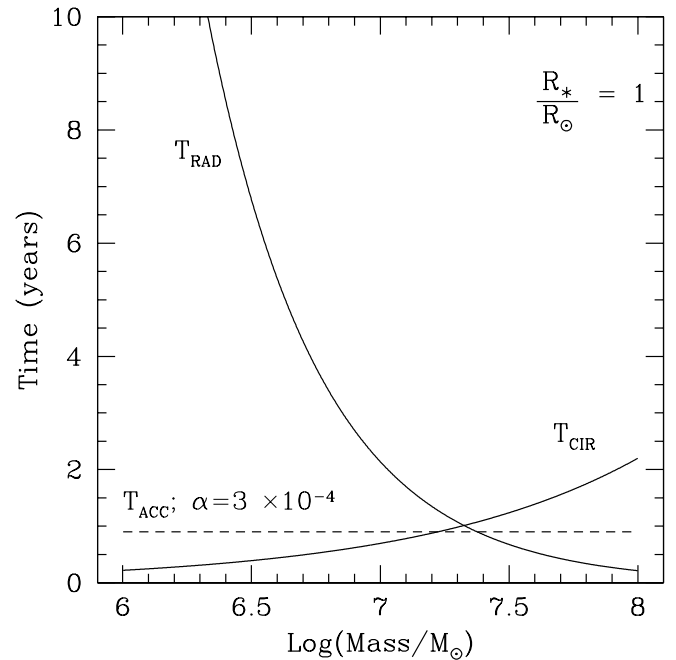


FIG. 2.—Timescales as a function of black hole mass for a solar-type star with a pericenter equal to the tidal radius. For a  $10^6 M_\odot$  black hole, thick disks can form, and bright flares can occur if  $t_{\text{acc}}$  is less than  $t_{\text{rad}}$ , which requires only that  $\alpha > 1.5 \times 10^{-5}$ . Above  $\sim 10^7 M_\odot$ , the Eddington time becomes shorter than the circularization time, so the disk cannot be thick, and the event will be less luminous.

fraction of the matter returns in times longer than the circularization time. Nevertheless, both estimates point to a transition mass around  $10^7 M_\odot$ . For more massive black holes, the return rate is lower, and the luminosity, which for slim to thin disks scales roughly linearly with mass accretion rate, will be lower than the Eddington limit.

The treatment above has been carried out for a solar-type star with a pericenter equal to the tidal radius. Let us next consider the results for a less massive, main-sequence star. These stars will generally be denser, with  $r \approx m^{0.8}$  (e.g., Kippenhahn & Weigert 1990). Then the transition mass will be  $\sim 2 \times 10^7 m^{0.8} M_\odot$ .

The last parameter that will be considered is the pericenter of the orbit that for disruptions may be small as  $2R_s$  in Schwarzschild coordinates. The upper or disruption radius,  $R_d$ , is the radius at which a star in a parabolic orbit is barely disrupted by the tidal field of a black hole. The disruption radius is of order the tidal radius and is given by

$$R_d = q R_t, \quad (16)$$

where the parameter  $q$  is between about 1 and 3 and hides a number of complexities, including relativistic corrections to the potential. Throughout the paper,  $q$  is taken as 2. Young et al. (1977), Luminet & Carter (1986), Khokhlov, Novikov, & Pethick (1993), Frolov et al. (1994), and Diener et al. (1997) contain discussions of this parameter with a range of analytical approximations and numerical methods. The demarcation of  $R_d$  is hazy, as stars still may lose a fraction of their mass at larger pericentric orbits.

According to the theory of loss cones (Frank & Rees 1976; Lightman & Shapiro 1977), there is a critical orbital radius outside of which a star's orbit diffuses in one orbit by

an angle greater than the loss-cone angle. If stars that diffuse onto the loss cone primarily come from inside this critical radius, there will be an excess of events with  $R_p \sim R_d$ . If stars diffuse primarily from large radii, where the diffusion angle is larger than the loss cone, one can use the “*n*σv approximation” that Hills (1975) used to calculate the tidal disruption rate to show that the distribution of lost stars is constant with  $\omega = R_p/R_t$  (gravitational focusing cancels the gain in cross section at larger  $\omega$ ). The true distribution of lost stars will be a combination of the two extremes mentioned, with an excess of events near  $R_p = R_d$  and a tail toward  $R_p = 2R_s$  corresponding to “plunge orbits” in a Schwarzschild metric (Novikov & Thorne 1973). The plunge orbits, in which the star is directly swallowed by the black hole, will not produce any tidal debris or accompanying flare. It is well known that above  $\sim 2 \times 10^8 M_\odot$  the plunge radius surpasses the disruption radius, so no tidal disruption events can be produced.

