Black Hole Formation in Fallback Supernova and the Spins of LIGO Sources

Sophie L. Schröder1,2, Aldo Batta1,2, and Enrico Ramirez-Ruiz1,2 ©

1 Niels Bohr Institute, University of Copenhagen, Blegdamsvej 17, DK-2100 Copenhagen, Denmark
2 Department of Astronomy and Astrophysics, University of California, Santa Cruz, CA 95064, USA

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Abstract

Here we investigate within the context of field binary progenitors how the the spin of the Laser Interferometer Gravitational Wave Observatory (LIGO) sources vary when the helium star-descendant black hole (BH) is formed in a failed supernova (SN) explosion rather than by direct collapse. To this end, we make use of 3d hydrodynamical simulations of fallback supernova in close binary systems with properties designed to emulate LIGO sources. By systematically varying the explosion energy and the binary properties, we are able to explore the effects that the companion has on redistributing the angular momentum of the system. We find that, unlike the mass, the spin of the newly formed BH varies only slightly with the currently theoretically unconstrained energy of the SN and is primarily determined by the initial binary separation. In contrast, variations in the initial binary separation yield sizable changes on the resultant effective spin of the system. This implies that the formation pathways of LIGO sources leading to a particular effective spin might be far less restrictive than the standard direct collapse scenario suggests.

Key words: binaries (including multiple): close – stars: black holes – supernovae: general

1. Introduction

The gravitational wave (GW) signals detected by the Laser Interferometer Gravitational Wave Observatory (LIGO; Abbott et al. 2016b, 2016c, 2017a, 2017b, 2017c) have uncovered a population of black holes (BHs) that is significantly more massive than the population known to reside in accreting binaries (Remillard & McClintock 2006). While there is significant debate in the community about how BH binaries are assembled (Zwart & McMillan 2000; Kalogera et al. 2007; Sadowski et al. 2008; Postnov & Yungelson 2014; Abbott et al. 2016a; Belczynski et al. 2016; de Mink & Mandel 2016; Rodriguez et al. 2016a; Gerosa & Berti 2017; Wysocki et al. 2017), the classical scenario (Tutukov & Yungelson 1993; Voss & Tauris 2003) remains one of the leading candidates. In this channel, a wide massive binary undergoes a series of mass transfer episodes leading to a tight binary comprised of a massive helium star ($M_*$) and a BH ($M_1$), prior to the formation of the second BH ($M_2$). It has been argued that LIGO observations of the mass-weighted angular momentum perpendicular to the orbital plane $\chi_{\text{eff}}$ provide constraints on this formation channel (Rodriguez et al. 2016b; Farr et al. 2017; Stevenson et al. 2017). This is because vital information on the mass transfer history of the binary and the spin of $M_*$ is imprinted on

$$\chi_{\text{eff}} = \frac{M_1 a_1 + M_2 a_2}{M_1 + M_2} \cdot \hat{L}.$$  \hspace{1cm} (1)

Here $a_1$ and $a_2$ are the dimensionless spins of the BHs and $\hat{L}$ is the direction of the angular momentum in the orbital plane.

The angular momentum of the secondary BH is intimately linked to that of the progenitor helium star, which in turn is determined by its mass-loss history and the torque exerted by the primary BH (Kushnir et al. 2016; Qin et al. 2018; Zaldarriaga et al. 2018). This torque can effectively drive synchronization of the stellar spin and the orbit in binaries tighter than $d_\star$, the maximum separation allowed for synchronization within the life of the helium star (Zaldarriaga et al. 2018). The final angular momentum of the star thus provides a reasonable estimate of $a_2$ when the mass and angular momentum losses from the final supernova explosion are ignored. All previous works have assumed that such effects are small based on the simple expectation that LIGO BHs are formed by direct collapse.

