

# Modeling the Complete Gravitational Wave Spectrum of Neutron Star Mergers

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In the context of neutron star mergers, we study the gravitational wave spectrum of the merger remnant using numerical relativity simulations. Postmerger spectra are characterized by a main peak frequency  $f_2$  related to the particular structure and dynamics of the remnant hot hypermassive neutron star. We show that  $f_2$  is correlated with the tidal coupling constant  $\kappa_2^T$  that characterizes the binary tidal interactions during the late-inspiral merger. The relation  $f_2(\kappa_2^T)$  depends very weakly on the binary total mass, mass ratio, equation of state, and thermal effects. This observation opens up the possibility of developing a model of the gravitational spectrum of *every* merger unifying the late-inspiral and postmerger descriptions.

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**Introduction.**—Direct gravitational wave (GW) observations of binary neutron stars (BNS), late-inspiral merger and postmerger by ground-based GW interferometric experiments, can lead to the strongest constraints on the equation of state (EOS) of matter at supranuclear densities [1–7]. There are two ways to set such constraints (GW observations of BNS mergers can also constrain the source redshift [8,9]): (I) measure the binary phase during the last minutes of coalescence using matched filtered searches [1,3–5] and (II) measure the postmerger GW spectrum frequencies using burst searches [6,7].

Method (I) relies on the availability of waveform models that include tidal effects and are accurate up to merger [4,5,10]. Here, “up to merger” indicates the end of the chirping signal in a precise sense that will be described below. Tidal interactions are significant during the late stages of coalescence at GW frequencies  $f_{\text{GW}} \gtrsim 400$  Hz (for typical binary masses), and they affect the phase evolution of the binary. The zero-temperature EOS is constrained by the measure of the quadrupolar tidal coupling constant  $\kappa_2^T$  (or equivalent or correlated parameters; see, e.g., Ref. [4]) that accounts for the magnitude of the tidal interactions [1,11].

Combining results from numerical relativity and the effective-one-body (EOB) approach to the general relativistic two-body problem [12–15], one can show that the merger dynamics of *every* irrotational binary is characterized by the value of  $\kappa_2^T$  [16]. At sufficiently small separations, the relevant dependency of the dimensionless GW frequency on the EOS, binary mass, and mass ratio is completely encoded in the tidal coupling constant. (The spin dependence is approximately linear for small spins aligned with the orbital angular momentum.) A tidal effective-one-body model compatible with numerical relativity data up to merger was introduced in Ref. [17], but no

prescription is available to extend the model to the postmerger.

Method (II) relies on the high-frequency GW spectrum and can, in principle, deliver a measure independent on method (I) [7]. Binary configurations with total mass  $M \leq M_{\text{thr}} \sim 2.9M_{\odot}$  are expected to produce a merger remnant composed of a hot massive or hypermassive neutron star. The merger remnant has a characteristic GW spectrum composed of a few broad peaks around  $f_{\text{GW}} \sim 1.8\text{--}4$  kHz. The key observation here is that the main peak frequencies of the postmerger spectrum are strongly correlated with properties (radius at a fiducial mass, compactness, etc.) of a zero-temperature spherical equilibrium star in an EOS-independent way [6,18]. Thus, a measure of the peak frequency constrains the correlated star parameter. Recently, there has been intense research on this topic, and various EOS-independent relations were proposed [6,18–24]. Most of the relations are constructed for equal-mass configurations and do not describe generic configurations for different total masses and mass ratios; see, e.g., Refs. [19,25]. Additionally, the postmerger GW spectrum might be influenced in a complicated way by thermal effects, magnetohydrodynamical instabilities, and dissipative processes.

In this Letter, we observe that the coupling constant  $\kappa_2^T$  can also be used to determine the main features of the postmerger GW spectrum in an EOS-independent way and for generic binary configurations, notably also in the unequal-mass case. The observation opens up the possibility of modeling the complete GW spectrum of neutron star mergers unifying the late-inspiral and postmerger descriptions. Geometrical units  $c = G = 1$  are employed throughout this Letter, unless otherwise stated. We use  $f$  for the spectrum frequencies and  $\omega$  for the instantaneous, time-dependent frequency.

