

Unified picture of the post-merger dynamics and gravitational wave emission in neutron star mergers

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We introduce a classification scheme of the postmerger dynamics and gravitational wave emission in binary neutron star mergers, after identifying a new mechanism by which a secondary peak in the gravitational wave spectrum is produced. It is caused by a spiral deformation, the pattern of which rotates slower with respect to the double-core structure in the center of the remnant. This secondary peak is typically well separated in frequency from the secondary peak produced by a nonlinear interaction between a quadrupole and a quasiradial oscillation. The new mechanism allows for an explanation of low-frequency modulations seen in a number of physical characteristics of the remnant, such as the central lapse function, the maximum density and the separation between the two cores, but also in the gravitational wave amplitude. We find empirical relations for both types of secondary peaks between their gravitational wave frequency and the compactness of nonrotating individual neutron stars, that exist for fixed total binary masses. These findings are derived for equal-mass binaries without intrinsic neutron star spin analyzing hydrodynamical simulations without magnetic field effects. Our classification scheme may form the basis for the construction of detailed gravitational wave templates of the postmerger phase. We find that the quasiradial oscillation frequency of the remnant decreases with the total binary mass. For a given merger event, our classification scheme may allow one to determine the proximity of the measured total binary mass to the threshold mass for prompt black hole formation, which can, in turn, yield an estimate of the maximum neutron star mass.

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I. INTRODUCTION

Neutron star (NS) mergers are strong emitters of gravitational waves (GWs) and thus among the prime targets for the upcoming GW detectors Advanced LIGO [1] and Advanced Virgo [2]. Future GW observations of such events [3] could reveal the properties of high-density matter and NSs (see e.g. Refs. [4–8] for reviews). The merger will likely result in the formation of a differentially rotating, strongly oscillating remnant [9–38]. Specifically, it has been shown that detections of its *dominant* postmerger GW frequency f_{peak} would strongly constrain the radius and the maximum mass of nonrotating NSs [28,29,33,39,40]. Additionally, there exist potentially detectable *secondary peaks* at lower frequencies [9,12,13,16,17,19,22,23,25–29,33,35–41], which could improve the constraints on NS properties if the mechanism by which they appear is understood.

In this study we find that *two different mechanisms* are at work for producing the low-frequency secondary peaks in the GW spectrum at frequencies below the main peak of the postmerger phase. Specifically, we show that apart from a *nonlinear combination frequency* [25], there exists also a distinct mechanism that generates a secondary GW peak by the *rotating pattern of a deformation of spiral shape*. This deformation is initially produced at the time of merging and is then sustained for a few rotational periods. The consideration of the two different mechanisms leads to a unified

classification scheme for the postmerger dynamics and GW emission. For high-mass binaries (relative to the threshold mass to prompt black-hole collapse), the nonlinear combination frequency dominates, while for low-mass binaries it is the spiral deformation that produces the strongest secondary peak. Both are simultaneously present and can produce peaks of comparable strength for intermediate binary masses. Hence, the secondary peaks cannot simply be explained by the dynamics of the double-core structure alone, which forms the inner part of the remnant, as in Refs. [36–38]. Identifying the importance of the rotating deformation, we explain for the first time the existence of a low-frequency modulation appearing in addition to the quasiradial oscillation mode in several physical quantities characterizing the remnant.

After clarifying the nature of the secondary peaks in the GW spectrum, we find mass-dependent empirical relations for the different peaks as a function of NS compactness. We rule out the existence of a universal, mass-independent and equation of state (EOS)-independent relation for the secondary peaks, a relation recently proposed in Refs. [36,37]. Such a relation can only be found within a small sample of EOSs for a very limited mass range (not the same for all EOSs). Our new relations can be used to optimize the constraints on the EOS in the case of the simultaneous detection of several postmerger frequencies. The understanding of the most prominent features of the GW spectrum is a prerequisite for constructing GW templates

for the postmerger phase, which could enhance the detection prospects compared to unmodeled searches [40,41] for the Advanced LIGO and Advanced Virgo detectors and their discussed upgrades [42–44]. For the planned Einstein Telescope [45], direct detections of secondary peaks are a viable prospect [36,37,40,41].

II. NATURE OF SECONDARY GW PEAKS

We investigate mergers of equal-mass, intrinsically non-spinning NSs with a 3D relativistic smoothed particle hydrodynamics (SPH) code, which imposes the conformal flatness condition on the spatial metric [46,47] to solve Einstein’s field equations and incorporates energy and angular momentum losses by a GW backreaction scheme [18,48] (see Refs. [12,18,28,29,49] for details on the code, the setup, resolution tests and model uncertainties). Comparisons to other numerical setups and also models with an approximate consideration of neutrino effects show an agreement in determining the postmerger spectrum within a few percent in the peak frequencies [27–29,33,36–38]. Magnetic field effects are negligible for not-too-high initial field strengths [24]. We explore a representative sample of ten microphysical, fully temperature-dependent equations of state (EOSs) (see Table I in Ref. [39] and Fig. 5 in this work for the mass-radius relations of non-rotating NSs of these EOSs) and consider total binary masses M_{tot} between $2.4 M_{\odot}$ and $3.0 M_{\odot}$. In this work we consider only NSs with an initially irrotational velocity profile, because known spin periods in observed NS binaries are slow compared to their orbital motion (see e.g. Ref. [50]), and simulations with initial intrinsic NS spin suggest an impact on the postmerger features of the GW signal only for very fast spins [19,35,38].

