Massive disc formation in the tidal disruption of a neutron star by a nearly extremal black hole

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Abstract
Black hole–neutron star (BHNS) binaries are important sources of gravitational waves for second-generation interferometers, and BHNS mergers are also a proposed engine for short, hard gamma-ray bursts. The behavior of both the spacetime (and thus the emitted gravitational waves) and the neutron-star matter in a BHNS merger depend strongly and nonlinearly on the black hole’s spin. While there is a significant possibility that astrophysical black holes could have spins that are nearly extremal (i.e. near the theoretical maximum), to date fully relativistic simulations of BHNS binaries have included black-hole spins only up to \( S/M^2 = 0.9 \), which corresponds to the black hole having approximately half as much rotational energy as possible, given the black hole’s mass. In this paper, we present a new simulation of a BHNS binary with a mass ratio \( q = 3 \) and black-hole spin \( S/M^2 = 0.97 \), the highest simulated to date. We find that the black hole’s large spin leads to the most massive accretion disc and the largest tidal tail outflow of any fully relativistic BHNS simulations to date, even exceeding the results implied by extrapolating results from simulations with lower black-hole spin. The disc appears to be remarkably stable. We also find that the high black-hole spin persists until shortly before the time of merger; afterward, both merger and accretion spin down the black hole.

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(Some figures may appear in colour only in the online journal)
1. Introduction

1.1. Motivation

Black hole–neutron star (BHNS) mergers are expected to be a major source for the upcoming advanced interferometric gravitational-wave detectors (Advanced LIGO, VIRGO, and KAGRA [1–3]), with an expected event rate of the order of 10 BHNS yr$^{-1}$ [4]. (Note however that this rate is poorly constrained: pessimistic models predict as little as an event every five years, while with optimistic assumptions, event rates of up to 300 BHNS yr$^{-1}$ are theoretically possible.) If the mergers leave a massive accretion disc around the black hole, they are also a promising setup for short-duration gamma ray bursts (SGRB) [5–7], which might be observed either directly or as an ‘orphan’ SGRB afterglow. If matter is ejected during the merger (either from the tidal tail or a disc wind), other signals are conceivable, including radio afterglows and/or ‘kilonovae’. (See [8] for a discussion of electromagnetic counterparts.)

It is to be expected that the parameters of BHNS systems will vary widely; the mass and spin of the black hole, in particular, can be substantially different from system to system. Population synthesis studies [9, 10] and mass measurements of stellar-mass black holes [11] both favor black-hole masses of 8–11$M_{\odot}$, but they are currently unable to constrain the spin of the black hole, because this is set mostly by the spin acquired at the hole’s formation [9]. This spin at formation will depend on the progenitor of the black hole and the complicated core collapse and accretion dynamics of the black hole’s birth; existing simulations can produce a wide range of spins [12, 13].

The black-hole spin can strongly affect the outcome of the BHNS merger. Spin decreases the innermost stable circular orbit (ISCO) radius and specific angular momentum for prograde circular orbits, and it increases them for retrograde orbits. It is expected that aligned black-hole spins will increase the mass of the post-merger accretion disc. Numerical simulations confirm this expectation, finding that the disc mass increases dramatically with black-hole spin [14–16]. Indeed, for BHNS systems with black-hole mass $\gtrsim 10M_{\odot}$, only those with high dimensionless black-hole spin appear capable of producing significant post-merger accretion discs [17–19]. For a low-mass system ($M_{\text{BH}} \sim 4.5M_{\odot}$), our previous simulations found that the mass remaining outside of the black hole after merger increased from 5% of the neutron-star mass for a system with an initially nonrotating black hole to 39% for a system with a pre-merger black-hole dimensionless spin of $\chi \equiv S/M^2 = 0.9$ [15].

It might appear that this previous study has already effectively covered the range of possible black-hole spins and the consequent post-merger behavior. However, $\chi$ turns out to be a poor measure of how close a black hole is to extremality. For example, a black hole with $\chi = 0.9$ has only about half the rotational energy of an extremal black hole of the same mass. More importantly for the accretion phase, the ISCO radius shrinks significantly between $\chi = 0.9$ and $\chi = 1$ (from 2.3$M_{\text{BH}}$ to $M_{\text{BH}}$ [20]). Thus, increasing the spin in this range can appreciably change the mass of nuclear matter that is able to avoid prompt accretion. It can also increase the efficiency of energy extraction from the accretion disc, $\eta \equiv L/\dot{M}c^2$, where $L$ is the luminosity of the disc, $\dot{M}$ is the accretion rate, and $c$ is the speed of light. Taking $\eta$ to be the specific binding energy at the ISCO, as for a thin disc with no ISCO shear, the efficiency for an extremal black hole is higher than that of a $\chi = 0.9$ black hole by a factor of 2.8. Finally, we note that for $\chi = 0.9$, the ISCO is still outside the ergosphere; the ISCO and, perhaps, significant amounts of disc matter, are only found inside the ergoregion for $\chi > 0.94$. It is conceivable that new behavior is possible for these ‘extreme’ cases of accreting $\chi > 0.94$ black holes.
Another important drawback of our previous study was that the disc was evolved for only a short time after its formation. A thick accretion disc of nuclear matter with a mass roughly 10% of the central black hole could show interesting dynamics. If the disc is sufficiently cool, self-gravity could make the disc unstable to spiral density modes or fragmentation. For thin discs, stability is determined by the Toomre parameter \( Q_T = \frac{\kappa c_s}{\pi G \Sigma} \) \([21, 22]\), where \( c_s \) is the sound speed of the disc, \( \kappa \) is the epicyclic frequency, \( \Sigma \) is surface density, and \( G \) is the gravitational constant. Instabilities in thick self-gravitating discs have been investigated numerically in both Newtonian \([23–27]\) and relativistic \([28–32]\) physics. The instabilities catalogued in these studies often are very sensitive to the angular momentum profile of the disc, especially the runaway \([33, 34]\) and the Papaloizou–Pringle instabilities \([35]\). Since the disc is thick and non-Keplerian, its angular momentum profile can only be known by modeling its formation—the BHNS merger, in our case.

