

Amplification of superkicks in black-hole binaries through orbital eccentricity

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We present new numerical-relativity simulations of eccentric merging black holes with initially anti-parallel spins lying in the orbital plane (the so-called *superkick* configuration). Binary eccentricity boosts the recoil of the merger remnant by up to 25%. The increase in the energy flux is much more modest, and therefore this kick enhancement is mainly due to asymmetry in the binary dynamics. Our findings might have important consequences for the retention of stellar-mass black holes in star clusters and supermassive black holes in galactic hosts.

I. INTRODUCTION

According to Einstein’s theory of general relativity (GR), gravitational waves carry energy, angular momentum, and linear momentum. In a binary black hole (BH) system the emission of energy and angular momentum causes the orbit to shrink, eventually leading to the merger of the two BHs. The emission of linear momentum imparts a recoil (or *kick*) to the merger remnant [1–3].

Calculations based on post-Newtonian (PN) theory found BH recoil speeds¹ of $\mathcal{O}(100)$ km/s [4–6]. Numerical relativity (NR) simulations, however, show that BH recoils can be more than an order of magnitude larger. This is because the vast majority of the linear momentum is emitted during the last few orbits and merger, where spin interactions are particularly prominent and analytic descriptions within the PN framework become inaccurate. In particular, in 2007 several groups realized that binary BHs with spins lying in the orbital plane and anti-parallel to each other might receive *superkicks* as large as ~ 3500 km/s [7–9]. Subsequent studies found that even larger kicks, up to ~ 5000 km/s, can be reached by further fine-tuning the spin directions [10–13]. Large kicks strongly affect the dominant mode of gravitational waveforms [14–16], and therefore it should be possible to directly measure their effect with future GW observations [17, 18]. Further studies targeted hyperbolic encounters [19] and ultrarelativistic collisions (which are not expected to occur in astrophysical settings) [20], where kicks can

reach 10^4 km/s. We refer to [21–23] for more extensive reviews on the phenomenology of BH recoils.

The occurrence of superkicks has striking astrophysical consequences for both stellar-mass and supermassive BHs. In particular, BH recoils predicted by NR simulations should be compared to the escape speeds of typical astrophysical environments [24].

The stellar-mass BH binaries observed by LIGO and Virgo may form dynamically in globular clusters [25], which present escape velocities in the range 10–50 km/s. These values are smaller even than typical recoil velocities of nonspinning BH binaries [26], which implies that a large fraction of stellar-mass BHs merging in those environments is likely to be ejected [27] (see [28] for a complementary study on intermediate-mass BHs in globular clusters). This may not be the case for environments with larger escape speeds such as nuclear star clusters [29] or accretion disks in active galactic nuclei [30, 31], which might therefore retain a majority of their merger remnants. If able to pair again, the BHs in such an environment can form “second generation” GW events detectable by LIGO and Virgo [32].

The supermassive BH mergers targeted by LISA and pulsar-timing arrays (PTAs) may also be significantly affected by large recoils. Superkicks of $\mathcal{O}(1000)$ km/s exceed the escape speed of even the most massive elliptical galaxies in our Universe. If supermassive BHs are efficiently ejected from their galactic hosts, this decreases their occupation fraction [33] and, consequently, LISA event rates [34, 35]. Spin-alignment processes of both astrophysical [36–39] and relativistic [40, 41] nature are commonly invoked to mitigate this effect.

Recoils are driven by asymmetries in the merging binary [42, 43]: no kick can be imparted if the emission of gravitational-wave energy is isotropic. For instance, an equal-mass nonspinning binary does not recoil by symmetry. Unequal masses or misaligned spins, however, introduce asymmetries in the GW emission. Orbital eccen-

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¹ Speeds are dimensionless in natural units ($c = G = 1$). Therefore, the recoil imparted to a BH does not depend on the total mass of the system.

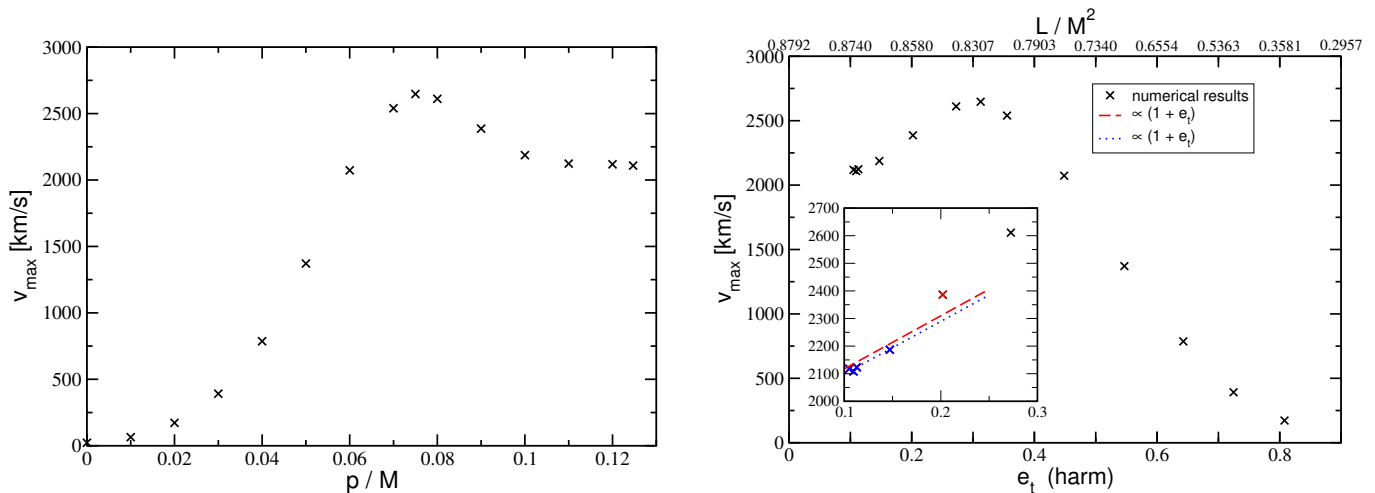


FIG. 1. Superkicks for eccentric binary BHs with equal masses and spins of magnitude $\chi_1 = \chi_2 = 0.596 \simeq 0.6$. *Left*: The maximum kick velocity v_{\max} as a function of the linear momentum parameter. The largest kicks correspond to moderate eccentricity and exceed the quasicircular value by about 25%. *Right*: The maximum kick velocity v_{\max} as a function of the eccentricity parameter e_t estimated in harmonic gauge. Labels on the upper horizontal axis display the corresponding initial orbital angular momentum L/M^2 of the binaries. The inset zooms in on the low-eccentricity regime and shows linear fits $v_{\text{kick}} \propto (1 + e_t)$ obtained from the first 4, blue data points (dotted curve) and also including the 5th, red data point (dashed curve). The increase of the recoil for small eccentricity is compatible with the $(1 + e)$ scaling from close-limit calculations [44].

tricity is a further natural ingredient to enhance the asymmetry of the binary and, consequently, the kick. Early PN estimates show that, for low eccentricities $e \lesssim 0.1$, the kick imparted to nonspinning BHs increases by about $\sim 10\%$, with a scaling proportional to $1 + e$ [44].