Motivated by the fact that the nature of BH-forming supernova (SN) explosions is not well understood (Fryer et al. 2012; Uglio et al. 2012; Pejcha & Thompson 2015), we explore in this Letter the effects on $\chi_{\text{eff}}$ when the second BH $M_2$ is formed instead by a fallback SN explosion (Moriya et al. 2010; Fryer et al. 2012; Dexter & Kasen 2013; Lovegrove & Woosley 2013; Perna et al. 2014; Batta et al. 2017; Fernández et al. 2018). For this purpose we make use of 3D hydrodynamical simulations of fallback SN in close binary systems with properties aimed at reproducing LIGO GW signals. The structure of this Letter is as follows. In Section 2 we describe the numerical formalism used to initiate the fallback SN explosion and compute the subsequent evolution of the binary. In Section 3 we describe the dynamics of the fallback material and its effect on the final spin of both BHs. Lastly, in Section 4 we present our key findings and relate them to the current population of LIGO sources.

2. Methods and Initial Setup

Here we follow the setup described in Batta et al. (2017) to study the evolution of the progenitor binary system after the birth of $M_2$. We make use of a modified version of the 3D smoothed particle hydrodynamics code GADGET2 (Springel 2005), with our initial setup consisting of a tidally locked 28$M_\odot$ helium star in orbit around a BH of $M_1 = 15M_\odot$. The reader is referred to Batta et al. (2017) for further details on
initial particle distribution and numerical accretion. We settled for a resolution of $5 \times 10^5$ particles, which showed convergence for $a_2$ and properly captured the dynamics of the ejecta and the binary system. Higher resolution studies using $2 \times 10^6$ particles were carried out to demonstrate convergence and accuracy for representative cases.

To study the interaction of the fallback material with the newly formed BH binary we explore three sets of simulations. Each set starts with the initial binary in a circular orbit with a separation $d = d_s R_*$, where $d_s = 2, 3, 5$ and $R_*$ is the pre-SN star’s radius. Then, for each orbital separation we run simulations with at least four different SN explosion energies. In all cases we assume $a_i = 0$ based on the results of Qin et al. (2018), and assume that synchronization of the stellar spin and the orbit has taken place, as is expected for the initial separations used in this analysis. Given the large uncertainties in BH natal kick estimations (Repetto & Nelemans 2015; Mandel 2016), we assumed the simplest scenario where no natal kick is applied to the recently formed BH. This, combined with the synchronization of the stellar spin and the orbit, translates into BH’s spins aligned with the orbital angular momentum.

The initial profile of the star was obtained from the 350C KEPLER model calculated by Woosley & Heger (2006) of a 28 $M_\odot$ pre-SN helium star with $R_* = 0.76 R_\odot$. We considered the innermost 3 $M_\odot$ of the pre-SN star to be the newly formed BH with $a_2(t = 0) = 0$, which we subsequently treat as a sink particle. After the removal of the inner core, we use a parameterized energy injection routine to mimic the supernova engine and derive the density and velocity profile of the expanding envelope. Specifically, the energy is parametrized as $E_{\text{SN}} = \zeta E_G$, where $E_G = 2.3 \times 10^{52}$ erg is the binding energy of the pre-SN star. This energy is then deposited instantaneously in a 1.5 $M_\odot$ mass shell located at the inner boundary between the BH and the stellar envelope.

The distribution that describes the ejecta is determined solely by the structure of the pre-SN star (Woosley & Weaver 1995; Matzner & McKee 1999) and is established by the hydrodynamics of the interaction. Initially, the shock propagates through the stellar material, pressurizing it and setting it into motion. Once the shock wave approaches the surface of the star, a rarefaction stage begins in which stellar material is accelerated by the entropy deposited by the shock. This stage terminates once the pressure ceases to be dynamically important and the material expands freely. Figure 1 shows the radial velocity profile of the envelope when the shock surfaces the stellar envelope for three different explosion energies: $\zeta = 0.1, 0.5$, and 0.9. The gray area shows the initial BH mass, while the dashed lines show $M_2(t)$ at the time the shock reaches $r = R_*$. Despite the complicated hydrodynamical interaction, the density and pressure approach steep power laws in velocity as the material enters the rarefaction stage and the ejecta begins to establish homologous expansion (Figure 1).

