The Hydrodynamic Evolution of Binary Black Holes Embedded within the Vertically Stratified Disks of Active Galactic Nuclei

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Abstract

Stellar-mass black holes can become embedded within the disks of active galactic nuclei (AGNs). Afterwards, their interactions are mediated by their gaseous surroundings. Here, we study the evolution of stellar-mass binary black holes (BBHs) embedded within AGN disks using three-dimensional hydrodynamic simulations and analytic methods, focusing on environments where the AGN disk scale height $H$ is $\gtrsim$ the BBH sphere of influence. We model the local surroundings of the embedded BBHs using a wind tunnel formalism and characterize different accretion regimes based on the local properties of the disk. We develop prescriptions for accretion and drag for embedded BBHs. Using these prescriptions with AGN disk models that can represent the Toomre-unstable outer regions of AGN disks, we study the long-term evolution of BBHs as they migrate through the disk. We find that BBHs typically merge within $\lesssim 1$–$30$ Myr, increasing their mass significantly in the process, allowing BBHs to enter (or cross) the pair-instability supernova mass gap. The BBH accretion rate often exceeds the Eddington limit, sometimes by several orders of magnitude. Many embedded BBHs will merge before migrating significantly in the disk. We also discuss possible electromagnetic signatures during and following the inspiral, finding that it is generally unlikely for the bolometric luminosity of the BBH to exceed the AGN luminosity.

Unified Astronomy Thesaurus concepts: Black holes (162); Astrophysical black holes (98); Active galactic nuclei (16); Binary stars (154); Hydrodynamics (1963); Hydrodynamical simulations (767)

1. Introduction

Direct observations of Sagittarius A*, the supermassive black hole (SMBH) at the center of our galaxy, indicate the existence of a population of massive stars within its sphere of influence (Schoedel et al. 2002; Ghez et al. 2003; Paumard et al. 2006; Bartko et al. 2009; Lu et al. 2009; Bartko et al. 2010). Furthermore, the growing observational sample of tidal disruption events suggests that relatively young stellar populations are ubiquitous around SMBHs (Law-Smith et al. 2017; French et al. 2020; Mockler & Ramirez-Ruiz 2021). While most active galactic nuclei (AGNs) harbor SMBHs with low accretion rates (Ptak 2001), a subset of AGN are commonly believed to harbor accretion disks that are abundant in cold, dense gas (Heckman & Best 2014). These disks remain “canonically” thin (Shakura & Sunyaev 1973) up to scales of $10^{-2}$–$1$ pc (Sirko & Goodman 2003) depending on the mass of the SMBH and, at larger radii, gradually transition between the gas-dominated accretion disk and the star-dominated galactic disk.

It has recently been realized that a subset of the stars around AGN comprise massive stellar binaries that evolve into binary black holes (BBHs). Over the past several years LIGO/Virgo has detected dozens of BBH mergers (Abbott et al. 2019, 2020a; The LIGO Scientific Collaboration et al., 2020), which are typically attributed to formation through isolated binary evolution (Belczynski et al. 2016; Giacobbo et al. 2018; Kruckow et al. 2018; Schröder et al. 2018; Baver et al. 2021) or dynamical interactions in dense stellar systems (O’Leary et al. 2009; Antonini & Perets 2012; Samsung et al. 2014; Samsung & Ramirez-Ruiz 2017; Antonini et al. 2017; Askar et al. 2017; Banerjee 2018; Fragione & Kocsis 2018; Kremer et al. 2018; Rodriguez et al. 2018; Samsung & D’Orazio 2018; Di Carlo et al. 2019; Kremer et al. 2019). However, a potentially significant fraction of all LIGO/Virgo events may be due to BBHs that merge within an AGN disk (McKernan et al. 2012, 2014; Bartos et al. 2017; Stone et al. 2017; Tagawa et al. 2020). Due to the high gas densities within the disk, accretion and gas drag will help the binary inspiral much faster than via three-body hardening and gravitational-wave emission alone (Antoni et al. 2019).

The possibility that BBHs may merge within an AGN disk provides the tantalizing prospect that some BBH mergers may produce an associated electromagnetic (EM) counterpart; at the moment of merger, the sudden mass loss and recoil of the product black hole (BH) may shock-heat the surrounding accretion flow, resulting in an optical–UV flare (Lippai et al. 2008; Corrales et al. 2010; Mink & King 2017; McKernan et al. 2019). While largely unproven, this scenario was reinvigorated with the recently claimed association of the BBH merger GW190521 with an AGN flare (Graham et al. 2020), although we note that the significance of this association has been questioned (Ashton et al. 2020). This scenario is particularly intriguing because the BHs that merged to generate GW190521 are sufficiently massive that they fall within the pair-instability supernova (PISN) mass gap (Abbott et al. 2020b), suggesting that GW190521 could not have been formed through standard binary evolution scenarios.
(Safarzadeh et al. 2020; Vigna-Gómez et al. 2021) unless our understanding of PISN is in significant error (e.g., Belczynski 2020). It is therefore worth exploring the possibility that certain BBH mergers, such as GW190521, may have been formed in an AGN disk.

Stellar-mass BBHs can end up orbiting in the midplane of an AGN disk via two methods. First, they can form in situ: in the outer reaches of AGN disks, the gas becomes Toomre unstable, forming generations of stars already embedded within the disk (Toomre 1964; Goodman 2003). The BHs that are birthed by this generation of stars couple to the surrounding gaseous disk and stellar population, causing them to inspiral toward the central SMBH—and, in the case of binary systems, cause their orbital separation to shrink (e.g., Stone et al. 2017). Second, BHs can form in the surrounding nuclear cluster (NC) then migrate into the AGN disk via dynamical interactions: the density of the NC is sufficiently high that its most massive constituents—the BHs—will mass segregate into the central regions, nearest the SMBH (Morris 1993; Rasio et al. 2004). Here, the gravitational influence of the AGN disk is strong, and the mass-segregated BHs will gradually have their inclinations and eccentricities damped until they occupy circular orbits embedded within the AGN disk (Ward & Hahn 1994; Tanaka & Ward 2004). Single, rather than binary, BHs may also become embedded by these same mechanisms and then efficiently pair up via gas-mediated single–single encounters (Tagawa et al. 2020).

Once embedded within the AGN disk, BBH evolution is driven by a combination of three-body and hydrodynamic interactions. Repeated single–double encounters increasingly harden binaries. It is estimated that, within \( \lesssim 10 \) encounters, a BBH can be sufficiently hardened such that gravitational waves can merge the binary within a Hubble time (Leigh et al. 2018). At the same time, gaseous torques from the disk influence the center-of-mass motion of the BBH. Depending on the assumed AGN disk profile, these gaseous torques can be negative far from the SMBH and positive close to the SMBH, causing embedded objects to become stuck in the disk (Toomre 1964; Goodman 2003). The BBH can be sufficiently hardened by the gas drag that exerts negative torques on the binary that helps it inspiral on short encounters.

\[ R_{\text{BH}} = D \left( \frac{q}{3} \right)^{1/3}, \]

where \( D \) is the distance from the SMBH to the BBH center of mass, and \( q = M/M_{\text{SMBH}} \) is the BBH–SMBH mass ratio. The second length-scale to consider is the Bondi radius,

\[ R_{\text{b}} = \frac{GM}{c_s^2}, \]

were \( c_s \) is the local sound speed at distance \( D \) in the disk. At radii larger than \( R_{\text{b}} \), the pressure support of the gas allows it to be unperturbed by the gravity of the BBH. The sphere of influence of the BBH will be limited by the smaller of these two length-scales. If \( R_{\text{H}} < R_{\text{b}} \), the ram pressure of the wind is more important than the gas pressure, and if \( R_{\text{b}} < R_{\text{H}} \), the opposite is true. To parameterize the relative strength of these two regimes, we start by writing down the velocity profile of the wind in the noninertial rest frame of the BBH,

\[ v(\delta r) = v_k(D + \delta r) - v_k(D), \]

where \( v_k \) is the Keplerian velocity, and \( \delta r \) is the relative distance from the BBH. Here, we have essentially assumed that, near \( (\pm \delta r) \) the BBH, the velocity profile has a Cartesian geometry. We also neglect the vertical velocity dependence of the flow. It is advantageous to linearize this velocity profile because it will later allow us to model embedded BBHs in a scale-free fashion,

\[ v(\delta r) \approx \frac{GM_{\text{SMBH}}}{\sqrt{D + \delta r}} - \frac{GM_{\text{SMBH}}}{D} \approx \sqrt{\frac{GM_{\text{SMBH}}}{D}} \left[ 1 - \frac{\delta r}{2D} \right] - \frac{GM_{\text{SMBH}}}{D} \approx -\frac{1}{2} \delta r \Omega_k, \]
where $\Omega_k = \sqrt{GM_{\text{BH}}/D}$ is the Keplerian frequency about the SMBH. The sphere of influence of the BBH is always limited by $R_H$, and at this radius, the linearized velocity deviates only marginally from the true velocity when $q \ll 1$. We define the wind-capture radius, $R_w$, as the radius at which a wind with the linearized velocity profile $v(\delta r)$ is marginally bound to the BBH,