TABLE I. BNS configurations and data. The columns are labeled as follows: name, EOS, binary total mass  $M$ , mass ratio  $q$ ,  $f_2$  frequency in kHz, dimensionless  $Mf_2$  frequency, and tidal coupling constant  $\kappa_2^T$ . Configurations marked with \* are stable MNS.

Name	EOS	$M [M_\odot]$	$q$	$f_2$ [kHz]	$Mf_2 [\times 10^2]$	$\kappa_2^T$
SLy-135135	SLy	2.70	1.00	3.48	4.628	74
SLy-145125	SLy	2.70	1.16	3.42	4.548	75
ENG-135135	ENG	2.70	1.00	2.86	3.803	91
SLy-140120	SLy	2.60	1.17	3.05	3.906	96
MPA1-135135	MPA1	2.70	1.00	2.57	3.418	115
SLy-140110	SLy	2.50	1.27	2.79	3.426	126
ALF2-135135	ALF2	2.70	1.00	2.73	3.630	138
ALF2-145125	ALF2	2.70	1.16	2.66	3.537	140
H4-135135	H4	2.70	1.00	2.50	3.325	211
H4-145125	H4	2.70	1.16	2.36	3.138	212
ALF2-140110	ALF2	2.50	1.27	2.38	2.931	216
MS1b-135135*	MS1b	2.70	1.00	2.00	2.660	290
MS1-135135*	MS1	2.70	1.00	1.95	2.593	327
MS1-145125*	MS1	2.70	1.16	2.06	2.740	331
2H-135135*	2H	2.70	1.00	1.87	2.561	439
MS1b-140110*	MS1b	2.50	1.27	2.08	2.487	441
MS1b-150100*	MS1b	2.50	1.50	1.87	2.303	461

*Numerical relativity GW spectra.*—The numerical relativity data used in this work were previously computed in Refs. [16,26]. In our simulations we solve Einstein equations using the Z4c formulation [27] and general relativistic hydrodynamics [28]. Our numerical methods are detailed in Refs. [26,29–33]. The binary configurations considered here are listed in Table I. In the following we summarize the main features of the GW radiation obtained by BNS simulations.

We consider equal and unequal mass configurations, different total masses, and a large variation of zero-temperature EOSs parametrized by piecewise polytropic fits [34]. Thermal effects are simulated with an additive thermal contribution in the pressure of a  $\Gamma$ -law form,  $P_{\text{th}} = (\Gamma_{\text{th}} - 1)\rho\epsilon$ , where  $\Gamma_{\text{th}} = 1.75$ ,  $\rho$  is the rest-mass density and  $\epsilon$  the specific internal energy of the fluid; see Refs. [32,35,36]. The initial configurations are prepared in quasicircular orbits assuming the fluid is irrotational.

Initial data are evolved for several orbits, during merger and in the postmerger phase for  $\gtrsim 30$  milliseconds. A detailed discussion of the merger properties determined by different EOSs, mass, and mass ratio is presented in Refs. [16,26]. The binary configurations in our sample do not promptly collapse to a black hole after merger, but form either a stable massive neutron star (MNS) or an unstable hypermassive neutron star (HMNS), which collapses on a dynamical time scale  $\tau_{\text{GW}} \lesssim \langle R \rangle^4 / \langle M \rangle^3 \approx 200$  ms [37]. Both HMNS and MNS remnants at formation are hot, differentially rotating, nonaxisymmetric, highly dynamical two-core structures (see, e.g., Refs. [35,38]).

