



A Shining Death of Unequal Supermassive Black Hole Binaries

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Abstract

In the Λ CDM scenario, small galaxies merge to produce larger entities. Since supermassive black holes (SMBHs) are found in galaxies of all sizes, SMBH binaries (SMBHBs) are generally expected to form during the amalgamation of galaxies. It is unclear what fraction of these binaries could eventually merge, but a general consensus is that initially the orbital decay is mediated by the surrounding gas and stars. In this Letter, we show that in active galactic nuclei (AGNs) the radiation field also causes the orbits of the accreting SMBHs to shrink. The corresponding mechanism, known as the “Poynting–Robertson drag” (PR drag), takes effect on a well-defined timescale CT_{Sal} , where T_{Sal} is the Salpeter timescale of the AGN, presumably coinciding with the primary SMBH, and $C = 4\xi^{-1}\epsilon^{-1}q^{1/3}(1+q)^{2/3}(1-\epsilon)$ is a constant determined by the radiative efficiency ϵ , the mass ratio q of the two black holes, and a parameter ξ characterizing the size of the circumsecondary accretion disk. We find that when $q \lesssim$ a few $\times 10^{-5}$, the PR drag is more efficient in shrinking the binary than many other mechanisms, such as dynamical friction and type-I migration. Our finding points to a possible new channel for the coalescence of unequal SMBHBs and the clearing of intermediate-massive black holes in AGNs.

Unified Astronomy Thesaurus concepts: [Accretion \(14\)](#); [Active galactic nuclei \(16\)](#); [Gravitational waves \(678\)](#); [Quasars \(1319\)](#); [Astrodynamics \(76\)](#)

1. Introduction

Almost all massive galaxies contain supermassive black holes (SMBHs) in their centers (Kormendy & Ho 2013). The current consensus is that such black holes (BHs) form in small galaxies in the early universe and grow to 10^6 – $10^{10} M_{\odot}$ by episodic accretion of gas (Soltan 1982; Yu & Tremaine 2002). This accretion phase, as is understood today, can be triggered by galaxy mergers (Kauffmann & Haehnelt 2000). During such a phase, a fraction of the gravitational energy of the gas is released in the form of radiation and the galaxy center becomes an active galactic nucleus (AGN). The luminosity could exceed the Eddington limit in the most extreme case (e.g., Wu et al. 2015).

Such a close relationship between galaxy merger and SMBH growth results in an inevitable consequence that pairs of SMBHs form in the nuclei of merging galaxies (Begelman et al. 1980). Such SMBH binaries (SMBHBs) are important astrophysical objects in the era of gravitational-wave astronomy. Merging SMBHBs are the major targets of the ongoing Pulsar Timing Arrays (Hobbs et al. 2010) and the planned Laser Interferometer Space Antenna (Amaro-Seoane et al. 2017).

However, there is a long-standing debate regarding the coalescence of SMBHBs. Earlier analysis of the dynamical evolution of the binaries revealed a bottleneck when the binaries shrink to a size of about 1 pc (Begelman et al. 1980). At this stage, the binaries become “hard” and start to slingshot the surrounding stars. In the simplest galaxy model, i.e., the stellar distribution is spherically symmetric and there is no gas, the replenishing of stars to the vicinity of the SMBHBs is inefficient (Makino 1997; Quinlan & Hernquist 1997; Milosavljević & Merritt 2001). Consequently, the evolution of the

binaries may stall. This theoretical prediction is inconsistent with the apparent scarcity of SMBHBs in galactic nuclei (Komossa 2006) and the conundrum is called “the final parsec problem.”

Real galaxies are more complicated than the idealized stellar systems based on spherical models. In general, merging galaxies are asymmetric. As a result, stars could be fed to the galaxy centers more efficiently so that the final parsec barrier may be avoided (Yu & Tremaine 2002; Zhao et al. 2002; Merritt & Poon 2004; Berczik et al. 2006). Moreover, since galaxy mergers often trigger gas inflow, it was realized early on that many SMBHBs may reside in gaseous environments and the interaction with gas may offer a potential solution to the above problem (Begelman et al. 1980; Ivanov et al. 1999; Gould & Rix 2000).