For a given progenitor structure, the ensuing ejecta will take similar ejecta distributions.

When $\zeta \lesssim 1$, energy injection fails to unbind the star such that a sizable fraction of its mass eventually falls back onto the newly formed BH. If this takes place in a binary system (Batta et al. 2017), a non-negligible fraction of the bound material can expand to a radius comparable or larger than the binary’s separation, thus immersing the BH companion in gas. The interaction of fallback material with the binary transfers orbital angular momentum to the gas, which upon accretion onto the orbiting BHs is finally transferred into spin angular momentum. Differences in $\zeta$ result in diverse accretion histories, which ultimately regulate the BHs’ final masses and spins. It is to this topic that we now draw our attention.

### 3. Fallback Supernova in Binaries and the Spins of LIGO Sources

#### 3.1. Spin Evolution of the Newly Formed BH

Figure 2 shows the gas column density in the equatorial plane of the binary for three different simulations (from top to bottom) and at three different stages (left to right). Evolutionary times in Figure 2 are measured in units of the dynamical time of the pre-SN star, $t_\text{dyn}$. All simulations have $\zeta = 0.4$ but differ on the initial separation of the binary: $d_a = 2$ (top panel), $d_a = 3$ (middle panel), and $d_a = 5$ (bottom panel). The frames are centered on the newly formed BH and the dashed circles in the left panels show the size of the pre-SN star.

Flow dynamics are similar in all three simulations shown in Figure 2. First, the envelope expands to rapidly engulf the companion BH. A bow shock is created as a result of this initial interaction. It is, however, only when the slower moving material reaches the companion that the resulting torque can supply the envelope gas with sizable angular momentum. This envelope material will remain bound to the system and will form a disk around $M_2$ if restricted to the region within which orbiting gas is gravitationally bound to the newly formed BH. A disk, albeit lighter, also forms around $M_1$, whose final mass depends sensitively on the initial separation.

The total mass bound to $M_2$ is the same in all simulations, yet the fraction of angular momentum accreted increases with decreasing separation. As a result, $a_2$ is higher for progressively more compact binaries despite the final mass of $M_2$ reaching similar values. This can be seen in the left panel of Figure 3, in which we show the evolution of $d_a$ as a function of the accreted mass in units of $M_\odot$. Initially, $a_2$ increases as envelope material is accreted directly onto the BH. The innate angular momentum in this initial phase is determined by tidal synchronization, which increases as the binary separation decreases (Kushnir et al. 2016, 2017; Zaldarriaga et al. 2018). A transition in the evolution of $a_2$ is observed in the left panel of Figure 3 when...
material that is effectively torqued by the binary is able to form a disk and is subsequently accreted onto $M_2$. This material has a higher specific angular momentum than the one initially set by tidal synchronization and, when accreted, is able to spin up the newly formed BH at a faster rate. The resultant change in slope observed in the left panel in Figure 3 due to the accretion of disk material is observed to occur earlier for smaller separations, which results in higher total spins values than those given by direct collapse of the same fallback material.

At a fixed $\zeta$, the spin of the newly formed BH depends sensitively on $d_*$. For $d_* \lesssim 5$, the resultant torque on the fallback material can be considerable and, as a result, $a_2$ can be appreciable larger than the one expected from tidal synchronization. In this case, the final mass of $M_2$ remains unchanged, while the final spin can vary drastically. The final mass of $M_2$ is, on the other hand, controlled by $\zeta$. The middle panel of Figure 3 shows the evolution of $a_2$ for a fixed separation $d_* = 3$ and changing $\zeta$. Initially, the spin evolution follows the trend expected from direct collapse. This is because the torque is unable to modify the original angular momentum of the promptly collapsing stellar material. A transition to disk accretion is seen in all cases, with the shift always occurring late in the mass accretion history of $M_2$. The resultant spin is similar in all cases due to the self-similarity of the mass distribution of the expanding ejecta (Matzner & McKee 1999), which results in a comparable mass ratio of directly falling