Tidal disruption events with smaller pericenters will have higher peak luminosities than those with  $R_p \sim R_t$  (see eq. [5]). Therefore, the actual transition mass discussed above is better regarded as a gradual rather than a sharp transition. In Figure 3 the fraction of tidal disruption events with peak accretion rates greater than the Eddington limit is shown for the case of a uniform distribution of  $R_p$  and stars with a Salpeter mass function between 0.3 and  $1 M_\odot$ . More realistic distributions of  $R_p$ , formed by integrating over a cusp, would provide a somewhat sharper transition.

An interesting result of these scaling arguments is that if the luminosity is roughly limited at the Eddington luminosity, there is a maximum luminosity that a tidal disruption event can produce. This luminosity occurs when the peak-mass return rate is the Eddington rate and the pericenter is the minimum ( $2R_s$ ) for a tidal disruption. The

maximum luminosity corresponds to the Eddington limit for a black hole mass  $\sim 5 \times 10^7 \epsilon_{0.1}^{1/3} M_\odot$ , i.e.,  $7 \times 10^{45}$  ergs  $s^{-1}$  and  $2 \times 10^{12} L_\odot$ . The maximum luminosity is probably somewhat lower, because the estimate assumes that all of the gravitational binding energy of the returning debris is released as it returns to pericenter, whereas the circularization process is probably slower. Such extreme events would be rare because of the small range of allowable impact parameters. The peak luminosity could be higher for tidal disruption events that occur in the orbital plane of a Kerr black hole.

Following Rees (1988) and Evans & Kochanek (1989), the duration of a flare may be calculated based on the time that the infall rate is super-Eddington, with the assumption that the circularization and accretion are rapid. The duration is

$$t_{\text{sup}} \approx t_{\text{min}} \left[ 18 \left( \frac{R_p}{R_t} \right)^{-9/5} r^{-9/10} m^{6/5} \epsilon_{0.1}^{3/5} M_6^{-9/10} - 1.5 \right] \quad (17)$$

$$\approx 1.9 \left( \frac{R_p}{R_t} \right)^{6/5} r^{3/5} m^{1/5} \epsilon_{0.1}^{3/5} M_6^{-2/5} \text{ yr} . \quad (18)$$

If the viscosity is high enough that the matter has been accreted roughly as fast as it circularizes, the emission should fall off after  $t_{\text{sup}}$ .

## 5. FLARE ENERGY SPECTRUM AND ABSOLUTE MAGNITUDES

The peak bolometric luminosity of the flares is determined from the scaling arguments given above. According to this scenario, the brightest flares are those in which accretion occurs in a thick disk (or in a slim disk), so a study of the spectrum for that case alone is the main focus. Cannizzo et al. (1990) have investigated the spectrum of a flare if the debris forms a thin disk. At the earliest and brightest times, the spectrum was found to peak at  $\sim 300 \text{ \AA}$ , with  $U$  and  $V$  absolute magnitudes of  $-16.7$  and  $-15.2$  for a  $10^7 M_\odot$  black hole.