Later numerical simulations of a smaller (secondary) BH embedded in the accretion disk of a bigger (primary) SMBH generally confirm the above picture and, furthermore, showed that the evolution is similar to the migration of a planet in a protoplanetary disk: the secondary opens a gap in the disk and migrates toward the primary on the viscous timescale (Armitage & Natarajan 2002; Cuadra et al. 2009). Equal-mass binaries could even clear out a cavity in the accretion disk and in this case the merger is driven mainly by the spiral arms in the circumbinary disk (MacFadyen & Milosavljević 2008). In hotter environments where the gas distribution is more or less isotropic, the SMBHBs could excite elongated structures (Escala et al. 2004; Dotti et al. 2006) or density wakes (Kim et al. 2008) which lag behind the major axes of the binaries. These structures impose a negative torque on the binaries, which could also accelerate the shrinking of the binary orbit.

However, recent hydrodynamical simulations revealed a more controversial picture. They show that the aforementioned

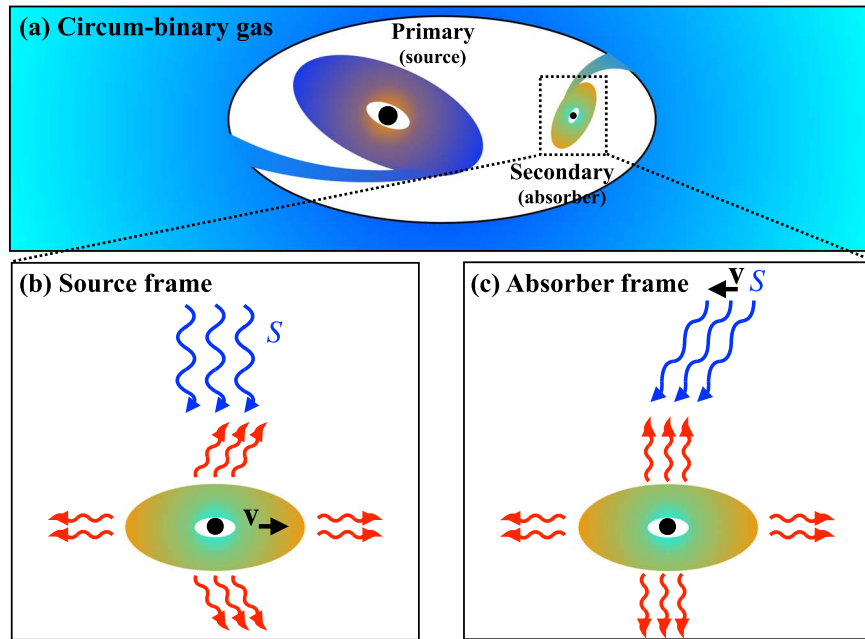


Figure 1. Physical picture of the PR drag. (a) An SMBHB embedded in a gaseous environment and both BHs are accreting. (b) In the rest frame of the light source (the circumprimary disk), the secondary BH is moving at a velocity of v . The light from the circumsecondary disk is beamed in the direction of motion, causing a linear-momentum loss. To conserve linear momentum, the circumsecondary disk and the embedded SMBH decelerates. (c) In the frame comoving with the absorber (the circumsecondary disk), the source is moving. The emission of the circumsecondary disk is now isotropic, but the light rays from the source is inclined due to the beaming effect. This inclined irradiation causes the circumsecondary disk and the embedded SMBH to drift relative to the comoving frame.

gap or cavity are not empty but filled with gas streams, which originate from the inner edge of the circumbinary disk and end up on both BHs (Hayasaki et al. 2007; Farris et al. 2014). These gas streams could exert a positive torque (Roedig et al. 2012) as well as directly deposit angular momentum onto the BHs (Hayasaki 2009; Shi et al. 2012). As a result, the binary orbit may even expand so that the final parsec problem remains (Miranda et al. 2017; Moody et al. 2019; Muñoz et al. 2019). It is worth noting that the generality of the expansion of the binary orbit deserves further investigation because the evolution is sensitive to the viscosity and the thermodynamical properties of the gas inside the gap and cavity (Tang et al. 2017).

