Light Curves of Tidal Disruption Events in Active Galactic Nuclei

Chi-Ho Chan1,2 □, Tsvi Piran1 □, and Julian H. Krolik3 □
1 Racah Institute of Physics, Hebrew University of Jerusalem, Jerusalem 91904, Israel
2 School of Physics and Astronomy, Tel Aviv University, Tel Aviv 69978, Israel
3 Department of Physics and Astronomy, Johns Hopkins University, Baltimore, MD 21218, USA

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Abstract

The black hole of an active galactic nucleus is encircled by an accretion disk. The surface density of the disk is always too low to affect the tidal disruption of a star, but it can be high enough that a vigorous interaction results when the debris stream returns to pericenter and punches through the disk. Shocks excited in the disk dissipate the kinetic energy of the disk interior to the impact point and expedite infall toward the black hole. Radiatively efficient disks with luminosity $\gtrsim 10^{-3}$ Eddington have a high enough surface density that the initial stream–disk interaction leads to energy dissipation at a super-Eddington rate. Because of the rapid infall, only part of this dissipated energy emerges as radiation, while the rest is advected into the black hole. Dissipation, infall, and cooling balance to keep the bolometric luminosity at an Eddington-level plateau whose duration is tens of days, with an almost linear dependence on stellar mass. After the plateau, the luminosity decreases in proportion to the disk surface density, with a power-law index between $-3$ and $-2$ at earlier times, and possibly a steeper index at later times.

Unified Astronomy Thesaurus concepts: Tidal disruption (1696); Hydrodynamical simulations (767); Active galactic nuclei (16); Accretion (14); Gravitation (661); Black hole physics (159)

1. Introduction

When a star of mass $M_*$ and radius $r_*$ approaches a black hole of mass $M_\bullet$ on an orbit whose pericenter $r_p$ is comparable to or smaller than the tidal radius $r_t \equiv r_*(M_*/M_\bullet)^{1/3}$ (Hills 1975), the tidal gravity of the black hole rips the star asunder (e.g., Rees 1988), leading to a tidal disruption event (TDE). The disruption takes place very close to the black hole: in units of the gravitational radius $r_g \equiv GM_\bullet/c^2$, the tidal radius is

$$r_t \approx 50 \frac{r_g}{M_\bullet} \left( \frac{M_*}{M_\bullet} \right)^{-2/3} \left( \frac{r_*}{r_g} \right)^{-1/3} \left( \frac{r_p}{r_g} \right)^{1/3}. \quad (1)$$

The mass return time, determined by the specific energy of the most bound debris, is the timescale on which this part of the debris returns to pericenter:

$$t_{\text{mb}} \sim 40 \text{ days} \times \left( \frac{M_\bullet}{10^6 M_\odot} \right)^{1/2} \left( \frac{M_*}{M_\odot} \right)^{-1} \left( \frac{r_*}{r_g} \right)^{3/2} \left( \frac{r_p}{r_*} \right)^3. \quad (2)$$

General relativistic simulations of TDEs starting from stars with realistic structures reveal that these quantities are accurate only within a factor of $\sim 2$ (Ryu et al. 2020).

The black hole of an active galactic nucleus (AGN) is girdled by an accretion disk. When a TDE happens in the vicinity of such a black hole, the passage of the star through the disk leaves the disk largely intact, and the disruption proceeds as if the disk were absent. However, upon return to the vicinity of the black hole, the debris of the disrupted star, being much more extended and dilute than the star, can interact with the disk in a more interesting way (Kochanek 1994; see also Kathirgamaraju et al. 2017).

As we shall argue later, the stream typically has much more inertia than the disk, so it acts as an immovable obstacle to disk rotation. Shocks form where the disk runs into the stream, and the shocks dissipate disk kinetic energy. When the disk surface density is large enough, the dissipation rate can be highly super-Eddington. This sort of shock dissipation could serve as the energy source for a bright flare different from the energy sources in ordinary TDEs, namely, accretion of any rapidly circularized debris (e.g., Rees 1988) and shocks due to stream self-interaction (Piran et al. 2015).