In this paper, we investigate for the first time how superkicks are affected by binary eccentricity using NR simulations of the merger. For this purpose, we consider equal-mass binaries with $M_1 = M_2 \equiv M/2$ with BH spins of equal magnitude pointing in opposite directions inside the orbital plane, $\mathbf{S}_1 = -\mathbf{S}_2$. Fixing the dimensionless spin $\chi_i \equiv |\mathbf{S}_i|/M_i^2$ to $\chi_1 = \chi_2 = 0.596$, we generate a sequence of increasing eccentricity by gradually reducing the initial orbital angular momentum L at fixed binding energy from the quasi-circular value to the head-on limit $L = 0$; in practice, we vary for this purpose the initial tangential momentum parameter p of each BH. For a given eccentricity (i.e. fixed L), the kick is known to depend sinusoidally on the initial angle of the two spins relative to the line connecting the BHs [14, 45, 46]. The maximum value of this sine function is the kick reported in Fig. 1 as a function of the linear momentum and of the eccentricity. The significant increase of the maximum kick from about 2100 km/s for approximately quasi circular binaries to 2600 km/s for moderate eccentricities $e_t \sim 0.3$ is the main finding of our study (where e_t is the eccentricity parameter of Ref. [47, 48]). We furthermore show that such an increase holds over a wider range of spin magnitudes and correspondingly raises the maximum superkick in BH binaries to about 4200 km/s, larger than the maximum of ~ 3700 km/s for negligible eccentricity.

The rest of this paper presents our methodology and results in more detail and is organized as follows. In

Sec. II we describe our NR runs; in Sec. III we present our recoil analysis; in Sec. IV we discuss the astrophysical relevance of our findings and possible directions for future work.

II. COMPUTATIONAL FRAMEWORK AND SET OF SIMULATIONS

A. Numerical-relativity setup

The BH binary simulations reported in this work have been performed with the LEAN code [49], which is based on the CACTUS computational toolkit [50, 51]. The Einstein equations are implemented in the form of the Baumgarte-Shapiro-Shibata-Nakamura-Oohara-Kojima (BSSNOK) formulation [52–54] using the method of lines with fourth-order Runge-Kutta differencing in time and sixth-order stencils in space for improved phase accuracy [55]. The wide range of lengthscales is accommodated through adaptive mesh refinement provided by CARPET [56, 57] and we compute apparent horizons with AHFINDERDIRECT [58, 59]. We start our simulations with puncture [60] data of Bowen-York [61] type computed with Ansorg’s spectral solver [62] inside CACTUS’ TWOPUNCTURE thorn. and evolve these using the moving puncture approach [63, 64].

The gravitational wave signal is extracted in the form of the Newman-Penrose scalar Ψ_4 computed from the grid variables [49].

B. Black-hole binary configurations

In this study we consider equal-mass BH binaries in the *superkick* configuration, i.e. the BHs have spins of equal magnitude pointing in opposite directions in the orbital plane.² In practice, we do not compute the dimensionless spins χ_i directly from the Bowen-York spin, because some angular momentum and energy is contained in the spurious radiation of the conformally flat initial data. This energy and momentum is partly accreted onto the BHs and partly radiated to infinity, leading to a brief period of spin adjustment. While negligible for slowly rotating BHs, this effect increases for larger spin parameters and ultimately leads to a saturation at $\chi \sim 0.928$ [65, 66]. In order to obtain a more accurate estimate of χ_i , we monitor the BH spins \mathbf{S}_i using the method described in Ref. [67] and compute the irreducible mass m_{ir} from the apparent horizon during the evolution. The dimensionless spin χ_i can then be computed according to [68]

$$M_i^2 = m_{\text{ir},i}^2 + \frac{|\mathbf{S}_i|^2}{4m_{\text{ir},i}^2}, \quad \chi_i = \frac{|\mathbf{S}_i|}{M_i^2}. \quad (1)$$

As expected from the above description, we observe a brief transient period in all simulations during which χ_i mildly decreases. Throughout this work we report the initial spin as the value at time $t_\chi = 20M$ measured from the beginning of the simulation. By this time χ_i has reached a nearly stationary value, so that the precise value of t_χ does not affect the results. We distinguish this estimate for the initial spin from the value directly obtained from the Bowen-York parameters, which we denote by $\chi_{\text{BY},i}$. The relation between χ_i and $\chi_{\text{BY},i}$ is shown in the fourth and fifth columns of Table I. All simulations presented in this paper have $\chi_1 = \chi_2$.

The net spin is zero in the superkick configurations, resulting in dynamics rather similar to those of nonspinning BH binaries; the main difference is a periodic motion of the orbital plane in the orthogonal (in our case z) direction. This motion of the binary orthogonal to the orbital plane results in a periodic blue and red shift of the gravitational radiation and the net effect of this beaming leads to asymmetric GW emission, especially in the $(\ell, m) = (2, 2)$ and $(2, -2)$ multipoles and, hence, net emission of linear momentum and the ensuing recoil of the post-merger remnant [14, 15]. For fixed initial position $(\pm x_0, 0, 0)$ of the BH binary, the periodic nature of the blue and red shifting of the gravitational radiation furthermore manifests itself in a sinusoidal dependence of the actual kick magnitude on the initial orientation of the spins in the orbital plane [14, 69]. We quantify this orientation in terms of the angle α between the initial

spin of the BH starting at $x > 0$ and the x axis, i.e. this BH has initial spin $\mathbf{S}_1 = S(\cos \alpha, \sin \alpha, 0)$ while the BH at $x < 0$ is initialized with $\mathbf{S}_2 = -\mathbf{S}_1$ [14, 23].