So far the models have ignored the impact of the radiation of the accretion disks on the evolution of the binaries. It is known that our Sun could induce a drag force on the dust particles in the solar system. The drag effect is caused by the asymmetry between the absorption and reemission of the solar irradiation (Poynting 1903). A fully relativistic treatment of the phenomenon further clarified that the drag force can be attributed to the light beaming effect: More light is reemitted in the direction of motion of the dust particles (Robertson 1937). Such an effect, also known as the ‘‘Poynting–Robertson effect’’ (PR effect, hereafter), in principle, also applies to an SMBHB system because (i) one of the SMBHs, by accreting gas, could radiate and (ii) the other BH surrounded by its own accretion disk could absorb and reemit this radiation. Here we study this effect and show that the drag force is indeed important for the orbital decay of SMBHBs.

2. The Poynting–Robertson Drag

We consider an SMBHB embedded in a gaseous environment, and both the primary and secondary BHs are accreting gas. The configuration is illustrated in panel (a) of Figure 1. We

note that the two accretion disks, namely the circumprimary and the circumsecondary disks, are not necessarily coplanar or aligned with the orbital plane of the SMBHB, as is shown in the numerical simulations of accreting SMBHB systems (Dotti et al. 2010; Nixon et al. 2013; Gerosa et al. 2015; Goicovic et al. 2016; Takakuwa et al. 2017) as well as the observations of the circumstellar disks in binary protostars (Takakuwa et al. 2017). Therefore, both accretion disks could be irradiated by the companion. For simplicity, we consider the circumprimary disk as the light source and the circumsecondary one as the absorber. In the following, we study the PR drag exerted on the secondary BH.

The physical picture of the PR drag is shown in the lower two panels of Figure 1. (b) In the rest frame of the irradiating source (the circumprimary disk), the secondary SMBH is moving in a direction perpendicular to the light rays from the source. An absorption of the light by the circumsecondary accretion disk does not change this perpendicular velocity. Meanwhile, the circumsecondary disk is re-emitting more light in the direction of the motion. This beamed emission carries momentum and, by the law of linear-momentum conservation, the circumsecondary disk and the embedded SMBH should recoil in the opposite direction. This recoiling effectively causes the PR drag. We note that the irradiation by an external source is crucial to the PR drag. Without it, although the emitted light from the small body is still beamed, it does not slow down the moving body because the light also takes away the rest mass, compensating the loss of the linear momentum (pointed out by Robertson 1937). (c) The same conclusion can be drawn in the rest frame of the absorber, i.e., the circumsecondary disk. In this frame, the emission from the circumsecondary disk is isotropic but the light rays from the source become inclined because the source is now moving. The irradiation imposes a pressure on the circumsecondary disk and because of the inclination, the pressure has a component

pointing in the direction of the motion of the source. This component forces the secondary disk to drift with respect to the comoving frame. This drift is equivalent to the recoiling effect seen in the source frame.

We now go back to the source frame and the drag force due to the PR effect can be calculated with Sv/c^2 (Robertson 1937), where, in our scenario, v is the orbital velocity of the secondary SMBH, c is the speed of light, and S denotes the energy flux that is incident on the circumsecondary disk. The above equation assumes that all the energy flux is absorbed by the disk, which is generally true for optically thick accretion disks. As a result of the PR drag, the secondary BH decelerates and migrates toward the primary on a timescale of