The typical GW signal computed in our simulations is shown in Fig. 1 for a fiducial configuration. We plot the real

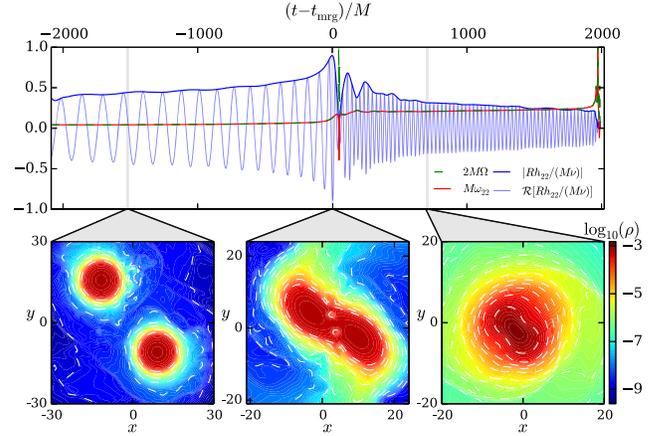


FIG. 1 (color online). Simulations of BNS and GWs. Top panel: Real part and amplitude of the GW mode  $Rh_{22}/(\nu M)$  and the associated dimensionless frequency  $M\omega_{22}$  versus the mass-normalized retarded time  $t/M$  for a fiducial configuration, H4-135135. The signal is shifted to the moment of merger,  $t_{\text{mrg}}$ , defined by the amplitude's peak (end of chirping). Also shown is (twice) the dynamical frequency  $M\Omega = \partial E_b / \partial j \sim M\omega_{22}/2$ . Bottom panels: Snapshots of  $\log_{10} \rho$  on the orbital plane, during the late inspiral (left panel), at simulation time corresponding to  $t_{\text{mrg}}$  (middle panel), and during the postmerger (right panel).

part and amplitude of the dominant  $\ell = m = 2$  multipole of the  $s = -2$  spin-weighted spherical harmonics decomposition of the GW,  $R(h_+ - ih_\times) = \sum_{\ell m} R h_{\ell m - 2} Y_{\ell m}(\theta, \phi)$ , versus the retarded time  $t$ . The figure's main panel also shows the  $\ell = m = 2$  instantaneous and dimensionless GW frequency  $M\omega_{22} = M d\phi/dt$ , where  $\phi = -\arg(Rh_{22})$ . The bottom panels show snapshots of  $\log_{10} \rho$  on the orbital plane, corresponding to three representative simulation times.

The waveform at early times is characterized by the well-known chirping signal; frequency and amplitude monotonically increase in time. The GW frequency reaches typical values  $\omega_{\text{GW}} = 2\pi f_{\text{GW}} \approx \omega_{22} \lesssim 0.1-0.2/M$ , i.e.,  $f_{\text{GW}} \sim 0.8-1.6$  kHz for an  $M = 2.7M_\odot$  binary. The chirping signal ends at the amplitude peak,  $\max |Rh_{22}|$ , which is marked in the figure by the middle vertical line. We formally define this time as the *moment of merger*,  $t_{\text{mrg}}$ , and refer to the signal at  $t > t_{\text{mrg}}$  as the postmerger signal. The GW postmerger signal is essentially generated by the  $m = 2$  structure of the remnant (see bottom right panel of Fig. 1). The frequency increases monotonically to  $M\omega_{22} \sim 0.2-0.5$  as the HMNS becomes more compact and eventually approaches the collapse. Assuming the remnant can be instantaneously approximated by a perturbed differentially rotating star [38], the  $f$ -mode of pulsation is strongly excited at formation and it is the most efficient emission channel for GWs.