$$R_w \equiv \frac{2GM}{v(R_w)^2} = 2Dq^{1/3}. \tag{5}$$

This quantity differs from the Hill radius only by a constant, i.e., $R_w = 24^{1/3}R_H$. We opt to sometimes use $R_w$ instead of $R_H$ because we are borrowing intuition from Bondi–Hoyle–Lyttleton accretion (e.g., Edgar 2004), where the “accretion radius” describes the region within which a supersonic gas can be gravitationally captured and is defined analogously to our wind-capture radius. Using Equation (4) with $\delta r = R_w$, we define the Mach number of the wind at the wind-capture radius as

$$M_w \equiv \frac{\Omega_k R_w}{2c_s}. \tag{6}$$

Using the relation $H/R = c_s/v_k$ (Pringle 1981), where $H/R$ is the disk aspect ratio, we can rewrite this expression as

$$M_w \approx q^{1/3}(H/R)^{-1}. \tag{7}$$

We have found that our accretion flow is most strongly dictated by $M_w$. This is illustrated in Figure 1, where we delineate three flow regimes defined by $M_w$. When $M_w < 1$, then $R_h < R_H$, and the flow accretes in a quasi-spherical, Bondi-like fashion. When $M_w > 1$, then the wind flowing through the BBH’s sphere of influence is supersonic, and has high specific angular momentum. Here, the flow is best characterized as being dominated by the ram pressure of the incident wind. In the intermediate case, such that $M_w \sim 0.1–1$, the flow still has a large amount of angular momentum, but the higher gas pressure allows the wind to thermalize and produce disk-like structures.

2.2. Defining Properties of the Accretion Flow

It is worth taking a moment to contextualize our wind tunnel framework in the landscape of other accretion flows. Our setup is similar to other local simulations of accretion flows, which have proven effective tools for understanding both single and binary BH accretion in a variety of environments, including the interstellar matter (Antoni et al. 2019; Kaaz et al. 2019, 2021; S. Schröder et al. 2022, in preparation) and common envelope evolution (MacLeod & Ramirez-Ruiz 2015a, 2015b; MacLeod et al. 2017; Murguia-Berthier et al. 2017; De et al. 2020; Everson et al. 2020). This family of accretion flows essentially considers the permutations of Bondi–Hoyle–Lyttleton accretion (for a review, see Edgar 2004) characterized by the gas properties of the surrounding medium.

The case of the AGN disk is distinguished from these other scenarios by the large angular momentum content available to the embedded binary. In the limiting case of an effectively infinite angular momentum reservoir, one might expect the accretion to resemble the commonly studied thin circumbinary disk (“CBD”; Muñoz et al. 2019, 2020), where BBHs are expected to expand rather than contract due to positive
gravitational torques from the surrounding gas. This is similar to the scenario explored by Baruteau et al. (2011) and more recently by Li et al. (2021), who performed global simulations of an AGN disk with an embedded binary system in two dimensions. In Li et al. (2021), they found that prograde embedded binaries also expand, as in the case of CBDs. We emphasize, however, that, for these works to globally resolve the AGN disk and the binary, they are required to use large values of \( q \) and model the system in two dimensions. In works such as Li & Lai (2022), who perform two-dimensional shearing box simulations and thus can achieve higher resolutions on the scale of the CBDs and circumsingle disks, they instead find that the flow behaves differently than in traditional CBDs. In particular, they find that the accretion flow is quasi-steady, exhibiting variability at twice the orbital frequency in the corotating frame. They also find that their binaries contract. However, in the two-dimensional shearing box simulations of Li et al. (2022), they find that the question of binary contraction versus expansion depends on the thermodynamics of the disk, and in the three-dimensional shearing box simulations of Dempsey et al. (2022), they find that that the semimajor axis of the binary is also important.

By using Equation (5), we can write \( R_w/H \approx 2q_{1/3}(H/R)^{-1} = 2M_w. \) When \( q \) is large and \( H/R \) is small, \( R_w \gg H \), which justifies two-dimensional approaches. However, as we will find in the following section, for much of the \( q-H/R \) parameter space, \( H \gg R_w \), and the resulting flow is inherently three-dimensional, supported partially by rotation and partially by pressure gradients. This suggests a dichotomy in the geometry of accretion flows surrounding embedded binaries: in thin accretion disks hosted by lower-mass (\( \lesssim 10^6 \)–\( 10^7 \) M\(_{\odot}\)) SMBHs, the embedded binary may be better represented by the more traditional two-dimensional accretion flows studied by Baruteau et al. (2011) and Li et al. (2021). In either thicker AGN disks (i.e., for high Eddington ratios, or in the outer Toomre-stabilized regions discussed in the next section) or for more massive (\( \gtrsim 10^7 \) M\(_{\odot}\)) SMBHs, the embedded binaries may be better represented by the vertically stratified accretion flows studied here.

2.3. Astrophysical Context

What value of \( M_w \) appropriately represents the flow surrounding an embedded binary? We can answer this by applying our \( M_w \) formalism to physical models of AGN disks. In particular, we consider Shakura–Sunyaev (S-S; Shakura & Sunyaev 1973) and Sirko–Goodman (S-G; Sirko & Goodman 2003) accretion disks. In the classical S-S disk, each disk annulus has Keplerian orbital energy that is gradually dissipated by some unspecified source of viscosity that is parameterized by the quantity \( \alpha \). While the S-S disk forms the basis for most modern disk models, it becomes gravitationally unstable at large distances (typically \( \gtrsim 10^6\)–\( 10^7 \) SMBH gravitational radii). The S-G model addresses these issues by including an unspecified pressure term that forces the disk to be marginally Toomre stable where it would otherwise be unstable. In reality, the disk likely continuously forms stars in this outer region, and the stellar feedback from these populations holds the disk in its marginally stable state.

We can explore how \( M_w \) depends on S-S and S-G AGN disk models by using Equation (7) to relate \( M_w \) to the mass ratio \( q \) and disk aspect ratio \( H/R \). The profile of \( M_w \) can depend sensitively on the parameters of the AGN. For simplicity, unless said otherwise, we will assume the host AGN disk has the following parameters for the remainder of this work:

1. a \( 10^8 \) M\(_{\odot}\) SMBH accretes at an Eddington ratio of \( \lambda_{\text{Edd}} = 0.5 \), with a radiative efficiency of \( \eta = 0.1 \).
2. viscosity is proportional to the total (gas + radiation) pressure, and \( \alpha = 0.01 \).
3. the embedded, equal-mass binary has a total mass \( M = 60 \) M\(_{\odot}\).

In the case of the S-G model, this is the same disk model used to study the orbital migration of embedded BHs in Secunda et al. (2019), who found that their migrating BHs became trapped at roughly \( \approx 10^3\)–\( 10^2 \) pc.

In the left panel of Figure 2, we plot \( M_w \) as a function of distance \( D \) in the AGN disk for both profiles. For the S-S profile, we show both a disk with \( H/R \) set by gas pressure
We notice a few features immediately. First, it is unusual for \( M_w \) to exceed unity. An exception is the gas pressure dominated S-S curve, but only at small radii where the disk is in fact radiation pressure dominated at this value of \( \alpha_{\text{SMBH}} \). The radiation pressure dominated S-S curve has large \( M_w \) at large radii, but here the flow is gas pressure dominated. For the S-G profile, which is more credible in the outer regions where the S-S disk is Toomre unstable, the flow is actually Bondi-like at regions beyond \( \approx 1 \) pc. We note the presence of the inflection point at \( \approx 0.05 \) pc in the S-G profile, which marks the transition between the inner disk and the Toomre-stabilized outer disk (see for reference Figure 2 of Sirko & Goodman 2003, which depicts the same profiles used here).