Extreme black-hole spin can also profoundly enhance the electromagnetic signal from the post-merger system. The ISCO binding energy is, as we have mentioned, a strongly nonlinear function of \( \chi \). Models of neutrino-dominated accretion flows have also shown that a high black-hole spin leads to a higher neutrino luminosity, presumably making more energy available for a gamma ray burst \([36–39]\). There is also much more rotational kinetic energy in the black hole that can be extracted via the Blandford–Znajek process (with the luminosity \( L_{BZ} \sim \chi^2 \)) \([40]\).

1.2. Summary and overview

To test the limit of high black-hole spin, high post-merger disc-mass BHNS mergers, we numerically model an extreme case. For the case that we consider, the pre-merger dimensionless black-hole spin is \( \chi = 0.97 \), aligned with the orbital angular momentum. The mass ratio is 3:1, corresponding to a low black-hole mass system. The neutron star has a low compaction \( M_{NS}/R_{NS} = 0.144 \), corresponding to a neutron-star radius \( R_{NS} \sim 14 \text{ km} \) for a 1.4\(M_\odot\) (\(R_{NS} \sim 12 \text{ km} \) for a 1.2\(M_\odot\) neutron star). This system is not intended to represent a typical BHNS binary but is a first step toward exploring the effects of nearly extremal black-hole spin in BHNS mergers. The configuration we consider here illustrates the dynamics possible at one extreme of the astrophysically allowable parameter space, and with this spin and compaction, we can compare to other results to isolate the effects of extreme spin, since the particular combination of mass ratio 3:1 and compaction 0.144 has been particularly well studied by numerical relativity for lower spins, from \( \chi = 0 \) to \( \chi = 0.9 \). The effects of high black-hole spin could be even more important in BHNS with higher mass ratios or with more compact neutron stars, since systems with high enough mass ratio or low enough compaction would need nearly extremal black-hole spin for the neutron star to tidally disrupt outside the horizon (and thus for the emission of an electromagnetic counterpart to the gravitational-wave signal) \([19]\). We intend to explore such configurations in future simulations.

We begin our evolution 5.5 orbits (\(\approx 21 \text{ ms}\)) before merger, when the binary is still in quasicircular orbit. We follow the system until \(\approx 27 \text{ ms}\) after merger; this is 13 orbital periods of the maximum density region of the post-merger accretion disc, long enough to let us clearly see the settling of the accretion disc to near equilibrium.

As expected, the extreme spin leads to a very large post-merger accretion disc. Indeed, a full 60% of the neutron-star matter is able to avoid falling into the black hole during the initial plunge and merger. The vast majority of this matter settles into a massive accretion disc. We note that the early disc mass is significantly larger (by about 50%) than what was found even for a \( \chi = 0.9 \) system, indicating that the disc mass is very sensitive to the black-hole spin when the spin is nearly extremal.
The matter evolution can be divided into phases, each described in more detail below. First, the neutron star is tidally disrupted, leading to an outgoing tidal tail and an ingoing accretion stream. Second, the accretion stream circles the black hole and collides with itself. This shock rapidly heats the matter and disrupts the accretion stream. Third, the disc settles into axisymmetric equilibrium, starting in the inner regions and proceeding outward. This requires redistribution of the orbital energy and angular momentum to allow fluid elements to settle into circular orbit. In the outer disc, the fluid is nonaxisymmetric and has sub-equilibrium average angular momentum. The disc accretes into the black hole slowly on a timescale of hundreds of milliseconds.

The rest of this paper is organized as follows. In section 2, we summarize the numerical methods we use to construct and evolve initial data for a black hole and neutron star in a circular inspiral, with the black hole having nearly extremal spin. Then, in section 3, we present and discuss the results of our simulation, including the emitted gravitational waves and the behavior of the black hole’s apparent horizon (section 3.1), the behavior of the accretion disc that forms following the neutron star’s tidal disruption (section 3.2), and the unbound material ejected by the merger (section 3.3). We briefly conclude in section 4.