The GW spectra are shown in Fig. 2 for a representative subset of configurations. Triangles mark frequencies  $f_{\text{mrg}}$  corresponding to  $t_{\text{mrg}}$ . Circles mark the main postmerger peak frequencies  $f_2 \sim 1.8-4$  kHz. The small frequency

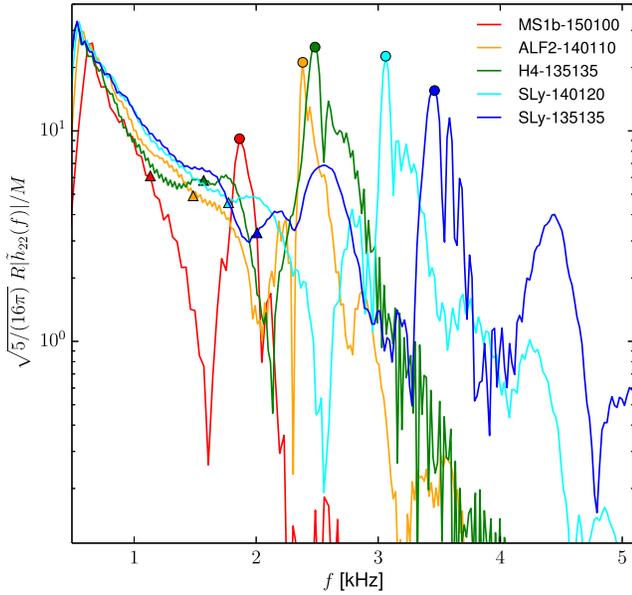


FIG. 2 (color online). GW spectra from BNS. The plot shows only a representative subset of the configurations of Table I. Triangles mark frequencies  $f_{\text{mrg}}$  corresponding to  $t_{\text{mrg}}$ ; circles mark  $f_2$  frequencies.

cutoff is artificial and related to the small binary separation of the initial data; physical spectra monotonically extend to lower frequencies. From the figure, one also observes that (i) there exists other peaks, expected by nonlinear mode coupling or other hydrodynamical interactions [23,35,38], (ii) peaks are broad, reflecting the nontrivial time evolution of the frequencies (see Fig. 1 and also the spectrogram in Ref. [39]), and (iii) secondary peaks are present in most of the configurations, and their physical interpretation has been discussed in Refs. [22,23,26,38]. We postpone the analysis of these features to future work. In the following we focus only on the  $f_2$  peak, which is the most robust and understood feature of the GW postmerger spectrum.

*Characterization of the postmerger GW spectra.*—Here, we show that  $f_2$  correlates with the tidal coupling constant  $\kappa_2^T$  that parametrizes the binary tidal interactions and waveforms during the late-inspiral merger. The relation  $f_2(\kappa_2^T)$  depends very weakly on the binary total mass, mass ratio, and EOS. We use a large data sample of 99 points including the data of Refs. [19,24].

Let us first briefly summarize the definition of  $\kappa_2^T$  and its role in the merger dynamics.

Within the EOB framework, tidal interactions are described by an additive correction  $A^T(r)$  to the radial, Schwarzschild-like metric potential  $A(r)$  of the EOB Hamiltonian [11]. The potential  $A(r)$  represents the binary interaction energy. In order to understand its physical meaning, it is sufficient to consider the Newtonian limit of the EOB Hamiltonian,  $H_{\text{EOB}} \approx Mc^2 + (\mu/2)\mathbf{p}^2 + (\mu/2)(A(r) - 1) + \mathcal{O}(c^{-2})$ , where  $\mu = M_A M_B / (M_A + M_B)$  is the binary reduced mass,  $\mathbf{p}$  the momenta, and

$A(r) = 1 - (2/r) = 1 - 2(GM/c^2 r_{AB})$ , with  $r_{AB}$  the relative distance between the stars (constants  $c$  and  $G$  are reintroduced for clarity). The tidal correction  $A^T(r)$  is parametrized by a multipolar set of relativistic tidal coupling constants  $\kappa_{(\ell)}^{A,B}$ , where  $A, B$  label the stars in the binary [1,11]. The leading-order contribution to  $A^T(r)$  is proportional to the quadrupolar ( $\ell = 2$ ) coupling constants,  $\kappa_2^A = 2k_2^A (X_A/C_A)^5 M_B/M_A$ , where  $M_A$  is the mass of star  $A$ ,  $C_A$  the compactness,  $X_A = M_A/M$ , and  $k_2^A$  the  $\ell = 2$  dimensionless Love number [40–43]. The total  $\ell = 2$  coupling constant is defined as  $\kappa_2^T = \kappa_2^A + \kappa_2^B$  and can be written as

