BOOSTED TIDAL DISRUPTION BY MASSIVE BLACK HOLE BINARIES DURING GALAXY MERGERS FROM THE VIEW OF N-BODY SIMULATION

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ABSTRACT

Supermassive black hole binaries (SMBHBs) are products of the hierarchical galaxy formation model. There are many close connections between a central SMBH and its host galaxy because the former plays very important roles on galaxy formation and evolution. For this reason, the evolution of SMBHBs in merging galaxies is a fundamental challenge. Since there are many discussions about SMBHB evolution in a gas-rich environment, we focus on the quiescent galaxy, using tidal disruption (TD) as a diagnostic tool. Our study is based on a series of numerical, large particle number, direct N-body simulations for dry major mergers. According to the simulation results, the evolution can be divided into three phases. In phase I, the TD rate for two well-separated SMBBHs in a merging system is similar to that for a single SMBH in an isolated galaxy. After two SMBHBs approach close enough to form a bound binary in phase II, the disruption rate can be enhanced by ~2 orders of magnitude within a short time. This “boosted” disruption stage finishes after the SMBHB evolves to a compact binary system in phase III, corresponding to a reduction in disruption rate back to a level of a few times higher than in phase I. We also discuss how to correctly extrapolate our N-body simulation results to reality, and the implications of our results to observations.

Key words: galaxies: evolution – galaxies: interactions – galaxies: kinematics and dynamics – galaxies: nuclei – methods: numerical

1. INTRODUCTION

In the Λ cold dark matter (ΛCDM) cosmology, supermassive black hole (SMBH) binaries (SMBHBs) are the direct descendants of hierarchical galaxy mergers (Begelman et al. 1980; Volonteri et al. 2003). If two merging galaxies are gas-rich with comparable masses, a large fraction of gas may be driven into the galactic center, prompting a burst of star formation and feeding the central SMBH (Barnes & Hernquist 1991; Mihos & Hernquist 1994; Hopkins et al. 2005; Springel et al. 2005a). As a result, an active galactic nucleus (AGN) forms in the center of the galaxy and the feedback provoked by the accreting SMBH may regulate the mass growth, both of the central black hole (BH) and the galaxy (Springel et al. 2005b; Sijacki et al. 2007; Di Matteo et al. 2008; Booth & Schaye 2009), which might be the physical mechanisms forming the observed tight correlations of the central SMBH and the host galaxy (Magorrian et al. 1998; Ferrarese & Merritt 2000; Gebhardt et al. 2000; Tremaine et al. 2002; Kormendy & Ho 2013). As a result of the progress on cosmological simulations in recent years, more details about BH growth and SMBH–host galaxy co-evolution have been revealed (Hirschmann et al. 2012; Choi et al. 2014; Kannan et al. 2015; Sijacki et al. 2015). Recent investigations suggest that minor or dry mergers without AGN activity play important roles in the formation and evolution of both late- and massive early-type galaxies (Naab et al. 2009; van Dokkum et al. 2010). Quiescent SMBHBs may form in all kinds of galaxy mergers.

To become gravitationally bound and finally coalesce, two SMBHs in merging galaxies have to lose effectively their large angular momentum and decrease their separation. It has been shown that the dynamical evolution of two SMBHs can be divided into several stages (Begelman et al. 1980). At the beginning of the merger, two SMBHs with galactic cores inspiral toward each other because of galactic dynamical friction (Chandrasekhar 1943). As the two SMBHBs become closer, increasing numbers of ambient stars around them are stripped away by the background potential. When the separation of the two SMBHs is approximately their influence radius, they become gravitationally bound and a SMBHB forms. During its subsequent evolution, ever more stars in the nuclear cores are scattered off the system through the three-body slingshot effect and dynamical friction becomes increasingly less efficient in hardening the SMBHB. When most of the stars bound to the SMBHBs are ejected away, the dynamical friction becomes inefficient. Then the evolution of the SMBHB is dominated by the slingshot effect scattering off stars, and depends on the replenishment of the scattered stars (Saslaw et al. 1974; Mikkola & Valtonen 1992; Quinlan 1996). If the replenishment of the scattered stars is dominated by a two-body relaxation process in a system of spherically symmetric distribution, the slingshot effect will be inefficient and the SMBHB may stall at the parsec scale for a timescale longer than the Hubble time, which is the so-called “final parsec problem” in SMBHB dynamical evolution (Begelman et al. 1980; Milosavljević & Merritt 2001; Berczik et al. 2005). However, both analytic calculations and N-body simulations in the past decade have suggested that the “final parsec problem” can be overcome in realistic systems of galaxy mergers due either to gas dynamics or to stellar dynamics, apart from spherical two-body relaxation (Gould & Rix 2000; Yu 2002; Merritt & Poon 2004; Berczik et al. 2006; Khan...
et al. 2011, 2013; Preto et al. 2011; Colpi 2014, and references therein). If an SMBHB can be driven to a separation of milliparsec scale, strong gravitational wave (GW) radiation would efficiently remove its orbital energy and angular momentum, and drive it to coalesce within the Hubble time (Peters 1964; Begelman et al. 1980). N-body numerical simulation results recently indicated that the coalescence timescales for this stage can be less than 1 Gyr, or even much shorter for SMBHBs with orbits with relatively high eccentricity (Khan et al. 2012; Sobolenko et al. 2015). This also has been confirmed under more general conditions, such as non-spherical galaxies and the collisionless limit using a novel Monte Carlo method instead of N-body simulation (Vasiliev et al. 2015), and in scattering experiments (Sesana & Khan 2015). Due to the anisotropy of GW emissions during coalescence, the remnant of the post-coalesced SMBHs can be kicked with a recoil velocity from a few hundreds of kilometers per second, as is usual, up to several thousands of kilometers per second in extreme cases (Peres 1962; Baker et al. 2006; Campanelli et al. 2007; Kopitz et al. 2007; Lousto & Zlochower 2011). The strong GW radiation from SMBHBs is the main target of the planned Laser Interferometer Space Antenna and the ongoing GW detection program Pulsar Timing Array (Seoane et al. 2013; Verbiest et al. 2016).

Some observational evidence for SMBHBs in gas-rich environments has been reported in the literature, ranging from resolved binary AGNs (Komossa et al. 2003; Hudson et al. 2006; Rodriguez et al. 2006; Fabbiano et al. 2011; Deane et al. 2014) and candidates for unresolved binary systems in AGNs with double-peaked broad lines (Tsamantza et al. 2011; Ju et al. 2013; Shen et al. 2013; Liu et al. 2014b), helical radio jets (Begelman et al. 1980; Conway & Wrobel 1995; Liu & Chen 2007; Roland et al. 2008; Kun et al. 2014), quasi-periodic variabilities (Sillanpaa et al. 1988; Liu et al. 1995, 2006; Liu & Wu 2002; Valtonen et al. 2011; Graham et al. 2015, and references therein), and X-shaped radio structures (Liu 2004), to candidates for the coalescence remnants of SMBHs in AGNs of interruption and recurrence of jet formation in double–double radio galaxies (Liu et al. 2003), and of characteristic signatures of recoiling SMBHs (Komossa et al. 2008; Civano et al. 2012; Liu et al. 2012; Lena et al. 2014). However, due to their native “black” property, SMBHBs in gas-poor environments are extremely hard to detect. A quiescent SMBH can be temporarily illuminated and investigated by tidally disrupting stars passing by (Hills 1975; Rees 1988; Evans & Kochanek 1989; Guillochon & Ramirez-Ruiz 2013), causing giant flares from γ-ray to radio bands (Gezari 2013; Komossa 2015, and references therein). It was suggested first theoretically and then confirmed by X-ray observations that a quiescent SMBHB in a galaxy can also be probed with observations of tidal disruption events (TDEs). Based on theoretical investigations, Liu et al. (2009) showed that when a TDE occurs in an SMBHB system, the gravitational perturbation of the companion SMBH would cause characteristic drops in the light curve of the TDE. A characteristic drop was observed recently in the X-ray light curve of the TDE candidate SDSS J120136.02+300305.5, precisely consistent with the prediction of the binary model (Liu et al. 2014a).

In the past decade, about 30–40 TDEs have been detected. According to analytical and numerical estimations, the TDE rate of single SMBH systems in isolated spherical galaxies is \(10^{-6} - 10^{-4} \, M_\odot \text{yr}^{-1}\) per galaxy, or a few times higher in flattened systems (Magorrian & Tremaine 1999; Syer & Ulmer 1999; Wang & Merritt 2004; Brockamp et al. 2011; Vasiliev & Merritt 2013; van Velzen & Farrar 2014; Zhong et al. 2014; Stone & Metzger 2016). This can be impacted if we take into account tidal disruption (TD) for giant stars (MacLeod et al. 2012). The TDE rate of gravitationally recoiling single SMBHs is several orders of magnitude smaller (Komossa & Merritt 2008; Stone & Loeb 2011, 2012; Li et al. 2012; O’Leary & Loeb 2012). The TDE rate of SMBHBs in galaxy mergers was investigated analytically and/or with scattering experiments. Chen et al. (2008) showed that the TDE rate of hard SMBHB systems is orders of magnitude lower than that in single SMBH systems if the SMBHB is in a spherical isotropic galactic core and two-body relaxation dominates. However, for a real SMBHB system, two SMBHs are bound by nuclear star clusters. Both analytic calculations and numerical scattering experiments indicate that the perturbations of the companion SMBH in a self-gravity bound SMBHB system would scatter a large fraction of the bound stars toward the central SMBH and significantly enhance the TDE rate by several orders of magnitude, up to one event per galaxy per year (Ivanov et al. 2005; Chen et al. 2009, 2011; Wegg & Nate Bode 2011). However, in their calculations, the contributions of the stars bound to the companion SMBH and the unbound/weakly bound stars due to two-body relaxation and/or triaxial gravitational potential were not taken into account. The analytic investigations showed that the TDE rate of stars by SMBHs in the early phases of galaxy mergers when galactic dynamical friction is dominant could also be enhanced by several orders of magnitude up to \(10^{-2}\) events per year per galaxy due to the perturbation of the companion galactic core and the triaxial distribution of the gravitational potential of the galactic nucleus (Liu & Chen 2013).