2. Methods

To construct initial data for a BHNS binary with nearly extremal black-hole spin, we follow the methods of [41], which are in part motivated by the methods for constructing binary-black-hole initial data with nearly extremal spins in [42]. We find that these methods, which had only been tested for spins up to \( \chi = 0.9 \) so far, can be applied without any modification to generate initial configurations with significantly higher black-hole spins—although obtaining the same accuracy as for lower spins requires the use of a finer numerical grid close to the black hole. Here, we consider a black hole with initial spin \( \chi = 0.970 \). Because the initial conditions are not perfectly in equilibrium, the spin slightly decreases during the initial relaxation, falling to \( \frac{a}{M_{\text{BH}}} = 0.967 \) by time \( t/M = 100M_\odot = 1.5 \text{ ms} \) after the beginning of the simulation (cf figure 2). We use a polytrope

\[
P = \kappa \rho^{\Gamma} + \tilde{T} \rho,
\]

where \( P \) is pressure, \( \rho \) is the rest-mass density, and \( \tilde{T} \) is a fluid variable related to the physical temperature. As in [15], we choose \( \Gamma = 2 \) and choose \( \kappa \) so that the compaction of the star is \( M_{\text{NS}}/R_{\text{SS}} = 0.144 \). Finally, we reduce the orbital eccentricity of the binary by iteratively solving for its instantaneous angular velocity and inspiral rate [43]. The last iteration, which is used throughout this paper, has eccentricity \( e = 0.005 \pm 0.001 \). The orbital trajectories of the black hole and neutron star before tidal disruption (i.e., during the first \( \approx 20 \text{ ms} \) of evolution after the start of the simulation, which corresponds to approximately 5.75 orbits), are shown in figure 1.

The coupled evolution of Einstein’s equations of general relativity and of the relativistic hydrodynamics equations is performed with the SpEC code [44], using our two-grid method [45]: Einstein’s equations are evolved in the generalized harmonic formalism [46], using pseudospectral methods and excision of the black-hole interior, while the neutron-star fluid is evolved on a separate finite-difference grid covering only the regions in which matter is present. The general relativistic equations of hydrodynamics are evolved in conservative form using a second-order finite volume scheme and high-order shock capturing methods: the approximate Riemann problem is solved on cell faces using the weighted essentially non-oscillatory reconstruction algorithm [47, 48] and Harten, Lax, and van Leer fluxes [49]. Both
Figure 1. The trajectory of the apparent horizon center (thin, black curve) and center of mass of the neutron star (thick, red curve) before tidal disruption (i.e., during the first \( \approx 20 \) ms after the start of the simulation). The orbital eccentricity in this case is \( e = 0.005 \pm 0.001 \).

Figure 2. The evolution of the black-hole spin \( \chi := S/M_{\text{sch}}^2 \) (top panel), Christodoulou mass \( M_{\text{chr}} \) (middle panel), and irreducible mass \( M_{\text{irr}} \) (bottom panel) as a function of time (with times \( t \) shown relative to time \( t_{0.8} \), the time when the baryonic mass outside the hole has fallen to 80\% of its initial value) for a BHNS merger with an initial black-hole spin of magnitude \( \chi = 0.97 \).

Sets of equations are evolved in time using the third-order Runge–Kutta method with adaptive choice of the time step.

To handle the black hole’s high spin during the evolution, we find that we can apply the same techniques that have previously been applied to achieve fully relativistic simulations of
Figure 3. The normalized Einstein constraint violation (as defined in equation (71) of [46]) as a function of time (with times $t$ shown relative to time $t_{0.8}$, the time when the baryonic mass outside the hole has fallen to 80% of its initial value) for low (dashed, black), medium (solid, red), and high (dotted, green) resolutions. In each simulation, the initial black-hole spin is aligned with the orbital angular momentum and has a magnitude of $\chi = 0.97$. The constraints rise around time $t - t_{0.8} = -14$ ms because the apparent horizon is moving farther from the excision surface. After modifying the simulation’s algorithm so that the excision surface and apparent horizon would nearly coincide (at approximately $t - t_{0.8} = -10$ ms), the constraint violation drops until the neutron star begins to disrupt.

We simulated three resolutions through inspiral and tidal disruption, evolving $\approx 27$ ms (low, medium resolution) and $\approx 11$ ms (high resolution) after merger. During the inspiral, these resolutions correspond to approximately $57^3$, $65^3$, and $72^3$ spectral gridpoints, respectively, and $100^3$, $120^3$, $140^3$ finite-difference gridpoints, respectively. After the neutron star begins to tidally disrupt, we change our domain decomposition, using low, medium, and high resolutions with approximately $98^3$, $111^3$, and $124^3$ spectral gridpoints, respectively, and $140^3$, $160^3$, $180^3$ finite-difference gridpoints, respectively. From then on, we use spectral adaptive mesh refinement to dynamically adjust the resolution of each spectral subdomain. Note that for all evolutions, since the black-hole spin is aligned with the orbital angular momentum, when solving the relativistic hydrodynamic equations we only evolve the region above the orbital plane, imposing a symmetry condition on the orbital plane and reducing the number of gridpoints in the direction normal to the orbital plane by a factor of 2.
We note that we do not find it necessary or advantageous to change to a damped harmonic gauge condition, as was done in the high-spin BHBH simulations in [50, 51]. Instead, we use the same gauge conditions used in previous two-grid BHNS evolutions at the same mass ratio [15].