First, we focus on a reference model for the moderately stiff DD2 EOS [51,52] with an intermediate binary mass of $M_{\text{tot}} = 2.7 M_{\odot}$. Figure 1 shows the x-polarization of the effective amplitude $h_{\text{eff},x} = \tilde{h}_x(f) \cdot f$ (with \tilde{h}_x being the Fourier transform of the waveform h_x) vs frequency f (reference model in black). Besides the dominant f_{peak} frequency [53], there are two secondary peaks at lower frequencies (f_{2-0} and f_{spiral}) with comparable signal-to-noise ratio. Both are generated in the postmerger phase, which can be seen by choosing a time window covering only the postmerger phase for computing the GW spectrum.

The secondary peak shown as f_{2-0} is a nonlinear combination frequency between the dominant quadrupolar f_{peak} oscillation and the quasiradial oscillation of the remnant, as described in Ref. [25]. We confirm this by performing additional simulations, after adding a quasiradial density perturbation to the remnant at late times. The frequency f_0 of the strongly excited quasiradial oscillation is determined by a Fourier analysis of the time evolution of the density or central lapse function and coincides with the

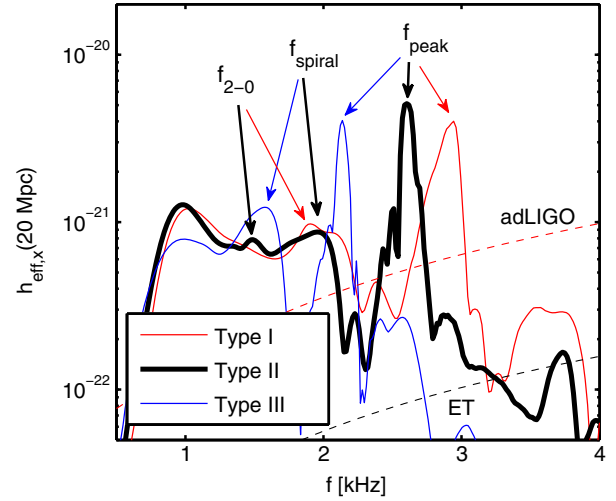


FIG. 1 (color online). GW spectra of $1.35\text{--}1.35 M_{\odot}$ mergers with the DD2 [51,52] (black), NL3 [51,54] (blue) and LS220 [55] (red) EOSs (cross polarization along the polar axis at a reference distance of 20 Mpc). Dashed lines show the anticipated unity SNR sensitivity curves of Advanced LIGO [1] (red) and of the Einstein Telescope [45] (black).

frequency difference $f_{\text{peak}} - f_{2-0}$. As in Ref. [25], the extracted eigenfunction at f_0 confirms the quasiradial nature.

The secondary f_{spiral} peak is produced by a strong deformation initiated at the time of merging, the pattern of which then rotates (in the inertial frame) slower than the inner remnant and lasts for a few rotational periods, while diminishing in amplitude. Figure 2 shows the density evolution in the equatorial plane, in which one can clearly identify the two antipodal bulges of the spiral pattern, which rotate slower than the central parts of the remnant. In this early phase the inner remnant is still composed of two dense cores rotating around each other (this is the nonlinear generalization of an $m = 2$ quadrupole oscillation producing the dominant f_{peak}). Extracting the rotational motion of the antipodal bulges in our simulations, we indeed find that their frequency equals $f_{\text{spiral}}/2$ producing gravitational waves at f_{spiral} (compare the times in the right panels in Fig. 2; recall the factor 2 in the frequency of the GW signal compared to the orbital frequency of orbiting point particles). In Fig. 2 the antipodal bulges are illustrated by selected fluid elements (tracers), which are shown as black and white dots, while the positions of the individual centers of the double cores are marked by a cross and a circle. (We define the centers of mass of the double cores by computing the centers of mass of the innermost 1000 SPH particles of the respective initial NSs and then following their time evolution.) While in the right panels the antipodal bulges completed approximately one orbit within one millisecond ($\approx \frac{2}{f_{\text{spiral}}}$), the double cores moved further ahead, i.e. with a significantly higher orbital frequency. Examining the GW