In Chan et al. (2019), we performed the first hydrodynamics simulations of the collision between the debris stream of a TDE and the pre-existing accretion disk of an AGN. We found that as long as the stream is much heavier than the disk, our simulation results are sensitive only to how dense the disk is, not how dense the stream is. This observation allows us to extrapolate our simulation results to even heavier streams and make quantitative predictions, even though our simulations did not explicitly cover that regime.

We begin by highlighting the most salient simulation results in Section 2. Based on these results, we estimate in Section 3 the bolometric light curve expected when a debris stream tears through the disk. In Section 4, we scrutinize the dependence of the light curve on black hole mass and disk properties, and we present sample light curves. This discussion is followed by a comparison with previous theoretical models and observations in Section 5. We end with our conclusions in Section 6.

2. Summary of Simulations

In Chan et al. (2019), we simulated the collision between the debris stream of a TDE and the pre-existing accretion disk of an AGN. As illustrated in Figure 1, a parabolic stream, representing the returning debris, strikes the disk perpendicularly from above as it reaches pericenter. To study late-time behavior, here we extend the simulations from the original duration of $\sim 12$ days to double that. The longer duration is still a fraction of $t_{\text{mb}}$, so it is reasonable to keep the stream conditions constant.

We adopt as fiducial parameters $M_* = M_\odot$, $r_* = R_\odot$, and $r_p = r_t$. Since our interest is in disk physics, the relevant timescale is the disk orbital time at the impact point, which is
The collision excites multiple spiral shocks in the disk inside the stream impact radius. The shocks remove angular momentum and dissipate kinetic energy, thereby clearing out the inner disk within several $t_{\text{orb}}$. Magnetohydrodynamic (MHD) stresses are ignored in the simulations, but they act much more slowly than shocks. The top left panel of Figure 2 tells us that the inflow time $t_{\text{infl}}$ or the time it takes shocks to redirect gas toward the black hole, gradually declines over the course of the event, from $\sim 1$ day at the start to $\sim 0.4$ day at $\gtrsim 10$ days.

As impact with the disk shaves off part of the stream, the inner disk mass is partially replenished with stream material. The inner disk mass is determined by the balance between inflow and replenishment. The longer simulations make clear this dynamic. The top center panel of Figure 2 demonstrates that the inner disk mass drops roughly in proportion to the disk aspect ratio at the impact point. The top right panel shows that $\gamma \equiv d \ln \Sigma / dt \ln t$ lies between $-3$ and $-2$ toward the end of the simulations, and the decrease of $\Sigma$ gradually steepens over time. We emphasize that the time evolution of $\Sigma$ is due to mechanisms wholly unrelated to those determining the power-law decay of the mass return rate at $t \gtrsim t_{\text{infl}}$.

The dissipated energy heats up the inner disk. As portrayed in the bottom left panel of Figure 2, heating raises the disk aspect ratio at the impact point to $H/r_p \sim 1/3$. The cooling time $t_{\text{cool}}$ is the time the inner disk takes to release its internal energy as radiation, including the time needed to convert internal energy to radiation, and the time needed for radiation to escape through free streaming, diffusion, or vertical advection due to convection and magnetic buoyancy. Because electron scattering dominates the opacity in these circumstances, the diffusion time is $\sim \tau H/c$, where $\tau = \Sigma \kappa T$ is the Thomson optical depth and $\kappa T$ is the cross section per mass for Thomson scattering. The diffusion time is long because $\tau$ drops rather slowly but $H/r_p \lesssim 1$ is much larger than in the unperturbed disk; in fact, the diffusion time is typically $\gtrsim t_{\text{infl}}$ at early times. The cooling time could be even longer because it also takes into account the radiation time. This means most of the radiation is trapped in the inflow toward the black hole and only a small fraction escapes.

3. Estimation of Bolometric Luminosity

The swift inflow in the perturbed inner disk could produce a flare. Similar to a steady-state disk, the energy of the flare ultimately comes from the release of the gravitational energy of the disk gas. The contrast with a steady-state disk is twofold: first, dissipation and inflow happen on much shorter timescales; second, mass resupply from the stream allows more energy to be released than the gravitational energy possessed by the unperturbed disk.