In order to assess the impact of the orbital eccentricity on the magnitude of the gravitational recoil, we have constructed a set of binary configurations guided by the second sequence of equal-mass, nonspinning BH binaries in Table I of Ref. [47]. This sequence starts with a quasicircular binary with initial separation $D/M = 7$ and a tangential linear momentum $p/M = 0.1247$ for each BH, resulting in an orbital angular momentum $L/M^2 = 0.8729$. These parameters determine the binding energy of the binary through $E_b \equiv M_{\text{ADM}} - M$, where M_{ADM} is the Arnowitt-Deser-Misner (ADM) mass [70] of the binary spacetime. We construct a sequence of configurations with increasing eccentricity by gradually reducing the initial linear momentum parameter while keeping the binding energy fixed at $E_b/M = -0.012$. For this choice, the gradual reduction of initial kinetic energy for larger eccentricity implies a larger initial separation, i.e. correspondingly less negative potential energy, and, thus, ensures an inspiral phase of comparable duration irrespective of the eccentricity.

The variation in the initial separation of the BHs requires a minor change in the setup of the computational grid for low- and high-eccentricity binaries. In the notation of Ref. [49] we employ a grid setup given in units of M by

$$\begin{aligned} &\{(256, 128, 64, 32, 16, 8) \times (2, 1), h\}, \\ &\{(256, 128, 64, 32, 16) \times (4, 2, 1), h\}, \end{aligned} \quad (2)$$

respectively, for binaries with $p/M \geq 0.8$ and those with $p/M < 0.8$. Here, the first line specifies a computational domain with six fixed outer grid components of cubic shape centered on the origin with radius 256, 128, 64, 32, 16, and 8, respectively, and two refinement levels with two cubic components each with radius 2 and 1 centered around either hole. The grid spacing is h on the innermost level and successively increases by a factor of 2 on each next outer level. The second line in (2) likewise specifies a grid with five fixed and three dynamic refinement levels.

Unless stated otherwise, we use a resolution $h = M/64$.

In order to accommodate the above mentioned sinusoidal variation of the kick velocity with the initial spin orientation α , we have performed for each value of the linear momentum parameter p a subset of 6 runs with $\alpha \in [0, 180^\circ)$. Due to the symmetry of the superkick configuration under a shift of the azimuthal angle $\phi \rightarrow \phi + 180^\circ$, the recoil will always point in the z direction with $v_x = v_y = 0$ [42, 43]. Furthermore, two binaries with initial spin orientations α and $\alpha + 180^\circ$ will generate kicks of equal magnitude but opposite direction, i.e. $v_z(\alpha) = -v_z(\alpha + 180^\circ)$ [14]. Kick velocities for $\alpha \geq 180^\circ$ can therefore be directly inferred through this symmetry from the simulations performed. For a few selected cases, we have performed additional simulations with $\alpha \geq 180^\circ$; the symmetry is confirmed with accuracy of $\mathcal{O}(0.1)$ % or better.

² We define here the orbital plane as the plane spanned by the initial position vector connecting the BHs and their initial linear momentum – in our case this is the xy plane, and the z axis points in the direction perpendicular to this plane.

C. Measuring the eccentricity

Our sequence of simulations is characterized by the variation of the orbital angular momentum at fixed binding energy. As discussed in detail in Ref. [47], there is no unambiguous way to assign an eccentricity parameter to BH binaries in the late stages of the inspiral. Motivated by the close similarity of the orbital dynamics of (equal-mass) superkick binaries and nonspinning binaries, we follow here the procedure used in [47] to obtain a PN estimate for nonspinning binaries. Specifically, we use Eqs. (20), (25) of Ref. [48], which provide the PN eccentricity parameter e_t for nonspinning binaries. This estimate needs to be taken with a grain of salt as it is only an approximation at the small binary separation during the last orbits before merger, and it ignores the effect of BH spins. Furthermore e_t exhibits an infinite gradient near the quasicircular limit when plotted as a function of the orbital angular momentum, leading to limited precision for values $e_t \lesssim 0.1$. Similarly, in the head-on limit the vanishing of L leads to a formal divergence of the eccentricity parameter and a Newtonian interpretation ceases to be valid (values $e_t > 1$ are possible in this regime). Nevertheless, e_t provides us with a rough estimate to quantify deviations from the quasicircular case and distinguish low-, moderate- and high-eccentricity configurations.

For all simulations, we have computed the following diagnostic variables. The energy, linear and angular momentum radiated in GWs are computed on extraction spheres of coordinate radius $r_{\text{ex}}/M = 30, 40, \dots, 90$ from the Newman-Penrose scalar according to the standard methods described, for example, in [71]. For the physical radiation reported in Table I we exclude the spurious radiation inherent in the initial data by considering only the wave signal starting at retarded time $u \equiv t - r_{\text{ex}} = 50 M$. We also compute the dimensionless spin of the post-merger BH from the apparent horizon [72]. We have confirmed these values using also conservation of energy and angular momentum, which yields agreement to within 0.5 % or better.

D. Numerical accuracy

Our numerical results for the GW emission and the recoil velocities are affected by two main sources of uncertainty: the discretization error and the finite extraction radii for the Newman-Penrose scalar.

We address the latter by extrapolating the GW signal to infinity using a Taylor series in $1/r$ as in Ref. [73]. The results reported are those extrapolated at linear order in $1/r$, and we estimate the error through the difference with respect to a second-order extrapolation. The magnitude of this error is $\sim 2\%$ or less.

In order to assess the error due to finite differencing, we have performed additional simulations of the configuration $p/M = 0.1247$, $\chi_i = 0.596$, $\alpha = 150^\circ$ using grid resolutions $h = M/48$ and $h = M/80$. Figure 2 shows

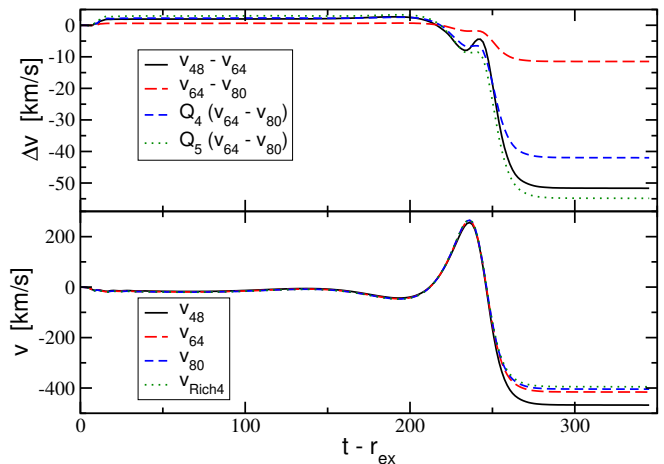


FIG. 2. Convergence analysis for the linear momentum radiated from a binary with $p/M = 0.1247$, $\chi_i = 0.596$, $\alpha = 150^\circ$. The linear momentum obtained for resolutions $h = M/48$, $h = M/64$ and $h = M/80$ is shown in the form of the kick velocity accumulated up to retarded time $u = t - r_{\text{ex}}$ in the bottom panel. The upper panel shows differences between various resolutions, together with rescaling according to fourth- and fifth-order convergence. We estimate the uncertainty from the more conservative fourth-order Richardson extrapolation, and we obtain a numerical error estimate of about 2 % for our standard resolution $h = M/64$.