Owing to the limitations of analytic calculations and scattering experiments, investigations of TDE rates in the literature have to be undertaken by splitting the evolution of SMBHBs in galaxy mergers into several stages based on the dominant physical processes. Other important issues, which are not always properly taken into account, are two-body relaxation between the stars around the SMBBH, the contribution of stars to the gravitational potential, and the evolution of the SMBBH with regard to its orbital parameters. This paper aims to resolve these issues; we investigate the dynamical evolution of SMBHBs in major galaxy mergers with direct N-body simulations and self-consistently calculate the TDE rates of SMBHBs from the early phase of galaxy mergers to the post-merger. A practical difficulty is that we are not yet able to accurately simulate a full galaxy with bulge and nucleus star-to-star due to computational limitations. An interesting method is to use expansion techniques on the Poisson equation to speed up the integration, which was recently proposed by Meiron et al. (2014). However, before being able to use that method in our model, several modifications have yet to be completed. As a result, we use the method of scaling from smaller particle number \(N\) and larger tidal radius to realistic values. Our results show that the TDE rates of SMBHBs in major galaxy mergers can be enhanced by up to two orders of magnitude, which is consistent with the preferential high detection rates of TDEs in E+A galaxies (Arcavi et al. 2014). This is very interesting because E+A galaxies are believed to be post-merger galaxies (Zabludoff et al. 1996; Goto 2005; Stone & Metzger 2016).
This paper is organized as follows. In Section 2, we briefly introduce the loss cone theory for galaxy mergers. In Section 3, we describe our N-body simulation method and the galaxy model. The numerical results for the TDE rate and its variations as a function of the dynamical evolution of SMBHBs in galaxy mergers are given in Section 4. We discuss and extrapolate our simulation results to a few typical realistic systems of galaxy mergers in Section 5, and provide a short discussion about the implications to observations in Section 6. Section 7 gives a short summary of our results.

2. TD IN A MERGING SYSTEM

Due to the strong tidal force from a BH, a star will be tidally disrupted if it can penetrate into the vicinity of a BH within $r < r_t$ (Hills 1975; Rees 1988)

$$ r_t \approx \mu r_h (M_{BH}/m_*)^{1/3}. \tag{1} $$

Here $\mu$ is a dimensionless parameter of order unity. $r_h$, $m_*$ and $M_{BH}$ are stellar radius, stellar mass, and BH mass, respectively. Since $r_t \propto M_{BH}^{1/3}$, while the Schwarzschild radius is $r_{sch} \propto M_{BH}$, a star will be directly swallowed instead of being tidally disrupted by a Schwarzschild SMBH with mass $M_{BH} \gtrsim 10^8 M_\odot$.

In a two-body system with only a single BH and a star, if the star has specific angular momentum smaller than a critical TD angular momentum $J_{k_c}$, which satisfies $J < J_{k_c} \approx (2GMBH/\mu)^{1/2}$, it will be tidally disrupted by the BH within one orbital period. Such stars always have their velocity vectors inside a so-called “loss cone” with opening angle (Lightman & Shapiro 1977; Merritt 2013, and references therein)

$$ \theta_{k_c} = \frac{J_{k_c}}{J_c}, \tag{2} $$

where $J_c$ is the specific circular angular momentum with same energy as the star. In a spherical isotropic stellar distribution, for stars inside the influence radius $r_h$ of the SMBH, $r < r_h$,

$$ \theta_{k_c}^2 \approx \frac{2r_t^3}{3r}, \tag{3} $$

while for stars at $r > r_h$,

$$ \theta_{k_c}^2 \approx \frac{2rhr_t}{3r^2}. \tag{4} $$

(Frank & Rees 1976; Baumgardt et al. 2004).

Due to the interaction with the environment, the angular momentum of a star changes with time. If $\theta_{D}(r)$ is the orbital averaged deflection angle of velocity and $\theta_D(r) \ll \theta_{k_c}(r)$, all of the stars inside the loss cone are in the diffusive regime and would be tidally disrupted within one orbital period (Frank & Rees 1976; Shapiro & Lightman 1976; Lightman & Shapiro 1977). The loss cone becomes empty because its replenishment through diffusion takes much longer than the stellar orbit period. In contrast, for a pinhole regime with $\theta_D(r) \gg \theta_{k_c}(r)$, the loss cone remains full because the star can be deflected in and out within one orbital period. Thus there is a critical radius $r_{crit}$, which gives $\theta_{D}(r_{crit}) \sim \theta_{k_c}(r_{crit})$. Detailed theoretical calculations and numerical simulations show that the TD rate of stars peaks at about $r_{crit}$ (Frank & Rees 1976; Lightman & Shapiro 1977; Amaro-Seoane et al. 2004; Zhong et al. 2014).

In a spherical system, the disruption rate contributed by the stars between $r$ and $r + dr$ is

$$ d\Gamma = \frac{4\pi r^2 dr \rho(r) \theta^2(r)}{t_\nu(r)} \tag{5} $$

(Frank & Rees 1976; Syer & Ulmer 1999; Liu & Chen 2013), where $t_\nu(r)$ is the dynamical timescale of the star, $\rho$ is stellar mass density, and $\theta$ is a dimensionless coefficient for stars depleting into the loss cone. Detailed analysis (Young 1977) suggests that

$$ \theta^2 = \min(\theta_{k_c}^2, \theta_{D}^2/\ln \theta_{k_c}^{-1}). \tag{6} $$

To investigate the TD rate in a merging system, we adopt the same scheme as Liu & Chen (2013) did, which writes the effective deflection angle as

$$ \theta_{k_c}^2 = \theta_2^2 + \theta_p^2 + \theta_c^2, \tag{7} $$

where $\theta_2$, $\theta_p$, and $\theta_c$ are, respectively, the contribution from two-body relaxation, the massive perturber and the triaxial gravitational potential. Based on our N-body simulation results presented in Section 4, we can analyze which mechanism dominates the TDE rate in different stages, and make analytical estimations. As we will present in Section 5, $\theta_2$ and $\theta_c$ can be well estimated by Equations (16) and (29), respectively, while $\theta_p$ is too complicated to be estimated. Detailed discussions can be found in Section 5.

The total disruption rate can be obtained by integrating Equation (5). That means we need to know $\rho$, $t_\nu$, and $\theta$, which are time-dependent in merging galaxies and evolving SMBHBs. For this reason, here we address these questions using direct N-body simulations.

3. DIRECT N-BODY SIMULATIONS OF THE TDE RATES OF SMBHS IN MERGING GALAXIES

The dynamical evolution of two SMBHs in merging galaxies has been investigated intensively using direct N-body simulations (Chatterjee et al. 2003; Berczik et al. 2006; Khan et al. 2011, 2013; Preto et al. 2011; Gualandris & Merritt 2012; Spurzem et al. 2012; Vasiliev et al. 2014; Wang et al. 2014), but none of them has taken into account the TD of stars of SMBHBs. For this reason, here we address these questions using direct N-body simulations.
number $N$ and the adopted value for the TD radius $r_\text{t}$ relative to our simulation units. Note that in our case we keep the ratio of the SMBH mass to the total mass of the system constant, motivated by observed correlations, so the number of stars $N$ is a proxy for the ratio between a single stellar mass and the SMBH mass. The goal is to provide several models in this two-dimensional framework of $N$ and $r_\text{t}$, and then extrapolate for realistic values of both parameters.

Several problems need to be addressed when using this procedure. It is quite common in physics, but for our self-gravitating systems with huge dynamical range in spatial and temporal scales it needs special attention. First, it is important that for every particle number $N$ the range of $r_\text{t}$ needs to be adjusted such that we get an appropriate loss cone: empty near the SMBH, going to a full loss cone near the critical radius in the vicinity of the gravitational influence radius of the SMBH (see e.g., Frank & Rees 1976 for the basis of this theory). If we ran a simulation with small $N$ and realistic (i.e., extremely small) $r_\text{t}$, all loss cones would be extremely small and full due to the short relaxation time, which is not the physical case. Therefore, for our range of particle numbers, only certain values of $r_\text{t}$ are physically relevant (see below) that can make the critical radius large enough to have the empty-to-full loss cone transition at the right location. Second, the growth of the BH per dynamical time becomes increasingly smaller for larger $N$ due to the increasing relaxation time. So, one may argue that our (relatively) small-$N$ simulations contain some artificial mass growth of the SMBH, which in turn affects the potential and all the dynamics in the central region. The alternative, to neglect the SMBH mass growth artificially, however, creates more severe difficulties; it leads to spurious mass relocations (e.g., the SMBH kicks out stars by slingshot which would otherwise be tidally accreted and disrupted, or if the stars are taken away the central gravitational potential is reduced artificially). We think that the best way to use the method of scaling is to include all physical processes (such as mass growth of the SMBH) for every choice of parameter, and then carefully analyze the scaling behavior with $N$ and $r_\text{t}$ (including also e.g., the mass growth in the scaling analysis). We will come back to this issue in Section 5.

In conclusion, we need to carry out a series of simulations with different $r_\text{t}$ and $N$ and then extrapolate the results to real systems. A similar strategy has been adopted in the literatures (Bauengardt et al. 2004; Brockamp et al. 2011; Li et al. 2012; Zhong et al. 2014, and references therein).