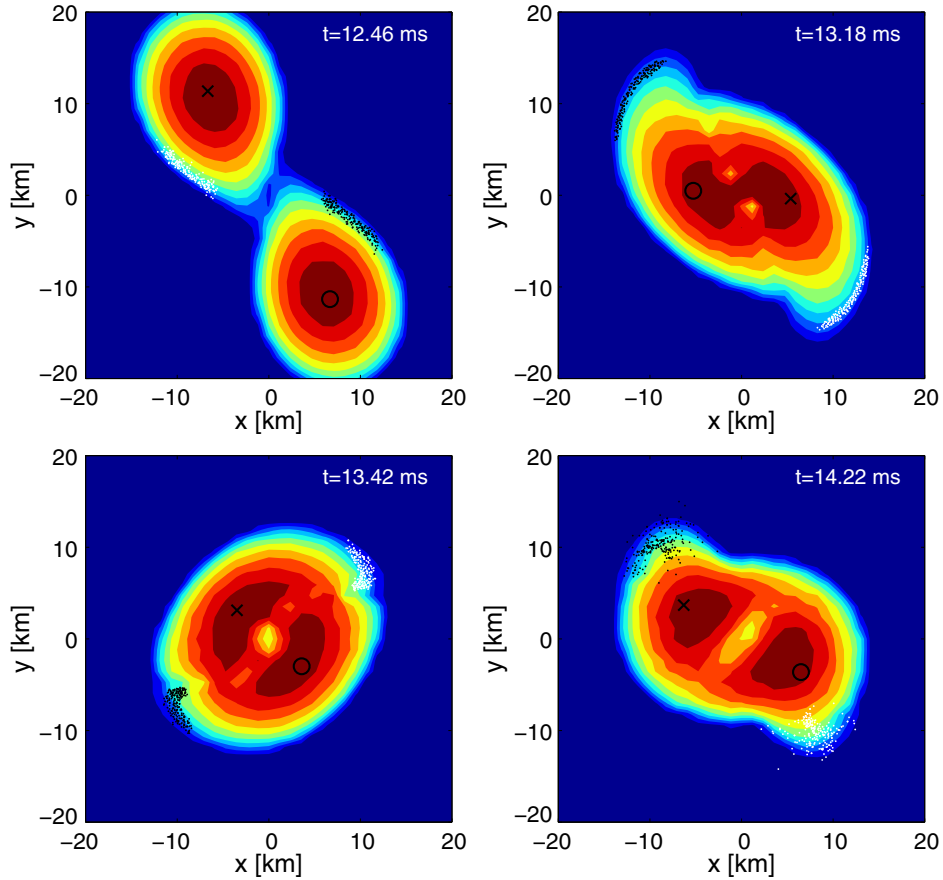


FIG. 2 (color online). Rest-mass density evolution in the equatorial plane for the $1.35\text{--}1.35M_{\odot}$ merger with the DD2 EOS (rotation counterclockwise). (The rest-mass density is shown with a variable linear scale relative to ρ_{\max} . A low number of contour levels is chosen for illustrative reasons; the underlying simulation data is smoother than it appears with the chosen color coding.) Black and white dots trace the positions of selected fluid elements of the antipodal bulges, which within approximately one millisecond complete one orbit (compare times of the right panels). The orbital motion of this pattern of spiral deformation produces the f_{spiral} peak in the GW spectrum at $2/(1\text{ ms})$ (Fig. 1). The cross and the circle mark the double cores, which rotate significantly faster than the antipodal bulges represented by the dots (compare times of the different panels).

spectrum and considering different time intervals, we find that the presence of the f_{spiral} peak agrees with the appearance and duration of the spiral deformation of the remnant.

In the upper-right panel of Fig. 2, the spiral deformation can be seen to initially reach deep inside the remnant. We approximately determine the amount of matter which belongs to the two antipodal bulges that are rotating more slowly compared to the double cores. This matter amounts to several tenths of M_{\odot} and is thus sufficient to explain the strength of the f_{spiral} GW peak. In addition, we find that the f_{spiral} GW peak can be roughly reproduced in a toy model, where the two bulges orbit as point particles around the central double-core structure for a duration of a few milliseconds. Note that this toy model differs significantly from the one in Ref. [37], which considers only the two cores to be contributing to the GW signal and considers only a single instantaneous orbital frequency of the system.

Furthermore, we take advantage of the quadrupole formalism to compute GW spectra considering only certain parts of the remnant, which are defined by using either a density or a spatial cutoff. In Fig. 3 we demonstrate that the dominant quadrupole f_{peak} frequency is generated mainly by those regions of the remnant which encompass densities exceeding 50% of the instantaneous maximum density ρ_{\max} . In contrast, most power of the f_{spiral} peak originates from densities below $0.5\rho_{\max}$, which corresponds to the outer parts of the remnant, where the two bulges form (see Fig. 2). Similar conclusions are reached when a spatial cutoff instead of a density cutoff is used.