convergence between fourth and fifth order resulting in a discretization error of about 2 % for the radiated linear momentum. A similar behavior is observed for the radiated energy E_{rad} . We use this value as an error estimate, but note that this is a conservative estimate for the *maximum* kick velocity at fixed eccentricity. The reason is that a considerable part of the numerical error consists in the inaccuracy of the inspiral phase of the binary. This phase error significantly affects the angle α_0 in Eq. (3) below, but has weaker repercussions on the maximum kick v_{max} . In other words, at lower resolution, we will obtain the maximum kick at a “wrong” phase angle α_0 , but still measure this maximum with decent precision. We have verified this expectation by generating a complete sequence for $p/M = 0.1247$, $\chi_i = 0.596$ at low, medium and high resolution. Applying the fit (3) to each of these gives us $v_{\text{max}} = 2098.1, 2108.3$, and 2109.7 km/s, respectively, for $h/M = 1/48, 1/64$, and $1/80$. Since we cannot entirely rule out fortuitous cancellation of errors in this excellent agreement, we keep in the remainder of this work the more conservative 2 % estimate from Fig. 2. Combined with the extrapolation procedure to $r_{\text{ex}} \rightarrow \infty$, we estimate our total error budget as $\sim 4\%$.

III. NUMERICAL RESULTS

The main results of our study are summarized in Table I. For each sequence with prescribed linear momentum p , we list there the initial separation D , orbital angular

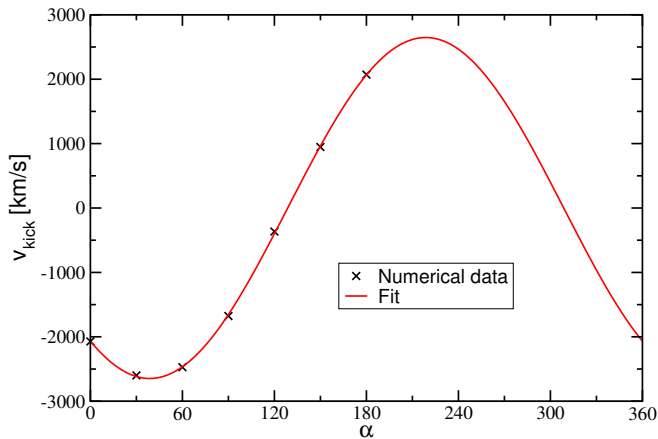


FIG. 3. The recoil velocity computed numerically from the GW signal for $p = 0.075$ is shown as \times symbols for different values of the spin orientation α . The red curve represents the fit obtained from these data according to Eq. (3). Only data for $\alpha < 180^\circ$ are needed to determine the fitting coefficients, due to the symmetry of the superkick configurations.

momentum L , the initial BH spins $\chi_{\text{BY},i}$ and χ_i , eccentricity estimates e_t obtained in ADMTT and harmonic gauge according to Eqs. (20), (25) of Ref. [48], the mean radiated energy E_0 , the maximum kick velocity v_{max} and the dimensionless spin χ_0 of the merger remnant.

A. Impact of the orbital eccentricity

The sinusoidal dependence of the kick magnitude on the initial spin orientation α is illustrated in Fig. 3 for the case $p/M = 0.075$, $\chi_i = 0.596$. The data are reproduced with high precision by a fit of the form

$$v_{\text{kick}} = v_{\text{max}} \times \cos(\alpha - \alpha_0), \quad (3)$$

where, for this specific series, $v_{\text{max}} = 2647$ km/s and $\alpha_0 = 218.7^\circ$. The radiated energy E_{rad} and the final spin, in contrast, vary only mildly (within the numerical uncertainties) with the angle α ; we report average values for these quantities. More specifically, we fit $E_{\text{rad}} = E_0 + E_1 \sin(2\alpha + \alpha_0)$ and report E_0 (and likewise χ_0).

The variation of the kick velocity with eccentricity is visualized in the left panel of Fig. 1, which shows v_{max} as a function of the linear momentum p . We clearly see that the largest kicks are not realized for quasicircular binaries but for moderate eccentricities. A similar effect is apparent for the radiated energy values of Table I, which closely resembles the observation in Table I of Ref. [47] for the nonspinning case. The increase in the recoil velocity, however, is much stronger: for $p/M = 0.75$, the maximum kick exceeds the quasicircular value by about 25% while the largest energy represents a meager 5% increase relative to the quasicircular case. This discrepancy shows that the enhanced kick is not merely due to increased radiation, but also to a higher degree of asymmetry in eccentric binaries.

An increase in the recoil at small eccentricities has already been noticed in the close limit calculations of Ref. [44, 74], which find a $(1 + e)$ proportionality for eccentricities $e \lesssim 0.1$. In the right panel of Fig. 1, we plot the maximum kick velocity as a function of the eccentricity parameter e_t in harmonic gauge (the ADMTT version of e_t would result in virtually the same figure). Due to the diverging gradient of e_t with respect to the orbital angular momentum [47], our data points are limited to $e_t \gtrsim 0.1$, but as shown in the inset of the figure, the data are compatible with the linear growth $\propto (1 + e_t)$ of the close-limit approximation. The two fits shown in the inset have been obtained using either the first 4 or the first 5 data points with the expression $v_{\text{max}} = v_0(1 + e_t)$. The numerical results suggest that above $e_t \approx 0.2$, v_{max} increases even more strongly with e_t before reaching the maximum at $e_t \approx 0.3$, and then decreases for yet higher eccentricity.

B. Impact of the spin magnitudes

The gravitational recoil in superkick configurations is known to increase approximately linearly with the spin magnitudes χ_i . Extrapolating numerical results to maximal spin $\chi_i = 1$ results in a maximal superkick of about 3680 km/s [69] for quasicircular binaries. We will now investigate to what extent nonzero eccentricity can increase this upper limit. In order to keep the computational costs manageable, we focus for this purpose on the $p/M = 0.75$ sequence which maximizes the recoil in our eccentricity analysis for $\chi_i = 0.596$. We cannot rule out that the “optimal” eccentricity maximizing recoil depends on the spin magnitude, so that our analysis should be regarded as a conservative estimate; the largest possible superkick in eccentric binaries may even exceed the value resulting from the analysis below.