In this paper we are focusing on a major merging system with two equal mass nuclei, which has a spherical stellar distribution with equal mass stars. For simplicity we choose a spherical Dehnen model to represent the nucleus for each galaxy (Dehnen 1993). Within this model, the space density profile can be written as

$$\rho(r) = \frac{3 - \gamma}{4\pi} \frac{M a}{r^\gamma (r + a)^{4-\gamma}},$$

where $a$ is the scaling radius, $M$ is the total mass of each nucleus, and the index $\gamma$ is a constant in the interval $[0, 3)$. Since there are two galaxies with equal mass, the total mass for the merging system is $M_{\text{tot}} = 2M$, and we can derive the cumulative mass for each nucleus by

$$M(r) = M \left( \frac{r}{r + a} \right)^{3-\gamma}. \tag{9}$$

Thus the influence radius $r_h$ which has been defined as $M_r (r \leq r_h) = 2M_{\text{BH}}$ is

$$r_h = \frac{a}{\left( \frac{M}{2M_{\text{BH}}} \right)^{\frac{\gamma}{3-\gamma}} - 1}, \tag{10}$$

and since $M \gg 2M_{\text{BH}}, r_h \approx a \left( 2M_{\text{BH}} / M \right)^{\frac{\gamma}{3-\gamma}}$.

For simplicity, we adopt model units with $G = M = a = 1$ hereafter, which is the same as in Li et al. (2012). Thus we have the same relations between our simulation units and physical quantities:

$$[T] = \left( \frac{GM}{a} \right)^{-1/2}$$

$$= 1.491 \times 10^6 \left( 2^{\frac{\gamma}{3-\gamma}} - 1 \right)^{3/2} \left( \frac{M}{10^{11}M_\odot} \right)^{-1/2} \times \left( \frac{r_{1/2}}{1 \text{ kpc}} \right)^{3/2} \text{ year}, \tag{11}$$

$$[V] = \left( \frac{GM}{a} \right)^{1/2}$$

$$= 655.8 \times \left( 2^{\frac{\gamma}{3-\gamma}} - 1 \right)^{-1/2} \left( \frac{M}{10^{11}M_\odot} \right)^{1/2} \times \left( \frac{r_{1/2}}{1 \text{ kpc}} \right)^{-1/2} \text{ km s}^{-1}, \tag{12}$$

$$[R] = a \left( 2^{\frac{\gamma}{3-\gamma}} - 1 \right) \left( \frac{r_{1/2}}{1 \text{ kpc}} \right) \text{ kpc}, \tag{13}$$

$$[M] = M/[T]$$

$$= 6.707 \times 10^4 \left( 2^{\frac{\gamma}{3-\gamma}} - 1 \right)^{-3/2} \left( \frac{M}{10^{11}M_\odot} \right)^{3/2} \times \left( \frac{r_{1/2}}{1 \text{ kpc}} \right)^{-3/2} \text{ M}_\odot \text{ yr}^{-1}. \tag{14}$$

Our fiducial galaxy model has been chosen as $\gamma = 1.0$, $M = 4 \times 10^{10} \text{ M}_\odot$, $M_{\text{BH}} = 4 \times 10^7 \text{ M}_\odot$, and half mass–radius $r_{1/2} = 1 \text{ kpc}$. Thus we can derive $[T] \approx 0.63 \text{ Myr}$, $[V] \approx 644 \text{ km s}^{-1}$ and $[R] \approx 0.4 \text{ kpc}$. For a solar-type star, the tidal radius is $\sim 10^{-5} \text{ pc}$ corresponding to $\sim 2 \times 10^{-8}$ in these simulation units. Obviously, comparing to $r_t \sim 10^{-4}$ we have adopted in our $N$-body simulations, this value is too small to have enough TDE records for a statistical study in integration.

Before the integrations for the merging system, each nucleus with a central SMBH has been dynamically relaxed through the same method used in Li et al. (2012). Similar to the procedure adopted by Preto et al. (2011), here we initially set two nuclei with distance $d \sim 20$, and a parabolic orbit with pericenter $\sim 1$. To accurately integrate the orbits of the particles, we use a parallel direct $N$-body $\varphi$-GRAPE/$\varphi$-GPU code with fourth-order Hermite integrator, which is the same code adopted by Li et al. (2012). For all of our integrations, the softening length has been set to $\epsilon = 10^{-5}$, and the integrations have been terminated at $t = 200$. All of the integrations are calculated on the laohu GPU cluster in the National Astronomical Observatories of China (NAOC).
| Model No. | $e_{ini}$ | $N/N_C$ | $r_t$ | $\gamma$ | $M_{BH}/M$ |
|-----------|----------|---------|-------|---------|------------|
| A01       | 1.0      | 125 K   | $5 \times 10^{-4}$ | 1.0 | 0.001 |
| A02       | 1.0      | 250 K   | $5 \times 10^{-4}$ | 1.0 | 0.001 |
| A03       | 1.0      | 500 K   | $5 \times 10^{-5}$ | 1.0 | 0.001 |
| A04       | 1.0      | 500 K   | $1 \times 10^{-4}$ | 1.0 | 0.001 |
| A05       | 1.0      | 500 K   | $5 \times 10^{-4}$ | 1.0 | 0.001 |
| A06       | 1.0      | 500 K   | $1 \times 10^{-3}$ | 1.0 | 0.001 |
| A07       | 1.0      | 500 K   | $5 \times 10^{-4}$ | 0.5  | 0.001 |
| A08       | 1.0      | 500 K   | $5 \times 10^{-4}$ | 1.5  | 0.001 |
| A09       | 1.0      | 500 K   | $5 \times 10^{-4}$ | 1.0  | 0.0001 |
| A10       | 1.0      | 500 K   | $5 \times 10^{-4}$ | 1.0  | 0.001 |
| A11       | 1.0      | 1 M     | $5 \times 10^{-4}$ | 1.0  | 0.001 |
| B01       | -isolated galaxy | 500 K | $5 \times 10^{-5}$ | 1.0 | 0.001 |
| B02       | -isolated galaxy | 500 K | $1 \times 10^{-4}$ | 1.0 | 0.001 |
| B03       | -isolated galaxy | 500 K | $5 \times 10^{-4}$ | 1.0 | 0.001 |
| B04       | -isolated galaxy | 500 K | $1 \times 10^{-3}$ | 1.0 | 0.001 |
| C01       | 0.3      | 500K    | $5 \times 10^{-4}$ | 1.0  | 0.001 |

**Note.** Column (1) Model sequence number. Column (2) Eccentricity of two nuclei at $t = 0$. Column (3) Particle numbers adopted in calculations for each nucleus. Column (4) TD radius $r_t$. Column (5) Density slope $\gamma$. Column (6) Initial BH mass. All the integrations keep $e = 10^{-5}$. For reference, we have four additional integrations in models B01–B04, with a single nucleus and SMBH. Model C01 has moderate initial eccentricity.

In order to investigate the dependency of our results on particle number $N$, numerical tidal radius $r_t$, density slope $\gamma$, and initial mass of the SMBH, we have run a series of simulations with the model parameters listed in Table 1. The first column is the sequence number of different simulations. Columns 2–6 are, respectively, initial orbital eccentricity $e_{ini}$ of the two galactic nuclei, particle number $N$ for each nucleus, numerical tidal radius $r_t$, initial stellar density slope $\gamma$, and initial mass of the SMBH. For comparison of the TDE rates, we have run four additional simulations for an isolated single nucleus in models B01–B04. In model C01, we have a tested a simulation for a merging galaxy with moderate initial orbital eccentricity $e_{ini} = 0.3$. All of the simulations except model C01 are terminated at $t = 200$, when the separation of the two SMBHs $r_B$ has $r_B \approx 400 r_h$. The simulation of model C01 is terminated at $t = 500$ because the dynamical evolution of the SMBHB is very slow for a moderate initial eccentric orbit. In the simulation, model A05 has been considered as our fiducial numerical model.

4. RESULTS

4.1. Evolution of Tidal Disruption

Figure 1 gives the evolution of the SMBHB orbit and corresponding TD accretion rates $\dot{M}$, in units of accreted mass per $N$-body time. For both panels, the red solid line and blue dashed line represent $r_{BH}$ and $\dot{M}$ respectively. The left panel corresponds to our fiducial model A05 with initial parabolic orbit, and the right panel is model C01 with initial orbital eccentricity $e = 0.3$ for comparison. Model C01 has been integrated to $t = 500$ because of the slower dynamical evolution of the two SMBHs. Here we only pick up the last 200 time units for easier comparison. According to Figure 1, the dynamical evolution of the two SMBHs in moderately eccentric orbits is slower than in parabolic orbits. Our integration result also shows that for A05 the eccentricity at the end of integration remains constant with $e \sim 0.7$, while in C01 it remains around 0.05. Actually, almost all of our integrations with initial parabolic orbits result in high eccentricities for the SMBHBs.

Figure 1 indicates that there is a significantly enhanced TD accretion rate as the two BHs approach close enough to form a bound binary. Thus we can empirically divide the entire evolution into three different phases. In phase I, $r_{BH}$, the separation of the two BHs is far larger than either’s influence radius, and $\dot{M}$ remains at a relatively constant low level. After the two BHs get close enough to significantly perturb each other, $\dot{M}$ increases significantly and the system evolves into phase II. In phase III, a compact SMBHB has formed and $\dot{M}$ drops down to a roughly constant lower level.

As shown in Figure 1, both models A05 and C01 have similar evolutions of TD accretion rates: $\dot{M}$ is roughly stable during phase I and III, and is significantly enhanced in phase II. This result can be easily understood. In phase I, the evolution of $\dot{M}$ is more or less the same as the case with a single SMBH in an isolated galaxy, because the perturbation from another nucleus is not significant. That corresponds to a diffusive regime for most large galactic nuclei. However phase II is a stage with violent relaxation. There may be two factors contributing to the enhancement of the TD accretion rate. First of all, stars bound to one SMBH suffer very strong perturbation from the companion SMBH, which will enhance the loss cone feeding. Second, due to triaxial stellar distribution and three-body interaction in phase II, more stars can be scattered into the TD loss cone. As a result, the loss cone feeding during phase II is very efficient, which leads to a significantly enhanced TD accretion rate.