$$\kappa_2^T = 2 \left( \frac{q^A}{(1+q)^5} \frac{k_2^A}{C_A^5} + \frac{q}{(1+q)^5} \frac{k_2^B}{C_B^5} \right), \quad (1)$$

assuming  $q = M_A/M_B \geq 1$ . The leading-order term of the tidal potential is simply  $A^T(r) = -\kappa_2^T r^{-6}$ .

A consequence of the latter expression for  $A^T(r)$  is that the merger dynamics is essentially determined by the value of  $\kappa_2^T$  [16]. All the dynamical quantities develop a nontrivial dependence on  $\kappa_2^T$  as the binary interaction becomes tidally dominated. The characterization of the merger dynamics via  $\kappa_2^T$  is “universal” in the sense that it does not require any other parameter such as EOS,  $M$ , and  $q$ . (There is, however, a dependency on the star spins.) For example, at the reference point  $t_{\text{mrg}}$ , the corresponding binary reduced binding energy  $E_b^{\text{mrg}}$ , the reduced angular momentum  $j^{\text{mrg}}$ , and the GW frequency  $M\omega_{22}^{\text{mrg}}$  can be fitted to simple rational polynomials [16],

$$Q(\kappa_2^T) = Q_0 \frac{1 + n_1 \kappa_2^T + n_2 (\kappa_2^T)^2}{1 + d_1 \kappa_2^T}, \quad (2)$$

with fit coefficients  $(n_i, d_i)$  given in Table II.

In view of these results, it appears natural to investigate the dependency of the *postmerger* spectrum on  $\kappa_2^T$ .

Our main result is summarized in Fig. 3, which shows the postmerger main peak dimensionless frequency  $Mf_2$  as a function of  $\kappa_2^T$  for a very large sample of binaries. Together with our data, we include those tabulated in Refs. [19,24]. The complete data set spans the ranges  $M \in [2.45M_\odot, 2.9M_\odot]$ ,  $q \in [1.0, 1.5]$ , and a large variation of EOSs. The peak location is typically determined within an accuracy of  $\delta f \sim \pm 0.2$  kHz (see also Ref. [18]). Each of the four panels of Fig. 3 shows the same data; the color code in each panel indicates different values of  $M$  (top left), EOS (top right),  $q$  (bottom left), and  $\Gamma_{\text{th}}$  (bottom right). The data correlate rather well with  $\kappa_2^T$ . As indicated by the colors and different panels, the scattering of the data does *not* correlate with variations of  $M$ , EOS,  $q$ , and  $\Gamma_{\text{th}}$ . The black solid line is our best fit to Eq. (2), where we set  $n_2 = 0$  and fit also for  $Q_0$  (see Table II). The 95% confidence interval of the fit is shown as a gray shaded area in Fig. 3.

TABLE II. Fit coefficients of different quantities at  $t_{\text{mrg}}$  and of  $Mf_2$  with the template in Eq. (2).

$Q(\kappa_2^T)$	$Q_0$	$n_1 [\times 10^2]$	$n_2 [\times 10^5]$	$d_1 [\times 10^2]$
$E_b^{\text{mrg}}$	-0.1201	+2.9905	-1.3665	+6.7484
$j^{\text{mrg}}$	+2.8077	+4.0302	+0.7538	+3.1956
$M\omega_{22}^{\text{mrg}}$	+0.3596	+2.4384	-1.7167	+6.8865
$Mf_2$	+0.053850	+0.087434	0	+0.45500