In models where f_{spiral} dominates over f_{2-0} (see our classification below), the presence of the two rotating antipodal bulges explains the appearance of a particular *low-frequency modulation* with $f_{\text{peak}} - f_{\text{spiral}}$ seen in the time evolution of the central lapse function (blue curve in Fig. 4), of ρ_{\max} , of the size of the central remnant and of the separation between the two cores. The same modulation

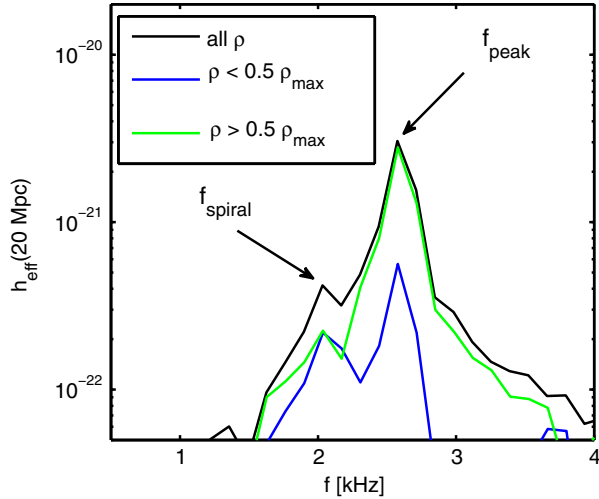


FIG. 3 (color online). Early postmerger spectrum of $1.35\text{--}1.35 M_{\odot}$ merger with the DD2 EOS (black) and GW spectra computed for fluid elements with densities above or below 50% of the maximum density.

occurs as a beat frequency in the time evolution of the GW amplitude (see e.g. Fig. 1 in Ref. [19]). The low-frequency modulation with $f_{\text{peak}} - f_{\text{spiral}}$ can coexist with the quasiradial oscillation, which has a higher frequency f_0 .

The above modulation frequency is associated with the rotating bulges and is explained as follows: the central remnant forms an elongated structure, around which the two bulges rotate, lagging behind. The characteristics of the remnant are modulated depending on the orientation of

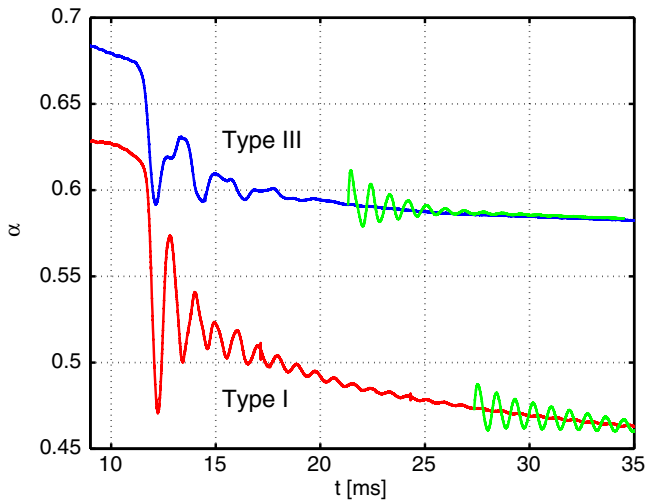


FIG. 4 (color online). Evolution of the central lapse function for $1.35\text{--}1.35 M_{\odot}$ mergers with the NL3 (blue) and LS220 (red) EOS (time-shifted to the same merging time). The evolution of models with an added quasiradial perturbation is shown (green). For type II and type III mergers, like the model with the NL3 EOS shown here, one recognizes a low-frequency oscillation in addition to the quasiradial mode, which is also strongly excited in the perturbed model at late times.

the antipodal bulges with respect to the double cores: the compactness is smaller, the central lapse function larger and the GW amplitude maximal when the bulges and the cores are aligned (lower-right panel in Fig. 2). Instead, when the bulges and the cores are orthogonal to each other (lower-left panel), the compactness is largest, the central lapse function smaller and the GW amplitude small.

In a frame corotating with the central remnant (i.e. a frame rotating with frequency $\sim f_{\text{peak}}/2$ with respect to the inertial frame), the two bulges counterrotate with a frequency of $\Delta f = (f_{\text{peak}} - f_{\text{spiral}})/2$ which corresponds to the $f_{\text{peak}} - f_{\text{spiral}}$ modulation in various quantities (notice the π -symmetry of the system, which compensates the factor $1/2$).

Our simulations also show that the rotational frequency of the centers of the double-core structure, although not constant, significantly exceeds $f_{\text{spiral}}/2$ at any time. Hence, variations in the angular frequency of the double-core structure alone are not sufficient to interpret the secondary peaks, and the simultaneous presence of both f_{peak} and f_{spiral} cannot simply be attributed to a single instantaneous angular frequency of the system, as suggested in Refs. [36–38]. We note that the simultaneous presence of two (orbital) frequencies (of the double-core structure and of the antipodal bulges) naturally explains strong time variations of the instantaneous GW frequency as seen in many simulations, e.g. Refs. [33,38].