$$T_{\text{PR}} = \frac{mv}{S(v/c^2)} = \frac{mc^2}{S}. \quad (1)$$

To derive the value of T_{PR} , we first express the energy flux using $S = \xi L r^2 / (4R^2)$, where L is the bolometric luminosity of the source, r denotes the radius of the circumsecondary disk, and R is the distance between the primary and secondary BHs. The coefficient ξ characterizes the cross section of the circumsecondary disk in the radiation field, and it is a function of the inclination of the disk relative to the light rays from the source. It is of order unity if the disks are misaligned. Even when the two disks are coplanar, the value of ξ does not vanish because accretion disks are not infinitely thin. The luminosity can be further written as $L = \eta L_{\text{Edd}}$, where L_{Edd} is the Eddington luminosity and η is the ‘‘Eddington ratio.’’ If the primary BH has a mass of M , the Eddington luminosity is $L_{\text{Edd}} \simeq 1.26 \times 10^{38} (M/M_{\odot}) \text{ erg s}^{-1}$. For the size of the circumsecondary disk (r), we notice that earlier numerical simulations found that it is comparable to the Roche radius $R_L := Rq^{1/3}(1+q)^{-1/3}$ (Lin & Pringle 1976; Artymowicz & Lubow 1994; Mösta et al. 2019), where $q := m/M$ is the mass ratio between the secondary BH and the primary one ($q \leq 1$ by definition).

In the derivation we assumed that the radiation field in the source frame is isotropic. It is known that the radiation from AGN accretion disks could be collimated when η is close to or exceeds the Eddington limit. In this case, if the radiation directly impacts the circumsecondary disk, the momentum flux would be greater than what we have estimated above. Moreover, accretion disks with high Eddington ratios also produce outflows or jets. These structures also carry momentum. If they strike the circumsecondary disk, the interaction could induce an additional drag force that is similar to the PR drag, only different in the sense that it is caused by the mass-momentum flux. We do not consider these additional effects in this work. Therefore, our PR timescale should be regarded as an upper limit.

With these considerations, we can rewrite Equation (1) as

$$T_{\text{PR}} = \frac{4mc^2}{L(r^2/R^2)} \simeq \frac{4q^{1/3}(1+q)^{2/3}}{\xi \eta} \frac{Mc^2}{L_{\text{Edd}}}. \quad (2)$$

Note that Mc^2/L_{Edd} is a constant independent of the BH mass or the distance between the two BHs. As a result,

$$T_{\text{PR}} \simeq 1.8 \times 10^9 \xi^{-1} \eta^{-1} q^{1/3} (1+q)^{2/3} \text{ yr}. \quad (3)$$

Therefore, despite the uncertainties in the hydrodynamics, the PR timescale is well determined by three parameters, namely, the mass ratio of the two BHs (q), the Eddington ratio for the

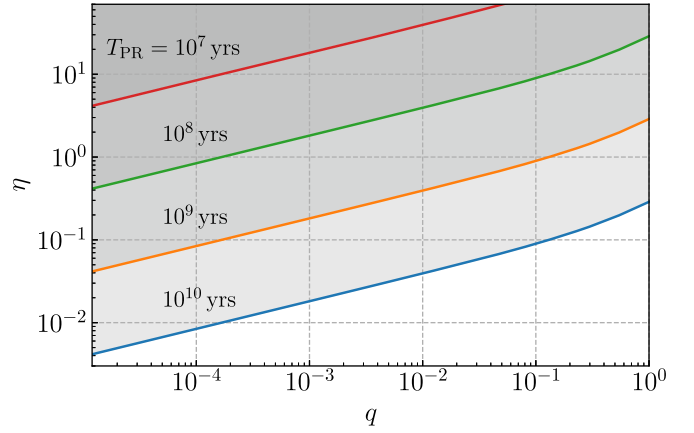


Figure 2. Dependence of the PR timescale (T_{PR}) on the mass ratio (q) of the SMBHB and the Eddington ratio (η) of the accreting primary BH. The curves represent the contours of constant T_{PR} .

primary BH (η), and the relative inclination of the two accretion disks (ξ).