In the right panel of Figure 2, we plot contours of \( M_w \) in the \( q-H/R \) plane (Equation (7)), and shade the same Bondi-like (orange) and gap-opening (blue) regions as in the left panel. The vertical axis in this panel is alternatively labeled in terms of \( M_{\text{SMBH}} \), given \( M = 60 M_\odot \). We have labeled \( M_w = 0.5, 1 \) and 2 isocountours, which correspond to the hydrodynamic simulations that we present in Section 3. The gap-opening criterion is satisfied only for particularly thin accretion disks with particularly high mass ratios. This point deserves emphasis, as there are conflicting claims in the literature regarding gap formation in AGN disks. In Baruteau et al. (2011), the gap formation happens primarily because of their choice of \( \alpha = 8 \times 10^{-4} \) (\( \alpha \approx 0.01-0.3 \) is more realistic for AGN disks, e.g., King et al. 2007), and the gap opening is clearly sensitive to this parameter as can be seen in Figure 2. In Bartos et al. 2017, they also invoke gap opening, but use a disk model that is thinner than those from both S-G and S-S models, causing \( M_w \) to be larger. In Stone et al. (2017), they also find that gaps are unlikely to open anywhere in the disk, and they use the AGN disk model from Thompson et al. (2005). This model is comparable to the S-G model because they also construct profiles that are marginally Toomre stable in the outer regions. Still, how significantly gap-opening affects the orbital evolution of embedded binaries remains unclear. For instance, in Li et al. (2021), they find that the gap depth weakly affects the evolution of the semimajor axis, particularly at larger \( H/R \). This is because the binary evolution is mainly dictated by gas within the Hill sphere(s), which can be readily replenished even when a gap is opened.

3. Hydrodynamics

3.1. Simulations

To test our \( M_w \) formalism, we performed three-dimensional hydrodynamic simulations of accretion onto embedded BBHs using version 4.5 of the grid-based, adaptive mesh refinement hydrodynamics code FLASH (Fryxell et al. 2000). We use a local, wind tunnel computational domain that is continuously replenished with a shearing gaseous wind (see the arrows in the inset panels of Figure 1). Physically, the velocity profile of the wind is an approximation derived from the Keplerian velocity about the SMBH in the noninertial rest frame of the BBH, which we acquire using the linearized velocity profile defined in Equation (4). For simplicity, we neglect gradients in the gas density and sound speed of the wind, setting them to ambient values \( \rho_\infty \) and \( c_\infty \) respectively. We evolve our simulations in a scale-free fashion, where \( c_\infty = R_\infty = 1 \), and \( \rho_\infty = 10 \). We relate the linearized velocity profile in Equation (4) to our scale-free setup by rewriting it as

\[ v(y) = \frac{y c_\infty M_y^4}{2 R_b} \]  

where we have substituted \( y \) for \( y \).

In this fiducial work, we solve the equations of inviscid hydrodynamics rather than the shearing box equations, and leave the latter for a later work. This means that we neglect centrifugal and coriolis accelerations from the SMBH potential. Our initial conditions in the absence of the BBH are setup to be in a steady state without these forces. When constructing the domain, we also neglect the vertical gradients in the wind, acknowledging that when \( M_w > 1 \) the AGN disk scale height \( H \) becomes comparable to the BBH sphere of influence. We also neglect the breaking of the symmetric velocity profile by pressure gradients in the AGN disk. These pressure gradients however modify the velocity profile on the order of \( (H/R)^2 \) (Pringle 1981), which is negligible for much of the parameter space (e.g., Figure 2).

These neglected effects are negligible when \( a \ll R_\text{H} \) (where \( a \) is the semimajor axis of the binary), but would affect our results when \( a \gtrsim R_\text{H} \). In a real system, binaries with \( a > R_\text{H} \) are unbound. Some of our simulation parameter choices correspond to this regime, but remain valuable for establishing hydrodynamic scaling relations within our wind tunnel setup. We also neglect the self-gravity of the gas; however, even if self-gravity is important for the equilibrium state of the Toomre-unstable regions of an AGN disk, within our domain the gravity of the binary dominates, and the self-gravity of the gas is dynamically unimportant. The width of our computational domain is \( 32 R_\infty \) in each Cartesian direction, with the wind oriented in the \( \pm x \) direction and generated in every time step at the \( x \) boundaries of the domain. For the \( y \)- and \( z \)-directions, we use outflow boundary conditions. Our simulation domain size is much larger than the BBH sphere of influence, and consequently the overall gas mass in the simulation domain remains roughly constant.

It is well known that hierarchical triple systems are sometimes unstable (Eggleton & Kiseleva 1995; Merline & Aarseth 2001; He & Petrovich 2018). The results of Eggleton & Kiseleva (1995) suggest that, for circular orbits, the stability of the inner binary is dictated by the semimajor axes of the inner (a) and outer (D) orbits and the mass ratio (q) of the inner binary to the outer massive object (the SMBH, in our case). So, while the hydrodynamics of our simulations are self-similar, they are not necessarily dynamically stable for all parameter choices. We acknowledge this as a limitation of our local simulations, but emphasize that the purpose of this work is to isolate the hydrodynamics of accreting, embedded BBHs. This
allows us to develop scaling relations that can then be applied to the dynamically stable parameter space.

We evolve the BBHs using the active sink particles module that was developed for FLASH by Federrath et al. (2010), and use the same implementation of it as Antoni et al. (2019). While we refer the reader to these works for complete descriptions of our sink prescription, the most salient details are as follows. The binary is initialized with its center of mass at \( x = y = z = 0 \) and occupies a circular orbit with semimajor axis \( a \) in the \( x-y \) plane. The \( +z \)-axis is parallel to the angular momentum vectors of the binary orbit and the AGN disk. We model the BBHs using Newtonian gravitational potentials with an absorbing boundary condition centered at the location of each binary companion. These absorbing boundaries have sink radii \( R_s = 0.0125 \ R_b \), and any gas that crosses this boundary is removed from the domain, with the mass and angular momentum of the removed gas being recorded and added to the corresponding binary member. We also check the convergence of our mass accretion and inspiral rates with the sink radius in Appendix A. While the rate of inspiral is converged, the mass accretion rate decreases with decreasing sink radius. The nonconvergence of the mass accretion rate with the sink radius is a known issue with the sink prescription in Bondi-like accretion flows (Xu & Stone 2019; De et al. 2020), and so the accretion rates derived from these simulations should be taken as upper limits. We adaptively refined around the region closest to the binary members, with smallest cell size \( \delta_{\text{min}} \approx 0.0039 \ R_b \). This means that there are \( R_s/\delta_{\text{min}} \approx 3.2 \) cells across a sink radius, a value that was tested for convergence in Antoni et al. (2019).

In total, we present six simulations, varying \( M_w \) and \( a/R_b \), where \( a \) is the BBH semimajor axis. We use values \( M_w = 0.5, 1, \) and \( 2 \), and \( a/R_b = 0.1 \) and 1. We use a gamma-law equation of state, and assuming that the gas in the AGN disk can cool efficiently, take \( \gamma = 1.1 \). In each simulation, the BBH occupies a circular, equatorial orbit that is prograde with the angular momentum of the wind. Each simulation was run until \( t = 50 \ R_b/\c_{\text{in}} \), well into the steady state as shown in Appendix A. In the following two subsections, we present the results of these simulations.

### 3.2. Flow

In Figure 3, we show the large-scale flow morphology of each simulation. In the top \( 2 \times 3 \) panel subfigure, we depict \( 12 \times 12 \ R_b \) equatorial slices of the gas density. When \( M_w = 0.5 \), the flow is structured and laminar, with an extended density profile that is wound in a direction prograde with the binary angular momentum. When \( M_w = 2 \), the flow becomes unwound, with tails that follow the direction of the wind. When \( M_w = 1 \), the flow shares characteristics of both the \( M_w = 0.5 \) and 2 simulations. It is worth taking a moment to compare these structures to those found in the global simulations of Li et al. (2021), Baruteau et al. (2011). In these works, the disk is cool and two-dimensional, and spirals around the mini-disks of the binaries are formed and become unwound as they transition into the global AGN disk. While our simulations share some of these features, particularly in the bottom left and middle panels of Figure 3, we generally find much less pronounced spiral patterns. We hypothesize that this is due to the following distinguishing features in our accretion flow. First, the disks in our simulations are hot and partially pressure-supported, making the angular velocity profile sub-Keplerian. This, in turn, will change the radial spacing of the Lindblad resonances, which is responsible for the formation of spiral density waves. Second, as seen most prominently at \( M_w \geq 1 \), our flow has a clumpy, turbulent density distribution, which will damp any spiral density waves that are formed. Disk-like structures are evident in our two-dimensional slices, and can be quantified with the circularization radius of the wind,

\[
R_{\text{circ}} = \frac{GM}{v^2(y = R_{\text{circ}})} \rightarrow R_{\text{circ}} = \frac{M_w}{2 \pi^2} \frac{\rho R_b}{M} \, (9)
\]

For \( M_w = 0.5, 1, \) and 2, the corresponding circularization radii are \( R_{\text{circ}} \approx 6.4, 1.6, \) and 0.4 \( R_b \). These values roughly correspond to the extent of the density structures in each simulation, particularly when \( a/R_b = 0.1 \).

The main difference between the \( a/R_b = 0.1 \) and 1 simulations is that, in the former, the accretion flow forms a high-density envelope shared by both binary companions, while in the latter each companion has its own envelope. This makes sense, because, when \( a/R_b = 0.1 \), the semimajor axis of the binary is much smaller than the sphere of influence, while, when \( a/R_b = 1 \), they are comparable.