Also, a single, initially strongly varying instantaneous frequency as explanation for the peaks in the spectrum is incompatible with the fact that the most pronounced peak in the spectrum occurs already at early times. Initially, the instantaneous frequency strongly oscillates around the frequency of the dominant peak. According to the picture of Ref. [38], peaks form at frequencies at which the instantaneous frequency spends most time, i.e. at the extrema of the instantaneous frequency. In this case, one would thus not expect that the dominant f_{peak} is strong at early times. In contrast, in our simulations the dominant f_{peak} has a substantial strength if one considers only the first few milliseconds.

III. CLASSIFICATION OF POSTMERGER DYNAMICS AND GW EMISSION

We have applied the above analysis tools (GW spectra, determination of f_0 from perturbed models, rotational frequency of the antipodal bulges and of the double cores, GW spectra with different cutoff densities) for a number of representative models, varying the binary mass and stiffness of the EOS. The results fully confirm the generic picture described above. Considering this larger set of models, we identify, based on the relative strength between f_{2-0} and f_{spiral} , three different types of postmerger dynamics and GW spectra for remnants which survive for more than several milliseconds.

- (i) *Type I*: When the total binary mass M_{tot} is not too far from the threshold mass for prompt quasiradial collapse of the remnant for a given EOS [32], the evolution of the central lapse function (and of ρ_{max}) is dominated by a very strong *quasiradial oscillation* of the remnant, see lower curve in Fig. 4. For such models the two initial NSs are more centrally condensed, and they merge with higher impact velocity (Fig. 3 in Ref. [56]). Because of the strongly excited quasiradial oscillation, f_{2-0} is the strongest secondary peak in the GW spectrum, while f_{spiral} is much weaker, likely because for more compact NSs the formation of the spiral pattern is less pronounced. There can be a partial overlap between f_{2-0} and f_{spiral} , see red curve in Fig. 1.
- (ii) *Type II*: For intermediate total binary masses, f_{2-0} and f_{spiral} have a *comparable strength* in the GW spectrum and the two types of secondary peaks are *well separated* (see black curve in Fig. 1). This is the generic type, which we discussed earlier in our reference model. In characteristic quantities such as the central lapse function, the quasiradial oscillation frequency f_0 as well as the low-frequency modulation with $f_{\text{peak}} - f_{\text{spiral}}$ are clearly noticeable.
- (iii) *Type III*: When the total binary mass M_{tot} is significantly below the threshold mass for quasiradial collapse, the time evolution of the central lapse function (as well as that of ρ_{max} and of the radius of the remnant) is dominated by the $f_{\text{peak}} - f_{\text{spiral}}$ modulation that we explained in the previous section as a result of the rotating spiral pattern with the two antipodal bulges. In the evolution of the central lapse function, this modulation typically has a smaller amplitude than type-I variations (Fig. 4). The quasiradial oscillation f_0 is also present, but with much smaller amplitude than the dominant modulation, because the smaller NS compactness implies a smaller impact velocity and thus a weaker excitation of the quasiradial oscillation. The smaller NS compactness also allows for a stronger spiral deformation. Consequently, the dominant secondary peak in the GW spectrum is f_{spiral} , while f_{2-0} is either very weak or hidden inside the background (see blue curve in Fig. 1).

For a given EOS, there is a continuous transition between the different types of postmerger dynamics depending on the total binary mass. Types I and III are the limiting cases of the more generic type II. Notice that since the threshold for quasiradial collapse is EOS dependent [32], the different types cover a different mass range for each EOS. For a total binary mass of $M_{\text{tot}} = 2.7 M_{\odot}$ all three types are possible depending on the EOS, where very soft EOSs yield type-I mergers and very stiff EOSs lead to type-III dynamics (see GW spectra in Fig. 1). This is also shown in

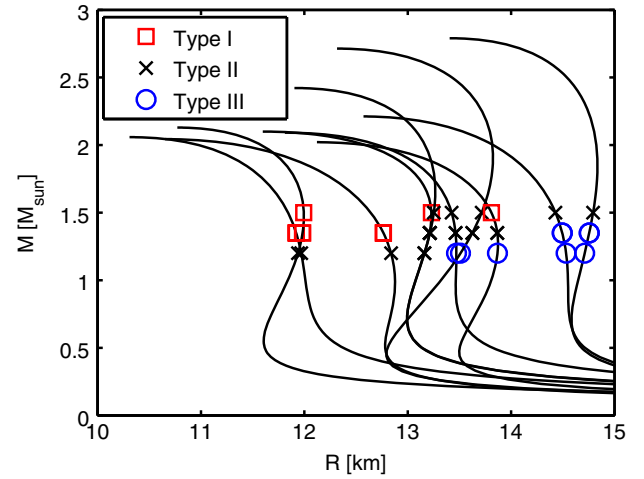


FIG. 5 (color online). Different types of postmerger dynamics and GW emission of the different merger models visualized by the mass-radius relations of nonrotating NSs of the EOSs considered in this work. The outcome of a given calculation with M_{tot} is shown as a symbol at $M_{\text{tot}}/2$ plotted on the mass-radius relation of the EOS employed in the simulation. Red squares indicate type I, black crosses stand for type II, and blue circles mark type III. See text for definitions of different types of postmerger dynamics and GW emission.