Liu & Chen (2013) have found that when the separation of two SMBHs is larger than $100 r_h$, the loss cone filling will be dominated by two-body relaxation and the TD accretion rate is identical to that of an isolated single SMBH, while as their separation shrinks to $d \sim 100 r_h$, the TD accretion rate induced by perturbation from another galaxy will be comparable to the contribution from two-body relaxation. We find, however, for our models A05 and C01, that TD accretion rates are enhanced by perturbations only at around $d \lesssim 10 r_h$ or below. To understand this difference, we ran their code with the same parameters (galaxy mass, density profile, and mass ratio of two galaxies) of our A05 model, but choosing a realistic tidal radius and particle number. The result based on their model shows that the contribution from perturbation becomes higher than two-body relaxation around $d \lesssim 400 r_h$. But if we rescale in their model the TD accretion rate obtained from two-body relaxation to the artificially small particle number and large tidal radius as used in our $N$-body simulation (using the scaling relations found in Sections 5.1 and 5.2) the effect of perturbation becomes important at $d \lesssim 8 r_h$, consistent with our results. We will come back to this issue in Section 5.

It is hard to distinguish which factor dominates the loss cone feeding in our simulation. Since we want to focus on the TD rate evolution in this paper, the discussion about the contributions from triaxial stellar distribution and perturbation will be postponed to a future paper. In phase III, the SMBHB has formed and evolved to a hard binary. Most of the stars bound to each BH have been ejected or disrupted by three-body interaction. The TD rate is dominated by stars outside the
SMBHB orbit. Because of a triaxial stellar distribution around SMBHB, the TDE rate should be higher than a relaxed spherical distribution system like a single galaxy. However, according to the results in the left panel of Figure 1, the separation of the two SMBHs is close to $r_\text{fl}$ by the end of the integration, which will not happen so early in real systems with much smaller $r_\text{fl}$. As a result, almost all of stars encountering the SMBHB in this stage will be disrupted. This will artificially suppress slingshot effects and enhance the TDE rate. For this reason, our results in the late phase III tend to overestimate the TDE accretion rates.

4.2. Dependence on Particle Number

As mentioned in Section 3, we cannot carry out a direct $N$-body simulation with the same particle number as in a real galaxy. For a stellar system with an embedded central SMBH in a real galaxy, the total number of stars can be easily as high as $10^9$, while direct $N$-body simulations, even if using very powerful high performance computers, can only currently manage $10^5$–$10^7$ particles. As a result, compared to real galaxies, the relaxation timescale in our model should be shorter. Besides, if we set the initial mass of a BH particle to be $10^{-3} M$ and assume every star particle has equal mass $M/N$, the mass ratio of the BH to the star in our simulation will be much smaller than reality. For this reason, it is necessary to investigate the dependence of our simulation results on particle number. The data are given in Table 2, with model index A01, A02, A05, and A11.

The left panel of Figure 2 shows the $N$ dependence for the evolution of TD accretion rate $M$, in units of accreted mass per $N$-body time. Since the two SMBHs have roughly identical growth, all of our calculations for SMBHB in this stage will be disrupted. This will artificially suppress slingshot effects and enhance the TDE rate. For this reason, our results in the late phase III tend to overestimate the TD accretion rates.

![Figure 1](image_url) Evolution of the separation for two SMBH and TD accretion rates. Here the $x$-axis and the left $y$-axis give evolved time and $r_{\text{BH}}$, the separation of SMBHs, respectively. The right $y$-axis gives the corresponding TD accretion rates $M$. The red solid line and blue dashed line represent $r_{\text{BH}}$ and $M$ respectively. Left panel: our fiducial model A05 with initial parabolic orbit. Right panel: model C01 with initial orbital eccentricity $e = 0.3$ for comparison. In order to compare to A05, here we only pick up the evolution for the last 200 time units. The $M$ in model C01 roughly remains constant in the first 300 time units.

| No. (1) | PI (2) | $M/\text{Galaxy}$ (3) | PII (4) | $M/\text{Galaxy}$ (5) | PIII (6) | $M/\text{Galaxy}$ (7) |
|--------|--------|------------------------|--------|------------------------|--------|------------------------|
| A01    | 0–83   | $3.74 \times 10^{-5}$   | 83–109 | $5.31 \times 10^{-5}$   | 109–200| $3.27 \times 10^{-5}$   |
| A02    | 0–84   | $2.75 \times 10^{-5}$   | 84–102 | $5.42 \times 10^{-5}$   | 102–200| $2.91 \times 10^{-5}$   |
| A03    | 0–87   | $4.31 \times 10^{-6}$   | 87–101 | $6.14 \times 10^{-6}$   | 101–200| $3.82 \times 10^{-6}$   |
| A04    | 0–86   | $7.12 \times 10^{-6}$   | 86–107 | $9.62 \times 10^{-6}$   | 107–200| $6.26 \times 10^{-6}$   |
| A05    | 0–83   | $2.14 \times 10^{-5}$   | 83–104 | $4.95 \times 10^{-5}$   | 104–200| $2.84 \times 10^{-5}$   |
| A06    | 0–80   | $3.38 \times 10^{-5}$   | 80–102 | $9.80 \times 10^{-5}$   | 102–200| $4.75 \times 10^{-5}$   |
| A07    | 0–108  | $6.69 \times 10^{-6}$   | 108–135| $1.37 \times 10^{-5}$   | 135–200| $1.00 \times 10^{-5}$   |
| A08    | 0–65   | $8.94 \times 10^{-5}$   | 65–82  | $1.90 \times 10^{-4}$   | 82–200 | $5.15 \times 10^{-5}$   |
| A09    | 0–92   | $9.53 \times 10^{-6}$   | 92–143 | $2.81 \times 10^{-5}$   | 143–200| $2.56 \times 10^{-5}$   |
| A10    | 0–73   | $4.05 \times 10^{-5}$   | 73–90  | $1.05 \times 10^{-4}$   | 90–200 | $2.44 \times 10^{-5}$   |
| A11    | 0–82   | $1.59 \times 10^{-5}$   | 82–102 | $5.00 \times 10^{-3}$   | 102–200| $2.62 \times 10^{-3}$   |

Note. Column (1) Model serial number in Table 1. Column (2) Period of PI. Column (3) TD accretion rate during PI. Column (4) Period of PII. Column (5) TD accretion rate during PII. Column (6) Period of PIII. Column (7) TD accretion rate during PIII.
red asterisks in the right panel, the averaged TD accretion rate in phase I declines with increasing particle number, and the rate in phase III with blue circles also shows a slight decrease. Phase II, represented by green squares, conversely maintains roughly a constant rate. This may because the loss cone feeding in phase I is dominated by two-body relaxation, which is $N$-dependent, while the loss cone feeding during phase II is dominated by the perturbative effect of the two SMBHs with their surrounding stars, the non-stationary potential, and the triaxial stellar distribution, none of which depends on $N$. Similar to phase II, both perturbation and the triaxial gravitational potential also contribute to the TD accretion rate during phase III, especially at the beginning. However, as the semimajor axis of the SMBHB is decreasing, more stars close to it will be ejected or disrupted. Thus the contribution from perturbation will diminish. At the same time, the contribution of the triaxial potential is also decreasing because the central region of the merger remnant is recovering to a spherical distribution. As a result, an ideal evolution of TD accretion rate in phase III will finally be dominated by two-body relaxation, which suggests a dependence on $N$ sooner or later. However, that may not happen in our simulations. Considering the large tidal radius we have adopted, the semimajor axis of the SMBHB will be smaller than $r_t$, at some time, which may give rise to artificial effects. For this reason, our integrations terminate at $t = 200$, when the average separation of the SMBHs is close to $r_t$. In brief, the evolution of the TD accretion rate in phase III is very complex, and we can only trace the evolution in its early stage.

### 4.3. Dependence on Tidal Radius

In order to explore the dependence of our results on $r_t$, we have tried series integrations with different $r_t$. All the data are given in Table 2 with model index A03, A04, A05, and A06. Figure 3 shows the influence of different $r_t$ on our results. In principle, larger $r_t$ correspond to higher $M$. However, as shown in the left panel, when $r_t$ is smaller than $5 \times 10^{-4}$, there is no significant $M$ peak in phase II. According to our analytical estimation in Section 5 and simulation results shown in the right panel of Figure 3, there is a roughly linear dependence of $M$ on $r_t$. Meanwhile, this dependence in phase II is a power law with index less than one. For this reason, with $r_t$ decreasing, $M$ in phase II drops faster than in phase I. For small enough $r_t$, the contribution of phase II will be comparable to phase I, and the peak can be washed out. This is an artificial effect due to limited particle resolution in our simulations. In reality, due to much larger $N$, $M$ contributed by two-body relaxation in phase I should be smaller than we have in the simulations. However, this effect is not crucial for our results. According to Table 2, all the averaged TD accretion rates in phase II have notable enhancement even if $r_t = 5 \times 10^{-5}$, which means the contribution from strong perturbation in phase II is still significant. A smaller $r_t$ may lead to equal accretion rates for phase I and II, but that does not mean our strategy is not suitable. For a proper scaling limit it is not sufficient just to keep one parameter fixed (e.g., $N$) and vary the other ($r_t$) to the realistically small value. In fact for every $N$ there is only a limited reasonable range of $r_t$. If $r_t$ is too small we will get too few events, while if it is too large we will get unphysically many events. The extrapolation can only be done simultaneously in $N$ and $r_t$, by carefully examining the scaling in both parameters, as we have done in Section 5. In phase III, it seems that $M$ always stays at a constant level which is slightly higher than in phase I.