Consider thick disk spectra to begin with. Figure 4 shows the spectra calculated using a thick disk code (Ulmer 1997). The spectra presented are typical thick disk spectra for a solar-mass star disrupted at  $\sim 2R_t$ , the disruption radius (eq. [16]), by a  $10^7 M_\odot$  black hole. Viewing angle will have a large effect on the spectra, because when seen edge-on the bright, hot funnel is obscured by the disk. Viewed face-on, the disk is very bright, radiating slightly above the Eddington limit. When viewed edge on, the bolometric luminosity is less by a factor of  $\sim 50$ . For the thick disks considered here, the luminosity is 2–3 times the Eddington limit. The absolute magnitudes for a disk around a  $10^7 M_\odot$  black hole radiating at 2.5 times the Eddington limit are, for face-on and edge-on disks,  $U: -21, V: -19.5$  and  $U: -19, V: -17.5$ , respectively. For a similarly bright disk around a  $10^6 M_\odot$  black hole, the absolute magnitudes are  $U: -17, V: -15.5$  and  $U: -15, V: -13.5$ . The bolometric corrections are  $\sim -7.5$  mag in  $V$ , where reddening, which may well be important in the center of a galaxy, and dimming from cosmological effects [although in a flat cosmology, very high redshift flares would brighten as  $(1+z)$  because of  $K$ -corrections] are neglected. The disks may radiate closer to the Eddington limit than do the disk models above, which radiate at 2–3 times the Eddington limit.

A thick disk may drive a strong wind. If a wind were dense enough, the effective photosphere would be moved

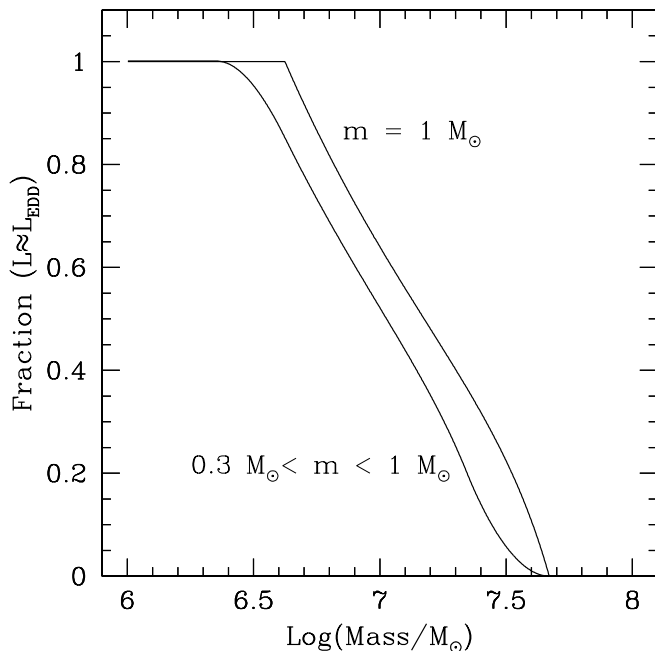


FIG. 3.—The fraction of tidal disruption events with peak accretion rates greater than the Eddington limit are shown (upper curve) for the case of solar-type stars with a uniform distribution of pericenters and (lower curve) stars with a Salpeter mass function between 0.3 and  $1 M_\odot$ , also with a uniform distribution of pericenters.

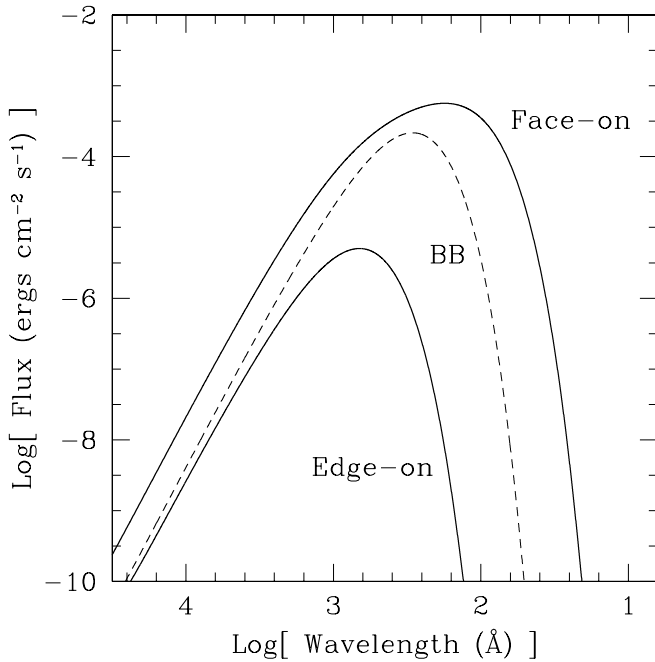


FIG. 4.—Predicted spectra of a tidal disruption flare around a  $10^7 M_\odot$  black hole. The flux is appropriate for a source at a distance of 100 Mpc. The two solid lines show calculated spectra from an optically and geometrically thick disk as viewed from face-on and edge-on. The former is much brighter and hotter because a large fraction of the energy is radiated in the funnel. The dashed line is a Planck curve with  $T = 10^5$  for comparison.