Assuming that the inclination angle between the two accretion disks is large ($\xi \sim 1$), we show in Figure 2 the dependence of T_{PR} on the other two parameters. Interestingly, T_{PR} is shorter than the Hubble time (10^{10} yr) in a large fraction of the parameter space. In particular, for equal-mass binaries ($q \simeq 1$), coalescing within a Hubble time requires that the Eddington ratio is greater than about 0.3. For unequal binaries, e.g., $q \lesssim 0.1$, the requirement becomes $\eta \gtrsim 0.1(q/0.1)^{1/3}$.

3. Comparison with Other Timescales

To understand the relative importance of the PR drag, we compare T_{PR} with the other timescales related to the formation and evolution of SMBHBs.

(1) We first consider the gravitational-wave radiation timescale. We calculate it with

$$T_{\text{gw}} = \frac{5}{64} \frac{R^4 c^5}{G^3 (M+m) M m} \quad (4)$$

$$\simeq 2.3 \times 10^7 q^{-1} (1+q)^{-1} M_8^{-3} \left(\frac{R}{0.01 \text{ pc}} \right)^4 \text{ yr} \quad (5)$$

assuming circular orbits (Peters & Mathews 1963) and denoting $M/(10^8 M_{\odot})$ with M_8 . The condition $T_{\text{PR}} < T_{\text{gw}}$ could be satisfied when

$$R \gtrsim 0.03 \xi^{-1/4} \eta^{-1/4} q^{1/3} (1+q)^{5/12} M_8^{3/4} \text{ pc}. \quad (6)$$

Therefore, the PR drag predominates at relatively large binary separation.

(2) Dynamical friction against the stellar background could also shrink the orbit of an SMBHB. Following Binney & Tremaine (2008), we calculate the dynamical-friction timescale with $T_{\text{df}} = \sigma_*^3 / (4\pi G^2 \rho_* m \ln \Lambda)$, where ρ_* is the mass density of the stellar background, $\ln \Lambda \simeq 6$ is the Coulomb logarithm, and in deriving the above equation we have assumed that the gravitational potential of the primary BH predominates so that the orbital velocity of the secondary, v , is comparable to the velocity dispersion of the background stars, σ_* . Now we evaluate this T_{df} within the gravitational influence radius of the primary BH, $R_{\text{inf}} \simeq GM/\sigma_*^2$, within which the gravity of the primary BH predominates. We have to consider this restriction

because the PR timescales derived above become invalid outside the influence radius. Noticing that (i) empirically $\sigma_*(R_{\text{inf}}) = 200M_8^{1/4} \text{ km s}^{-1}$ (Tremaine et al. 2002), (ii) $\sigma_*^2(R) \simeq GM/R$ within the influence radius, (iii) the stellar mass enclosed in the binary orbit is about $M_* \sim 4\pi\rho_*R^3/3 \propto R^{3-\gamma}$ where γ denotes the power-law index of the density profile ($\rho_* \propto R^{-\gamma}$), and (iv) $M_*(R_{\text{inf}}) \sim M$ according to the definition of the influence radius, we deduce that

$$T_{\text{df}} \sim \frac{1}{3 \ln \Lambda} \left(\frac{M^2}{mM_*} \right) \left(\frac{R}{\sigma_*} \right) \simeq \frac{3000M_8^{1/4}}{q} \left(\frac{R}{R_{\text{inf}}} \right)^{\gamma-3/2} \text{ yr.} \quad (7)$$

Comparing it with Equation (3), we find that $T_{\text{PR}} < T_{\text{df}}$ when $q \lesssim 5 \times 10^{-5} (\xi\eta)^{3/4} M_8^{3/16} (R/R_{\text{inf}})^{3\gamma/4-9/8}$. This result suggests that the PR drag is more important than dynamical friction for unequal binaries.