**Figure 2.** Gas column density (code units) in the equatorial plane of the binary for three different initial separations (from top to bottom) and at three different evolutionary stages (left to right). The frames are centered on the newly formed BH (black circle) and the dashed circles in the left panels show the initial size of the pre-SN star. The cyan circle shows the companion BH. Times are measured in units of $t_\infty$. All simulations have $\zeta = 0.4$ but differ on the initial separation of the binary: $d_* = 2$ (top panel), $d_* = 3$ (middle panel), and $d_* = 5$ (bottom panel). In the last frame of the bottom panel, the resulting large binary separation places the companion outside of the frame.
stellar material to disk material for different values of $\zeta$. For example, this fraction varies from 5.3% for $\zeta = 0.5\%$–7.6% for $\zeta = 0.9$ (see the middle panel of Figure 3 for $d_* = 3$). This mass ratio is mainly responsible for determining the final spin of the BH and varies only slightly with $\zeta$.

We have discussed, in the context of the classical scenario, the effects that the binary separation and the energy of the SN have on the resulting spin of the newly formed BH. The right panel of Figure 3 provides a clear summary of our findings as it shows the final spin of $M_2$ as function of $d$ and $\zeta$. The final mass of the newly formed BH is also shown by the size of the symbols. The masses for $M_2$ range from 22.5 $M_\odot$ for $\zeta = 0.1$ ($M_2/M_\odot \approx 0.8$) to $8.2 M_\odot$ for $\zeta = 0.9$ ($M_2/M_\odot \approx 0.3$).

Together with the results from our simulations we also plot the expected spin obtained from the direct collapse of the pre-SN stellar profile. This formalism makes use of the KEPLER model and assumes solid body rotation determined by tidal synchronization. Then, by assuming the spherical collapse of the star, we obtain the BH’s spin $a_2$ for different fractions $f = 0.3, 0.5, 0.8, 1$ of the collapsed stellar mass $M_\odot$. If the entire star was to collapse directly onto a BH, this will give a final spin $a_2$ that is solely dependent on $d$, as predicted by the dashed line in Figure 3, labeled $M_\odot$.

When $\zeta$ is small and a significant fraction of the material is promptly accreted by the BH, the simple direct collapse formalism provides an accurate description of the final spin of the newly formed BH. This can be seen by comparing the dashed line in Figure 3 labeled $0.8 M_\odot$ with the simulation results obtained for $\zeta = 0.1$, which give BHs with $M_2 \approx 0.8 M_\odot$ and final spins that closely resemble the direct collapse ones. By contrast, when $f M_\odot \lesssim M_\odot$, the final spin is significantly higher than the one predicted by direct collapse of the same enclosed material. This is because in such cases the fallback material is effectively torqued by the BH companion, which results in disk formation and consequently higher final spin values. Binary BH formation in the classical scenario depends critically on the currently poorly constrained energy of the resulting SN, which for fallback-mediated remnant growth results in faster-spinning BHs than what would have been attainable for a single star progenitor.

**Figure 3.** Dependence of $a_2$ on $\zeta$ and $d$. Left panel: the spin parameter of $M_2$ as a function of the accreted mass in units of $M_\odot$ for the three simulations shown in Figure 2. Here, all of the simulations have $\zeta = 0.4$ and, as a result, the total accreted mass is similar. The solid lines correspond to the values derived from the simulations, while the dashed lines indicate the expected mass and angular momentum accretion from material that remains in the disk (Bardeen 1970; Thorne 1974) when the simulation ends at $t = 80 t_\odot$. Middle panel: the dependence of $a_2$ on $\zeta$. All of the simulations shown have the same separation ($d = 3 R_\odot$), but $\zeta = 0.1, 0.4, 0.5, 0.7$, and 0.9. The dashed line represents the spin expected from direct collapse (i.e., $\zeta = 0$). Right panel: the dependence of $a_2$ on $d_*$ and $\zeta$. The size of the symbols show the final BH mass, ranging from 8.2 $M_\odot$ to 22.5 $M_\odot$. The vertical dashed line shows $d_*$, the radius for effective tidal synchronization (Zaldarriaga et al. 2018), and the vertical solid line shows $d_*$, the binary separation required for merging within a Hubble time ($28 M_\odot + 15 M_\odot$). Dashed lines depict the spin expected from the direct collapse of different mass fractions, $f M_\odot$, of the pre-SN star with $f = 0.3, 0.5, 0.8, 1$.