In the bottom \( 1 \times 3 \) panel subfigure of Figure 3, we show the three-dimensional gas density contours from each \( a/R_b = 1 \) simulation. As also seen in the two-dimensional subfigure, when \( a/R_b = 1 \), each binary companion has its own high-density envelope, which overlap more when \( M_w \) is smaller. As \( M_w \) increases, the peak density remains similar, but each density envelope becomes truncated. While in the two-dimensional subfigure, the flow appears disk-like; we see from the three-dimensional panel that the density envelopes are as extended in the vertical directions as they are in the equatorial directions. This suggests that the flow structures surrounding embedded BBHs are partially Bondi-like, and partially disk-like.

### 3.3. Accretion and Inspiral

The orbits of our simulated BBHs evolve due to a combination of gas drag and the accretion of mass and angular momentum. To compare to our numerical results, we will begin by writing down the analytic binary inspiral rate, \( \dot{a}/a \). The full expression for \( \dot{a}/a \) is

\[
\frac{\dot{a}}{a} = \frac{2}{L} \left( \frac{L}{L}_{\text{accr}} - \frac{3M}{M} \right) + \left( \frac{a}{a}_{\text{drag}} \right), \tag{10}
\]

which assumes that the specific angular momentum of the binary is Keplerian (\( \ell = \sqrt{GM a} \)). We have split the angular momentum evolution into two terms for accretion and gas drag, i.e., \( 2L/L = 2(L/L)_{\text{accr}} + 2(L/L)_{\text{drag}} \), where we have written \( 2(L/L)_{\text{drag}} = \left( \frac{a}{a}_{\text{drag}} \right) \) to be consistent with our analytic estimate of drag, which we will introduce shortly. The contribution from gas drag is found to always be negative in our work. This is in contrast to hydrodynamic simulations of thin, gap-forming circumbinary disks where the drag term is often found to be positive, e.g., Muñoz et al. (2019). We are

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6 In cases where the accretion rate is super-Eddington, the flow is radiation dominated, and \( \gamma = 4/3 \) would be more appropriate, but we leave the explicit effects of radiation to be considered in a later work.
mainly interested in the latter two terms for mass accretion rate and drag. Note, however, we do include a time series for Equation (10) in Figure 8. To determine $\mathcal{M}$ analytically, we integrate over streamlines entering the gravitational sphere of influence of the BBH. We start by defining the length-scale,

$$R_{bw} = \frac{R_b}{1 + \frac{1}{2} \mathcal{M}_w^2},$$

which smoothly transitions between $R_b$ when $\mathcal{M}_w \ll 1$, and $R_w$ when $\mathcal{M}_w \gg 1$. This allows us to continuously characterize the sphere of influence of the BBH across flow regimes analytically. The integral for the mass accretion rate is

$$M = f \times \rho_\infty \int_{r < R_{bw}} v(y) dA$$

$$= f \times 4\pi R_{bw}^2 \rho_\infty c_\infty \left( \frac{\mathcal{M}_w^3}{8 + 4\mathcal{M}_w^2} + 1 \right),$$

where we have included an ad hoc numerical prefactor $f$. We do this because these calculations assume that all streamlines entering the BBH sphere of influence are accreted; though in reality some fraction of this material is advected away due to gas and ram pressure, such that $f < 1$. To estimate $\frac{a}{a}_{\text{drag}}$, we

\textbf{Figure 3.} Upper panel: here, we highlight the large-scale flow morphology for each of our simulations by plotting $12 \times 12 R_b$ slices of gas density at late times. Each column from left to right represents increasing local Mach number $\mathcal{M}_w$, where at high $\mathcal{M}_w$ the ram pressure of the wind prevents stable flow structures from forming. In each row, we plot different values of the binary semimajor axis $a$. At small $a$, the binary shares a density enhancement, whereas at large $a$ each binary member has its own density enhancement. Lower panel: we plot three-dimensional gas density isocontours for each of our $a = 1 R_b$ simulations at the same times as the upper panel. As $\mathcal{M}_w$ increases, the density enhancements surrounding each BH become smaller.
follow the approach of Antoni et al. (2019), who determined the inspiral rate for a BBH embedded in a Bondi–Hoyle wind tunnel (e.g., Edgar 2004; Blondin & Raymer 2012). In Appendix B, we modify their derivation for our wind tunnel setup and derive the expression

\[
\left(\frac{\dot{a}}{a}\right)_{\text{drag}} = s \times \frac{8\pi G a \rho_{\infty}}{(GM/4a + c_{\infty}^2 + M_w c_{\infty}^2)^{1/2}} \left(1 + \frac{2R_{\text{desc}}}{a}\right)^p.
\]  

(13)

Here, \(s\) and \(p\) are numerically determined parameters that we will fit for. This calculation assumes that the orbital energy of the binary is dissipated by a gravitational wake induced by each component of the binary as they orbit through a dense gaseous envelope. These wake-induced gravitational torques are different from those that dictate the orbital angular momentum evolution in thin circumbinary disks, which is mainly driven by Lindblad torques from the inner and outer disks. Our main argument for neglecting the study of Lindblad torques is that our accretion flows are pressure dominated, rather than rotation dominated, and we do not have any circumbinary disks. It is possible that there is some transition that occurs at higher \(M_w\), where the angular momentum evolution of the binary transitions to a structure that better resembles the usual circumbinary disk.

In Figure 4, we assess how well our analytic estimates compare to our simulated values for the accretion and drag rates as a function of \(M_w\). Our simulated values of mass accretion rate are calculated by taking the time derivative of the mass of the binary, which is updated each time step during our simulation. The drag rate is determined by recording the net drag force that the gas exerts on the binary at each time step. We express our rates in units of \(\dot{M}_{\text{Bondi}} = \frac{c_{\infty}^3}{4\pi G M_{\text{Bondi}} \rho_{\infty}}\), which is the Bondi accretion rate per unit mass (Bondi 1952). We have numerically determined that \(f = 0.58\) (Equation (12)), \(s = 0.35\), and \(p = -0.86\) (Equation (13)). Using these adjustments, Figure 4 shows that our analytic estimates for our mass accretion and drag rates reasonably reproduce our simulated results. In general, the binaries accrete and inspiral on similar timescales, with inspiral occurring faster (slower) at larger (smaller) separations. This is in contrast to what is seen in Bondi–Hoyle accretion onto binaries, e.g., Figure 11 of Antoni et al. (2019), where inspiral happens about 3–4 times faster than accretion. We note that, when \(M_w \gg 1\), our local wind tunnel approximation breaks down because the BBH can begin influencing the AGN disk, so care should be taken in extrapolating to larger values of \(M_w\). In those cases, the flow geometry is likely better represented by the results of Baruteau et al. (2011). In contrast, extrapolating our mass accretion prescription to \(M_w = 0\) provides no issue, as the flow approaches a Bondi solution in that limit.

4. Evolution of Embedded BBHs

4.1. How Does the Binary First Become Embedded?

To determine the evolution of embedded BBHs, it is important to understand how they are first captured. The BBH is either born in the outer regions of the disk or, having previously existed in the NC when the disk was formed, it gradually aligns with the disk. The initial position of the embedded BBH is determined by how it became embedded, which strongly impacts the binary’s subsequent evolution.

![Figure 4](image-url)
Hailey et al. 2018). However, the formation of this cusp requires that the stellar population has relaxed around the central SMBH (Bahcall & Wolf 1976). The relaxation timescale for a $10^6 M_\odot$ SMBH can range from 0.1–10 Gyr (O’Leary et al. 2009), and can exceed a Hubble time for more massive SMBHs. So, it may be that heavier SMBHs harbor no BH cusp during the AGN phase, which makes the viability of the dynamical capture channel less certain. For this reason, we favor the in situ formation of embedded BBHs, which should begin their journey in the AGN disk at larger ($\gtrsim 1$ pc) distances.