Fig. 5, which provides an overview of the types of the different models considered in this study. The type of a given simulation is indicated by a symbol plotted at the mass of the individual inspiraling stars on the mass-radius relation of the EOS which was used in the calculation (e.g. the results of the 1.2–1.2 M_{\odot} binaries are displayed at the radii of NSs with 1.2 M_{\odot}). As described, there is a continuous transition between the different classes, which is why one should consider the transitions between the different types in Fig. 5 as being tentative, and a slightly different classification may be possible at the borders between the different types. More quantitative definitions of the types may be useful in the future. Still, one can clearly identify a diagonal band of type-II mergers for intermediate binary masses, and also the binary setups leading to the limiting cases of type I or type III are seen to form roughly diagonal bands.

For $2.4 M_{\odot} \leq M_{\text{tot}} \leq 3.0 M_{\odot}$, we find that f_{spiral} typically ranges between $f_{\text{peak}} - 0.5$ kHz and $f_{\text{peak}} - 0.9$ kHz, while f_{2-0} ranges between $f_{\text{peak}} - 0.9$ kHz and $f_{\text{peak}} - 1.3$ kHz. This property will be useful for identifying either f_{2-0} or f_{spiral} (or both) in future GW observations. Furthermore, we find that $f_{\text{peak}} - f_{2-0} (= f_0)$ decreases with increasing M_{tot} in all models for which f_{2-0} is clearly present, in agreement with the fact that the quasiradial frequency decreases near the threshold to collapse. This observation may be useful to estimate the proximity to prompt gravitational collapse. Very near the threshold, one thus may expect $f_{2-0} \rightarrow f_{\text{peak}}$. In contrast, $f_{\text{peak}} - f_{\text{spiral}}$ typically increases with increasing M_{tot} , and

above the threshold to collapse a spiral pattern during the dynamical collapse could still produce a weak peak in the GW spectrum, as in Ref. [57].

IV. EMPIRICAL RELATIONS FOR DOMINANT AND SECONDARY PEAK FREQUENCIES

For our sample of EOSs, Fig. 6 shows f_{peak} , f_{spiral} and f_{2-0} as functions of the compactness M/R of the *non-spinning, individual NSs* (at infinite separation) for $M_{\text{tot}} = 2.7 M_{\odot}$ (with the compactness in units of $c = G = 1$). We find strong correlations that can be described by the following quadratic fits:

$$f_{\text{peak}}[\text{kHz}] = 199(M/R)^2 - 28.1(M/R) + 2.33, \quad (1)$$

$$f_{\text{spiral}}[\text{kHz}] = 358(M/R)^2 - 82.1(M/R) + 6.16, \quad (2)$$

$$f_{2-0}[\text{kHz}] = 392(M/R)^2 - 88.3(M/R) + 5.95. \quad (3)$$

The maximum deviations of the data used for these fits are 140 Hz, 86 Hz and 153 Hz for f_{peak} , f_{spiral} and f_{2-0} , respectively. If the compactness is determined from a measured frequency by inverting Eqs. (1)–(3), these maximum deviations imply errors of 3%, 3% and 4% in the compactness for f_{peak} , f_{spiral} and f_{2-0} , respectively. (Note that the slope in the relation for f_{spiral} is somewhat flatter.) Thus, the accuracy of these empirical relations is similarly good for f_{peak} and f_{spiral} and only slightly worse for f_{2-0} . Even if one assumes an error of 10% uncertainty in the frequency determinations from an actual GW observation (which at least for f_{peak} is too pessimistic [40]), this would add 6%, 8% and 5% error, respectively, in the compactness

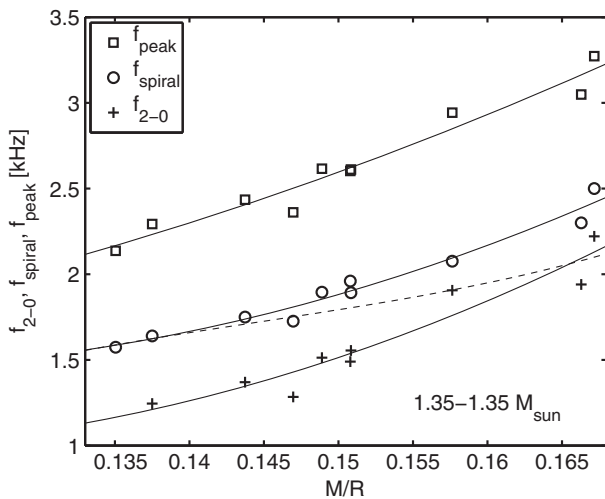


FIG. 6. f_{peak} , f_{spiral} and f_{2-0} for mergers with ten different EOSs and $M_{\text{tot}} = 2.7 M_{\odot}$ vs the compactness M/R for non-rotating, single NSs. Solid lines show empirical relations. The dashed line is taken from Ref. [37] (see text for explanations).