(3) Salpeter timescale (T_{Sal}) characterizes how fast an accreting body increases its mass by one e-folding. It can be shown that T_{Sal} is closely related to T_{PR} . Given a radiative efficiency of ϵ for an accretion disk ($\epsilon \sim 0.1$), the Salpeter timescale can be calculated with

$$T_{\text{Sal}} := \frac{\epsilon Mc^2}{(1-\epsilon)L} = \frac{\epsilon}{(1-\epsilon)\eta} \frac{Mc^2}{L_{\text{Edd}}} \quad (8)$$

$$\simeq 4.53 \times 10^8 \epsilon \eta^{-1} (1-\epsilon)^{-1} \text{ yr.} \quad (9)$$

Using the relationship between T_{Sal} and Mc^2/L_{Edd} , we can rewrite Equation (2) as

$$T_{\text{PR}} \simeq \frac{4q^{1/3}(1+q)^{2/3}(1-\epsilon)}{\xi\epsilon} T_{\text{Sal}}. \quad (10)$$

The two timescales become comparable when $q \simeq q_{\text{cri}} := (\xi\epsilon/4)^3$. For $\xi = 1$ and $\epsilon = 0.1$, we find that $q_{\text{cri}} \simeq 1.6 \times 10^{-5}$.

When $q < q_{\text{cri}}$, the PR timescale becomes shorter than the Salpeter timescale. In this case, an SMBH could clear away the surrounding small BHs before it grows by one e-folding.

For an SMBHB with $q > q_{\text{cri}}$, the PR timescale is longer than the Salpeter timescale, i.e., $T_{\text{PR}} > T_{\text{Sal}}$. In this case, to drive the SMBHB to coalescence, the primary BH must grow by more than one e-folding. This scenario applies to those BHs in the early universe, where they have to accrete enough gas to grow from a mass of 10^2 – $10^5 M_\odot$ to the current 10^6 – $10^9 M_\odot$ (Volonteri 2010). The growth, in fact, amounts to nine e-foldings. Therefore, if we take $T_{\text{PR}} < 9T_{\text{Sal}}$ as the criterion for binary coalescence, adopting the assumption that $\epsilon = 0.1$ and $\xi = 1$, we find that $q \lesssim 0.016$.

(4) Duty cycle (T_D) is another important timescale. Conventionally, it is designed to characterize the lifetimes of AGNs. In our problem, it provides an estimation of the total duration of the accretion episodes of an SMBHB. Only when $T_{\text{PR}} \lesssim T_D$ is the PR drag efficient enough to affect the dynamical evolution of the binary. Observations of luminous AGNs suggest that T_D is a decreasing function of η (Hopkins & Hernquist 2009; Shankar et al. 2009). For $0.1 \lesssim \eta \leq 1$, T_D is typically 10^8 yr, and for $0.01 \lesssim \eta \lesssim 0.1$, T_D increases to 10^9 yr. These results in general agree with the Salpeter timescale as is derived in Equation (9). Using these values for the duty cycle, we find that the condition $T_{\text{PR}} < T_D$ is satisfied when $q \lesssim 1.7 \times 10^{-4}$ according to Equation (3). We note that this requirement applies mainly to those SMBHBs in the local

universe, because the duty cycles used in our analysis are derived based on relatively low-redshift AGNs.

(5) Small objects on inclined orbits with respect to an accretion disk could be ground down into the disk due to the mutual collisions (Syer et al. 1991). We adopt the formula in Ivanov et al. (1999) and calculate the ‘‘ground-down’’ timescale with $T_{\text{gd}} = m/(\Sigma A \Omega)$, where Σ is the surface density of the accretion disk at the point of collision, $\Omega = v/R$ is the angular velocity of the secondary BH, and A is the effective cross section of collision. The above equation is derived in the approximation that the relative velocity of the BH-disk collision is of the order of v .