We proposed in Chan et al. (2019) a crude estimate for the bolometric luminosity $L_{\text{c}}(t)$ of the collision-induced flare. Here we use the modified form

$$L_{\text{c}} \sim Q \min(1, t_{\text{infl}}/t_{\text{cool}}),$$

where $Q(t)$ is the energy dissipation rate. This equation encapsulates a competition between the inner disk releasing its internal energy as radiation, represented by $t_{\text{cool}}$, and the inflow sweeping this radiation into the black hole, represented by $t_{\text{infl}}$. When $t_{\text{infl}}/t_{\text{cool}} \lesssim 1$, the ratio $t_{\text{infl}}/t_{\text{cool}}$ estimates the fraction of dissipated energy the inner disk manages to radiate.
away; when $t_{\text{infl}}/t_{\text{cool}} \gtrsim 1$, all the dissipated energy is radiated away before the gas is accreted, so $L_c \sim Q$. Equation (4) is only an estimate; the true luminosity must be determined by performing time-dependent radiative transfer on the actual distribution of heating and opacity.

The top left panel of Figure 2 and Equation (3) together tell us that

$$t_{\text{infl}} \sim 4t_{\text{orb}}.$$  

Gas does not plunge into the black hole; rather, it spirals inward along trajectories that shift from quasicircular near the impact radius to more eccentric at smaller radii. When disk material strikes the obstacle posed by the stream, it loses a large part of both its angular momentum and its energy; when it is deflected by shocks at smaller radii, the fractional loss of angular momentum is greater. For this reason, the heating rate as a function of time is reasonably estimated by the orbital kinetic energy of gas orbiting at the impact radius per inflow time:

$$Q \sim \frac{\pi GM_b \Sigma r_p}{t_{\text{infl}}} \sim \frac{\tau}{8\pi} \left( \frac{r_p}{r_g} \right)^{-1/2} \frac{t_{\text{orb}}}{t_{\text{infl}}} L_E,$$  

where $L_E = 4\pi GM_b c \kappa_T$ is the Eddington luminosity. The bottom center panel of Figure 2 demonstrates that, apart from an initial adjustment phase, this estimate is accurate to within a factor of $\sim 2$.

In estimating $t_{\text{cool}}$, we ignore the contributions of radiation time and vertical advection for simplicity. In fact, finite radiation time slows cooling but vertical advection speeds it, hence the two effects partially cancel. The inner disk starts out optically thick, so

$$t_{\text{cool}} \sim \frac{\tau H}{c} \sim \frac{\tau}{4\pi} \left( \frac{r_p}{r_g} \right)^{-1/2} \frac{t_{\text{orb}}}{t_{\text{infl}}}.$$  

where we used $H/r_p \sim \frac{1}{2}$ in the second step. The inner disk eventually becomes so depleted that $\tau \ll 1$ and radiation can escape freely, at which point $L_c \sim Q$. This free-streaming limit is reached only when $t_{\text{infl}}/t_{\text{cool}} \gtrsim 1$, since $t_{\text{infl}} \gtrsim t_{\text{orb}}$ always and $t_{\text{cool}} \lesssim t_{\text{orb}}$ in this limit. Considering that $L_c$ is capped by the minimum function in Equation (4) long before the free-streaming limit kicks in, the limit is irrelevant and we can use Equation (7) for all $\tau$.

We see from Equation (7) that, as long as the inner disk maintains a high enough $\Sigma$ to make

$$\tau \gtrsim 4\pi \left( \frac{r_p}{r_g} \right)^{1/2} \frac{t_{\text{infl}}}{t_{\text{orb}}},$$  

we have $t_{\text{infl}}/t_{\text{cool}} \lesssim 1$. Combining Equations (4), (6), and (7) in this regime yields

$$L_c/L_E \sim \frac{1}{2}.$$  

This is our main result: the bolometric light curve of a TDE in an AGN starts off with an Eddington-level plateau. Our simulations confirm the validity of Equation (9): $Q_{\text{infl}}/t_{\text{cool}}$ in the bottom right panel of Figure 2 quickly stabilizes to $\sim L_E$. The remarkable constancy of $L_c$ is due to the cancellation of $t_{\text{infl}}$ and $\tau$ in the derivation of Equation (9) (see also...
Krolik 2010). When the inner disk is cleared out to the point that Equation (8) is violated and $t_{\text{int}}/t_{\text{cool}} \gtrsim 1$, the luminosity falls as $L_\nu \sim Q \propto \Sigma$. Both the plateau duration and the post-plateau $L_\nu$ depend on the time evolution of a single parameter, $\Sigma$.