determination. Similar empirical relations hold, with varying accuracy, for each total binary mass. As already pointed out and explained for f_{peak} in Ref. [29], even better empirical relations are obtained for $1.35\text{--}1.35 M_{\odot}$ binaries when the above three frequencies are examined as functions of the radius of a nonrotating NS with $1.6 M_{\odot}$. Notice that all three empirical relations follow similar trends.

Because the frequency of the f_{spiral} peak agrees with the rotational frequency of the antipodal bulges near the surface of the remnant, a scaling of f_{spiral} with the compactness is not unexpected, which explains the similarities with the behavior of f_{peak} . In fact, for the models with $M_{\text{tot}} = 2.7 M_{\odot}$ one finds very tight relations between f_{spiral} and f_{peak} , and between f_{2-0} and f_{peak} .

A measurement of f_{peak} is already sufficient to accurately constrain the radii of nonrotating NSs (and thus the EOS) [28,29,39] if the total mass has been obtained accurately from the inspiral signal, as is likely to be the case for distances within which f_{peak} has the required signal-to-noise ratio to be detectable with second-generation interferometers (e.g. Refs. [40,58]). A detection of the weaker secondary peaks, with similar M/R dependence as f_{peak} , could further optimize the constraints on the EOS.

Figure 7 displays f_{spiral} as a function of the compactness M/R of the nonspinning, individual NSs for different EOSs and for different M_{tot} in the range $2.4\text{--}3.0 M_{\odot}$ (the most likely range of total binary masses [59]). The same figure, but for f_{2-0} , is shown in Fig. 8. Especially for f_{spiral} , we discover that there exist tight relations between the compactness and the secondary frequencies for fixed binary masses.

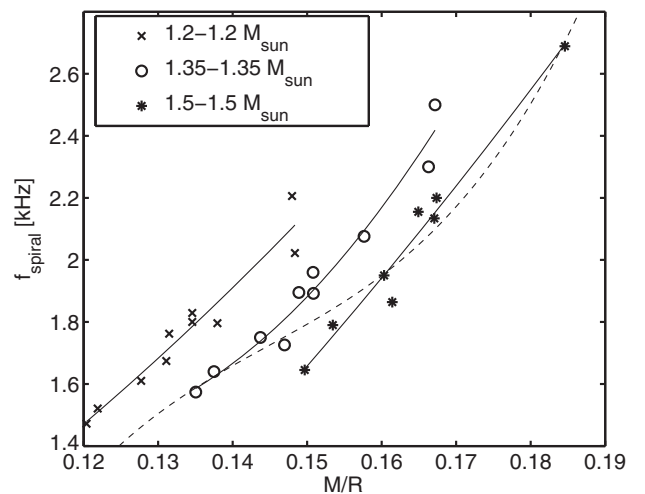


FIG. 7. f_{spiral} vs the compactness, but for different binary masses. Solid lines show empirical relations. The dashed line is taken from Ref. [37] (see text for explanations).

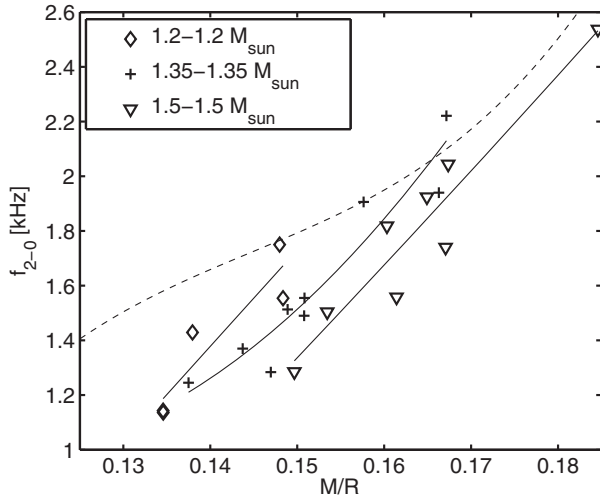


FIG. 8. f_{2-0} vs the compactness, but for different binary masses. Solid lines show empirical relations. The dashed line is taken from Ref. [37] (see text for explanations).