### 4.2. Binary Evolution

Once BBHs become embedded in the AGN disk, they will simultaneously inspiral, migrate, and grow in mass. In this section, we model these processes by evolving a system of coupled ordinary differential equations for $D$ (migration), $M$ (accretion), and $\dot{a}$ (inspiral). To guide our intuition, we begin by estimating the timescales associated with each of these processes. We use an S-G AGN disk model with $M_{\text{SMBH}} = 10^8$, $\alpha = 0.01$, $\lambda_{\text{Edd}} = 0.5$, and assume that the viscosity is proportional to the total (gas + radiation) pressure. At a characteristic distance of 1 pc for a $60 M_\odot$ binary, the migration timescale is

$$t_{\text{migr}} \approx 61 \text{ Myr} \left( \frac{0.06}{H/R} \right)^2 \left( \frac{60 M_\odot}{M_{\text{BBH}}} \right)^{-1} \left( \frac{10^8 M_\odot}{M_{\text{SMBH}}} \right)^{3/2} \left( \frac{0.5 \text{ pc}}{D} \right)^{-1/2} \left( \frac{1540 \text{ g cm}^{-2}}{\Sigma} \right)^{-1},$$

(14)

which we acquire from Equation (10) of Paardekooper (2014). This characteristic timescale is the comparable lifetime of the host disk ($\approx 50–150$ Myr), suggesting that embedded binaries may migrate significantly (e.g., Secunda et al. 2019, 2020). This equation is for the static torque usually associated with Type I migration, which is valid under our assumption that the embedded BBH only negligibly affects the evolution of the host AGN disk. This expression also neglects self-gravity of the host disk and radiative torques, which may enhance the migration rate (Kley & Nelson 2012).

At large semimajor axes, before gravitational-wave emission dominates, the binaries inspiral due to the combined influence of mass accretion and drag. At distances of $\gtrsim 0.5$ pc, the accretion flow is Bondi-like (e.g., Figure 2), for which the mass doubling time ($\equiv M/M_{\text{Bondi}}$) is characteristically

$$t_{\text{Bondi}} \approx 148 \text{ Kyr} \left( \frac{M}{60 M_\odot} \right)^{-1} \left( \frac{\rho}{8 \times 10^{-15} \text{ g cm}^{-3}} \right) \left( \frac{c_s}{54 \text{ km s}^{-1}} \right)^{-3}.$$  

(15)

This is extremely rapid, which at first glance suggests the embedded binaries should easily balloon in mass. The associated drag timescale ($\equiv$ the inverse of Equation (13), where for simplicity we take $M_a = 0$, and $a = 100$ au) is $\approx 18$ Kyr, increasing to $\approx 72$ Kyr at $a = 10$ au. Together, these three timescales form a hierarchy: the binary inspirals first (mainly due to drag, but also due to mass accretion), then increases in mass, then migrates.

However, the mass accretion rate associated with Equation (15) is super-Eddington throughout much of the disk.

This can be seen in Figure 5, where we plot the Eddington ratio of the embedded binary, $\lambda_{\text{Edd}}$, as a function of distance in the AGN disk. This figure shows that the embedded binary can be supplied gas at rates of up to $\sim 10^4–10^5$ the Eddington ratio. These profiles peak at radii of $\sim 0.05–0.07$ pc, where there is a transition to the star-forming outer regions of the host AGN disk. In practice, we limit the mass accretion rate to its Eddington-limited value,

$$\dot{M}_{\text{Edd}} = \frac{L_{\text{Edd}}}{\eta c^2}.$$  

(16)

Simulations of super-Eddington accretion have shown that the mass accretion rates can significantly exceed the Eddington rate (McKinney et al. 2014, 2015; Dai et al. 2018), which is essentially set by how small the radiative efficiency ($\eta$) is. Thin disks have typical values of $\eta = 0.1$, which lowers to roughly $\eta = 0.01$ near the Eddington limit, and can be even smaller at significantly super-Eddington rates. For simplicity, we adopt a constant value of $\eta = 0.01$. The associated Eddington accretion timescale ($\equiv M/M_{\text{Edd}}$) is independent of mass,

$$t_{\text{Edd}} \approx 4.5 \text{ Myr} \left( \frac{\eta}{0.01} \right).$$  

(17)

Even with Eddington-limiting our mass accretion rate, it still occurs on a shorter timescale than that in migration. We must make a few modifications for the inspiral rate in our binary evolution prescriptions, as well. First, we include a term for the gravitational-wave inspiral rate of an equal-mass binary (Peters 1964). Second, we neglect the $(L/L_{\text{ISCO}}$ term in Equation (10), because in real systems this term can only contribute marginally ($O(Ml_{\text{ISCO}})$, where $l_{\text{ISCO}}$ is the specific angular momentum at the innermost stable circular orbit) to the binary’s evolution. We also make the assumption that whenever the mass accretion rate is Eddington limited, so is the drag rate, because an Eddington-limited disk will be depleted of the gas that drives the binary to inspiral.
We evolve our binaries in Figure 6, where we show how the binary separation, mass, and AGN disk position change in time. We initialize our evolved binaries with masses $M = m_1 + m_2 = 20$ (dashed lines) and $60 M_\odot$ (solid lines), and initial distances $D = 0.5, 1, 2$ pc. We initialize our binary semimajor axes such that they are marginally stable to the hierarchical triple instability outlined in Eggleton & Kiseleva (1995). We use $\alpha_{\text{Edd}}^{\text{SMBH}} = 0.5$ on the left subfigure, and $\alpha_{\text{Edd}} = 0.05$ on the right. Each integration is terminated when the semimajor axis reaches $a = 0.1$ au, after which the inspiral becomes rapid due to gravitational-wave decay. In all cases, the full evolution occurs on timescales of $\lesssim 1 - 30$ Myr. The relative hierarchy of timescales follows our expectations from our analytic predictions; the binary inspirals the quickest, and is driven to inspiral by a combination of drag, gravitational-wave radiation, and mass accretion. Second, the masses begin to increase and, in some cases, double. For each profile, the binaries migrate only marginally before they merge. To highlight uncertainties in our binary evolution prescriptions, we reran each profile with the prescribed mass accretion rate reduced or increased by a factor of 10, the results of which are shown by the error bars in our middle panel.

In general, the evolution of the binary happens well within the lifetime of the AGN disk ($\approx 50 - 150$ Myr), labeled $\tau_{\text{AGN}}$ in Figure 6). However, a potentially important caveat in this evolution is that we keep the AGN disk static. In reality, the AGN disk can “flicker” on short timescales ($\approx 0.1 - 10$ Myr, labeled $\delta \tau_{\text{AGN}}$ in Figure 6; e.g., Schawinski et al. 2015). A similar phenomena is seen in the simulations of Alcacer et al. (2020), who found that the gas supplying luminous quasars can restructure and change the orientation on timescales of $\approx 0.1 - 1$ Myr. If the binary becomes embedded, but then within $\delta \tau_{\text{AGN}}$, the disk restructures and reorients itself; the BBH may once again require its inclination to be damped with respect to the restructured disk for a gas-assisted inspiral to proceed.

5. Discussion

5.1. Implications for Gravitational-wave Signals

The results of Figure 6 have profound implications for the gravitational-wave signals observed by LIGO and Virgo. If a BBH becomes embedded within an AGN disk, it can be driven to merger on timescales $\lesssim 1 - 30$ Myr. While the rates at which BBHs form within AGN disks are unclear, our calculations suggest that a gas-assisted inspiral is a viable formation channel for BBH merger events. Furthermore, the BHs in these binaries accrete at significant rates, suggesting that a pair of $30 + 30 M_\odot$ BHs can potentially evolve into the pair-instability supernova mass gap. Such a system was recently observed by LIGO/Virgo (Abbott et al. 2020, 2020), with a reported possible association with an EM transient event in an AGN (Graham et al. 2020). If further analysis confirms the association or if additional associations are identified with future LIGO events, such signals would amount to a confirmation of the merger channel studied here.

A caveat is that the accreting BHs in our paradigm should significantly spin up and become aligned, leading to a large effective spin parameter, $\chi_{\text{eff}} \gtrsim 0.9$. In contrast, the GW190521 gravitational-wave signal was associated with a much smaller value of $\chi_{\text{eff}} = 0.08$. This suggests that the BHs either had low spins individually or had large spins that were misaligned with the orbit. If GW190521 did come from the AGN disk formation channel, then there must be some difference in the evolutionary process that we characterized in Section 4.2. A few plausible scenarios are as follows:

1. While we focus on embedded BBHs, individual BHs will migrate and accrete mass from the disk in an analogous way. The inner parsec surrounding the host SMBH may be replete with these single BHs, each of which can potentially grow to sufficiently large masses that they enter the PISN mass gap. These BHs could then pair and merge dynamically later in their life. In this scenario, the BH spins would be randomly oriented, potentially leading to a low $\chi_{\text{eff}}$, regardless of the individual BH spins.

2. The scenario presented here involves no dynamical interactions while the BBH evolves within the disk. We can calculate the interaction time between a BBH and a
t = 1 \frac{n}{\dot{M}_w} \\ \simeq 2 \left( \frac{n}{10^3 \text{pc}^{-3}} \right)^{-1} \left( \frac{a}{100 \text{ au}} \right)^{-2} \left( \frac{v_{k}}{650 \text{ km s}^{-1}} \right)^{-1} \text{Myr},

where we have adopted a cross section, \( \sigma \), derived from the orbital separation, and an interaction velocity \( V \) corresponding to the Keplerian orbit of a BBH at a distance of \( \sim 1 \text{ pc} \) from a \( 10^8 M_\odot \) SMBH. For a stellar density of \( 10^3 \text{ pc}^{-3} \) and an orbital separation of \( 10^2 \text{ au} \), we find an interaction time of \( \sim 2 \text{ Myr} \), shorter than a merger time of \( \sim 10 \text{ Myr} \). These strong dynamical encounters can fundamentally alter the BBHs orbit, reorienting the orbital angular momentum vector, altering the orbital eccentricity, and even evading companions. However, from Figure 6, note that most of the mass accreted by a BH occurs when the orbital separation shrinks to \( \lesssim 10 \text{ au} \), with a much longer interaction time.