These simulations are computationally expensive. The low, medium, and high resolutions cost approximately $9.8 \times 10^4$, $1.4 \times 10^5$, and $1.5 \times 10^5$ CPU hours, corresponding to approximately 70, 105, and 116 days of wallclock time. (Note that because we did not continue the high-resolution simulation for as long after disruption, its expense is only slightly higher than that of the medium resolution.)

3. Results

In this section, we present our main numerical results. First, in section 3.1, we show results for the spacetime’s evolution. Then, we discuss the properties of the resulting accretion disc (section 3.2) and ejecta (section 3.3) after the black hole tidally disrupts the neutron star.

3.1. Spacetime

To characterize the behavior of the curved spacetime, we begin by examining the mass and spin of the black hole as measured on the apparent horizon. Figure 2 shows, for each resolution, the black-hole dimensionless spin \( \chi := \frac{S}{M_{\text{ch}}^2} \) (top panel), Christodoulou mass \( M_{\text{ch}} := \sqrt{M_{\text{irr}}^2 + \frac{S^2}{4M_{\text{irr}}^2}} \) (middle panel), and irreducible mass \( M_{\text{irr}} \) (lower panel) as functions of time. Following the initial relaxation, the masses increase slightly (from \( M_{\text{irr}} = 3.31 \) and \( M_{\text{ch}} = 4.20 \) to \( M_{\text{irr}} = 3.33 \) and \( M_{\text{ch}} = 4.21 \)) while the dimensionless spin decreases slightly (from \( \chi = 0.970 \) to \( \chi = 0.967 \)). The masses and spins then remain constant in time to within our numerical accuracy until approximately 20 ms after the start of the simulation, when the neutron star begins to disrupt. During the next \( \approx 1 \) ms, the masses sharply increase, while the dimensionless spin sharply decreases. The hole’s masses then continue to slowly increase as its dimensionless spin slowly decreases.

While all three resolutions exhibit the same qualitative behavior, we were unable to demonstrate convergence of the black hole’s mass and spin through the end of the simulation. Shortly after disruption begins, in the high-resolution simulation the black hole experiences an anomalously large decrease in the dimensionless spin; through experimentation, we found this decrease to be sensitive to the details of dynamic regridding (i.e., the algorithm used to resize the finite-difference grid as the region containing matter expands—cf section II B 1 of [15]). This could be a consequence of the extremely large amount of mass remaining in the disc. Apparently, there are certain special distributions of matter on our grid that, upon a regrid, result in unexpectedly large interpolation error near the excision surface and a corresponding small but abrupt change in the evolution. Since our primary concern in this paper is the behavior of the disc and ejecta, and since this \( O(\%) \) effect does not qualitatively change our results,

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7 We are confident that this anomaly does not qualitatively affect our results or our conclusions, for the following reasons: (i) cf the top panel of figure 10, which shows the result for all three \( \chi = 0.97 \) resolutions: all three resolutions show the same qualitative behavior. (ii) The \( O(\%) \) anomalous behavior seen in the high-resolution black-hole quantities around time \( t_8 \approx 1 \) ms coincides with a period of the evolution when the finite-difference grid rapidly, repeatedly regrids. The effect of these regrids can also be seen in small discontinuous changes in disc properties such as the baryon rest mass. Disabling regridding did eliminate the anomalous behavior in the high-resolution run. Unfortunately, long evolutions cannot be done without regridding. The frequent regridding behavior that seems to cause the problem is not seen in the low and medium resolution runs. (iii) The low and medium resolutions, which do not display the frequent regridding problem, show good quantitative agreement with each other.
we have chosen to leave our efforts to resolve the black-hole spin after disruption (perhaps by improving our regridding method) for future work.

Figure 3 further demonstrates our difficulty in obtaining convergence after tidal disruption. Before disruption (i.e. during the first 20 ms of the simulation), the normalized constraint violation appears to be convergent, as expected. The constraints do rise around $t - t_{0.8} \approx -14$ ms as the excision surface falls far inside the apparent horizon, but this trend reverses after a modification of our method to keep the excision surface close to the apparent horizon\footnote{Because of the high cost of our simulations (cf section 2), we chose to adjust our algorithm during the simulations—at a time early enough that the constraint violation is still small—rather than to repeat the simulations from the beginning. We observe no significant effect corresponding to this adjustment in our results; therefore, we are confident that this adjustment does not affect our conclusions. However, had we continued to allow the excision surface to remain far inside the apparent horizon, constraint growth would eventually (before the time of merger) have caused the simulations to fail.}. After interpolating to our higher resolution domain (around $t - t_{0.8} \approx -2$ ms), the constraints sharply fall at first, but then sharply spike during the merger. The subsequent bump in the constraints corresponds to the small bump in the dimensionless spin and the small drop in irreducible mass seen around $t - t_{0.8} \approx 7$ ms in figure 2.