outward from the surface of the disk, and the spectra would appear redder. Some variable stars, namely luminous blue variables, show what are believed to be wind-induced color changes resulting in  $V$ -band brightening of 2 mag (e.g., Humphreys & Davidson 1994). In super-Eddington thick disks, strong winds are to be expected on the purely phenomenological grounds that in stars, winds increase as a star's luminosity approaches the Eddington limit. Additionally, instabilities in the funnel of thick disks may drive strong winds as well as jets (e.g., Jaroszyński, Abramowicz, & Paczyński 1980; Nityananda & Narayan 1982). Alternatively, a strong wind could be driven before the debris completely settles into a disk. Rees (1988) observed that the total energy required to eject the majority of the bound material is small (see eq. [2]),  $\sim M_\odot c^2 R_\star R_s/R_p^2$ .

Limits on the strength of the wind are now examined. The mass-loss rate is limited by the energy required to lift the material from the disk and carry it to infinity. In the context of neutron stars, Paczyński & Prószyński (1986) have considered energy balance limits on winds. Following Paczyński & Prószyński, the mass-loss rate is

$$\dot{m} \approx \frac{(L_b - L_{ph})2R_b}{c^2 R_s} \quad (19)$$

$$\lesssim \frac{L_{\text{Edd}} 2R_b}{c^2 R_s}, \quad (20)$$

where  $R_b$  is the base radius from which the matter is lifted into the wind by the radiation, and  $L_b$  and  $L_{ph}$  are the luminosity at the base radius and at the photosphere. If the base radius is taken to be  $\sim 10R_s$ ,  $\dot{m} \lesssim 0.05 M_\odot M_6 \text{ yr}^{-1}$ .

The wind may, therefore, carry away a sizable fraction of the disk mass for a  $10^7 M_\odot$  black hole.

Given a maximal mass-loss rate, the maximum photospheric radius may be estimated as follows. Near the photosphere,  $\dot{m} = 4\pi v \rho R_{ph}^2$  and  $\rho R_{ph} \sim m_p/\sigma_T$ , where  $m_p$  is the proton mass, and  $\sigma_T$  is the Thompson cross section. If  $v \sim v_{\text{esc}} \sim \sqrt{GM_{\text{BH}}/R_{ph}}$ , then

$$R_{\text{phot}} \lesssim \frac{\dot{m}^2 \sigma_T^2}{32\pi^2 GM_{\text{BH}} m_p^2} \quad (21)$$

$$\lesssim 120 R_s \left( \frac{v}{v_{\text{esc}}} \right)^{-2}. \quad (22)$$

The minimum effective temperature is expected to be  $\sim 1.1 \times 10^5$  K, which is still very hot. The velocity of the outgoing wind could be lower than the escape velocity, and if the luminosity is very close to the Eddington limit at the photosphere, the escape velocity would be reduced by  $1 - L/L_{\text{Edd}}$  from the expression given above, creating a larger photosphere. Including a photosphere with  $T \sim 10^5$  K, which radiates at the Eddington limit in the thick disk spectra calculated above, damps the high-energy spectrum but has little effect on the  $U$ - and  $V$ -band luminosities. The optical bands are weakly affected, even though a larger fraction of the light goes into optical, primarily because the total luminosity of the source is diminished as energy is carried away by the wind, and because a significant fraction of the optical light comes from part of the disk outside the photosphere at a few hundred  $R_s$ . The optical luminosities of a thick disk and of a thick disk with a wind are nearly the same. The main effects of a strong wind are the reduction of the very high energy spectrum of the flare and a possible shortening of the flare by loss of the stellar debris to the wind. Such strong winds would greatly reduce the expected number of soft X-ray tidal disruption flares (Sembay & West 1993) accessible, for instance, to *ROSAT*.