3. The accretion flow studied here is initially laminar, while AGN disk models typically assumed some \( \alpha \) viscosity that parameterizes the role of isotropic, turbulent eddies that permeate the AGN disk. The characteristic length-scale of an eddy is \( \sim \alpha H \), and depending on the value of \( \dot{M}_w \), this can be comparable to \( R_w \). If so, the evolution of the embedded binary could be driven more strongly by the accretion of eddies with randomly distributed angular momenta than by the angular momentum supplied from the AGN disk velocity profile. If this is the case, mass accretion might occur similarly, but the spin evolution of the binaries would be stochastic. This is, however, a different flow geometry than that studied here and deserves attention in its own right.

While the relative importance of AGN disks as a formation scenario for BBH mergers may remain unresolved even after future LIGO/Virgo observing runs, the space-based gravitational-wave detector LISA may be sensitive to BBHs in AGN, a prospect that we will consider now. Although many different factors can impact the overall rate of systems observed by LISA, at sufficiently small orbital separations where gravitational-wave radiation dominates the orbital evolution of point masses such as BBHs, general relativity makes a clear prediction for the distribution of BBH orbital separations, \( P(a) \); since the strength of gravitational-wave radiation increases with decreasing orbital separation, \( (\dot{a}/a)_{GR} \sim a^{-4} \) (Peters 1964), a population of equal-mass BBH binaries will have \( P(a) \sim a^{-3} \). However, at sufficiently large separations, the gas drag forces will dominate. Taking \( p \approx 1 \), Equation (13) indicates that \( (\dot{a}/a)_{\text{drag}} \propto a^{5/2} \), when \( a \ll R_w \), producing an orbital separation distribution \( P(a) \sim a^{-7/2} \). By equating these two forces, we can find the critical orbital separation, \( a_{\text{crit}} \), indicating the transition between the two regimes:

\[
a_{\text{crit}} \approx 25.7 \left( \frac{m_1}{30 M_\odot} \right)^{9/13} \left( \frac{c_s}{20 \text{ km s}^{-1}} \right)^{-4/13} \times \left( \frac{\rho_s}{10^{-11} \text{ g cm}^{-3}} \right)^{-2/13}. \tag{18}
\]

This corresponds to an orbital period \( \approx \text{days} \) or a gravitational-wave frequency \( \approx \mu \text{Hz} \). This is, unfortunately, well below the LISA sensitivity curve, which extends down to \( \approx \text{mHz} \). Additionally, the strong dependence on \( a \) of both the gravitational-wave inspiral rate and our drag rate makes the dependence of \( a_{\text{crit}} \) on the gas density and sound speed extremely weak. The binaries may also grow significantly by accretion during their inspiral, which would cause \( a_{\text{crit}} \) to increase further. Equation (18) also assumes a purely hydrodynamic drag rate, while in reality the accretion flow is highly super-Eddington, which likely further decreases the efficiency of drag. Thus, if our \( a \) scalings predicted in Equation (13) hold, we find it unlikely for LISA to be able to detect embedded BBHs.

5.2. Electromagnetic Signatures

In recent years, the possibility that BBH mergers may be accompanied by EM signatures has intrigued the astrophysics community. There have been many proposed mechanisms for producing a signature, including the rapid accretion of relic disks post-merger (Perna et al. 2016; Schröder et al. 2018; S. Schröder et al. 2021, in preparation), the shock-heating of circumbinary disks due to the post-merger recoil (Corrales et al. 2010; Mink & King 2017), and mergers following a single-star progenitor (D’Orazi & Loeb 2018). In Bartos et al. (2017), they also studied the fates of embedded BBHs, and argued that these binaries could produce observable signatures by reaching super-Eddington luminosities via relativistic, beamed outflows. Interest in this scenario has been reinvigorated by the claimed association of an AGN flare with BBH merger GW190521. While this claim is tenuous (Atson et al. 2020), it highlights one of the difficulties of this scenario: even if an embedded BBH produces a luminous signal, it must be distinguished from the dominant AGN emission. We note the recent work, Graham et al. (2020), which compared AGN flaring activity with BBH merger candidates observed by LIGO/Virgo in O3, and found 9 possible associations. These associations depend on the nature of EM signatures from merging embedded BBHs, the definition of what typical AGN flaring is, and the large localization windows from gravitational-wave observations. Given these considerations, we expect that the identification of EM signatures from embedded BBH mergers will only be made robust when we better understand what an associated flare would look like.

Radiation may be produced by an embedded BBH either by accretion during the inspiral phase or transiently immediately following the merger. We turn our attention, first, to the case of a steady-state luminosity. In many cases, the embedded BBH is supplied\(^7\) gas at rates near the Eddington limit or significantly above it (as pointed out by other authors, e.g., Bartos et al. 2017; Stone et al. 2017). However, even in cases where the mass accretion rate can exceed the Eddington limit, the luminosity itself will still be limited because radiation becomes trapped within the flow. In this case, the accretion flow will be a geometrically and optically thick advection-dominated accretion flow (“ADAF”; Narayan & Yi 1994), producing an unusually soft blackbody spectrum. For an Eddington-limited pair of stellar-mass BHs to be observable, they must compete with the host AGN. The relative luminosity of the BBH

\(^7\) We intentionally avoid the term accreted, here, because while we can confidently estimate the gas supply rate via Equation (12), the actual fraction of accreted material on subgrid scales is uncertain.
accretor to the AGN is

\[
\frac{L_{\text{BBH}}}{L_{\text{AGN}}} \sim 10^{-7} \left( \frac{\lambda_{\text{Edd}}^{(\text{SMBH})}}{0.1} \right) \left( \frac{M}{60 \, M_\odot} \right) \left( \frac{M_{\text{SMBH}}}{10^8 \, M_\odot} \right)^{-1}.
\]  

(19)

For typical values of \(\lambda_{\text{Edd}}^{(\text{SMBH})}, M_{\text{BBH}},\) and \(M_{\text{SMBH}},\) this value is extremely small, and the accreting BBH will be indistinguishable from the host AGN unless the two are resolved.

Clearly, if the EM signatures from an embedded BBH are to be observed, it must shine brighter than the Eddington luminosity. The most natural way for this to happen is by producing a collimated, relativistic jet that has its emission beamed toward us (e.g., MacLeod et al. 2014). We can estimate an upper limit on the jet luminosity by assuming the flow is magnetically arrested, producing powerful Blandford–Znajek jets due to large poloidal magnetic fields (Blandford & Znajek 1977; Tchekhovskoy et al. 2011). Assuming this, we take

\[
L_{\text{jet}} = \eta_{\text{MAD}} M c^2
\]

where \(\eta_{\text{MAD}}\) is the jet efficiency for a magnetically arrested disk (MAD). The jet efficiency in MADs commonly reaches \(\sim 200\%–300\%\) by tapping the rotational energy of the BH (Tchekhovskoy et al. 2011), so we assume a characteristic \(\eta_{\text{MAD}} = 2\). Assuming that the radiative efficiency of the accretion flow is \(\eta \approx 0.1\), then the total jet luminosity is

\[
L_{\text{jet}} \approx 20 \left( \frac{\eta_{\text{MAD}}}{2} \right) \left( \frac{\eta}{0.1} \right)^{-1} L_{\text{Edd}}
\]

(21)

If this jet is directed toward us, the emission will be beamed within an opening angle \(\theta \sim 1/\Gamma\), where \(\Gamma\) is the Lorentz factor of the jet. If we measure the flux from this jet at Earth, then the inferred isotropic luminosity is

\[
L_{\text{iso}} \sim \frac{4\pi}{\theta^2} \frac{L_{\text{jet}}}{2} = \Gamma^2 2\pi L_{\text{jet}}
\]

(22)

If we compare this boosted luminosity to the host AGN luminosity, then we find that

\[
\frac{L_{\text{iso}}}{L_{\text{AGN}}} \sim 10^{-3} \Gamma^2 \left( \frac{2}{\eta_{\text{MAD}}} \right) \left( \frac{\eta}{0.1} \right)^{-1} \left( \frac{\lambda_{\text{Edd}}^{(\text{SMBH})}}{0.1} \right)
\]

\[
\times \left( \frac{M}{60 \, M_\odot} \right) \left( \frac{M_{\text{SMBH}}}{10^8 \, M_\odot} \right)^{-1}.
\]