We conclude this section by examining the dominant mode of the gravitational waveform (figure 4). Because we plot $\Psi^4 = (d^2/dt^2)(h_z - ih_x)$ instead of the wave amplitude $h$, the spurious ‘junk’ gravitational radiation emitted during the initial relaxation is clearly visible. Each resolution shows a small, secondary burst of gravitational waves after the primary wave has terminated. The phase of this secondary burst cannot be resolved in the simulation: as is typical for numerical simulations of BHNS or binary-neutron-star mergers, the phase error in the post-merger waveform is on the order of a few radians. At all resolutions, the burst carries an energy $E_{\text{burst}} \sim 10^{-4} M_{\odot} c^2$ in a time $t_{\text{burst}} \sim 0.75$ ms and at frequencies $f_{\text{burst}} \sim 2$–3 kHz. There are no easy ways to associate this burst in the gravitational-wave signal extracted at large radius with a specific feature of the forming accretion disc. Additionally, sharp features due to shocks will only converge at first order with numerical resolution, and their effects are likely to be poorly resolved. Nonetheless, an estimate from the quadrupole formula shows

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Figure 4. The real part of the dominant $\ell = m = 2$ mode of the gravitational waveform $rM_{\odot} \Psi^4$, where $r$ is the radius of the sphere on which the gravitational waveform is extracted (here the outermost sphere, with a radius of $r = 1540M_{\odot}$), $M_{\odot} = 1M_{\odot}$, and $\Psi^4$ is the Newman–Penrose scalar corresponding to outgoing gravitational waves. The inset zooms in on the part of the waveform enclosed by the dashed box.
Figure 5. The evolution of the baryonic mass $M_b$ as a function of time $t$ for the medium resolution BHNS merger with an initial aligned black-hole spin of magnitude $\chi = 0.97$. The solid, green line is the total baryonic mass $M_b$ outside the black hole, the dashed, black line is the total rest mass $M_{\text{grid}}$ remaining on the finite-difference computational grid, and their difference $M_{\text{out}}$ is the dotted, red line. All masses are normalized by the baryonic mass $M_{b,0}$ at the start of the simulation. Times are shown relative to time $t_{0.8}$, the time when the baryonic mass outside the hole has fallen to 80% of its initial value.

that these numbers are consistent with emission from an overdense region with excess mass $M_{\text{emitter}} \sim 0.01\text{–}0.03 M_{\odot}$, orbiting in the inner region of the disk ($R \sim 30$ km). For the massive disc observed here, the presence of such an asymmetry immediately after merger is quite natural. We thus expect these features to be qualitatively correct, even though the detailed properties of the waveform are not numerically resolved.

3.2. Accretion disc

After tidal disruption, the nuclear matter forms an outgoing tidal tail and an ingoing accretion stream. At a time $t = t_{0.8} + 10$ ms (10 ms after the time when 80% of the initial baryonic mass remains outside the hole), $\approx 23\%$ of the initial baryonic mass has moved outward and left the finite-difference grid, and $\approx 31\%$ of the initial baryonic mass remains on the finite-difference grid (figure 5). Of the mass that has left, $\approx 17\%$ appears to be unbound and quickly leaves the grid (see section 3.3), while another $\approx 6\%$ is weakly bound but extends far enough from the black hole to leave the finite-difference grid. This material would have fallen back onto the disc at later times ($\sim 10^2$ ms) and its loss is one reason we limit our simulations to the early to middle ($\leq 30$ ms) post-merger evolution.

During this tidal disruption phase, the matter remains cold and degenerate. The infalling stream of matter circles around the black hole and collides with itself, producing an intense shock that travels through and disrupts the incipient accretion flow. Nearly all of the entropy generation in the evolution occurs during this crucial millisecond, which could be labeled the shock phase of the evolution. Shock heating renders the matter completely nondegenerate, and hot matter flows to the outer-disc region ($r > 100$ km) over the next $\sim 10$ ms. By the end of this time, a clear hierarchy of velocities establishes itself through the disc out to 200 km. (See figure 6.) Inside this region, the azimuthal velocity dominates and is roughly twice the sound speed throughout ($\Omega r \approx 2c_s$). The orbital period sets the dynamical timescale; we may define
Figure 6. A break-up of the components of the fluid velocity, shown at three times after merger, as a fraction of the speed of light. Also shown is the sound speed $c_s$. The extents on the horizontal and vertical axes are the same in all three panels. The radial velocity itself is negative (infall).