A third possibility, and the most optically luminous one, is that an extended, static envelope could be formed around the disk (Loeb & Ulmer 1997). Such an envelope might form either if the initial collisions of the debris streams or subsequent evolution of the disk stream expels much of the matter to large radii and strong radiation pressure disperses this marginally bound gas into a quasi-spherical configuration. In this scenario the object is spherically symmetric with Eddington luminosity and effective temperature,

$$T_{\text{eff}} \approx 2.3 \times 10^4 \text{ K} \left( \frac{M_{\text{bh}}}{10^7 M_{\text{env}}} \right)^{1/4}, \quad (23)$$

where  $M_{\text{env}}$  is the mass in the envelope. In this case, the bolometric corrections would be much less ( $\sim 2.5$ ), and the objects would be as bright as Seyferts.

In addition to observing the very blue continuum of the thick disk, it may be possible to observe a flare in emission lines. The disk may radiate in emission lines but the extent is uncertain because the emission would depend on the temperature and density structure both of the disk photosphere and of any wind. As was discussed in Roos (1992) and Kochanek (1994), prior to circularization the bound material generally has densities higher than  $10^{12}$  atoms  $\text{cm}^{-3}$ , so most of the energy the material radiates (or absorbs and reradiates) will be in the continuum and not in emission lines. Moreover, such debris has a relatively small

covering fraction. Recombination lines, e.g., H $\alpha$ , might be radiated from the recombination of the debris or from the ionization and subsequent recombination of a neutral surrounding environment. The possibility that the disk may radiate in recombination lines was suggested by Eracleous et al. (1995) and Storchi-Bergmann et al. (1995) to explain observations of variable H $\alpha$  lines in spectrum of the Seyfert NGC 1097. However, the disk is probably extremely optically thick even at thousands of  $R_s$ , with typical surface density  $0.5 M_\odot / \pi (1000 R_s)^2$  of  $3000 M_\odot^{-2} \text{ g cm}^{-2}$ , so any emission lines would have to be formed close to the surface of the disk and would be suppressed from optically thin estimates (Storchi-Bergmann et al. 1995, eq. 2).

## 6. LONG-TERM EMISSION

Because much of the flare's energy is radiated between 10 and 100 eV, the possibility that the H II region created by the ionizing radiation of a flare could be observed is intriguing (e.g., Rees 1997). Assuming that  $\sim 5\%$  of the rest mass of the star is radiated in the flare and using the thick disk spectra discussed above, the total number of photons between  $\sim 100$  and  $1000 \text{ \AA}$  is  $N_\gamma \sim 10^{63}$ . The ISM surrounding the black hole will be in a steady state, because the recombination time,  $t_{\text{rec}} \sim 10^6 n^{-1} T_4^{1/2} \text{ yr}$ , is longer than the time between disruptions ( $t_{\text{dis}} \sim 10^4 \text{ yr}$ , where  $n$  is the number of atoms per  $\text{cm}^3$ ). The ionizing photon rate is therefore  $\sim 3 \times 10^{51} \text{ s}^{-1}$ , or  $\sim 1\%$  the total ionizing photon rate from the Galaxy (Mezger 1978; Bennett et al. 1994). Direct detection of H $\alpha$  from tidal disruption ionization is therefore unlikely.

At higher photon energies, the tidal disruption events may (in a time-averaged sense) contribute the majority of photons. The production rate of He II ionizing photons by tidal disruption events is  $\sim 3 \times 10^{50} \text{ s}^{-1}$ , which, if a large fraction (e.g., 10%) of the photons ionize He II and  $\sim 20\%$  of these ionization produce He II  $\lambda 4686$ , then the luminosity in He II  $\lambda 4686$  (Pacschen line) would be  $\sim 2 \times 10^{37} \text{ ergs s}^{-1}$ . The number of He II photons would be reduced by a factor of  $\sim 50$  in the case of the strongest winds described above. For densities of  $1 \text{ atom cm}^{-3}$ , the radius of the nebula would be

$$R_{\text{H II}} \approx \left( \frac{3N_\gamma}{4\pi n^2 \alpha_B} \right)^{1/3} \approx 1.5 \times 10^{21} n^{-2/3} T_4^{1/6} \text{ cm}, \quad (24)$$