Because the stream feeds the inner disk, the plateau duration is not limited by the initial inner disk mass and can be $\gg t_{\text{inf}}$; the plateau duration will be estimated in Section 4. A plateau may not appear if radiation time, convection, or magnetic buoyancy modifies the functional form of $t_{\text{cool}}$. Our calculations give us only a crude estimate of the bolometric luminosity in some time-averaged sense. Accurate predictions of the light curve and the spectrum should be based on detailed radiation MHD simulations.

4. Light Curves

We consider two disk models for the unperturbed disk, depending on how the unperturbed accretion rate $M_\nu = L_\nu/(\eta \dot{\nu})$ compares with the Eddington accretion rate $M_{\nu,\text{E}} = L_{\nu,\text{E}}/(\eta \dot{\nu})$, where $\eta = 0.1$ is the fiducial radiative efficiency of a radiatively efficient disk. Any time-steady disk satisfies $M_\nu \sim 4\pi R \Sigma_0 v_R$, where $\Sigma_0$ and $v_R$ are the unperturbed surface density and radial velocity respectively. The value of $\Sigma_0$ is fixed by $M_\nu$; therefore, a choice of disk model boils down to a choice of $v_R$. The inflow mechanism in the unperturbed disk is the outward transport of angular momentum by internal stresses, but once the stream strikes the disk, inflow is driven by spiral shocks instead.

If $M_\nu/M_{\nu,\text{E}} \gtrsim 10^{-3}$, we assume the disk is geometrically thin, optically thick, and radiatively efficient (Shakura & Sunyaev 1973) with fiducial stress parameter $\alpha = 0.1$. Both the disk aspect ratio and $v_R/v_0$ at the impact point are $\ll 1$ in this model, $v_0$ being the orbital velocity. As a result, $\Sigma_0$ is large, and the initial energy dissipation rate $Q_0$, obtained by substituting the unperturbed surface density $\Sigma_0$ into Equation (6), is large as well.

If $M_\nu/M_{\nu,\text{E}}$ is any lower, the disk could be a geometrically thick, optically thin, radiatively inefficient accretion flow (RIAF) (e.g., Ichimaru 1977; Rees et al. 1982; Narayan & Yi 1994). Parameterizing the radial velocity in a geometrically thick disk as $v_R \sim \alpha' v_0$, we write

$$\kappa T \Sigma_0 \sim \frac{1}{\alpha' \eta'} \frac{L_\nu}{L_{\nu,\text{E}}} \left(\frac{\rho_{\text{crit}}}{\dot{\nu}}\right)^{1/2}.$$  

We emphasize that $\alpha'$ is merely a parameterization; for simplicity we take $\alpha' = 0.1$. Ryan et al. (2017) performed general relativistic MHD simulations of RIAFs assuming that electrons are heated by Coulomb scattering off ions at a rate derived empirically from solar wind measurements (Ressler et al. 2015). They found that the effective radiative efficiency is $\eta' \sim 10^{-2}$ for $M_\nu/M_{\nu,\text{E}} \sim 10^{-3}$, suggesting that accretion may proceed at efficiencies close to radiatively efficient values even at very low accretion rates. We adopt this value of $\eta'$ below.

The much larger $v_R/v_0$ in an RIAF compared to a radiatively efficient disk means that $\Sigma_0$ and hence $Q_0$ are much smaller.

The top panel of Figure 3 displays $Q_0$ for a radiatively efficient disk, and Figure 4 does the same for an RIAF. We end Figure 3 at $M_\nu/M_{\nu,\text{E}} = L_\nu/L_{\nu,\text{E}} = 10^{-4}$ and begin Figure 4 at

![Figure 3](image-url)