We vary the initial spin magnitude χ_i while keeping all other parameters, including the eccentricity e_t , fixed. A convergence analysis for $\chi_i = 0.9$ yields a similar order as in Fig. 2, but demonstrates that higher resolution is needed for these configurations. We use $h = M/80$ for the simulations discussed in this subsection, which results in a discretization error of about 4%. As before, we cover the range of the initial spin orientation by evolving 6 binaries with $\alpha \in [0, 180^\circ)$ for each value of χ_i and fit the resulting v_{kick} according to the sinusoidal function of Eq. (3). The results for these simulations are listed in the lower block of Table I. As expected, the maximum recoil velocity v_{max} increases with the spins χ_i . We display v_{max} as a function of χ_i in Fig. 4, together with a linear fit to model the leading-order dependence of the maximum recoil velocity v_{max} on the spin magnitude χ_i [8, 69]. This fit is given by

$$v_{\text{max}} = [(243 \pm 122) + (4020 \pm 163) \chi_i] \text{ km/s}, \quad (4)$$

and predicts a maximum kick of 4263 ± 285 km/s for extremal spins $\chi_i = 1$. This value exceeds the maximal

p/M	D/M	L/M^2	$\chi_{\text{BY},1} = \chi_{\text{BY},2}$	$\chi_1 = \chi_2$	$e_t(\text{ADMTT})$	$e_t(\text{harm})$	$10^2 E_0/M$	$v_{\text{max}} [\text{km/s}]$	χ_0
0.1247	7.000	0.8729	0.6	0.596	0.1095	0.1096	3.687	2108	0.6815
0.12	7.278	0.8734	0.6	0.596	0.1049	0.1052	3.678	2118	0.6810
0.11	7.932	0.8725	0.6	0.596	0.1130	0.1130	3.664	2123	0.6798
0.10	8.678	0.8678	0.6	0.596	0.1480	0.1472	3.757	2187	0.6808
0.09	9.529	0.8576	0.6	0.596	0.2040	0.2020	3.862	2387	0.6884
0.08	10.493	0.8394	0.6	0.596	0.2758	0.2725	3.656	2611	0.6999
0.075	11.018	0.8264	0.6	0.596	0.3166	0.3124	3.368	2647	0.7010
0.07	11.571	0.8100	0.6	0.596	0.3608	0.3555	3.069	2540	0.7021
0.06	12.754	0.7652	0.6	0.596	0.4567	0.4485	2.258	2073	0.6905
0.05	14.013	0.7007	0.6	0.596	0.5603	0.5467	1.452	1371	0.6539
0.04	15.288	0.6115	0.6	0.596	0.6681	0.6428	0.833	786	0.5862
0.03	16.487	0.4946	0.6	0.596	0.7835	0.7247	0.429	391	0.4839
0.02	17.488	0.3498	0.6	0.596	1.0122	0.8078	0.203	172	0.3467
0.01	18.162	0.1816	0.6	0.596	3.0771	2.0975	0.100	64	0.1813
0	18.398	0	0.6	0.596	∞	∞	0.071	22	0
0.075	11.018	0.8264	0.6	0.596	0.3166	0.3124	3.368	2647	0.7010
0.075	11.018	0.8264	0.65	0.645	0.3166	0.3124	3.383	2849	0.7002
0.075	11.018	0.8264	0.7	0.694	0.3166	0.3124	3.368	3019	0.6990
0.075	11.018	0.8264	0.75	0.742	0.3166	0.3124	3.386	3166	0.6969
0.075	11.018	0.8264	0.8	0.789	0.3166	0.3124	3.330	3479	0.6976
0.075	11.018	0.8264	0.85	0.834	0.3166	0.3124	3.233	3583	0.6960
0.075	11.018	0.8264	0.9	0.876	0.3166	0.3124	3.167	3776	0.6950

TABLE I. Each sequence of simulations is characterized by the linear momentum parameter p and the initial BH separation D (which determine the orbital angular momentum L and the eccentricity of the binary), as well as the initial spins, given here both in the form of the pristine Bowen-York parameters $\chi_{\text{BY},i}$ and of the more accurate horizon estimate χ_i . The remaining columns list: estimates of the eccentricity e_t obtained from PN relations in the ADMTT and harmonic gauge, respectively; the mean radiated GW energy E_0 ; the maximum kick velocity v_{max} ; and the mean spin χ_0 of the remnant BH.

superkick for quasicircular binaries of about 3680 km/s [8, 69] by about 16 %, but falls short of the 5000 km/s maximum for the hang-up kicks reported in Ref. [10]. To the best of our knowledge, the effect of eccentricity on these hang-up kicks has, not yet been explored. The results reported here and the findings of Ref. [44] hint that yet larger recoils may be possible in bound BH binary systems.

IV. CONCLUSIONS

Orbital eccentricity amplifies superkicks. We have presented an extensive series of numerical simulations of merging BHs with spin vectors of magnitude ~ 0.6 in the orbital plane and initially antialigned with each other. We then vary the initial linear momentum of the holes for fixed binding energy, which is equivalent to modifying the initial eccentricity. We find that orbital eccentricity can boost the final recoil by up to $\sim 25\%$. The binaries that receive the largest kick of ~ 2600 km/s have moderate

eccentricity $e_t \sim 0.3$ [47, 48]. For comparison, the maximal kick imparted to a quasicircular binary with the same parameters is ~ 2100 km/s. Our results suggest that the enhanced radiation of linear-momentum is mainly due to the more pronounced asymmetry in the binary's GW emission rather than the mere consequence of a larger energy flux.

An additional series of simulations with fixed eccentricity and varying spin magnitudes allows us to extrapolate these results to maximally rotating BHs. We predict a maximum superkick of at least ~ 4300 km/s, compared to the quasicircular result ~ 3700 km/s. We stress that this estimate is conservative because (i) we did not explore the optimal value of the eccentricity as a function of the spin magnitude, and (ii) we constrain the spins to the orbital plane; partial alignment is known to generate larger recoils [10, 11]. The impact of orbital eccentricity on these *hang-up* kicks with partial spin alignment is a complex task that we leave for future work: the recoil has a more complicated dependence on the eccentricity and the initial spin orientations because of spin precession.