The right panel shows averaged TD accretion rates in different phases for different $r_t$. The correlation of TD accretion rate and $r_t$ in phase II can be well fitted by a linear correlation, implying that the loss cone is fed efficiently. In the reference $N$-body simulations for an isolated galaxy with a single SMBH, the TD accretion rates represented by the black plus line is similar to the results of phase I represented by the red asterisk line. This is consistent with the expectation that the TDE rate evolution in phase I is the same as for a single BH in an isolated spherical galaxy. Therefore, the TDE loss cone in
phase I is empty and the loss cone feeding is dominated by two-body relaxation. Another interesting thing is that the dependence of averaged TD accretion rate on \( r_t \) in phase III, represented by blue circles, lies beside the results of phase I and II. It actually nearly overlaps the pink triangles, the averaged result for all three phases. The dependence of averaged TD accretion rate on \( r_t \) in phase III is neither a linear relation as in phase II, nor a relation close to the single BH case, which indicates that the loss cone feeding in phase III is more efficient than in phase I, but not as full as in phase II.

As shown in the left panel, the TD accretion rates for sufficiently large \( r_t \) in phase I are slightly lower than those that in phase III while significantly less than those in phase II. In other words, we have low TDE rates in phase I and slightly higher rates in phase III with relatively long period, and very high rates in phase II with very short period. Considering of the short period of phase II, the contribution from phase II to the averaged TD accretion rates for all three phases should not be dominant. For example, in model A05, the number of disrupted particles in phase II is only 1040, compared to a total of 5547 in the whole simulation. The contribution of phase I, II, and III to the total TD rate is \( \sim 32\% \), \( \sim 19\% \), and \( \sim 49\% \), respectively. According to the estimation in Section 6, the contribution of phase II is more than a quarter for our fiducial model. As a result, the rate in phase III is closer to the averaged result for all of three phases.

4.4. Dependence on Density Profile and Initial BH Mass

As discussed before, the numerical calculation capability limits the particle resolution, which cannot provide a realistic high-mass ratio of SMBHs to stars in a direct \( N \)-body simulation. Here we consider heavier SMBHs in the simulations to improve the situation. In addition to the initial BH mass, it is also important to investigate the dependence of TDE rates on the initial stellar density profile. Therefore we have carried out three integrations with \( M_{\text{BH}} = 0.0001, 0.001, 0.01 \) and another three with \( \gamma = 0.5, 1.0, 1.5 \). Here we adopt the fiducial values for other model parameters. We do not have any tests for \( \gamma > 1.5 \) because all those integrations are prohibitively time consuming.

The results are given in Figures 4 and 5. The left panels of Figures 4 and 5 show that the enhancement of TD accretion rates and the starting time of phase II strongly depend on density distribution and initial BH mass. The larger \( M_{\text{BH}} \) and \( \gamma \) are, the more significant the enhancement of the TD accretion rates, and the earlier the staring time. The duration of phase II also depends both on initial SMBH mass and density profile index \( \gamma \). A system with more massive SMBHs or larger \( \gamma \) usually corresponds to a shorter duration of phase II. More details can be found in Table 2. It shows that model A07, A05, and A08 have different \( \gamma \), and A09, A05, and A10 have different initial \( M_{\text{BH}} \). The influence of initial \( M_{\text{BH}} \) and \( \gamma \) can be understood as follows. In principle, a more massive SMBH corresponds to a larger influence radius \( r_h \) and gives a stronger perturbation to its companion. Thus, SMBHs with larger initial \( M_{\text{BH}} \) form a bound system at larger separation and the formed SMBHB also evolves faster. For a system with larger \( \gamma \), the stellar distribution is more concentrated toward the center and the SMBH, and its TDE rates are enhanced much more significantly, leading to faster growth of the SMBH masses. In particular, at the transient time between phase I and II, the differences in BH mass between models A05 (\( \gamma = 1.0 \)) and A07 (\( \gamma = 0.5 \)) and between A08 (\( \gamma = 1.5 \)) and A05 are, respectively, \( \sim 1.6 \) and \( \sim 2.5 \) times. Heavier SMBH mass gives a stronger perturbation to the stars around the companion SMBH, leading to faster ejection and more frequent TDEs, and thus shorter duration of phase II. But this effect may not be significant in reality because the mass growth of the SMBHs in our \( N \)-body simulation is artificially enlarged. It is interesting to notice that for the heaviest initial SMBH mass \( M_{\text{BH}} = 0.01 \) and the largest density profile index \( \gamma = 1.5 \), the TD accretion rate evolution in phase I has a slight decrease until the start of phase II. This may be due to a problem with the set-up of the initial density distribution. Because the systems for models A10 and
A08 have much longer relaxation timescales in the bound regions and at the influence radius of SMBH, our limited computer resources prohibit us from making the systems well relaxed before the integrations start.

The right panels of Figures 4 and 5 give the averaged TD accretion rates at different phases as a function of initial $M_{\text{BH}}$ and $\gamma$, respectively. The averaged TD accretion rates in phase I and II significantly correlate with the initial BH mass and stellar distribution of the galactic nucleus. In phase II, the TD accretion rate for the model with $\gamma = 1.5$ (model A08) is about 14 times higher than that with $\gamma = 0.5$ (model A07). It is expected that the TD accretions for $\gamma > 1.5$ are even higher. In addition, the rate in phase II strongly depends on the initial mass of the SMBH, and the rate for $M_{\text{BH}} = 0.0001$ (model A09) is about four times that for $M_{\text{BH}} = 0.0001$ (model A09). In phase III, the averaged TD accretion rate weakly depends on $\gamma$ and is nearly independent of $M_{\text{BH}}$.

5. EXTRAPOLATION TO REAL GALAXIES

To apply our simulation results in real systems of galaxy mergers, we have to extrapolate. To obtain TDE rates, one needs to integrate Equation (5). Since the TDE rate is dominated by the contribution of stars around the critical radius $r_{\text{crit}}$, we can estimate it as

$$\Gamma \sim \frac{r^3 \rho^2 d\rho}{t_d m_\star} \bigg|_{r = r_{\text{crit}}}$$  (15)

(Frank & Rees 1976; Baumgardt et al. 2004), where for simplicity all the stars are assumed to have equal mass ($m_\star = M_\star$). Now the key point is to estimate $r_{\text{crit}}$. As we will show in the following sections, $r_{\text{crit}}$ changes with the dynamical evolution of SMBHs during galaxy mergers. For the order of magnitude, we estimate the averaged TDE rates with typical critical radius $r_{\text{crit}}$ for each phase.

5.1. Extrapolation for Phase I

In phase I, the separation of two SMBHs is much larger than the influence radius of the SMBH (typically $r_{\text{BH}} \gg r_{\text{h}}$). As was shown by Liu & Chen (2013), the perturbation and torque of the companion galactic core is small in this stage. The typical TDE rate can be estimated as to that of an isolated galaxy with a single central SMBH. For a spherical isotropic system, two-body relaxation dominates the diffusion process and $\theta_d \sim \theta_2$ in Equation (7) with

$$\theta_2 \approx \frac{t_d}{t_r}$$  (16)

(Frank & Rees 1976), where the dynamical time of the star is $t_d \sim \sqrt{r^3/(GM_{\text{BH}})}$ within the influence radius and the two-body relaxation timescale $t_r$ is

$$t_r \approx \frac{0.34 \sigma^3}{G^2 m_\star \ln \Lambda}$$  (17)

(Binney & Tremaine 2008). Here $\ln \Lambda$ is the Coulomb logarithm and $\sigma$ is the one-dimensional velocity dispersion of the stars. For $r \lesssim r_{\text{h}}$, we have $\sigma \approx \sqrt{GM_{\text{BH}}/r}$. The Coulomb logarithm is approximately

$$\ln \Lambda \approx \ln \left(\frac{M_{\text{BH}}}{2m_\star} \right)$$  (18)

(Preto et al. 2004).

From Equations (3), (16), and (17) and the assumption $\rho(r) = \rho_0 (r/r_0)^{-\gamma}$, we have the typical critical radius of Phase I

$$r_{\text{crit}} \approx \left(0.23M_{\text{BH}} \frac{r_{\text{h}}}{m_\star \rho_0 r_0^2 \ln \Lambda} \right)^\frac{1}{\gamma - 1}$$

$$\propto M_{\text{BH}}^2 \left(\frac{N}{\ln \Lambda}\right)^\frac{1}{\gamma - 1} r_{\text{h}}^\frac{1}{\gamma - 1}.$$  (19)
Thus, Equation (15) gives the mass accretion rate due to TD

\[ \dot{M} = \Gamma m_{\bullet} \sim \frac{r^3 \theta_c^3 \rho(r)}{t_d} \bigg|_{r=r_{\text{crit}}} \]  

(20)

and

\[
\dot{M} \approx 2.94 \times 0.23^{2.6-0.23} G^2 \left( \rho_0 \rho_0^2 \right)^{\gamma/2} M_{\text{BH}}^{6.54} \left( \frac{N}{\ln \Lambda} \right)^{2-\gamma} \frac{1}{r_{\text{h}}^{0.54}} \]

\[ \propto M_{\text{BH}}^{6.54} \left( \frac{N}{\ln \Lambda} \right)^{2-\gamma} \frac{1}{r_{\text{h}}^{0.54}}. \]  

(21)

We assumed that \( \rho_0 \rho_0^2 \) is independent of \( M_{\text{BH}}, N, \) and \( r_1 \). This is different from the assumption had been used by Frank & Rees (1976), which assumed that \( \rho = \rho_0 \) and \( r_0 = r_1 \), with \( \rho_0 \) representing the density at influence radius. \( r_1 \) depends on \( M_{\text{BH}} \) by its definition \( r_1 \simeq G M_{\text{BH}} / \sigma_{\text{v}}^2 \). Here \( \sigma_{\text{v}}^2 \) denotes one-dimensional stellar velocity dispersion, which they considered to be constant outside \( r_1 \). As a result, they argued that the TD accretion rate should sensitively depend on BH mass. But \( \sigma_{\text{v}} \) is a strong function of radius in cuspy galaxy models such as the Dehnen model we adopted here (Tremaine et al. 1994; Merritt et al. 2007). Therefore, we simply consider \( \rho \rho_0^2 \) as an undetermined coefficient, and estimate the TD accretion rate in reality by fitting and extrapolating our simulation results. Our results based on direct \( N \)-body simulation should be more reliable. For this reason, we choose Equation (21) as our extrapolation solution.