Figure 3. Color contours of three properties of the perturbed inner disk, as functions of the black hole mass $M_\nu$ and the Eddington ratio $L_\nu/L_{\nu,\text{E}}$ of the unperturbed disk, with other parameters fixed at their fiducial values (Sections 2 and 4). This figure portrays the case in which the unperturbed inner disk is a radiatively efficient disk, while Figure 4 illustrates the case of an RIAF. The top panel shows the initial energy dissipation rate $Q_0$, divided by the Eddington luminosity $L_{\nu,\text{E}}$. The middle panel shows the initial value of the ratio of inflow to cooling time $t_{\text{inf}}/t_{\text{cool}}$. The bottom panel shows the plateau duration in units of days (Appendix). The blue regions outside the dashed–dotted lines have only an upper limit, while the empty regions have no plateau at all because $t_{\text{inf}}/t_{\text{cool}} \gtrsim 1$ initially. The crosses are the values of $(M_\nu, L_{\nu,\text{E}})$ for which light curves are shown in Figure 5. In all panels, the lower dashed line at $L_{\nu,\text{E}} \sim 1.7 \times 10^{39}$ erg s$^{-1}$ separates two opacity regimes of the unperturbed disk: opacity is dominated by free–free absorption below and electron scattering above. The upper dashed line at $L_{\nu,\text{E}} \sim 5 \times 10^{41}$ erg s$^{-1}$ separates two pressure regimes of the unperturbed disk: pressure is dominated by gas below and radiation above. The thick gray chevron is the contour for $M_\nu/M_{\nu,\text{E}} = 1$, and the two thin ones above it are for $M_\nu/M_{\nu,\text{E}} = 10$ and $M_\nu/M_{\nu,\text{E}} = 100$, respectively. Our approximation that the stream has greater inertia than the disk holds for $M_\nu/M_{\nu,\text{E}}$ above a few tenths (Chan et al. 2019). The horizontal dotted line marks the value of $L_{\nu,\text{E}}$ below which the radiatively efficient disk may become an RIAF.
The Eddington-level luminosity plateau discussed here is radiatively efficient. The level of the initial plateau is given by Equation (9), and the plateau duration by the bottom panel of Figure 3. That panel also shows, with crosses, $(M_\text{in}/M_\odot, L_*/L_\odot)$ of the depicted light curves. The plateau duration depends weakly on $M_\text{in}$ and $L_*/L_\odot$, and the dependence changes sign when disk pressure transitions from gas- to radiation-dominated (Appendix). The longest plateau duration for radiatively efficient disks across all $M_\text{in}$ and $L_*/L_\odot$ is $\sim 23$ days (Figure 3); the light curves here come close to that limit.

\[ \frac{M_\text{in}}{M_\odot} = (\eta/\eta')^c (L_*/L_\odot) = 10^{-2} \]

in view of the uncertain $M_\text{in}/M_\odot$, marking the transition between the two disk models. The rise of $Q_0$ follows from the fact that, in a radiatively efficient disk, $\Sigma_0$ reaches its maximum when disk pressure switches from gas- to radiation-dominated. Because $Q_0 \ll \dot{L}_\text{E}$ generally in a radiatively efficient disk, there is enough dissipation to power Eddington-level luminosity. By contrast, $v_\text{R}$ is much larger and $\Sigma_0$ much smaller in an RIAF with the same $M_\text{in}/M_\odot$, so $Q_0 \ll \dot{L}_\text{E}$. A noticeable flare is unlikely, so we drop the RIAF from further consideration.

The middle panel of Figure 3 plots the initial value of $t_\text{fall}/t_\text{cool}$ for the radiatively efficient disk, with $t_\text{fall}$ and $t_\text{cool}$ from Equations (5) and (7), respectively. For most of the parameter space plotted, $t_\text{fall}/t_\text{cool} \lesssim 1$, so $L_c \sim \frac{1}{2} \dot{L}_\text{E}$ irrespective of $M_\text{in}$ and $L_*/L_\odot$. With the gradual depletion of the inner disk, $t_\text{fall}/t_\text{cool} \propto 1/\Sigma$ rises but $Q \propto \Sigma t_\text{fall}$ falls, and the two changes offset each other to sustain a constant $L_c$.