**Figure 4.** Dependence of $a_1$ and $a_2$ on both $\zeta$ and $d$. The size of the symbols depicts the value of the final mass of the BH. Top panel: $a_2$ as a function $\zeta$ for initial $d_* = 2$, $d_* = 3$ and $d_* = 5$. Bottom panel: $a_1$ as a function $\zeta$ for initial $d_* = 2$, $d_* = 3$ and $d_* = 5$, under the assumption that $M_1$ had no spin before the final SN explosion $a_1(t = 0) = 0$.

### 3.2. Spin Evolution of the Orbiting BH

Figure 4 shows the dependence of $a_1$ and $M_1$ on $\zeta$ and $d$. In contrast to $M_2$, the final mass of the companion BH is only weakly altered by changes in $\zeta$. The reason is that a comparatively small mass can be effectively restricted to the region within which the expanding envelope material is gravitationally bound to $M_1$. This bound material forms a disk whose final mass depends on both $d$ and $\zeta$. The resultant changes in $a_1$, under the assumption of $a_1(t = 0) = 0$, are observed to be more pronounced when the initial binary separation changes. Although, as expected, no sizable changes...
take place at large separations given that only a tiny fraction of the companion’s envelope can be under the gravitational influence of $M_1$. We note here that values of $a_1 \lesssim 0.08$, involving the accretion of a small number of particles, are not converged at the resolution used in this study. The final value of $a_1$ shows a modest variation with SN energy with a small preference for $\zeta \approx 0.5$ at small separations. This indicates that although the ejecta distributions are similar for changing values of $\zeta$, the fraction of bound material to $M_1$ is largest for this particular explosion energy, although its exact value is likely to change for different pre-SN progenitors.

4. Discussion

In this Letter we have explored within the classical binary scenario how the spin of LIGO sources vary when the remnant BH is formed in weak SN explosions instead of direct collapse. Our key findings are summarized below.

1. The final mass of the newly formed BH depends on the explosion energy. Its mass varies from $M_2 \approx 0.8 M_*$ for $\zeta = 0.1$ to $M_2 \approx 0.3 M_*$ for $\zeta = 0.9$ (Figure 1).
2. At a fixed SN energy, the final spin increases significantly with decreasing $d$ as a larger fraction of the fallback material is torqued by the companion. This results in similar mass BHs but with widely different spins (see the left panel in Figure 3).
3. Due to the self-similarity of the mass distribution of the expanding ejecta, the final spin of the BH varies only slightly with $\zeta$. This results in BHs with a wide range in masses but similar spins (see the middle panel in Figure 3).
4. In the presence of a companion, the final spin of a BH formed by a fallback SN explosion can be significantly higher than the one predicted by direct collapse of the same stellar material (see the right panel in Figure 3).
5. The spin of the BH companion, on the other hand, depends on both $\zeta$ and $d$. This is because its accretion history is determined by the amount of fallback material that it is able to seize (Figure 4).

In Figure 5 we present a comparison of our results (upper panel) in the context of both direct collapse solutions and current LIGO observations of binary BHs. Shown are $\chi_{\text{eff}}$ as a function of the chirp mass, $M_\chi$, of the resulting BH binary system. The shaded quadrilateral regions (upper and lower panels) show systems produced by the direct collapse of pre-SN helium stars of varying masses, whose structures have been taken from the KEPLER models of Woosley & Heger (2006). The final spin of $M_2$ is calculated using the radial stellar profile and assuming the rigid body rotation of the tidally synchronized SN progenitor. The pre-SN helium stars ($M_*$) are assumed to be orbiting around a BH with $M_1 = q M_*$ and $a_1 = 0$. The dependence of $\chi_{\text{eff}}$ with $M_\chi$ is obtained by varying $q$ from 0.53 to 1 in all cases, while the dependence of $\chi_{\text{eff}}$ at a fixed $M_\chi$ is obtained by changing $d$ from $2 R_* \rightarrow 5 R_*$ at constant $q$. To facilitate comparisons, we plot as shaded ellipses the 90% credibility intervals of the GW signals measured so far (Abbott et al. 2016b, 2016c, 2017a, 2017b, 2017c).