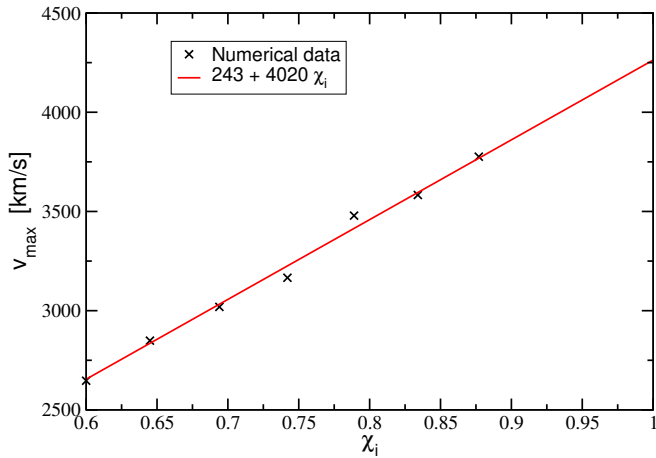


FIG. 4. The maximum recoil velocity v_{\max} for $p/M = 0.075$ as a function of the initial spin magnitude χ_i . The curve represents the linear fit (4).

The amplification of superkicks due to orbital eccentricity may have important consequences for the modeling of GW sources. For the stellar-mass BHs targeted by ground-based interferometers, a nonnegligible eccentricity at merger would be a powerful signature of strong and recent interactions with external bodies (cf. e.g. [75–81]). If BH binaries coalescing in dynamical environments are indeed eccentric, our findings further limit the ability of stellar cluster to retain their merger remnants [32]. For instance, Refs. [82, 83] found that dynamical interactions in globular clusters are a viable formation mechanism to explain multiple generations of eccentric BH mergers. The calculation of the retention fraction, however, does not take into account the significant kick enhancement due to eccentricity that we have found in this work. Given the low escape speed of globular clusters, this amplification may considerably reduce the predicted number of second-generation BH mergers.

For the case of supermassive BH binaries, eccentric sources are commonly invoked to explain current PTA limits. Orbital eccentricity shifts some of the emitted power to higher frequencies, causing a turnover in the predicted spectrum [84–87]. The presence of this feature allows current astrophysical formation models calibrated

on galaxy counts to more easily accommodate the measured upper limits. Our work highlights that kicks may be higher than currently assumed, further reducing the merger rate and the predicted stochastic GW background.

Numerical relativity simulations now provide a thorough understanding of the properties of the BH remnants left behind following mergers of BHs on quasicircular orbits. Efficient and accurate models for final mass, spin, and kick are available and routinely implemented in astrophysical predictions. For eccentric orbits, the additional dimensionality of the parameter space increases the computational resources required to accurately predict waveforms and remnant properties. Comparatively few numerical studies have focused on the eccentric regime in the past [47, 88, 89], but more recently systematic efforts in GW modeling have expanded into the eccentric regime [90]. We hope that our findings have further demonstrated the fertile ground of this class of binaries and that they will spark future work in this direction.

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- [1] W. B. Bonnor and M. A. Rotenberg, *Proceedings of the Royal Society of London Series A* **265**, 109 (1961).
 - [2] A. Peres, *Physical Review* **128**, 2471 (1962).
 - [3] J. D. Bekenstein, *Astrophys. J.* **183**, 657 (1973).
 - [4] M. J. Fitchett, *Mon. Not. Roy. Astron. Soc.* **203**, 1049 (1983).
 - [5] M. Favata, S. A. Hughes, and D. E. Holz, *Astrophys. J.* **607**, L5 (2004), [astro-ph/0402056](#).
 - [6] L. Blanchet, M. S. S. Qusailah, and C. M. Will, *Astrophys. J.* **635**, 508 (2005), [astro-ph/0507692](#).
 - [7] J. A. González, M. Hannam, U. Sperhake, B. Brügmann, and S. Husa, *Phys. Rev. Lett.* **98**, 231101 (2007), [gr-qc/0702052](#).
 - [8] M. Campanelli, C. O. Lousto, Y. Zlochower, and D. Merritt, *Phys. Rev. Lett.* **98**, 231102 (2007), [gr-qc/0702133](#).
 - [9] W. Tichy and P. Marronetti, *Phys. Rev. D* **76**, 061502(R) (2007), [gr-qc/0703075](#).