The luminosity plateau continues until $t_\text{fall}/t_\text{cool} \sim 1$. The plateau duration is calculated by applying Equation (8), the condition for a plateau. To connect $\tau$ in that equation to $\Sigma = \tau/\kappa T$ in the simulations, we approximate $\Sigma(t)$ in the top right panel of Figure 2 by

\[ \Sigma/\Sigma_0 \sim 0.1(t/t_\text{dec})^{\gamma} \quad \text{for } t \gtrsim t_\text{dec}, \]

with $t_\text{dec} = 10$ days as the start of the approximately power-law decay and $\gamma = -2.5$ as the power-law index. If $t_\text{fall}/t_\text{cool} \gtrsim 0.1$ initially, this equation allows us to determine only an upper limit of $\lesssim 10$ days to the plateau duration. The plateau duration is shown in the bottom panel of Figure 3, and explicit expressions can be found in the Appendix. The plateau typically lasts tens of days. After the plateau, the inner disk fades as $L_c \sim Q \propto \Sigma$. For a small part of the parameter space, $t_\text{fall}/t_\text{cool} \gtrsim 1$ even at the beginning; these inner disks do not have a luminosity plateau.

The plateau duration is comparable to but typically smaller than $t_\text{ab}$. Although the stream mass current decreases during the plateau, the disk mass current would decrease by an even greater amount, according to Equation (11). Therefore, the stream would become heavier and heavier relative to the disk as the event progresses, so the assumption that energy dissipation in the disk dominates that in the stream remains valid.

For a radiatively efficient unperturbed disk, Figure 3 shows that the maximum plateau duration of $\sim 23$ days is found at $M_\text{in} \sim 4 \times 10^6 M_\odot$ and $L_*/L_\odot \sim 10^{-3}$. The disk giving rise to the longest plateau is marginally radiatively efficient, and its pressure at the stream impact radius is on the cusp between gas and radiation dominance; such a disk has the greatest $\Sigma_0$ and thus the longest $t_\text{cool}$. The numbers here are for the fiducial case of a Sun-like star, but the argument holds for other main-sequence stars as well. As detailed in the Appendix and illustrated in Figure 6, the maximum plateau duration and the corresponding $M_\text{in}$ increase with $M_\text{in}$ as $\propto (M_\text{in}/M_\odot)^{-1/12}$ and $\propto (M_\text{in}/M_\odot)^{7/11}$, respectively.

Figure 5 displays the estimated bolometric light curves for four radiatively efficient unperturbed disks. For all four, the luminosity plateau lasts $\sim 20$ days to within a factor of $\sim 1.5$. After the plateau, $L_c \sim Q \propto \Sigma \sim \tau^\gamma$. The power-law index of the falloff is steeper than $-5/3$, the power-law index of the mass return rate (Phinney 1989), and the two are completely unrelated.

5. Discussion

The Eddington-level luminosity plateau discussed here is qualitatively different from the plateaus in other TDE models. The classical picture (e.g., Rees 1988) describes a TDE around an isolated black hole. The debris, returning to pericenter at super-Eddington rates for up to several years (e.g., Evans & Kochanek 1989), quickly forms a circular disk. The disk bolometric luminosity tracks the mass return rate, but instead of going super-Eddington, it reaches a plateau (e.g., Loeb & Ulmer 1997; Krolik & Piran 2012). It should be noted that the classical picture has been called into question by simulations (Shiokawa et al. 2015) and their implications for observations (Piran et al. 2015; Krolik et al. 2020). In particular, the predicted plateau is not seen in...
the light curves of most optical TDEs (van Velzen et al. 2020), and its existence is under debate.

The stream–disk interaction we discussed applies to a TDE around a black hole with a pre-existing accretion disk. In contrast with the classical picture, the plateau is produced when the debris slams into the disk, and the principal parameter determining whether and for how long our plateau can be observed is the accretion rate of the unperturbed disk. Moreover, the plateau is shorter than in the classical picture, lasting only tens of days. One respect in which our mechanism does resemble the classical picture, however, is that the plateaus in both cases arise from radiation trapping, that is, a situation in which inflow is faster than radiation can escape by diffusion.

Saxton et al. (2019) reported a TDE candidate in a quiescent galaxy with a \( \sim 10^{7} M_\odot \) black hole. The 0.2-to-2 keV flux held steady for \( \sim 90 \) days at a \( \sim 10^{-3} \) Eddington level before diminishing as a power law. If the galaxy were to harbor a weakly accreting black hole with no recognizable AGN signatures, then the disruption of a \( \sim 4 M_\odot \) main-sequence star could produce a plateau with the observed duration. The dimness of the plateau may be explained by most of the energy being radiated at other frequencies, or by dust extinction.