Figure 7. The density profile of the accretion disc, shown at times $t = 8$, 17, and 26 ms after merger. For each time, two lines are shown. The top line is the vertically integrated surface density $\Sigma$ at each radius, averaged over all angles. The bottom line is the rms deviation from the average due to nonaxisymmetry. A dot on the bottom line indicates the location of $r_{\text{settle}}$ at this time. In this and all other radial plots, the radius shown is constructed from the proper circumference at the equator.

a ‘settling radius’ $r_{\text{settle}}(t)$ as the radius at which gas has had time to complete exactly one orbital period since merger. Since orbits will not fully circularize after one period, this radius provides only a rough sense of which parts of the disc have had time to settle. By the end of our evolution (30 ms post-merger), $r_{\text{settle}} = 190$ km. Outside $r_{\text{settle}}$, the disc does not have time to reach equilibrium. The radial velocity is a non-negligible fraction of the azimuthal velocity in the outer disc (see figure 6), but the radius at which this infall becomes supersonic recedes with $r_{\text{settle}}$.

In figure 7, we show the surface density $\Sigma$ at three times during the settling phase. The density peaks at a circumferential radius of around 30 km; the location of the peak
remains fixed, while the overall density slowly decreases everywhere as the mass is depleted by accretion. Also shown in this plot is the RMS deviation in $\Sigma$ at each radius due to deviations from axisymmetry. We see that the inner disc is quite axisymmetric, but that deviations from axisymmetry reach unity at around $r_{\text{settle}}(t)$. A Fourier analysis of $\Sigma$ at late times shows that the deviation from axisymmetry primarily subsists in trailing spiral features, although localized twists in which a mode is leading over a small radial range are occasionally seen in some modes. The dominant mode in the outer disc is $m = 1$. Only in the inner region can these modes be regarded as linear perturbations of an equilibrium system; here they provide a natural way to understand the extraction of angular momentum from the inner disc which must occur if the low-angular-momentum outer disc (see below) is to settle. In addition to these smooth modes, equilibrium in $r_{\text{settle}}/2 < r < r_{\text{settle}}$ is sometimes disturbed by weak, localized sharp features, presumably weak shocks where gas on eccentric orbits hits the more settled disc. These features are most clearly seen in the pressure force (see below).

In figure 8, we plot the specific orbital energy ($E = u_t$) and angular momentum ($L = -u_{\phi}/u_t$) on the equator as a function of radius. For comparison, we plot the expected curves for circular orbit geodesics in this spacetime metric. The geodesic $E$ and $L$ are the expected orbital parameters for a stationary disc taking into account the disc’s self-gravity (which is included in the numerical metric) but not pressure forces. Except in the very inner disc, we see significant deviations from geodesic orbits, as would be expected for a thick disc. Figure 8 also includes $E$ and $L$ for equilibrium circular orbits, i.e. stationary circular orbits taking into account the pressure force for the given fluid profile. These agree with the actual $E$ and $L$ up to about 100 km. Beyond this, the energy curves continue to agree (accounting for deviations due to localized shocks), but the angular momentum profiles flatten more rapidly than the equilibrium curve (although the latter seems to become somewhat noisy) and become
Figure 9. The entropy profile of the accretion disc, shown at times $t = 8, 17, \text{and } 26 \text{ ms after merger.}$ The entropy for our $\Gamma$-law equation of state is $S = \log \left( \frac{P(\rho, T)}{P(\rho, 0)} \right)$. For each radius, a density-weighted average is carried out over the $\phi$ and $z$ coordinates.

sub-equilibrium. The gas at these radii was ejected (from the tidal tail or the post-merger shock) into eccentric orbits, as seen from the $L$ deficit and the nonzero radial velocity (cf figure 6). Equivalently, the nonaxisymmetric modes carry net negative angular momentum. Pressure support becomes especially important in the outer disc, as can be seen from the deviation between geodesic and equilibrium curves. We also have observed that the disc thickness, height $H$ divided by radius $r$, increases with radius from about $H/r \approx 0.2$ in the inner disc to about $H/r \approx 0.35$ in the outer disc.

There are no clear signs of instability or turbulence in the disc, a fact that itself should be explained. The lack of obvious global corotation instability of the kind found in some other massive disc simulations is a natural consequence of the angular momentum profile. Instability is only expected for $L \propto r^n, n < 2 - \frac{\sqrt{3}}{3}$ [35] with growth rate decreasing rapidly as $n$ increases above zero, and our disc has $n \approx 0.3$ in the equilibrated region. This region is also stable against the effects of self-gravity, with Toomre parameter no lower than $Q_T \approx 400$. This stability is a consequence both of the shear and the thermal pressure. Radiative cooling may later remove some of the disc’s heat, but even removing all of it would leave $Q_T \sim 10$. Shear and thermal energy also conspire to protect the disc from convective instability. In Newtonian physics, the Solberg–Hoiland stability criteria are

$$r^{-3} \nabla^2 L^2 - (C_p \rho)^{-1} \nabla P \cdot \nabla S > 0$$

(2)