- [10] C. O. Lousto and Y. Zlochower, *Phys. Rev. Lett.* **107**, 231102 (2011), [arXiv:1108.2009 \[gr-qc\]](#).
- [11] C. O. Lousto and Y. Zlochower, *Phys. Rev. D* **87**, 084027 (2013), [arXiv:1211.7099 \[gr-qc\]](#).
- [12] Y. Zlochower and C. O. Lousto, *Phys. Rev. D* **92**, 024022 (2015), [Erratum: *Phys. Rev. D* **94**, 2, 029901 (2016)], [arXiv:1503.07536 \[gr-qc\]](#).
- [13] C. O. Lousto and J. Healy, *Phys. Rev. D* **100**, 104039 (2019), [arXiv:1908.04382 \[gr-qc\]](#).
- [14] B. Brügmann, J. A. González, M. Hannam, S. Husa, and U. Sperhake, *Phys. Rev. D* **77**, 124047 (2008), [arXiv:0707.0135 \[gr-qc\]](#).
- [15] F. Pretorius, in *Physics of Relativistic Objects in Compact Binaries: From Birth to Coalescence*, edited by M. Colpi *et al.* (Springer, New York, 2009) [arXiv:0710.1338 \[gr-qc\]](#).
- [16] J. D. Schnittman, A. Buonanno, J. R. van Meter, J. G. Baker, W. D. Boggs, J. Centrella, B. J. Kelly, and S. T. McWilliams, *Phys. Rev. D* **77**, 044031 (2008), [arXiv:0707.0301 \[gr-qc\]](#).
- [17] D. Gerosa and C. J. Moore, *Phys. Rev. Lett.* **117**, 011101 (2016), [arXiv:1606.04226 \[gr-qc\]](#).
- [18] J. Calderón Bustillo, J. A. Clark, P. Laguna, and D. Shoemaker, *Phys. Rev. Lett.* **121**, 191102 (2018), [arXiv:1806.11160 \[gr-qc\]](#).
- [19] J. Healy, F. Herrmann, I. Hinder, D. M. Shoemaker, P. Laguna, and R. A. Matzner, *Phys. Rev. Lett.* **102**, 041101 (2009), [arXiv:0807.3292 \[gr-qc\]](#).
- [20] U. Sperhake, E. Berti, V. Cardoso, F. Pretorius, and N. Yunes, *Phys. Rev. D* **83**, 024037 (2011), [arXiv:1011.3281 \[gr-qc\]](#).
- [21] J. Centrella, J. G. Baker, B. J. Kelly, and J. R. van Meter, *Reviews of Modern Physics* **82**, 3069 (2010), [arXiv:1010.5260 \[gr-qc\]](#).
- [22] U. Sperhake, in *Gravitational Wave Astrophysics*, Astrophysics and Space Science Proceedings, Vol. 40, edited by C. F. Sopuerta (2015) p. 185.
- [23] D. Gerosa, F. Hébert, and L. C. Stein, *Phys. Rev. D* **97**, 104049 (2018), [arXiv:1802.04276 \[gr-qc\]](#).
- [24] D. Merritt, M. Milosavljević, M. Favata, S. A. Hughes, and D. E. Holz, *Astrophys. J.* **607**, L9 (2004), [astro-ph/0402057](#).
- [25] M. J. Benacquista and J. M. B. Downing, *Living Rev. Rel.* **16**, 4 (2013), [arXiv:1110.4423 \[astro-ph.SR\]](#).
- [26] J. A. González, U. Sperhake, B. Brügmann, M. Hannam, and S. Husa, *Phys. Rev. Lett.* **98**, 091101 (2007), [gr-qc/0610154](#).
- [27] J. Morawski, M. Giersz, A. Askar, and K. Belczynski, *Mon. Not. Roy. Astron. Soc.* **481**, 2168 (2018), [arXiv:1802.01192 \[astro-ph.GA\]](#).
- [28] K. Holley-Bockelmann, K. Gültekin, D. Shoemaker, and N. Yunes, *Astrophys. J.* **686**, 829 (2008), [arXiv:0707.1334 \[astro-ph\]](#).
- [29] F. Antonini and F. A. Rasio, *Astrophys. J.* **831**, 187 (2016), [arXiv:1606.04889 \[astro-ph.HE\]](#).
- [30] N. C. Stone, B. D. Metzger, and Z. Haiman, *Mon. Not. Roy. Astron. Soc.* **464**, 946 (2017), [arXiv:1602.04226 \[astro-ph.GA\]](#).
- [31] I. Bartos, B. Kocsis, Z. Haiman, and S. Márka, *Astrophys. J.* **835**, 165 (2017), [arXiv:1602.03831 \[astro-ph.HE\]](#).
- [32] D. Gerosa and E. Berti, *Phys. Rev. D* **100**, 041301(R) (2019), [arXiv:1906.05295 \[astro-ph.HE\]](#).
- [33] D. Gerosa and A. Sesana, *Mon. Not. Roy. Astron. Soc.* **446**, 38 (2015), [arXiv:1405.2072 \[astro-ph.GA\]](#).
- [34] L. Blecha and A. Loeb, *Mon. Not. Roy. Astron. Soc.* **390**, 1311 (2008), [arXiv:0805.1420 \[astro-ph\]](#).
- [35] A. Sesana, *Mon. Not. Roy. Astron. Soc.* **382**, L6 (2007), [arXiv:0707.4677 \[astro-ph\]](#).
- [36] T. Bogdanović, C. S. Reynolds, and M. C. Miller, *Astrophys. J.* **661**, L147 (2007), [astro-ph/0703054](#).
- [37] G. Lodato and D. Gerosa, *Mon. Not. Roy. Astron. Soc.* **429**, L30 (2013), [arXiv:1211.0284 \[astro-ph.CO\]](#).
- [38] M. C. Miller and J. H. Krolik, *Astrophys. J.* **774**, 43 (2013), [arXiv:1307.6569 \[astro-ph.HE\]](#).
- [39] D. Gerosa, B. Veronesi, G. Lodato, and G. Rosotti, *Mon. Not. Roy. Astron. Soc.* **451**, 3941 (2015), [arXiv:1503.06807 \[astro-ph.GA\]](#).
- [40] M. Kesden, U. Sperhake, and E. Berti, *Astrophys. J.* **715**, 1006 (2010), [arXiv:1003.4993 \[astro-ph.CO\]](#).
- [41] E. Berti, M. Kesden, and U. Sperhake, *Phys. Rev. D* **85**, 124049 (2012), [arXiv:1203.2920 \[astro-ph.HE\]](#).
- [42] L. Boyle, M. Kesden, and S. Nissanke, *Phys. Rev. Lett.* **100**, 151101 (2008), [arXiv:0709.0299 \[gr-qc\]](#).
- [43] L. Boyle and M. Kesden, *Phys. Rev. D* **78**, 024017 (2008), [arXiv:0712.2819 \[astro-ph\]](#).
- [44] C. F. Sopuerta, N. Yunes, and P. Laguna, *Astrophys. J.* **656**, L9 (2007), [astro-ph/0611110](#).
- [45] L. E. Kidder, *Phys. Rev. D* **52**, 821 (1995), [arXiv:gr-qc/9506022 \[gr-qc\]](#).
- [46] M. Campanelli, C. O. Lousto, Y. Zlochower, and D. Merritt, *Phys. Rev. Lett.* **98**, 231102 (2007), [arXiv:gr-qc/0702133 \[GR-QC\]](#).
- [47] U. Sperhake, E. Berti, V. Cardoso, J. A. González, B. Brügmann, and M. Ansorg, *Phys. Rev. D* **78**, 064069 (2008), [arXiv:0710.3823 \[gr-qc\]](#).
- [48] R.-M. Memmesheimer, A. Gopakumar, and G. Schäfer, *Phys. Rev. D* **70**, 104011 (2004), [gr-qc/0407049](#).
- [49] U. Sperhake, *Phys. Rev. D* **76**, 104015 (2007), [gr-qc/0606079](#).
- [50] Allen, G. and Goodale, T. and Massó, J. and Seidel, E., in *Proceedings of Eighth IEEE International Symposium on High Performance Distributed Computing* (IEEE Press, , 1999).
- [51] Cactus Computational Toolkit homepage: [www.cactuscode.org](#).
- [52] T. Nakamura, K. Oohara, and Y. Kojima, *Prog. Theor. Phys. Suppl.* **90**, 1 (1987).
- [53] M. Shibata and T. Nakamura, *Phys. Rev. D* **52**, 5428 (1995).
- [54] T. W. Baumgarte and S. L. Shapiro, *Phys. Rev. D* **59**, 024007 (1998), [gr-qc/9810065](#).
- [55] S. Husa, J. A. González, M. Hannam, B. Brügmann, and U. Sperhake, *Class. Quant. Grav.* **25**, 105006 (2008), [arXiv:0706.0740 \[gr-qc\]](#).
- [56] E. Schnetter, S. H. Hawley, and I. Hawke, *Class. Quant. Grav.* **21**, 1465 (2004), [gr-qc/0310042](#).
- [57] Carpet Code homepage: [www.carpetcode.org](#).
- [58] J. Thornburg, *Phys. Rev. D* **54**, 4899 (1996), [gr-qc/9508014](#).
- [59] J. Thornburg, *Class. Quant. Grav.* **21**, 743 (2004), [gr-qc/0306056](#).
- [60] S. Brandt and B. Brügmann, *Phys. Rev. Lett.* **78**, 3606 (1997), [gr-qc/9703066](#).
- [61] J. M. Bowen and J. W. York, Jr., *Phys. Rev. D* **21**, 2047 (1980).
- [62] M. Ansorg, B. Brügmann, and W. Tichy, *Phys. Rev. D* **70**, 064011 (2004), [gr-qc/0404056](#).