Our simulations did not account for magnetic fields. Magnetic-field loops in the outer disk can protrude into the depleted inner disk (Noble et al. 2011), enhancing magnetic stresses to the degree that the outer disk may replenish the inner disk on a timescale as fast as a few days. Any resupply from the outer disk complements resupply from the stream, keeping \( \Sigma \) at a high level and prolonging the plateau.

Our calculations do not capture the gradual increase of the stream mass current to its peak value in realistic TDEs. During the early parts of this rising phase when the stream is lighter than the disk, our predictions here do not apply, but shocks excited by the light stream can still deflect orbiting disk material toward the black hole. Therefore, at later times when the stream is heavier than the disk and our predictions do apply, the disk will have a surface density smaller than its unperturbed value, suggesting that any luminosity plateau that may appear would be shorter. However, if the net mass loss from the inner disk during the rising phase is small due to resupply from the stream or the outer disk, the plateau might be extended. Exactly how a time-varying stream mass current affects the light curve is the subject of future simulations.

Even though predictions about the spectrum of the inner disk are of great astrophysical interest, we refrain from making such predictions because the spectrum depends sensitively on the structure of the inner disk and its cooling mechanisms, but we have little knowledge of either. The comparable magnitudes of \( t_{\text{infl}} \) and \( t_{\text{cool}} \) mean that gas and radiation energy densities in the inner disk vary on similar timescales, and the vigorous energy dissipation in the inner disk implies that radiation pressure may be instrumental in shaping the inner disk. A careful treatment of the inner disk therefore requires computationally expensive radiative MHD simulations that can evolve gas and radiation in unison. Radiative transfer must be performed in three dimensions on account of the lack of symmetry: shocks are localized and heating is uneven; the disk structure shown in Figure 1 is highly irregular, both along the midplane and in the vertical direction; and the stream arches over and partly occludes the disk. It is also unclear to what degree free–free and Compton processes bring the disk to thermal equilibrium. In light of all these considerations, we content ourselves with an estimate of the light curve, which already has a very distinctive shape.

6. Conclusions

The black hole of an AGN is surrounded by an accretion disk. The debris stream of a TDE around such a black hole runs into the disk near pericenter. The stream delivers a much stronger mass current to the impact point than the disk does, so the dynamical evolution of the former is largely unaffected by the latter. On the other hand, the heavy stream prevents the disk from rotating freely. Shocks are formed where the disk crashes into the stream; disk gas is deflected inward and drives shocks against gas closer in. As a result, the kinetic energy of the disk interior to the impact point is dissipated into internal energy at a rate high enough to power a bright flare.

If the unperturbed disk is radiatively efficient, that is, if its luminosity is \( \gtrsim 10^{-3} \) Eddington, then the disk surface density \( \Sigma \) is large enough to make the initial energy dissipation rate \( \dot{Q} \propto \Sigma \) super-Eddington. Conversely, if the unperturbed disk is an RIAF, \( \Sigma \) is too small and \( \dot{Q} \) too low to beget any visible flare.

Because gas in the inner disk falls into the black hole in only a few orbits, not all the energy dissipated issues forth as radiation. We estimate the bolometric luminosity as \( L_\text{bol} \sim Q \min(1, t_{\text{infl}}/t_{\text{cool}}) \), where \( t_{\text{infl}} \) is the time for a gas packet to flow inward from the impact point to the black hole, and \( t_{\text{cool}} \) is the cooling time of the inner disk. The luminosity at early times is regulated to an Eddington-level plateau by the facts that \( Q \propto \Sigma/t_{\text{infl}} \) and \( t_{\text{cool}} \propto \Sigma \).

The luminosity plateau ends when \( t_{\text{infl}}/t_{\text{cool}} \) falls to \( \sim 1 \), which occurs after the disk has been depleted by shock-driven inflows. Resupply from the stream and the outer disk could keep both \( \Sigma \) and \( t_{\text{cool}} \) large, delaying the end of the plateau. If only stream resupply acts, as in our simulations, then the plateau duration is largely determined by the initial \( \Sigma \), which is greatest for marginally radiatively efficient disks in which gas and radiation contribute comparably to the pressure. The maximum plateau duration is \( \sim 23 \) days \( \times (M_\odot/M_\ast) \), produced by a TDE around a \( \sim 4 \times 10^9 M_\odot \times (M_\ast/M_\odot)^{1/3} \) black hole.