$$\nabla \cdot P (\nabla L^2 \nabla S - \nabla L^2 \nabla S) < 0,$$

(3)

where $\rho$ is the density, $S$ the entropy, $P$ the pressure, and $C_p$ the heat capacity. The relativistic version of these equations is slightly more complicated but basically similar [53]. Figure 9 shows the average specific entropy as a function of radius during the settling phase. The entropy has a minimum close to the density maximum, while the gas on the edges is hotter. Thus $\nabla P \cdot \nabla S < 0$ except near a small ring near $r = 40 \text{ km}$, and this feature is stabilized by the strong shear. In fact, the epicyclic frequency is much greater than the Brunt–Väisälä frequency $N$ out to $r \approx 80 \text{ km}$. Beyond this, the $S$ gradient steepens, the $L$ gradient flattens, and buoyancy becomes an important restoring force. Vertical convection is also suppressed by a strong positive $\nabla_s S$. The disc will be unstable to the magnetorotational instability, since
\[ \frac{dQ}{\ln r + N^2} < 0 \] throughout, but this does not appear in our simulations, which do not include magnetic fields.

The disc evolves slowly due to the settling of the outer disc and accretion of gas in the inner disc (the two perhaps connected by transfer of angular momentum). The late-time accretion rate is \( M \approx 2M_\odot \text{s}^{-1} \), giving a lifetime of \( \tau = M/M \sim 200 \text{ ms} \). Before this time, we expect magnetorotational effects to become important. The most unstable modes will grow on the dynamical timescale (ms), and then magnetic turbulence will transport angular momentum. Assuming this acts like an \( \alpha \) viscosity with \( \alpha = 0.01–0.1 \), the effective viscous timescale will be of order \( r^2/(\alpha c_s H) \sim 10–100 \text{ ms} \) for radius \( r \), height \( H \), and sound speed \( c_s \) characteristic of the central torus. Accretion increases the spin and mass of the black hole, but since the accreted matter has low angular momentum (\( L \approx 10 \), see figure 8), the total effect on the dimensionless spin \( \chi = S/M^2 \) will be small even for accretion of most of the disc mass. (At late times in our evolution, \( \chi \) is slowly decreasing.)

Our simulation does neglect several effects that could influence the accretion disc’s later evolution. A magnetic field could produce a jet, powered either by the accretion flow or the high black-hole spin. It would certainly render our disc magnetorotationally unstable. Although we find that MRI turbulence is not needed for strong early-time accretion, it will still alter the disc evolution in important ways. Second, neutrino radiation will cool the gas and make it less neutron-rich. Assuming a simple ideal gas plus radiation pressure equation of state, we infer that the temperature in the disc is of order 10 MeV. (In this model, the disc is gas pressure-dominated everywhere except the inner edge.) Using a simple diffusion approximation, one can estimate a luminosity \( L_\nu \sim 10^{52} \text{ erg s}^{-1} \), which could deplete the disc’s thermal energy on a timescale of \( \sim 10^2 \text{ ms} \). More likely, the cooling will balance the dissipative heating introduced by magnetic turbulence. Neutrino radiation could itself induce unstable entropy or composition gradients and drive convection. Finally, our grid limitations have forced us to allow gravitationally bounded, outward-moving material to leave the finite-difference grid. This matter should eventually fall back onto the disc, perturbing and perhaps shocking it. Tracking the fallback matter will be a computational challenge for future simulations that wish to evolve discs like this one further in time than was done here.

We should note that the amount of matter remaining outside of the black hole after merger is not only larger than in any previous studies of BHNS mergers: it is also significantly above the predictions obtained by extrapolating results for lower spin black holes in the high-spin regime. The remnant mass for lower spin systems is generally well approximated by \[ M_{\text{rem}} M_\odot = 0.288 \left( \frac{3M_{\text{BH}}}{M_{\text{NS}}} \right)^{1/3} \left( 1 - 2 \frac{M_{\text{NS}}}{R_{\text{NS}}} \right) - 0.148 \frac{R_{\text{ISCO}}}{R_{\text{NS}}}, \] with \( R_{\text{ISCO}} \) the radius of the ISCO around the black hole. Figure 10 shows the predicted remnant mass for BHNS mergers at \( q = 3, M_{\text{NS}}/R_{\text{NS}} = 0.144 \), together with numerical results at \( \chi = (0, 0.5, 0.75, 0.9, 0.97) \) [14, 18]. We see that the analytical predictions perform well for low-spin black holes (\( \chi < 0.9 \)), while strong deviations are visible for the higher spin configurations. This is not surprising: the analytical results are only supposed to be valid for relatively low-mass remnants (\( M_{\text{rem}} < 0.2M_\odot \)). But it does emphasize the need for further studies of BHNS mergers at high black-hole spin in order to properly model the characteristics of the post-merger remnant in that regime—especially considering that the most energetic SGRBs observed might require the presence of such high-mass accretion discs [54].