Some points should be emphasized. The current LIGO observations are inconsistent with the direct collapse of pre-SN helium stars in close binaries (Kushnir et al. 2016; Belczynski et al. 2017; Hotokezaka & Piran 2017; Zaldarriaga et al. 2018). When the assumption of direct collapse is relaxed, the mass of $M_2$ can be altered by small changes to the explosion energy $\zeta$, while $a_2$ and $a_1$ (to a lesser extent) depend primarily on $d$ (Figure 4). For the specific $28 M_\odot + 15 M_\odot$ system studied here, we show that changes in $\zeta$ alone can produce systems like LVT151012 ($\zeta = 0.1$, $d_\star = 5$) or GW170608 ($\zeta = 0.9$, $d_\star = 5$). For a fixed SN energy of $\zeta = 0.9$, changes in the initial separation can, on the other hand, yield systems like GW170608 ($d_\star = 5$) or GW151226 ($d_\star = 2$).

Irrespective of the exact progenitor system, the processes discussed here imply that the formation pathways of LIGO binary BHs are more complicated than the standard scenario suggests. But the effects are especially interesting for weak SN explosions taking place in close binary systems. Future LIGO observations can offer clues to the nature of the SN explosion leading to the formation of BHs, which is currently not well understood (Perna et al. 2014; Sukhbold et al. 2016; Raithel et al. 2018). For instance, GW170608 could be indicative of the weak SN explosion of a more massive pre-SN progenitor system, while GW151226 might arise due to direct collapse of a lighter, yet more compact, progenitor system.

The properties of LIGO sources in the $(\chi_{\text{eff}}, M_\chi)$ plane are diverse. One appealing aspect of the classical scenario is that the great variety of binary and explosion parameters can probably help explain this diversity. Given the need for a large helium core mass in progenitors, BH formation may be favored not only by slow rotation but also by low metallicity (Izzard et al. 2004). Larger-mass helium cores might have less energetic explosions, but this is currently highly uncertain. Many massive stars may produce supernovae by forming neutron stars in spherically symmetric explosions, but some
may fail during neutrino energy deposition, forming BHs in the center of the star (Fryer et al. 2012; Ugliano et al. 2012; Pejcha & Thompson 2015) and possibly a wide range of weak SN explosions (Moriya et al. 2010; Fryer et al. 2012; Dexter & Kasen 2013; Lovegrove & Woosley 2013; Battat et al. 2017; Fernández et al. 2018). Asymmetric SN explosions, which might be a natural consequence of BH formation, could change the results presented in this study. However, Chan et al. 2018 recently showed that the structure of the fallback material can only be significantly modified by asymmetries when the explosion is strong ($\zeta \gtrsim 1$) and, as such, we expect our results to be representative of weak SN models. One expects various outcomes ranging from very massive BHs with low spins (GW150914) to lighter and faster-spinning BHs (GW151226).

The number density of binary BHs of different masses and spins would provide a natural test to distinguish between different stellar explosion avenues.

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ORCID iDs
Sophie L. Schröder © https://orcid.org/0000-0003-1735-8263
Enrico Ramirez-Ruiz © https://orcid.org/0000-0003-2558-3102