As mentioned in Section 3, our model has artificially enhanced BH accretion, which may be not accurate enough to estimate the TD accretion rate in phase I. Considering the analysis above and our simulation results, we believe this effect is not significant. Our results show that, by the end of phase I, \( M_{\text{BH}} \) and \( r_1 \) have been increased up to \( \sim 2.6 \) and \( \sim 1.6 \) times, respectively. For most of our models, at the beginning, the relaxation timescale around \( \sim r_1 \) for each galaxy is longer than the duration of phase I. For simplicity, we can assume that this relation could persist during the entire phase I. That means the mass growth of the SMBHs will only impact the very central region within phase I, which cannot efficiently change the TD accretion rate. Combined with a non-growing tidal radius adopted in the simulation, the artificial mass growth of the SMBH does not have enough time to have a serious effect during phase I. In order to check this speculation, we have calculated the stellar dispersion and density evolution of model B03, and two additional similar integrations with \( N = 250 \) K and \( N = 125 \) K. The comparisons between \( t = 0 \) and \( t = 80 \) show that there is no significant change for these two parameters around \( r_1 \). Our analysis of \( \dot{M} - M_{\text{BH}} \) dependence also indicates that, along with the growth of the SMBH, \( M \) remains roughly constant in phase I. For this reason, we believe that the mass growth of the SMBH cannot significantly change the TDE rate in phase I. However, it does not mean the influence of the initial \( M_{\text{BH}} \) is negligible, because the initial stellar distribution sensitively depends on the initial \( M_{\text{BH}} \).

For fixed \( M_{\text{BH}} \) and \( N \), we have the relation

\[ \dot{M} \propto \frac{1}{r_1^{0.54}}. \]  

(22)

Although Equations (21) and (22) are valid for \( r_{\text{crit}} \approx r_1 \), it can be expected that the power-law relationship between mass accretion rate and \( r_1 \) or \( N \) should be a good approximation even for \( r_{\text{crit}} \approx r_1 \). Figure 6 shows the simulated TD accretion rates of phase I as functions of \( r_1 \) (upper left panel) and particle number \( N \) (bottom right panel), and the fits with power-laws as suggested by Equations (21) and (22). The TD accretion rates, obtained by direct \( N \)-body simulations, as functions of \( r_1 \) and \( N \) are well fitted by power-laws \( \dot{M} \approx 3.92 \times 10^{-3} r_1^{-0.6867} \) and \( \dot{M} \approx 6.76 \times 10^{-3} (N / \ln \Lambda)^{-0.3051} \), respectively. The accretion rate is in our simulation units \([M]\). With these fitting results, as well as Equations (21), we can derive \( \gamma \approx 1.33 \) and \( \gamma \approx 1.67 \), respectively. We consider these two results as consistent, given the approximations used here.
Actually, according to Equation (21), with fixed $M_{\text{BH}}$, we can derive
\[ \ln M \approx \ln A - \frac{2\gamma - 1}{8 - 2\gamma} \ln \left( \frac{N}{\ln \Lambda} \right) + \frac{9 - 4\gamma}{8 - 2\gamma} \ln \tau, \tag{23} \]
where
\[ A = 2.94 \times 0.23^{\frac{9-4\gamma}{8-2\gamma}} G^2 (\rho_0 r_0^3) \frac{1}{r_{1/2}} M_{\text{BH}}^{6-4\gamma}. \tag{24} \]
Based on Equation (23) we can make a two-dimensional fitting, which gives results consistent with those we obtained above
\[ M \sim 1.26 \times \left( \frac{N}{\ln \Lambda} \right)^{-0.506} \tau^{0.686}. \tag{25} \]

With Equation (14), we can finally derive the TD accretion rate for a real galaxy
\[ \dot{M} \sim 0.387 \times \left( \frac{M}{10^{11} M_\odot} \right)^{3/2} \left( \frac{r_{1/2}}{1 \text{ kpc}} \right)^{-2.186} \times \left( \frac{N}{\ln \Lambda} \right)^{-0.506} \left( \frac{\tau}{10^{-6} \text{ pc}} \right)^{0.686} M_\odot \text{ yr}^{-1}. \tag{26} \]

Now we can easily extrapolate our simulation results of phase I to real galaxies. For our fiducial galaxy model, the averaged TD accretion rate in phase I should be $\sim 7 \times 10^{-6} M_\odot \text{ yr}^{-1}$. For comparison, we also have calculated the averaged TDE rates for galaxies similar to M32 or M105. For an M32-like galaxy, according to observations, we adopt $r_{1/2} \sim 0.1 \text{ kpc}$ and $M_{\text{BH}} \sim 3 \times 10^6 M_\odot$ (Wang & Merritt 2004; Cappellari et al. 2006). Here we assume that the half-mass radius of the galaxy is roughly equal to the effective radius. In order to maintain consistency with our simulation model, we set $\gamma = 1.0$. As a result, our extrapolation shows that the averaged TDE rate in phase I is $\sim 4 \times 10^{-3} M_\odot \text{ yr}^{-1}$. For a more massive galaxy, similar to M105, with $r_{1/2} \sim 2 \text{ kpc}$, $M_{\text{BH}} \sim 10^8 M_\odot$, $M \sim 10^{11} M_\odot$ and $\gamma = 1.0$, the averaged TDE rate in phase I is $5 \times 10^{-6} M_\odot \text{ yr}^{-1}$.

Since the averaged TDE rates in phase I of the merger should be similar to a normal galaxy with a single SMBH, it will be interesting to compare our results with the estimations for normal galaxies. For a merging system composed of two galaxies similar to M32, the averaged TDE rate in phase I is an order of magnitude lower than that argued by Wang & Merritt (2004). That may be because they used a density slope index of $\gamma = 2$, which is larger than we have for $\gamma = 1$. As Figure 5 and Equations (21) and (22) show, TDE rates depend on $\gamma$. By fitting the dependence of averaged TD accretion rates on $\gamma$, which is given in the right panel of Figure 5, we get a power-

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**Figure 6.** Fitting results for $\tau$, $N$, and $r$ dependence. Here the red solid squares represent the integration results in our simulations, and the red solid lines represent the fitting results according to formulae (21), (22), and (31), or a power-law relation. The upper left, upper right, and bottom left panels show the $\tau$ fitting result for phase I, II, and III, individually. The bottom right panel gives the $N$ fitting result for phase I.

The Astrophysical Journal, 834:195 (15pp), 2017 January 10

Li et al.
law index \(\sim 2.29\). After simply scaling to \(\gamma = 2\), we have \(M \simeq 2 \times 10^{-4} M_\odot \, \text{yr}^{-1}\) for M32. Considering that there are differences for the results obtained through different methods, this result is consistent with their estimation (Alexander 2012).

5.2. Extrapolation for Phase II

In phase II, the separation of the two SMBHs evolves from \(\sim 20 r_h \) to \(\sim 3 \times 10^{-2} r_h\), where \(r_h \approx 0.1\) for an initial stellar distribution with \(\gamma = 1.0\). In a real galaxy merger, that corresponds to the evolution of two SMBHs from formation of a bound SMBHB to hard binaries. Because of the strong perturbation by the companion galactic core and of the triaxiality in stellar distribution, we have \(\theta_0^2 \gg \theta_2^2\), \(\theta_2^2 \gg \theta_5^2\) and \(\theta_5^2 \sim \theta_2^2\) in Equation (7).

As was shown by Liu & Chen (2013), the companion SMBH together with the bound star cluster would exert a strong tidal torque to the stars around the main SMBH with distance \(r\). If \(M_p\) is the total mass of the perturber and \(r_{\text{BH}}\) is the distance between the two SMBHs, the tidal torque is approximately

\[
T_p \simeq \frac{GM_p r^2}{r_{\text{BH}}} \tag{27}
\]

for \(r \ll r_{\text{BH}}\). Thus the changes of the angular momentum of a star can be up to a timescale \(t_{\text{c}} = \min(t_{\text{fH}}, t_{\text{pec}})\), with \(t_{\text{fH}}\) and \(t_{\text{pec}}\) respectively, the dynamical timescale of the perturber and the stellar apsidal precession timescale. For \(t > t_{\text{c}}\), we can estimate the average change of the angular momentum of the star on a dynamical timescale \(t_{\text{d}}(r)\) (Liu & Chen 2013).

\[
J_p^2 = T_p^2 t_{\text{d}}(r). \tag{28}
\]

Because \(\theta_p^2 = J_p^2 / L_p^2\), we have

\[
\begin{align*}
\theta_p^2 &= \frac{J_p^2}{(r^2 \sigma^2)} \quad \text{for } r \gtrsim r_h, \\
\theta_p^2 &= \frac{J_p^2}{GM_p r^2} \quad \text{for } r < r_h.
\end{align*} \tag{29}
\]

Where we have assumed a constant stellar velocity dispersion for \(r \gtrsim r_h\) and \(\sigma^2 = GM_{\text{BH}} / r_h\). From Equations (28) and (29) and with the assumption of \(r \gtrsim r_h\), we obtain the critical radius

\[
r_{\text{crit}} \simeq 0.93 G^{-1/11} M_{\text{BH}}^{3/11} M_p^{4/11} t_{\text{d}}^{12/11}/r_h^{1/11}. \tag{30}
\]