To compare T_{gd} with T_{PR} , we first calculate the collisional cross section with $A = \pi(Gm/v^2)^2$. We also note that in the standard accretion-disk model, Σ is related to the accretion rate \dot{M} as $\dot{M} = 3\pi\nu\Sigma$, where ν is the viscosity and it is related to the viscosity parameter α and the disk scale height H as $\nu = \alpha\Omega H^2$ (Frank et al. 2002). From these relations, we find that

$$T_{\text{gd}} \simeq 3\alpha h^2 q^{-1} (1-\epsilon) T_{\text{Sal}}, \quad (11)$$

where $h := H/R$ is the aspect ratio of the disk. It is now clear that when $q \lesssim 8 \times 10^{-4}$, the PR timescale would be shorter than the ground-down timescale, if we adopt the typical parameters $\alpha = h = \epsilon = 0.1$ and $\xi = 1$. In this case, the binary would have shrunk significantly due to the PR drag before its orbit becomes coplanar the accretion disk.

(6) Even after the secondary BH has been ground down into the accretion disk, the PR drag may not vanish completely because the primary disk could be warped, e.g., due to the Bardeen–Petterson effect (Bardeen & Petterson 1975). It is well known that an embedded secondary would excite density waves in the accretion disk (Goldreich & Tremaine 1980). If the mass of the secondary is small, the disk surface density is not significantly perturbed and the interaction between the secondary and the disk can be calculated in a linear approximation. The interaction would result in a radial migration of the secondary, which is known as the type-I migration. The migration timescale can be calculated with

$$T_I = \frac{fh^2 M}{q\Sigma R^2 \Omega} \simeq 3\pi\alpha h^4 f q^{-1} (1-\epsilon) T_{\text{Sal}}, \quad (12)$$

where f is a parameter depending on the temperature and density profiles of the disk near the secondary’s orbit (e.g., Paardekooper et al. 2011), and we have applied the relation $\dot{M} = 3\pi\nu\Sigma$ in the second equation. Using our fiducial parameters and Equation (10), we find that $T_{\text{PR}} < T_I$ when $q \lesssim 6 \times 10^{-5}$. It is worth noting that the direction of the type-I migration could be inward or outward depending on the sign of f , which in turn depends on the exact temperature and density profiles of the accretion disk. Meanwhile, the PR drag always leads to an inward drift.

A small, embedded secondary BH could also accrete from the accretion disk. The increase in mass leads to an inward migration of the secondary because of the conservation of angular momentum. To estimate the corresponding timescale, we calculate the deceleration due to accretion with $\dot{v} = \pi\rho R_B^2 \Delta V^2/m$, where $\rho \sim \Sigma/(2H)$ is the surrounding gas density, ΔV is the relative velocity between the secondary and the surrounding gas, which is of the order of the sound speed c_s , and $R_B = Gm/\Delta V^2$ is the Bondi radius. Using the

condition $c_s = \Omega H$ for hydrostatic equilibrium, we find that the migration timescale is

$$T_M := v/|\dot{v}| \simeq 6\alpha h^5 q^{-1}(1 - \epsilon)T_{\text{Sal}}. \quad (13)$$

Compared to T_I , T_M is more sensitive to h , so the condition $T_{\text{PR}} < T_M$ requires an even smaller q , i.e., $q \lesssim 8 \times 10^{-6}$ in our fiducial model.

If the secondary is massive enough, an annular gap could be opened in the disk around the orbit of the secondary. In this case, the secondary will be locked in the gap and migrate on a timescale correlated with the viscous timescale of the disk, $t_{\text{vis}} = 2R^2/(3\nu)$. Following Lin & Papaloizou (1986), we calculate the timescale of this type-II migration with $T_{II} = t_{\text{vis}}(m/M_d)$, where $M_d \simeq \pi R^2 \Sigma$ is the disk mass enclosed in the orbit of the secondary and we have assumed that $M_d \ll m$. For a standard thin disk, we find that

$$T_{II} \simeq 2q(1 - \epsilon)T_{\text{Sal}}, \quad (14)$$

and, in fact, it is shorter than T_{PR} for any q if we adopt the fiducial parameters of $\eta = \epsilon = 0.1$. This result indicates that if the secondary becomes massive enough to open a gap in the accretion disk, the later evolution would be dominated by type-II migration and the PR drag is relatively unimportant.