We argue that the observed postmerger correlation with  $\kappa_2^T$  is a direct consequence of the merger universality. Although an analytical or approximate description of the postmerger dynamics is not available, the gauge-invariant  $E_b(j)$  curves contain, in analogy to the merger case, significant information about the system dynamics [17]. Specifically, we interpret  $E_b(j)$  as being generated by some Hamiltonian flow that continuously connects merger and postmerger. In terms of this Hamiltonian evolution, the values  $[E_b^{\text{mrg}}(\kappa_2^T), j^{\text{mrg}}(\kappa_2^T)]$  provide initial conditions for the dynamics of the MNS/HMNS; it is then plausible to assume that the postmerger correlation follows from these initial conditions by continuity. In order to assess this conjecture, we define the frequency given by the equation  $M\Omega = \partial E_b / \partial j$ , notice that  $\Omega^{\text{mrg}} = \Omega^{\text{mrg}}(\kappa_2^T)$ , and show that  $\Omega$  is the relevant dynamical frequency for both inspiral merger and postmerger. Recalling that the standard quadrupole formula predicts that a generic source with  $m = 2$  geometry and rotating at frequency  $\Omega$  emits GWs at a frequency  $2\Omega$ , we plot the latter in Fig. 1 and indeed observe that it corresponds to the main emission channel  $\omega_{22}$  during the whole evolution. In practice, the gauge-invariant  $\Omega$  can be interpreted as the orbital frequency during the inspiral and the angular frequency of the MNS/HMNS during postmerger. Furthermore, since merger remnants from larger  $\kappa_2^T$  binaries are less bound and have larger angular momentum support at formation,  $\Omega(\kappa_2^T)$  (and thus  $f_2$ ) must be a monotonically decreasing function of  $\kappa_2^T$ , which is what one can observe in Fig. 3.

The frequency evolution is also expected to depend on angular momentum dissipation due to magnetic field instabilities (see, e.g., Refs. [44–46]), cooling, and shear viscosity [37]. However, the available literature indicates these physical effects are negligible in the first approximation, and we argue that they might result in frequency shifts  $\Delta f_2 \lesssim \delta f_2$ . The star rotation can instead play a relevant role via spin-orbit coupling effects: Stars with dimensionless spin parameters  $\gtrsim 0.05$ – $0.1$  can give frequency shifts  $\gtrsim \delta f_2$  [39].

*Outlook.*—The result of this work, coupled with the modeling of the merger process given in Refs. [16,17], indicates the possibility of modeling the late-inspiral-merger-postmerger GW spectrum in a consistent way using  $\kappa_2^T$  as a main parameter. In particular, an accurate late-inspiral-merger GW spectrum is given by a suitable frequency-domain representation,  $\tilde{h}(f) = A(f) \exp[-i\Psi(f)]$ , of the

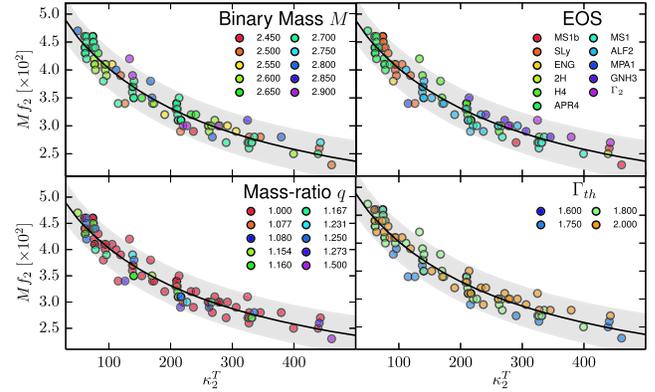


FIG. 3 (color online).  $Mf_2$  dimensionless frequency as a function of the tidal coupling constant  $\kappa_2^T$ . Each panel shows the same data set; the color code in each panel indicates the different values of binary mass (top left), EOS (top right), mass ratio (bottom left), and  $\Gamma_{\text{th}}$  (bottom right). The black solid line is our fit [see Eq. (2) and Table II]; the gray area marks the 95% confidence interval.