Table 2 and the discussions in Section 4.2 imply that the typical value of the separation \(r_{\text{BH}}\) in phase II depends neither on \(N\) nor \(t_d\). Because \(M_p\) depends on \(r_{\text{BH}}\), we have \(r_{\text{crit}} \propto r_{\text{BH}}^{1/11}\), which is independent of \(N\). From Equations (3), (20), and (30), we have

\[
\begin{align*}
\dot{M} &\approx \frac{2}{3} \rho_0 r_0 \left\{ G M_{\text{BH}} \right\}^{1/2} r_{\text{crit}}^{-1/2} \gamma t_{\text{fH}} \\
&\approx \frac{2}{3} \times 0.93^{1-2\gamma} \rho_0 r_0 G^{4/11} \frac{M_{\text{BH}}}{M_p}^{4/11} \frac{t_{\text{d}}^{12/11}}{r_{\text{BH}}^{12/11}} \frac{M_{\text{BH}}^{7/3}}{M_{\text{BH}}^{7/3} \times r_{\text{BH}}^{12/11}} \\
&\propto M_{\text{BH}}^{7-3\gamma} r_{\text{BH}}^{12-2\gamma}.
\end{align*} \tag{31}
\]

Equation (31) implies that the TD accretion rate in phase II is approximately a power-law function of \(M_{\text{BH}}\) and \(r_{\text{BH}}\). The upper right panel of Figure 6 shows that our N-body simulation results for \(\dot{M} \sim r_{\text{BH}}\) in phase II can be well fitted with a power-law function with \(M = 6.55 \times 10^{-2} \left(\frac{r_{\text{BH}}}{10^6 \text{ pc}}\right)^{0.9456}\), which is very consistent with Equation (31) for \(\gamma = 1.0\). We did not subtract the contribution of two-body relaxation in phase II. According to Figure 2 and our analysis above, the TD accretion rate in phase II does not sensitively depend on \(N\), therefore we conclude that the contribution of two-body relaxation here is negligible. Similar to Section 5.1, we can estimate the TD accretion rate for a real galaxy in phase II:

\[
\dot{M} \sim 1.17 \times 10^{-4} \left(\frac{M_{\odot}}{10^{11} M_{\odot}}\right)^{3/2} \left(\frac{r_{\text{BH}}}{1 \text{ kpc}}\right)^{-2.4456} \times \left(\frac{r_{\text{BH}}}{10^6 \text{ pc}}\right)^{0.9456} M_{\odot} \text{ yr}^{-1}. \tag{32}
\]

With this fitted power-law relation, we can extrapolate our simulation results in phase II to a real system of galaxy mergers. For our fiducial model of SMBH mass \(M_{\text{BH}} = 4 \times 10^7 M_{\odot}\), the averaged TD accretion rate in phase II is \(\sim 2 \times 10^{-4} M_{\odot} \text{ yr}^{-1}\) with a peak TD accretion rate \(\sim 5.5 \times 10^{-4} M_{\odot} \text{ yr}^{-1}\). The TD accretion rate for phase II is about 30 times for the averaged and about 80 times at the peak higher than the rate in phase I. For an M32-like galaxy with \(M_{\text{BH}} = 3 \times 10^6 M_{\odot}\), the averaged TD accretion rate for phase II is \(5 \times 10^{-4} M_{\odot} \text{ yr}^{-1}\), and the peak rate is \(\sim 1.4 \times 10^{-3} M_{\odot} \text{ yr}^{-1}\), which is about 13 times for the averaged or 35 times at the peak higher than the rate in phase I. For a more massive galaxy, similar to M105 with \(M_{\text{BH}} \sim 10^8 M_{\odot}\), the TD accretion rate is \(2 \times 10^{-4} M_{\odot} \text{ yr}^{-1}\) with peak \(\sim 5.4 \times 10^{-4} M_{\odot} \text{ yr}^{-1}\).

The TD rates given here are averaged over the range of the SMBH separation from about \(20 r_h \) to \(3 \times 10^{-2} r_h\). Liu & Chen (2013) analytically investigated the TD rates as a function of SMBH separation in galaxy mergers for \(r_{\text{BH}} \sim 10^2 r_h - 2 r_h\) with \(r_h \sim r_{\text{BH}}\), which overlap with the early stage of phase II. Their Figure 4 shows that the TD rates can be enhanced by up to \(\sim 200\) at \(r_{\text{BH}} \sim 2 r_h\) over the rates at \(r_{\text{BH}} \gg 100 r_h\) for \(\gamma = 1.75\), \(M_{\text{BH}} = 10^7 M_{\odot}\) and SMBH mass ratio \(q = 1\). These are, respectively, about seven times and three times larger than the averaged and peak values for our fiducial model. The difference may be because they used a deeper density profile with \(\gamma = 1.75\). As was shown in Section 4.4, a larger \(\gamma\) usually corresponds to a larger TDE rate. Similar to the \(\gamma\) extrapolation for phase I, if we fit the averaged TD accretion rates for different \(\gamma\) in phase II with a power law of index \(\sim 2.34\), the enhancement of the averaged TDE rates for our fiducial values would become \(\sim 100\), which is roughly consistent with those given by Liu & Chen (2013). For an M32-like galaxy with central SMBH mass \(M_{\text{BH}} \sim 3 \times 10^6 M_{\odot}\), the enhancement of the TD rates in phase II over phase I is about \(\sim 66\) for the averaged or \(\sim 180\) for the peak TD rates for \(\gamma = 2.0\). It is also roughly consistent with the results given by Liu & Chen (2013).

Using numerical scattering experiments, Chen et al. (2009, 2011) investigated the TD rates of hard SMBHBs with extremely unequal masses at dynamical evolution corresponding to the late half-stage of phase II. They showed that the TD rate of an SMBHB with mass \(10^7 M_{\odot}\) and mass ratio \(q = 0.01\) can be enhanced by up to three or four orders of magnitude higher than that for a single SMBHB fed by two-body relaxation. This is about one to two orders of magnitude higher than the peak rates obtained from our fiducial model. One of the reasons is that they assumed \(\gamma = 2\) in their calculations. When a \(\gamma = 1.5\) is adopted for the same system, the peak TD rates decrease to about a few times \(10^{-3}\) events per year as
shown in Figure 20 of Chen et al. (2011), which is only a few times higher than the peak TDE rate $\sim 1.4 \times 10^{-3}$ per year estimated by extrapolating the values of our fiducial model to $\gamma = 1.5$ with the fitted power-law extrapolation. Another reason for the differences may be that the enhancement of TDE rates is less pronounced for less-unequal-mass binaries (Chen et al. 2009). Our results are based on $q = 1$, which is much larger than the largest mass ratio $q = 0.1$ investigated by Chen et al. (2011). In any case, our $N$-body simulations indicate that the enhancement effect for equal mass binaries is still significant and can be as large as a factor of a few hundred times for SMBH systems with $M_{\text{BH}} \sim 10^6 M_\odot$ and $\gamma = 2.0$.

5.3. Extrapolation for Phase III

In phase III, SMBHBs are hard binaries with averaged separations $r_{\text{BH}} \sim 3 \times 10^{-3} r_h - 3 \times 10^{-2} r_h$. As was indicated in Section 4, it is difficult analytically to obtain a physically meaningful extrapolation relation for phase III, because the tidal loss cone is neither empty nor full. Since $r_h \sim 0.1$ in phase III, the separation corresponds to $3 \times 10^{-4} - 3 \times 10^{-3}$, which is very close to the tidal radius $r_t = 5 \times 10^{-4}$ in our fiducial model. That means in the late stage of phase III, stars are tidally accreted in our simulation which would be ejected by slingshot, if $r_t$ were much smaller. As mentioned by Chen et al. (2008), if we only consider unbound stars or bound stars with eccentricity close to unity, the TDE rate is proportional to the geometrical cross section between the stars and the SMBHs. This cross section, by approximation, is proportional to $r_t$. On the other hand, the rate of stars ejected through three-body interaction is proportional to the semimajor axis of the SMBHB. Thus when $r_{\text{BH}}$ is close to $r_t$, almost all of the stars interacting with the SMBHB will be tidally disrupted. This effect happens in reality much later than in our simulation, due to the much smaller $r_t$ in real systems. Due to our rather large $r_t$ we may underestimate the rate of slingshot ejections and overestimate the rate of TD. So, one should be aware that in phase III our extrapolation method is more uncertain than in phases I and II.

Here we fit the $N$-body simulation results with a power law and estimate the TDE rates for a real system of galaxy mergers by extrapolating the simulation results with the fitted power-law relation. The relation between TD accretion rates $\dot{M}$ and tidal radius $r_t$ is given in the bottom left panel of Figure 6. It shows that the simulation results for different $r_t$ can be well fitted by a presumed power law of $\dot{M} \propto 1.90 \times 10^{-2} r_t^{0.8634}$, which is not seriously different from the linear relation predicted by simple cross-section analysis. We also fit the numerical results $\dot{M}$ as a function of particle number $N$ with a power law and obtain $\dot{M} \propto 1.03 \times 10^{-4} N^{0.10}$. The weak dependence of TD accretion rates on particle number $N$ implies that the tidal loss cone in phase III is indeed not empty but close to full.

We estimated the TDE rates for a few typical real systems of galaxy mergers by extrapolating the numerical results with fitted power-law relations. For our fiducial model, the separation of the SMBHB in phase III ranges from about 0.1–1 pc and the TD accretion rate is $\dot{M} \sim 7 \times 10^{-5} M_\odot \text{yr}^{-1}$, which is about an order of magnitude higher than that of the two separated single SMBHs fed by two-body relaxation in Phase I. This is inconsistent with the results of hard SMBHBs given in the literature. Chen et al. (2008) showed that because of the three-body slingshot effect, the TDE rate in a hard SMBHB system is more than an order of magnitude smaller than that in a single SMBH system if the loss cone refilling is dominated by two-body interaction. This inconsistency can be understood as follows. First, as discussed above, the extrapolation based on our $N$-body simulations tends to overestimate TDE rate in phase III. Second, when an SMBHB becomes hard, the nuclear stellar system is out of equilibrium and evolving with time. That leads to significantly higher tidal loss cone refilling rates comparing to the predictions based on two-body interaction.