(23)

If we require that \(L_{\text{iso}} \gtrsim L_{\text{AGN}},\) then the corresponding \(\Gamma\) for which this is achieved is \(\Gamma \approx 316\). This is in the upper range of \(\Gamma\) values associated with gamma-ray bursts (Lithwick & Sari 2001; Gehrels et al. 2009), and corresponds to an opening angle \(\theta \approx 0.018\). Even if embedded BBHs could produce jets with this Lorentz factor, we would only see \(\times 0.0002\%\) of them due to the inclination-dependence of the emission. Furthermore, while Bartos et al. (2017) suggest that embedded BBHs form gaps, allowing jets to escape unimpeded, Figure 2 suggests most binaries will not form gaps, and Figure 6 suggests most binaries merge in the outer regions (\(\gtrsim 0.5–1\) pc) of the disk. In these regions, \(M_{\text{inj}} \approx 0.1\), and \(H\) is roughly 2 orders of magnitude larger than the binary sphere of influence. This provides a large column density of material for the jet to traverse, likely resulting in significant mass-loading, which would lower the resulting \(\Gamma\). This suggests that it is difficult to detect a jet associated with an embedded BBH unless it occurs at frequencies separate from the dominant AGN emission. We note that the EM signatures associated with these jets have been studied in detail by two recent papers: Zhu et al. (2021), for the case of embedded neutron stars, and Perna et al. (2021), for jet propagation from both BHs and neutron stars. In both works, they emphasize the issue of distinguishing the transient emission from the AGN emission, and find that the transient can more easily outshine the AGN in the X-ray, infrared, and radio bands.

Finally, we consider the possibility of transient EM signatures that occur during or immediately following the merger. There are two possibilities; the first possibility is that the binary experiences a transient spike in its accretion rate and thus luminosity post-merger (Milosavljević & Phinney 2005). However, these studies were focused on thin circumbinary disks with carved out cavities, and in the thick flows that accompany these binaries, no such spike is likely to exist. We therefore ignore this possibility. The other possibility is that the sudden decrease in gravitational potential at the onset of the merger rapidly shocks the surrounding gaseous envelope, producing an intense photospheric emission.

The sudden mass loss accompanying the BBH merger is associated with a drop of potential energy in the surrounding high-density envelope,

\[
\Delta E = \int_{R_{\text{inj}}} R_{\text{inj}} G M_{\text{merger}} \rho(r) dV.
\]

(24)

We assume that within the BBH sphere of influence, characterized by \(R_{\text{inj}}\) (Equation (11)), the density profile is Bondi-like, i.e., \(\rho(r) \sim \rho_{\infty} (r/R_{\text{inj}})^{-3/2}\) (Shapiro & Teukolsky 1983). This yields a change in potential energy,

\[
\Delta E = 4\pi \rho_{\infty} G M_{\text{merger}} R_{\text{inj}}^2.
\]

(25)

If we assume the flow is Bondi-like in the region of interest and that roughly \(\approx 50\%\) of the BBH mass is lost upon merger, then this change in energy can be rewritten as

\[
\Delta E \approx 8.1 \times 10^{40} \left( \frac{M}{60 \, M_\odot} \right)^3 \left( \frac{\rho_{\infty}}{10^{-18} \, \text{g cm}^{-3}} \right)
\]

\[
\times \left( \frac{c_{\infty}}{25 \, \text{km s}^{-1}} \right)^{-4}.
\]

(26)

If we crudely assume that this energy is lost within a sound-crossing time \(\Delta t = R_{\text{inj}}/c_{\infty},\) then the luminosity produced by the shock (\(\approx \Delta E/\Delta t\)) is

\[
L_{\text{shock}} \approx 1.6 \times 10^{32} \, \text{erg s}^{-1} \left( \frac{M}{60 \, M_\odot} \right)^2 \left( \frac{\rho_{\infty}}{10^{-18} \, \text{g cm}^{-3}} \right)
\]

\[
\times \left( \frac{c_{\infty}}{25 \, \text{km s}^{-1}} \right)^{-1}.
\]

(27)

Since a near-Eddington, \(10^8 \, M_\odot\) host SMBH should have a luminosity of roughly \(L_{\text{AGN}} \approx 10^{45}$–$10^{46}$ erg s$^{-1},\) it is implausible for the luminosity of the post-merger shock to exceed that of the AGN.

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5.3. Feedback

One of the uncertainties in the accretion flow embedding the binaries is the role of feedback. Feedback can be deposited into the gas supply either radiatively or mechanically, through winds and relativistic jets, and can affect the mass accretion rate and efficiency of drag.

If an accreting binary produces a jet, it is possible for the jet to carve out a cavity and suppress the gas supply (e.g., Ramirez-Ruiz et al. 2002). If the force exerted on the ambient medium by a relativistic jet \( \sim L_{\text{jet}}/c \), then it will exert a ram pressure \( \sim L_{\text{jet}}/c^2 \theta^2 \), where \( \theta \) is the opening angle of the jet. By assuming that, as the jet plunges through the medium, it becomes subrelativistic and spreads laterally, we can take \( \theta^2 = 4\pi \). We can set the jet ram pressure equal to the gas pressure of the ambient medium to determine the radius of the jet-medium interface (denoted \( r_j \)),

\[
\frac{L_{\text{jet}}}{4\pi r_j^2 c} = \rho_{\infty} c^3, \\
\Rightarrow r_j = \sqrt{\frac{L_{\text{jet}}}{4\pi \rho_{\infty} c^3 c}}. \tag{28}
\]

To compare this to the relevant physical scale, we can write \( r_j \) in terms of the Bondi radius and physical values,

\[
r_j/R_b \approx 56 \left( \frac{\chi_{\text{jet}}}{20} \right)^{1/2} \left( \frac{M}{60 M_\odot} \right)^{-1/2} \\
\times \left( \frac{\rho_{\infty}}{10^{-18} \text{ g cm}^{-3}} \right)^{-1/2} \left( \frac{c_{\infty}}{25 \text{ km s}^{-1}} \right)^{-1}. \tag{29}
\]

In this relation, we have expressed the jet luminosity in terms of its Eddington ratio, \( \chi_{\text{Edd}} \). For the assumed ambient conditions (typical of the AGN disk at \( \approx 1 \) pc) and Eddington ratio (for an upper limit, see Equation (21)), \( r_j \) is much larger than \( R_b \). This suggests that if a powerful jet is produced by the binary, it should easily carve out a large cavity and impede the gas supply. This could at least transiently diminish the accretion rate and may impose a duty cycle on the growth of the BBH. Alternatively, we can use \( r_j < R_b \) as the condition required for a cavity not to form. Then, the maximum Eddington ratio of the jet is

\[
\chi_{\text{jet}}^{\text{Edd}} < 0.0064 \left( \frac{M}{60 M_\odot} \right) \left( \frac{\rho_{\infty}}{10^{-18} \text{ g cm}^{-3}} \right) \left( \frac{c_{\infty}}{25 \text{ km s}^{-1}} \right)^{-2}. \tag{30}
\]

This is an extremely low jet efficiency and would likely require a BH with a very small spin (Narayan et al. 2003). However, closer to \( \approx 10^{-1} \)–\( 10^{-2} \) pc, the gas density can reach \( \approx 10^{-14} \) g cm\(^{-3}\) (Sirko & Goodman 2003), which would result in a minimum Eddington ratio of 64. This suggests that the ability of a jet to impede the gas supply strongly depends on where the BBH is in the disk; at larger distances, it likely induces a duty cycle in the accretion flow, while at smaller distances the jet will be quickly extinguished.

5.4. Caveats

The results of this work, particularly the evolutionary tracks laid out in Section 4.2, have uncertainties that deserve emphasis. Here are a few essential ingredients required to properly understand the evolution of embedded BBHs:

1. Understanding how drag and mass accretion rates change in the super-Eddington regime is critical to understanding the evolution of embedded BBHs. The ratio of the accretion to inspiral timescale is what determines the final mass of the binary upon merger and will be altered if radiation is realistically taken into account. Additionally, the radiative efficiency affects both the mass accretion rate and the evolutionary timescale. While we use a constant radiative efficiency of \( \eta = 0.01 \), it will be higher in the sub-Eddington regime, but can be even lower in the highly super-Eddington regime. While the details of the evolutionary tracks depend on how the drag in the super-Eddington regime is prescribed, two main features are clear: that embedded BBHs can grow significantly, possibly entering the PISN mass gap before merging, and that the time to merger is \( \lesssim 1 \)–\( 30 \) Myr.