4. Discussion

In our problem, the object receiving the PR drag is quite different from a solid dust particle. Nevertheless we calculated the drag using the formula derived for dust particles. Such a simplification deserves justification.

(i) In our problem the drag force is imposed directly on the circumsecondary accretion disk. However, the gravitational coupling between the disk and the secondary BH is so strong that the drag force can be imparted to the BH almost immediately. One can see this by calculating the dynamical timescale of the circumsecondary disk, which characterizes how quickly the accretion disk responds to a perturbation. We can calculate it with $2\pi(Gm/R_L^3)^{-1/2}$, which equals $2\pi(GM/R^3)^{-1/2}$. The dynamical timescale is the longest when the binary is at a distance of R_{inf} , but even in this case it is approximately $R_{\text{inf}}/\sigma_* \simeq 5 \times 10^3 M_8^{1/4}$ yr. We see that the timescale is indeed much shorter than the PR timescale. Therefore, any offset between the BH and the surrounding accretion disk will be damped relatively quickly.

(ii) Our emitter, i.e., the accretion disk, is rotating but it should not significantly affect the calculation of the drag force. This is so because the rotation velocity is typically $v_c \sim (Gm/R_L)^{1/2}$, and we find that $v_c/v \sim q^{1/3}$. Therefore, when $q \ll 1$, most of the gas in the circumsecondary disk is rotating at a velocity smaller than the orbital velocity, so the rotation can be neglected. Even when $q \sim 1$ so that $v_c \sim v$, the previous calculation of the PR drag force is approximately correct, because an axisymmetric rotation does not break the symmetry of the reemission and hence does not contribute to the PR drag. The orbital motion induces asymmetry due to aberration, and hence is the main source of the PR drag.

(iii) Our Figure 2 includes a region where the luminosity is super-Eddington ($\eta > 1$). Such large luminosity should not destroy the circumsecondary disk by blowing it off. This is because the conventional Eddington luminosity, which is used in this work, is derived assuming an opacity of $\kappa_e \simeq 0.4 \text{ cm}^2 \text{ g}^{-1}$, dominated by electron scattering. However,

accretion disks are normally optically thick, so that $\kappa_e \Sigma \gg 1$. As a result, the effective opacity is $1/\Sigma$ and it is much smaller than κ_e . Since Eddington luminosity is inversely proportional to the opacity, the corresponding effective Eddington luminosity, to blow away the circumsecondary accretion disk, is much higher than the conventional one.

(iv) Since the dynamical friction timescale T_{df} is inversely proportional to q , it seems that very small BH cannot come from outside the influence radius and be delivered to the vicinity of an SMBH. However, small BH can form in situ, as the remnants of massive stars, or be brought in by massive star clusters. How fast these channels populate the galactic nuclei with small BH is out of the scope of this work and deserves further investigation.

5. Conclusion

In this Letter, we investigate the impact of a new drag force, induced by the Poynting–Robertson effect, on the dynamical evolution of SMBHBs. We find that for a mass ratio of $q \lesssim a \text{ few} \times 10^{-5}$, the PR drag could predominate the dynamical evolution and lead to a fast coalescence of the BHs. The relevant systems include stellar-mass BHs of $\mathcal{O}(10) M_\odot$ around 10^6 – $10^7 M_\odot$ SMBHBs (the mergers are known as the “extreme-mass-ratio inspirals”), as well as intermediate-massive BHs (10^3 – $10^5 M_\odot$) around 10^8 – $10^{10} M_\odot$ SMBHBs. Unlike the dynamical-friction or type-I/II migration timescales, which are sensitive to the properties of the stellar and gas distribution around SMBHBs and hence are uncertain, the PR timescale is determined by fewer parameters, essentially only the mass ratio q and the Eddington ratio η . Our work made it possible to implement the PR drag in the future hydrodynamic and cosmological simulations so that we can better understand the evolution of unequal SMBHBs in galactic nuclei.

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