waveform of Ref. [17]. The leading-order tidal contribution of such a spectrum reads  $\Psi^T(f) = -39/4\kappa_2^T x^{5/2}$  with  $x(f) \propto f^{2/3}$ ; see Ref. [1] for  $\Psi^T(f)$  at 2.5 post-Newtonian order. A simple template for the postmerger spectrum for binaries with  $M \leq M_{\text{thr}}$  is then given by a single-peak model and our fit for  $f_2$ . The precise construction of such a complete spectrum will be the subject of future work. As mentioned in our discussion, it will be particularly important to include spin effects (see, e.g., Refs. [5,39]).

The performance of the proposed model in a GW data-analysis context will be carefully evaluated in a separate study. In this respect, we suggest that an optimal strategy to constrain the EOS could be to combine the late-inspiral measurement of type (I) with measurement of type (II). The inclusion of the postmerger model might lead to an improved estimate of  $\kappa_2^T$ , for the same number of observed events [3–5].

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[1] T. Damour, A. Nagar, and L. Villain, *Phys. Rev. D* **85**, 123007 (2012).

[2] J. S. Read, L. Baiotti, J. D. E. Creighton, J. L. Friedman, B. Giacomazzo, K. Kyutoku, C. Markakis, L. Rezzolla, M. Shibata, and K. Taniguchi, *Phys. Rev. D* **88**, 044042 (2013).