2. The true nature of AGN disks, particularly in the outer Toomre-unstable region, remains an open question. While we favor the profiles of Sirko & Goodman (2003), they use an unspecified pressure source to stabilize this region. This provides a useful model for making predictions but remains a crude estimate, which should be explored further. Additionally, as discussed in Section 4.2, the time-dependence of AGN disks is a major uncertainty. As found in the simulations of Angles-Alcazar et al. (2020), at parsec scales, the disk can reorient itself on \( \approx \) Myr timescales, which may alter the occupation fraction of BHs in the AGN disk or limit the total amount of time the binaries have to evolve in the disk.

3. We assumed that the embedded BBH is allowed to evolve through the AGN undisturbed, but in reality it dynamically couples to the surrounding stellar population. Three-body scattering may merge the BBH faster (Stone et al. 2017) or may reorient the binaries with respect to the surrounding flow, which can alter the mass and spin evolution of the binaries prior to merger. Additionally, as discussed in Section 4.1, a proper understanding of the BH distribution near the SMBH will allow us to better understand how many BHs can be captured dynamically by the AGN and where their initial position in the disk is, which can strongly affect their subsequent evolution.

4. We have neglected multiple effects in our hydrodynamic model. Importantly, these include the noninertial forces of coriolis and centrifugal forces, which will likely modify the flow. Additionally, properly modeling the vertical stratification of the flow is also likely to be important. Even given a more realistic domain, there are large uncertainties pertaining both to the size of the sink radius (which will affect the convergence of the mass accretion rate, e.g., Appendix A) and on the neglected physics. In particular, a proper treatment of the thermodynamics (i.e., via a cooling prescription or better yet radiation transport) will affect the mass accretion and drag rates (e.g., Li et al. 2022), and the feedback
processes (such as relativistic jets or outflows launched on subgrid scales) may also alter the flow substantially.

5.5. Summary

To summarize, we describe the chronology of embedded BBHs as they evolve hydrodynamically through the host AGN disk:

1. Birth. A pair of BHs either are born in the disk (typically \( \gtrsim 0.1-1 \) pc) or are dynamically captured. If the NC harbors a BH cusp, it is possible for the BHs to be initially captured at much smaller radii. As discussed in Section 4.1, it is unclear whether or not a BH cusp will have been formed during the AGN phase of the SMBH, particularly for higher-mass SMBHs where the relaxation timescale of the surrounding stellar population can be longer than the age of the universe. For this reason, we find it more likely that BBHs first become embedded in the outer regions of the AGN disk.

2. Accretion. As embedded BBH orbits prograde with the AGN disk, it will become engulfed in a gaseous wind that has an asymmetric velocity profile in the rest frame of the BBH. This is depicted in Figures 1, 2, and 3, where the characteristics of the accretion flow are primarily defined by \( M_w \), the Mach number of the wind in the BBH rest frame. At large \( M_w \), the ram pressure of the wind dominates the accretion flow, and at small \( M_w \), the accretion flow becomes Bondi-like. In the outer, Toomre unstable regions of an AGN disk, \( M_w \) is often small enough that Bondi accretion provides an adequate approximation (Figure 2). When \( M_w \) is large enough (\( \gtrsim 1 \)), the accretion flow becomes planar, likely resembling a pair of thin mini-disks and potentially opening an annular gap in the AGN disk (Baruteau et al. 2011; Li et al. 2021). Our analysis is invalid in this regime, which likely holds for AGN disks that are geometrically thin and are hosted by lighter (\( \sim 10^5 M_\odot \)) SMBHs. This makes our results most salient for binaries embedded within the disks of heavier (\( \gtrsim 10^7 M_\odot \)) SMBHs. For most regions in the disk, a BBH will accrete at super-Eddington rates (Section 4.2); if the BBH accretes at the Eddington limit for \( \gtrsim 5 \) Myr, it can double its mass.

3. Inspiral. The inspiral of the BBH is most strongly governed by mass accretion and drag, with the drag generally contributing more than the mass accretion (Figure 4). In Section 3.3, we provide analytic expressions with best-fit numerical coefficients that reasonably reproduce the inspiral rate measured in our hydrodynamic simulations. We find that typical BBHs, beginning their journey anywhere between 0.1 and 2 pc in the AGN disk, will merge within \( \lesssim 1-30 \) Myr.

4. Migration. Concurrent to accretion and inspiral, the embedded BBHs will migrate in the AGN disk (typically inwards). For the initial distances we considered (0.5–2 pc), we found that our BBHs migrated only marginally before merging.

5. Multimessenger signatures. In Section 4, we found that BBHs can often double their mass during their inspiral. This can alter the resulting BH mass spectrum detected by LIGO, allowing embedded BBHs to enter the PISN mass gap, as discussed in Section 5.1. In Section 5.2, we studied the possible EM signatures accompanying an accreting BBH or occurring transiently post-merger. In general, it is difficult for the BBH to outshine its host AGN. If the BBH produced a relativistic jet, it would have to have a large Lorentz factor to be detectable, leaving only a very small fraction of these sources detectable due to the inclination-dependence of the beamed emission. Shock-heating in the AGN disk due to the post-merger recoil can produce EM transients, but with luminosities far less than that of the AGN.

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Appendix A

Simulation Convergence

In our simulations, we represented the boundaries of the BHs with an absorbing boundary condition of radius \( R_s = 0.0125 R_h \). In Figure 7, we show the convergence of our accretion and inspiral rates as a function of \( R_s \) by comparing them at \( R_s = 0.0625 \), 0.0125, and 0.025 \( R_h \). A full simulation run at \( R_s = 0.0625 R_h \) is computationally expensive, so we opted to restart the \( R_s = 0.0125 R_h \) simulation, at \( t = 40 R_h/c_s \) (well into steady state) and run it for 10 accretion timescales. For \( R_s = 0.025 R_h \), we also increased the maximum adaptive mesh refinement level, so the number of cells across a sink radius would remain constant. We used this restart approach for \( R_s = 0.025 R_h \) as well, and also compared it to running the full simulation duration at that sink radius, and found no difference in the result. In the specific case of Figure 7, we used the simulation with \( a = 0.1 R_h \) and \( M_w = 0.5 \). We find that the accretion rate typically decreases with the sink radius and that the inspiral rate is relatively constant albeit with higher variability at smaller sink radii. This is expected; as the sink radius decreases, more streamlines are deflected before reaching the boundary, decreasing the accretion rate (e.g., Xu & Stone 2019). On the other hand, gravitational drag results from the cumulative influence of all gas in the binary’s sphere of influence, and so it is less sensitive to the sink radius. This suggests that our results regarding the drag in Section 3.3 should be relatively robust, but the mass accretion rate should be taken as an upper limit.
Figure 8. We plot the accretion and inspiral rates for several of our simulations for the majority of the simulation run-time. Since these rates have strong variability, we averaged both quantities in \(4 \frac{R_s}{c_\infty}\) bins, represented by each scatter point. The plots depict times \(> 20 \frac{R_s}{c_\infty}\) after which the flow is in quasi-steady state.

In Figure 8, we plot the accretion and inspiral rates for the majority of the simulation run-time for each simulation (excluding the intermediate \(M_w = 1\) simulations). Here, both mass and inspiral rates are evolved on the fly in our simulation, and account for all terms in Equation (13). In general, both rates are highly variable, so we opted to plot the averages over both quantities in \(t = 4 \frac{R_s}{c_\infty}\) bins. We depict both rates at times \(> 20 \frac{R_s}{c_\infty}\), after transient spikes in the accretion and inspiral rates have died off, and the flow reaches a steady state.

Appendix B
Derivation of Drag Prescription

Here, we derive the drag formula given by Equation (13), where we follow analogously from Appendix A of Antoni et al. (2019). As a binary companion of mass \(m_1\) orbits, it will capture gas into a dense posterior wake, to which we associate a characteristic radius,

\[
R_{BH,1} = \frac{2Gm_1}{V_1^2 + c_\infty^2 + M_w c_\infty^2}.
\]  

(B1)

This is essentially the Bondi–Hoyle accretion radius, with terms in the denominator for the orbital velocity of the binary companion \((V_1)\), the sound speed \((c_\infty)\), and the characteristic velocity of the shearing wind \(M_w c_\infty\). We can relate this to a characteristic energy dissipation rate, which is defined as the kinetic energy flux entering the accretion cross section \(\pi R_{BH,1}^2\).

\[
E_{BH,1} = \frac{1}{2} \pi R_{BH,1}^2 \rho (V_1^2 + c_\infty^2 + M_w c_\infty^2)^{3/2}.
\]  

(B2)

Since we are considering an equal-mass binary, the total energy dissipation rate of the binary is double this. We also take \(V_1 = \sqrt{GM/m} = \sqrt{GM/\alpha/2}\), where \(M = m_1 + m_2\) is the total
we expect the overall scaling of this expression to hold.

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