For a galaxy like M32 with SMBH mass $M_{\text{BH}} \sim 3 \times 10^6 M_\odot$, the separation of the SMBHB in phase III ranges from about 0.01 pc to 0.1 pc and the extrapolated TD accretion rate is $\sim 2 \times 10^{-4} M_\odot \text{yr}^{-1}$, which is about five times higher than that in phase I. For a galaxy like M105, the extrapolated TD accretion rate is $6 \times 10^{-5} M_\odot \text{yr}^{-1}$, about 12 times higher than that in phase I. The TD accretion rate in phase III also depends on $\gamma$. It can be well fitted by a power law of index $\sim 1.49$. This gives an estimate of TD accretion rate for an M32-like galaxy with $\gamma = 2.0$, $\sim 6 \times 10^{-4} M_\odot \text{yr}^{-1}$, which is about three times higher than in Phase I. Therefore, the enhancement of TD accretion rates between phases I and III only weakly depends on $\gamma$.

6. IMPLICATIONS FOR OBSERVATIONS

As our results have shown, the TDE rate of a merging galaxy can be boosted considerably. For our fiducial model, it can be boosted up to 80 times, which means the TDE detection in merging galaxies should be easier. For a dry major merger like we discussed here, the typical evolution time is $\sim 1 \text{ Gyr}$ (Colpi 2014). For our extrapolated fiducial model, the period for phase II is $\sim 13 \text{ Myr}$, corresponding to $\sim 2600 \text{ TDEs}$. If we simply assume that the TD rate for the rest of the time is equal to that in phase I, the corresponding TDEs should be $\sim 6900$. That means more than a quarter of the TDEs in a merging galaxy are prompted in phase II, which occupies only $\sim 1\%$ of the entire evolution time. However, this result is only based on major mergers, which are relatively scarce in the evolution history of a galaxy. According to the results by Casteels et al. (2014), the major merger rate for the local universe is only a few percent galaxy$^{-1} \text{Gyr}^{-1}$. Considering the upcoming optical transient survey, the Large Synoptic Survey Telescope6, which might detect several hundreds of TDEs every year (Wegg & Nate Bode 2011), several TDEs corresponding to a merging galaxy in phase II should be found. For minor mergers which are more common in the lifetime of a galaxy (Hopkins et al. 2010), as mentioned above, it is hard to make predictions. In this paper we have only studied major mergers, for which we found a significant boost of their TDE.

Boosted TDE rates in phase II correspond to two close SMBHs with strong perturbation to each other. A detected TDE from one of these two SMBHs in this stage should be observed as an offset flare. But the displacement from the center of the galaxy should be of order 1–10 pc at maximum, which is very hard to detect in distant galactic nuclei. On the other hand, such events should occur with a high velocity relative to the local rest frame in the galaxy (could be several thousand kilometers per second), so they might be detectable through a blue- or redshift of their spectral features. This kind of observational signature is similar to a TDE from a recoiling SMBH, but the latter has a much more smaller event rate (Komossa & Merritt

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6. http://www.lsst.org/lsst
2008; Li et al. 2012; Stone & Loeb 2012). As suggested by Liu & Chen (2013), it is possible to distinguish the TDE from merging galaxies and a recoiling remnant.

7. SUMMARY

In this paper, we investigated the dynamical evolution of SMBHBs in gas-poor systems of galaxy mergers with equal mass. By using direct N-body simulations with resolution up to two million particles on the special many-core hardware "laohu" GPU cluster at NAOC, we self-consistently calculated the TDE rates of stars of the SMBHs and their variations with the dynamical evolution of SMBHBs. The typical separation of SMBHs in galaxy mergers ranges from $r_{BH} \sim 200r_h$ to $\sim 2 \times 10^{-3} r_h$. We found that the evolution of TDE rates of SMBHBs in galaxy mergers can be divided into three stages, which tightly correlate with their dynamical evolution. In phase I, the typical separation of two SMBHs is $r_{BH} \sim 200r_h$ and their dynamical evolution is dominated by dynamical friction. At this stage, the tidal loss cone is empty and its refilling is dominated by two-body relaxation. The TDE rate of each SMBH is about that estimated for a single SMBH feeding by two-body relaxation and strongly depends on the mass of the central SMBH and the stellar density profile index $\gamma$. In phase II, the two SMBHs rapidly evolve from two strongly interacting but unbound nuclear systems to a self-bound and finally hard SMBHB system, with typical separation $r_{BH} \sim 20r_h$ to $\sim 3 \times 10^{-2} r_h$. In phase II, the tidal loss cone is full because of the triaxial stellar distribution and the strong perturbations to the star excited by the companion. The TDE rates of SMBHBs can be enhanced by up to two orders of magnitude relative to the TDE rates in phase I. The enhancement effect depends on the mass of the central SMBH but only weakly on the stellar distribution parameter $\gamma$.

Our conclusions about phases I and II are generally consistent with the results given in the literature (Ivanov et al. 2005; Chen et al. 2009, 2011; Liu & Chen 2013). However, our results for phase III are inconsistent with those for hard SMBHBs by Chen et al. (2008). In phase III, SMBHBs are hard and evolve slowly from a separation about $3 \times 10^{-2} r_h$ to $2 \times 10^{-3} r_h$. During this stage, the tidal loss cone is neither full nor completely empty and the loss cone feeding depends weakly on the particle number. The averaged TDE rate during phase III is about or weakly enhanced by a factor of a few relative to that for phase I. This is in contrast with the conclusions given by Chen et al. (2008), in which the TDE rates for hard SMBHBs are found to be by about an order of magnitude smaller than those estimated for single SMBHS feeding by two-body relaxation in spherical isotropic stellar systems.

We conclude that the TDE rates in galaxy major mergers are strongly enhanced by up to two orders of magnitude relative to the estimation of a single SMBH feeding by two-body relaxation, which is consistent with the high detection rate of TDEs in E+A galaxies (Arcavi et al. 2014), a type of galaxy which is probably a post-merger of galaxies (Zabludoff et al. 1996; Goto 2005; Stone & Metzger 2016). For minor mergers, which are more common compared to major mergers, it is hard to make predictions without detailed numerical simulations. This issue is beyond the scope of this paper.

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REFERENCES
Alexander, T. 2012, in Tidal Disruption Events and AGN Outbursts, Vol. 39, ed. R. Saxton & S. Komossa (Madrid: EPJ Web of Conferences), 05001
Amaro-Seoane, P., Freitag, M., & Spurzem, R. 2004, MNRAS, 352, 655
Arcavi, I., Gal-Yam, A., Sullivan, M., et al. 2014, ApJL, 793, 38
Baker, J. G., Gentrelle, J., Choi, D.-I., et al. 2006, ApJL, 653, L93
Barnes, J. E., & Hernquist, L. E. 1991, ApJL, 370, L65
Baumgardt, H., Makino, J., & Ebisuzaki, T. 2004, ApJ, 613, 1133
Begelman, M. C., Blandford, R. D., & Rees, M. J. 1980, Natur, 287, 307
Berczik, P., Merritt, D., & Spurzem, R. 2005, ApJ, 633, 680
Berczik, P., Merritt, D., Spurzem, R., & Bischof, H.-P. 2006, ApJL, 642, L21
Binney, J., & Tremaine, S. 2008, in Galactic Dynamics, ed. J. Binney & S. Tremaine (2nd ed., Princeton, NJ: Princeton Univ. Press), 587
Booth, C. M., & Schaye, J. 2009, MNRAS, 398, 55
Brockamp, M., Baumgardt, H., & Kroupa, P. 2011, MNRAS, 418, 1308
Campanelli, M., Lousto, C., Zlochower, Y., & Merritt, D. 2007, ApJL, 659, L5
Cappellari, M., Bacon, R., Bureau, M., et al. 2006, MNRAS, 366, 1126
Casteels, K. R. V., Conselice, C. J., Bamford, S. P., et al. 2014, MNRAS, 445, 1157
Chandrasekhar, S. 1943, ApJ, 97, 255
Chatterjee, P., Hernquist, L., & Loeb, A. 2003, ApJ, 592, 32
Chen, X., Liu, F. K., & Magorrian, J. 2008, ApJ, 676, 54
Chen, X., Madau, P., Sesana, A., & Liu, F. K. 2009, ApJL, 697, L149
Chen, X., Sesana, A., Madau, P., & Liu, F. K. 2011, ApJ, 729, 13
Choi, E., Naab, T., Ostriker, J. P., Johansson, P. H., & Moster, B. P. 2014, MNRAS, 442, 440
Civano, F., Elvis, M., Lanzuisi, G., et al. 2012, ApJ, 752, 49
Colpi, M. 2014, SSRv, 183, 189
Conway, J. E., & Wrobel, J. M. 1995, ApJ, 439, 98
Deane, R. P., Paragi, Z., Jarvis, M. J., et al. 2014, Natur, 511, 57
Dehnen, W. 1993, MNRAS, 265, 250
Di Matteo, T., Colberg, J., Springel, V., Hernquist, L., & Sijacki, D. 2008, ApJL, 676, 33
Evans, C. R., & Kochanek, C. S. 1989, ApJL, 346, L13
Fabbiano, G., Wang, J., Elvis, M., & Risaliti, G. 2011, Natur, 477, 431
Ferrarese, L., & Merritt, D. 2000, ApJL, 539, L9
Frank, J., & Rees, M. J. 1976, MNRAS, 176, 633
Gebhardt, K., Bender, R., Bower, G., et al. 2000, ApJL, 539, L13
Gezari, S. 2013, BrJph, 43, 351
Goto, T. 2005, MNRAS, 357, 937
Gould, A., & Rix, H.-W. 2000, ApJL, 532, L29
Graham, M. J., Djorgovski, S. G., Stern, D., et al. 2015, Natur, 518, 74
Gualandris, A., & Merritt, D. 2012, ApJ, 744, 74
Guillochon, J., & Ramirez-Ruiz, E. 2013, ApJL, 767, 25
Hills, J. G. 1975, Natur, 254, 295