References
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016a, ApJL, 818, L22
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016b, PhRvL, 116, 241103
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2016c, PhRvL, 116, 061102
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017a, PhRvL, 118, 221101
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017b, ApJL, 851, L35
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017c, PhRvL, 119, 141101
Bardeen, J. M. 1970, Natur, 226, 64
Batta, A., Ramirez-Ruiz, E., & Fryer, C. 2017, ApJL, 846, L15
Belczynski, K., Holz, D. E., Bulik, T., & O’Shaughnessy, R. 2016, Natur, 534, 512
Belczynski, K., Klencki, J., Meynet, G., et al. 2017, arXiv:1706.07053
Chan, C., Müller, B., Heger, A., et al. 2018, ApJL, 852, L19
de Mink, S. E., & Mandel, I. 2016, MNRAS, 460, 3545
Dexter, J., & Kasen, D. 2013, ApJ, 772, 30
Farr, W. M., Stevenson, S., Miller, M. C., et al. 2017, Natur, 548, 426
Fernández, R., Quataert, E., Kashiyama, K., & Coughlin, E. R. 2018, MNRAS, 476, 2166
Fryer, C. L., Belczynski, K., Wiktorowicz, G., et al. 2012, ApJ, 749, 91
Gerosa, D., & Berti, E. 2017, arXiv:1703.06223
Hořezekáza, K., & Piran, T. 2017, ApJ, 842, 111
Izzard, R. G., Ramirez-Ruiz, E., & Tout, C. A. 2004, MNRAS, 348, 1215
Kalogera, V., Belczynski, K., Kim, C., O’Shaughnessy, R., & Willems, B. 2007, PhR, 442, 75
Kushnir, D., Zaldarriaga, M., Kollmeier, J. A., & Waldman, R. 2016, MNRAS, 462, 844
Kushnir, D., Zaldarriaga, M., Kollmeier, J. A., & Waldman, R. 2017, MNRAS, 467, 2146
Lovegrove, E., & Woosley, S. E. 2013, ApJ, 769, 109
Mandel, I. 2016, MNRAS, 456, 578
Matzner, C. D., & McKee, C. F. 1999, ApJ, 510, 379
Moriya, T., Tominaga, N., Tanaka, M., et al. 2010, ApJ, 719, 1445
Pejcha, O., & Thompson, T. A. 2015, ApJ, 801, 90
Perna, R., Duffell, P., Cantiello, M., & MacFadyen, A. I. 2014, ApJ, 781, 119
Postnov, K. A., & Yungelson, L. R. 2014, LRR, 17, 3
Qin, Y., Frags, T., Meynet, G., et al. 2018, arXiv:1802.05738
Raihels, C. A., Sukhbold, T., & Ozel, F. 2018, ApJ, 856, 35
Remillard, R. A., & McClintock, J. E. 2006, arXiv:0606352
Repetto, S., & Nelemans, G. 2015, MNRAS, 453, 3341
Rodriguez, C. L., Haster, C.-J., Chatterjee, S., Kalogera, V., & Rasio, F. A. 2016a, ApJL, 824, L8
Rodriguez, C. L., Zevin, M., Pankow, C., Kalogera, V., & Rasio, F. A. 2016b, ApJL, 832, L2
Sadowski, A., Belczynski, K., Bulik, T., et al. 2008, ApJ, 676, 1162
Springel, V. 2005, MNRAS, 364, 1105
Stevenson, S., Berry, C. P. L., & Mandel, I. 2017, arXiv:1703.06873
Sukhbold, T., Ertl, T., Woosley, S. E., Brown, J. M., & Janka, H.-T. 2016, ApJL, 821, 38
Thorne, K. S. 1974, ApJ, 191, 507
Tutukov, A. V., & Yungelson, L. R. 1993, MNRAS, 260, 675
Ugliano, M., Janka, H.-T., Marek, A., & Arcones, A. 2012, ApJ, 757, 69
Voss, R., & Tauris, T. M. 2003, MNRAS, 342, 1169
Woosley, S. E., & Heger, A. 2006, ApJ, 637, 914
Woosley, S. E., & Weaver, T. A. 1995, ApJS, 101, 181
Wysocki, D., Gerosa, D., O’Shaughnessy, R., et al. 2017, arXiv:1709.01943
Zaldarriaga, M., Kushnir, D., & Kollmeier, J. A. 2018, MNRAS, 473, 4174
Zwart, S. F. P., & McMillan, S. L. W. 2000, ApJL, 528, L17