- [63] J. G. Baker, J. Centrella, D.-I. Choi, M. Koppitz, and J. van Meter, *Phys. Rev. Lett.* **96**, 111102 (2006), [arXiv:gr-qc/0511103 \[gr-qc\]](#).
- [64] M. Campanelli, C. O. Lousto, P. Marronetti, and Y. Zlochower, *Phys. Rev. Lett.* **96**, 111101 (2006), [arXiv:gr-qc/0511048 \[gr-qc\]](#).
- [65] G. B. Cook and J. W. York, Jr., *Phys. Rev. D* **41**, 1077 (1990).
- [66] S. Dain, C. O. Lousto, and R. Takahashi, *Phys. Rev. D* **65**, 104038 (2002), [gr-qc/0201062](#).
- [67] M. Campanelli, C. O. Lousto, Y. Zlochower, B. Krishnan, and D. Merritt, *Phys. Rev. D* **75**, 064030 (2007), [gr-qc/0612076](#).
- [68] D. Christodoulou, *Phys. Rev. Lett.* **25**, 1596 (1970).
- [69] M. Campanelli, C. Lousto, Y. Zlochower, and D. Merritt, *Astrophys. J.* **659**, L5 (2007), [gr-qc/0701164](#).
- [70] R. Arnowitt, S. Deser, and C. W. Misner, in *Gravitation an introduction to current research*, edited by L. Witten (John Wiley, New York, 1962) pp. 227–265, [gr-qc/0405109](#).
- [71] M. Ruiz, M. Alcubierre, D. Núñez, and R. Takahashi, *Gen. Rel. Grav.* **40**, 1705 (2008), [arXiv:0707.4654 \[gr-qc\]](#).
- [72] W. G. Cook, D. Wang, and U. Sperhake, *Class. Quant. Grav.* **35**, 235008 (2018), [arXiv:1808.05834 \[gr-qc\]](#).
- [73] U. Sperhake, B. Brügmann, D. Müller, and C. F. Sopuerta, *Class. Quant. Grav.* **28**, 134004 (2011), [1012.3173 \[gr-qc\]](#).
- [74] C. F. Sopuerta, N. Yunes, and P. Laguna, *Phys. Rev. D* **74**, 124010 (2006), [Erratum: *Phys. Rev. D* **78**, 049901 (2008)], [astro-ph/0608600](#).
- [75] F. Antonini and H. B. Perets, *Astrophys. J.* **757**, 27 (2012), [arXiv:1203.2938 \[astro-ph.GA\]](#).
- [76] J. Samsing, *Phys. Rev. D* **97**, 103014 (2018), [arXiv:1711.07452 \[astro-ph.HE\]](#).
- [77] L. Gondán, B. Kocsis, P. Raffai, and Z. Frei, *Astrophys. J.* **860**, 5 (2018), [arXiv:1711.09989 \[astro-ph.HE\]](#).
- [78] J. Samsing, A. Askar, and M. Giersz, *Astrophys. J.* **855**, 124 (2018), [arXiv:1712.06186 \[astro-ph.HE\]](#).
- [79] M. Zevin, J. Samsing, C. Rodriguez, C.-J. Haster, and E. Ramirez-Ruiz, *Astrophys. J.* **871**, 91 (2019), [arXiv:1810.00901 \[astro-ph.HE\]](#).
- [80] G. Fragione and A. Loeb, *Mon. Not. Roy. Astron. Soc.* **486**, 4443 (2019), [arXiv:1903.10511](#).
- [81] J. Samsing, D. J. D’Orazio, K. Kremer, C. L. Rodriguez, and A. Askar, (2019), [arXiv:1907.11231 \[astro-ph.HE\]](#).
- [82] C. L. Rodriguez, P. Amaro-Seoane, S. Chatterjee, and F. A. Rasio, *Phys. Rev. Lett.* **120**, 151101 (2018), [arXiv:1712.04937 \[astro-ph.HE\]](#).
- [83] C. L. Rodriguez, M. Zevin, P. Amaro-Seoane, S. Chatterjee, K. Kremer, F. A. Rasio, and C. S. Ye, *Phys. Rev. D* **100**, 043027 (2019), [arXiv:1906.10260 \[astro-ph.HE\]](#).
- [84] M. Enoki and M. Nagashima, *Progress of Theoretical Physics* **117**, 241 (2007), [astro-ph/0609377](#).
- [85] A. Sesana, *Astrophys. J.* **719**, 851 (2010), [arXiv:1006.0730 \[astro-ph.CO\]](#).
- [86] E. A. Huerta, S. T. McWilliams, J. R. Gair, and S. R. Taylor, *Phys. Rev. D* **92**, 063010 (2015), [arXiv:1504.00928 \[gr-qc\]](#).
- [87] S. R. Taylor, E. A. Huerta, J. R. Gair, and S. T. McWilliams, *Astrophys. J.* **817**, 70 (2016), [arXiv:1505.06208 \[gr-qc\]](#).
- [88] I. Hinder, B. Vaishnav, F. Herrmann, D. M. Shoemaker, and P. Laguna, *Phys. Rev. D* **77**, 081502(R) (2008), [arXiv:0710.5167 \[gr-qc\]](#).
- [89] I. Hinder, F. Herrmann, P. Laguna, and D. Shoemaker, *Phys. Rev. D* **82**, 024033 (2010), [arXiv:0806.1037 \[gr-qc\]](#).
- [90] A. Ramos-Buades, S. Husa, G. Pratten, H. Estellés, C. García-Quirós, M. Mateu, M. Colleoni, and R. Jaume, (2019), [arXiv:1909.11011 \[gr-qc\]](#).