3.3. Ejecta

In addition to the formation of an accretion disc, the tidal disruption of the neutron star can lead to the ejection of a significant amount of unbound material. This material is initially
neutron-rich, and its radioactive decay can lead to detectable optical afterglows ('kilonovae’ [55, 8]), as well as contribute to the formation of heavy elements (r-process nucleosynthesis).

The deceleration of the ejecta in the interstellar medium can also cause emission in the radio band [8]. For binary neutron stars [56] and low-spin BHNS [57] binaries, general relativistic simulations show that only a small fraction of the neutron-star material (<1%) is unbound, with kinetic energies of $10^{49}$–$10^{51}$ erg [56].

Measurements of the unbound mass are fairly inaccurate in our code, especially for very energetic and massive ejecta. This is due to the fact that the fluid equations are only evolved up to $20M_{BH}–30M_{BH}$ from the center of mass of the system, while measurements of the properties of the ejecta are best done far away from the black hole and in low-density tidal tails (the condition $u_t < -1$, used to determine whether material is unbound, is only valid for a stationary metric and pressureless fluid). Additionally, the accuracy of the evolution far away from the black hole is lower than in the forming disc, where most gridpoints are concentrated. The development of techniques allowing us to efficiently evolve the fluid equations at larger distances without losing accuracy close to the black hole (e.g. adaptive mesh refinement) will be required to avoid these issues. In a recent paper [58], we showed that the main source of error when measuring the mass of the ejecta is generally the large grid spacing in the far zone.

The same holds in this simulation: we find

$$M_{ej} = 0.26M_\odot \pm 0.16M_\odot$$  

(5)
assuming second-order convergence between the medium and high resolution. This confirms that BHNS mergers with rapidly rotating black holes are significantly more favorable to the ejection of neutron-rich material. In [58], we found that for higher mass ratio systems ($q = 7$) and lower spins $\chi = 0.9$, a few per cent of a solar mass is likely to be unbound, already a noticeable improvement compared to low-spin systems. The ejecta in the case studied here, which is of course an extremely favorable configuration, is an order of magnitude larger—thus showing that for rapidly spinning systems, it is conceivable that a large fraction of the mass remaining outside of the black hole at late time is unbound.

Measurements of the kinetic energy of the ejecta are even less reliable than its mass. Indeed, the kinetic energy is dominated by the most relativistic parts of the ejecta, which are also the most poorly resolved. From measurements of $u_\text{e}$ as the outward-moving material is leaving the finite-difference grid, we estimate that the ejected material has a median velocity $v \sim 0.5c$, and kinetic energy $E_{\text{ej}} \sim 10^{52}–10^{53}$ erg. These numbers should, however, be considered only as order of magnitude estimates. Nonetheless, this is an extremely large amount of energy, which would cause the emission of a detectable radio signal for a large fraction of the mergers within the range of Advanced LIGO as the ejecta slows down in the interstellar medium.

4. Conclusion

We have simulated the merger of a black hole–neutron star system with a premerger black-hole spin parameter of $\chi = 0.97$, the highest spin yet attempted for modeling such systems. Given the strongly nonlinear dependence of many aspects of the merger on $\chi$ (black-hole rotation energy, ISCO location, accretion efficiency), this simulation carries the numerical exploration of BHNS mergers into a distinctly new regime not sampled even by our previous simulations with $\chi = 0.9$. These nonlinearities manifest themselves clearly in our results.

Upon disruption of the neutron star, less than half of the nuclear matter is promptly accreted into the black hole. Almost 20% appears to be ejected from the system in an unbounded tidal tail outflow, and the rest settles into a massive accretion disc. Both the ejecta and disc masses are higher than any previous fully relativistic BHNS simulation. This in itself is unsurprising; both theoretical considerations and previous simulations would lead us to expect that in higher $\chi$ systems more matter should evade prompt accretion. What is notable is that the disc mass exceeds even the expectations for this spin based on extrapolating the trends of lower spin systems. If such systems occur in nature, they must be particularly spectacular events even by the standard of BHNS mergers.

We have closely followed the behavior of the horizon and of the accretion disc. Both merger and accretion have the net effect of decreasing the dimensionless spin $\chi$ of the black hole. While the disc is thick and self-gravitating, it appears to be quite stable, settling to an axisymmetric quasistationary state and evolving only slowly under the influences of outer-disc settling and accretion-induced mass decrease. Its subsequent evolution will be driven in part by physical processes not included in these simulations.

That existing numerical relativity techniques can successfully treat such an extreme system essentially unaltered is encouraging. Further studies of the extreme-spin regime of BHNS parameter space can now be attempted; particularly interesting will be the exploration of systems with higher black-hole mass. These are expected to be more common astrophysically, and they may also be more intriguing from a relativist’s point of view, since the disruption event happens nearer to the horizon, as measured by the disruption separation divided by $M_{\text{BH}}$. The investigation of the accretion system studied in this paper should also be completed by incorporating the remaining crucial physics, particularly the magnetic turbulence.
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