Following the end of the plateau, \( L_\text{bol} \propto \Sigma \), which in our simulations declines as a power law in time with index between \( -3 \) and \( -2 \) at earlier times, and possibly a steeper index at later times.

Lastly, the simulations of Chan et al. (2019) show that part of the stream drills through the disk and fans out on the other side into a dilute plume, a fraction of which crashes back and collides inelastically with the disk. The resulting luminosity enhancement, which could also be a fraction of Eddington at early times, will be characterized in future work.

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Software: Athena++, (Stone et al. 2008), Python, NumPy, Matplotlib.

Appendix
Plateau Duration

We present explicit expressions, applicable to all stars, for the plateau duration when the unperturbed disk is radiatively
...we generalize $g_i$.

Our $M_i = -g_i$ approximation that the stream has greater inertia than the disk holds for constant sequence stars with masses as indicated in the top right corner of each panel. The plateau duration is Thomson scattering opacity at the stream impact radius, the plateau duration is $t_{\text{dec}}$.

If the unperturbed disk is dominated by gas pressure and Thomson scattering opacity, the plateau duration is

$$t_{\text{dec}} = 10 \text{ days} \times \left( \frac{M_*}{M_\odot} \right)^{-1/2} \left( \frac{r_*}{R_{\odot}} \right)^{3/2} \left( \frac{r_{\text{p}}}{R_{\odot}} \right)^{-1/2}.$$

(A1)

If the unperturbed disk is dominated instead by radiation pressure and Thomson scattering opacity, the plateau duration is

$$\sim 18 \text{ days} \times \left( \frac{M_*}{M_\odot} \right)^{-1/2+1/(3\gamma)} \left( \frac{r_*}{R_{\odot}} \right)^{3/2-1/\gamma} \times \left[ \frac{M_h}{10^3 M_\odot} \right]^{1/(1+1/3\gamma)} \left( \frac{L_{aE}}{10^{-3}} \right)^{1/(1+1/3\gamma)} \left( \frac{\alpha}{0.1} \right)^{1/(1+1/3\gamma)} \left( \frac{\eta}{0.1} \right)^{1/(1+1/3\gamma)}.$$

(A3)

Because Equation (11) is valid only for $t > t_{\text{dec}}$, a plateau duration shorter than that should be interpreted as an upper limit at $t_{\text{dec}}$.

Equations (A2) and (A3) are plotted in the bottom panel of Figure 3 for the fiducial case of a Sun-like star, and in Figure 6 for other main-sequence stars obeying $r_* \sim M_*$. The maximum plateau duration is

$$\sim 23 \text{ days} \times \left( \frac{M_*}{M_\odot} \right)^{-1/2+1/(2\gamma)} \left( \frac{r_*}{R_{\odot}} \right)^{3/2-1/(8\gamma)} \times \left[ \frac{M_h}{10^3 M_\odot} \right]^{1/16} \left( \frac{r_{\text{p}}}{R_{\odot}} \right)^{21/16}.$$

(A4)

which is attained at

$$M_h \sim 4 \times 10^6 M_\odot \times \left( \frac{M_*}{M_\odot} \right)^{-7/16} \left( \frac{r_*}{R_{\odot}} \right)^{21/16} \left( \frac{r_{\text{p}}}{R_{\odot}} \right)^{21/16}.$$

(A5)

ORCID iDs

Chi-Ho Chan https://orcid.org/0000-0001-5949-6109
Tsvi Piran https://orcid.org/0000-0002-7964-5420
Julian H. Krolik https://orcid.org/0000-0002-2995-7717

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Figure 6. Same as the bottom panel of Figure 3, but for different main-sequence stars with masses as indicated in the top right corner of each panel. The plateau is longer for more massive stars. The gray chevrons are contours of constant $M_i/M_\odot$; they increase by factors of 10 toward the top left, and the thick one is for $M_i/M_\odot = 1$. The contours slide upward with rising $M_*$. Our approximation that the stream has greater inertia than the disk holds for $M_i/M_\odot$ above a few tenths (Chan et al. 2019).