where  $\alpha_B$  is the case B recombination coefficient for hydrogen. A prediction of tidal disruptions is therefore a highly ionized nebular emission in the inner kpc of galaxies. Such nebular emission may be produced by a variety of sources. He II  $\lambda 4686$  lines, for instance, may be produced by Of stars (Bergeron 1977) or Wolf-Rayet stars (e.g., Schmutz, Leitherer, & Gruenwald 1992), although such lines would generally be broad ( $\sim 3\text{--}10 \text{ \AA}$ ), whereas nebular lines would be narrow. Possible competing ionizing sources include shocks (Shull & McKee 1979), X-ray binaries (Garnett et al. 1990), and Wolf-Rayet stars. While all of these objects are concentrated toward galactic centers, it may be possible to distinguish between a tidal disruption nebula and most other ionizing sources. The fossil nebulae may be most readily observed in galaxies where the background ionizing level is small, corresponding to a small star formation rate, but since ample gas is required to absorb the radiation, the best candidates for detection of the nebulae may be Sa galaxies. Such nebulae may also be visible in radio lines.

Harder photons,  $\sim 10 \text{ keV}$ , may possibly be produced in a tidal disruption event, especially if the funnel drives a low-density jet, as opposed to a high-density wind, as was discussed above. Vainshtein & Sunyaev (1980) suggested that scattering of such hard photons, particularly in the K $\alpha$  lines of Fe, S, and Si, could allow detection of a transient X-ray source long after it had faded. Recently, Sunyaev, Markevitch, & Pavlinsky (1993) used such techniques to study the X-ray brightness history of our Galactic center and concluded that the center has not been significantly brighter in the last 400 yr.

The idea that the unbound debris may produce a supernova remnant has been discussed by Khokhlov & Melia (1996). For typical disruptions, the unbound material is emitted in a fan of small opening angle,  $\Omega < 0.1 \text{ rad}^2$ , with an energy of  $\sim 10^{50} \text{ ergs}$ . Rare events with pericenters close to the Schwarzschild radii may provide  $10^{52} \text{ ergs}$ . Such events may produce supernova remnants that could be observed in external galaxies. For remnants that sweep up 100 times their initial mass, the remnants would be displaced from the black hole by

$$R_{\text{disp}} \approx 40 \Omega_{0.1}^{-1/3} n^{-1/3} \text{ pc}. \quad (25)$$

A supernova remnant lasts  $\sim 10^4 \text{ yr}$ , corresponding to one such remnant per galaxy.

## 7. CONCLUSIONS

Tidal disruption flares occur in two classes: those that are sub-Eddington and those that radiate near the Eddington limit. Flares from black holes above  $\sim 2 \times 10^7 M_\odot$  will generally not radiate above the Eddington limit.

The brightest flares will occur around  $\sim 5 \times 10^7 M_\odot$  black holes, but they will be relatively rare, as they require a finely tuned impact parameter. The majority of the brightest events will occur around  $\sim 2 \times 10^7 M_\odot$  black holes. The corresponding timescale for the decay of such flares (eq. [17]) is about 1 month. For all of the spectral scenarios considered, the optical luminosity is high, with peak apparent magnitudes in  $U$  and  $V$  of  $-21$  and  $-19.5$  for a  $10^7 M_\odot$  black hole and  $-19$  and  $-17.5$  for a  $10^6 M_\odot$  black hole, neglecting reddening. Both high- and low-redshift supernovae searches may observe the flares. As tidal disruption events occur at a rate of  $\sim 100$  times less frequently than supernovae, many supernovae must be observed before a flare is expected to be seen. Furthermore, if the flares come from the centers of galaxies, they may be harder to detect against the bright nuclei. The extreme blueness of the flares may help separate them from the background. In fact, because the spectra are so blue,  $K$ -corrections may actually brighten the flares in optical bands.

The He II ionizing radiating produced in the flares may dominate that which is produced in the centers of quiescent galaxies, creating a steady state, highly ionized diffuse nebula with an extent of  $\sim 1 \text{ kpc}$  that may be observable in recombination lines in all galaxies with both tidal disruption events and gas near the nucleus.

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