- [3] W. Del Pozzo, T. G. F. Li, M. Agathos, C. V. D. Broeck, and S. Vitale, *Phys. Rev. Lett.* **111**, 071101 (2013).
- [4] B. D. Lackey and L. Wade, *Phys. Rev. D* **91**, 043002 (2015).
- [5] M. Agathos, J. Meidam, W. Del Pozzo, T. G. F. Li, M. Tompitak, J. Veitch, S. Vitale, and C. Van Den Broeck, *Phys. Rev. D* **92**, 023012 (2015).
- [6] A. Bauswein and H.-T. Janka, *Phys. Rev. Lett.* **108**, 011101 (2012).
- [7] J. Clark, A. Bauswein, L. Cadonati, H.-T. Janka, C. Pankow, and N. Stergioulas, *Phys. Rev. D* **90**, 062004 (2014).
- [8] C. Messenger and J. Read, *Phys. Rev. Lett.* **108**, 091101 (2012).
- [9] C. Messenger, K. Takami, S. Gossan, L. Rezzolla, and B. Sathyaprakash, *Phys. Rev. X* **4**, 041004 (2014).
- [10] M. Favata, *Phys. Rev. Lett.* **112**, 101101 (2014).
- [11] T. Damour and A. Nagar, *Phys. Rev. D* **81**, 084016 (2010).
- [12] A. Buonanno and T. Damour, *Phys. Rev. D* **59**, 084006 (1999).
- [13] A. Buonanno and T. Damour, *Phys. Rev. D* **62**, 064015 (2000).
- [14] T. Damour, P. Jaranowski, and G. Schäfer, *Phys. Rev. D* **62**, 084011 (2000).
- [15] T. Damour, *Phys. Rev. D* **64**, 124013 (2001).
- [16] S. Bernuzzi, A. Nagar, S. Balmelli, T. Dietrich, and M. Ujevic, *Phys. Rev. Lett.* **112**, 201101 (2014).
- [17] S. Bernuzzi, A. Nagar, T. Dietrich, and T. Damour, *Phys. Rev. Lett.* **114**, 161103 (2015).
- [18] A. Bauswein, H. Janka, K. Hebeler, and A. Schwenk, *Phys. Rev. D* **86**, 063001 (2012).
- [19] K. Hotokezaka, K. Kiuchi, K. Kyutoku, T. Muranushi, Y.-i. Sekiguchi, M. Shibata, and K. Taniguchi, *Phys. Rev. D* **88**, 044026 (2013).
- [20] A. Bauswein, T. Baumgarte, and H. T. Janka, *Phys. Rev. Lett.* **111**, 131101 (2013).
- [21] A. Bauswein, N. Stergioulas, and H.-T. Janka, *Phys. Rev. D* **90**, 023002 (2014).
- [22] K. Takami, L. Rezzolla, and L. Baiotti, *Phys. Rev. Lett.* **113**, 091104 (2014).
- [23] A. Bauswein and N. Stergioulas, *Phys. Rev. D* **91**, 124056 (2015).
- [24] K. Takami, L. Rezzolla, and L. Baiotti, *Phys. Rev. D* **91**, 064001 (2015).
- [25] A. Bauswein, N. Stergioulas, and H.-T. Janka, *Phys. Rev. D* **91**, 124056 (2015).
- [26] T. Dietrich, S. Bernuzzi, M. Ujevic, and B. Brügmann, *Phys. Rev. D* **91**, 124041 (2015).
- [27] S. Bernuzzi and D. Hilditch, *Phys. Rev. D* **81**, 084003 (2010).
- [28] J. A. Font, *Living Rev. Relativity* **11**, 7 (2007).
- [29] D. Hilditch, S. Bernuzzi, M. Thierfelder, Z. Cao, W. Tichy, and B. Brügmann, *Phys. Rev. D* **88**, 084057 (2013).
- [30] S. Bernuzzi, A. Nagar, M. Thierfelder, and B. Brügmann, *Phys. Rev. D* **86**, 044030 (2012).
- [31] S. Bernuzzi, M. Thierfelder, and B. Brügmann, *Phys. Rev. D* **85**, 104030 (2012).
- [32] M. Thierfelder, S. Bernuzzi, and B. Brügmann, *Phys. Rev. D* **84**, 044012 (2011).
- [33] B. Brügmann, J. A. Gonzalez, M. Hannam, S. Husa, U. Sperhake, and W. Tichy, *Phys. Rev. D* **77**, 024027 (2008).
- [34] J. S. Read, B. D. Lackey, B. J. Owen, and J. L. Friedman, *Phys. Rev. D* **79**, 124032 (2009).
- [35] M. Shibata, K. Taniguchi, and K. Uryu, *Phys. Rev. D* **71**, 084021 (2005).
- [36] A. Bauswein, H.-T. Janka, and R. Oechslin, *Phys. Rev. D* **82**, 084043 (2010).
- [37] T. W. Baumgarte, S. L. Shapiro, and M. Shibata, *Astrophys. J.* **528**, L29 (2000).
- [38] N. Stergioulas, A. Bauswein, K. Zagkouris, and H.-T. Janka, *Mon. Not. R. Astron. Soc.* **418**, 427 (2011).
- [39] S. Bernuzzi, T. Dietrich, W. Tichy, and B. Brügmann, *Phys. Rev. D* **89**, 104021 (2014).
- [40] T. Hinderer, *Astrophys. J.* **677**, 1216 (2008).
- [41] T. Damour and A. Nagar, *Phys. Rev. D* **80**, 084035 (2009).
- [42] T. Binnington and E. Poisson, *Phys. Rev. D* **80**, 084018 (2009).
- [43] T. Hinderer, B. D. Lackey, R. N. Lang, and J. S. Read, *Phys. Rev. D* **81**, 123016 (2010).
- [44] R. Ciolfi, S. K. Lander, G. M. Manca, and L. Rezzolla, *Astrophys. J.* **736**, L6 (2011).
- [45] B. Zink, P. D. Lasky, and K. D. Kokkotas, *Phys. Rev. D* **85**, 024030 (2012).
- [46] K. Kiuchi, K. Kyutoku, Y. Sekiguchi, M. Shibata, and T. Wada, *Phys. Rev. D* **90**, 041502 (2014).