In Refs. [36,37] the existence of a single, universal, mass-independent relation (dashed line in Figs. 6 to 8) between the frequency of the strongest secondary GW peak (denoted there as f_1) and M/R was proposed (there was no distinction of two different secondary peaks, as we find here). However, this result was based on using a limited set of five EOSs of soft or moderate stiffness (with corresponding maximum masses of nonrotating NS only up to $2.2 M_\odot$), as well as on different chosen mass ranges for each EOS with a spread of only $0.2 M_\odot$ in the total binary mass.

In contrast to Refs. [36,37], within our larger sample of ten EOSs (that includes stiff EOSs with maximum masses reaching up to $2.8 M_\odot$) and for a more representative total binary mass range of $2.4\text{--}3.0 M_\odot$ (the same for all EOSs), *such a mass-independent, universal relation does not exist* (dashed curve in Figs. 6 to 8). Figures 7 and 8 show that there is a large spread of up to 500 Hz both in the f_{spiral} vs M/R relation and in the f_{2-0} vs M/R relation for an expected range of total binary masses, and that the data cannot be described by a single function (dashed line). Even if one consistently chooses the strongest secondary peak among f_{spiral} and f_{2-0} in each case, there does not exist a mass-independent universal relation with compactness. The absence of a mass-independent, universal relation implies that the procedure for EOS constraints and binary mass determinations as described in Refs. [36,37] is not valid. Notice also that in Fig. 6, the relation proposed in Refs. [36,37] describes either f_{spiral} at low compactness or f_{2-0} at high compactness, which is consistent with the expectations from our classification scheme describing which of the secondary peaks dominates the GW spectrum. However, for intermediate values of the compactness the merger will be of type II, and both types of secondary peaks can be present with comparable amplitude, which further

complicates the definition of a single f_1 frequency, as was assumed in Refs. [36,37]. We also point out that detections of the secondary peaks may possibly be less accurate since the peaks are broader in comparison to the main peak and sometimes do not stand out clearly.

Instead of a universal mass-independent relation, we find that there exist useful empirical relations only for fixed binary masses (shown as thin line segments in Figs. 7 and 8), such as the case shown in Fig. 6. For merger events sufficiently close to allow a detection of postmerger GW peaks, an accurate determination of the binary masses is expected (e.g. Ref. [58]). With the upcoming GW detectors, the dominant postmerger oscillation frequency has been shown to be detectable with high precision for mergers within several Mpc using a morphology-independent burst analysis [40]. Significant improvements in the detectability are possible for more sophisticated search algorithms, e.g. matched filtering, which, however, require detailed modeling of the expected signal features to which our present study should contribute by clarifying the relation between dominant and secondary GW peaks. Depending on the highly uncertain binary NS merger rate, a detection may succeed with Advanced LIGO/Virgo [40] or with discussed upgrades [42–44] (the latter may yield a sensitivity increase at high frequencies of a factor 3). Since the strength of the secondary features is somewhat lower than that of the dominant peak (in relation to the anticipated instrument noise curves), direct detections of the secondary GW peaks can be expected for the planned Einstein Telescope [36,37,40,41,45], unless the merger rate is on the more optimistic side as defined in Ref. [3], which may enable an earlier observation. The exact detection rate of secondary peaks is hard to quantify not only because of the uncertain merger rate, but also because of the varying strength and prominence of the secondary peaks depending on the exact model.

In this study we focus on equal-mass binaries. To assess the impact of the binary mass ratio, we perform additional simulations for $1.3\text{--}1.4 M_\odot$ binaries with the DD2 EOS [51,52], LS220 EOS [55] and NL3 EOS [51,54]. The GW spectra in comparison to the symmetric case exhibit qualitatively the same features. The frequencies of the three different GW peaks are very similar to the equal-mass case with only small deviations of at most 2%, 4% and 6% for f_{peak} , f_{spiral} and f_{2-0} , respectively, compared to the results of the $1.35\text{--}1.35 M_\odot$ binaries. Also, the evolution of the central lapse function shows qualitatively the same behavior. This suggests that, at least for moderately asymmetric binaries, our classification scheme applies as well. It also implies that the frequency dependencies and their implications discussed above hold quantitatively with a good accuracy, although the secondary peaks show somewhat larger deviations, which may impede their use for EOS constraints if the mass ratio is not measured accurately from the GW inspiral signal.

V. SUMMARY AND OUTLOOK

We present a unified picture of the postmerger dynamics and GW emission in binary NS mergers. Aside from the secondary GW peak produced by a nonlinear coupling between quadrupole and quasiradial oscillations, we identify a new mechanism by which another secondary peak in the GW spectrum appears: it is caused by a spiral deformation, the pattern of which rotates more slowly with respect to the double-core structure of the inner remnant. Based on the presence and strength of these two secondary peaks, we introduce a classification scheme of the postmerger dynamics and of the GW spectrum. Moreover, the new mechanism allows for an explanation of low-frequency modulations seen in a number of physical characteristics of the remnant, such as the central lapse function, the maximum density, the separation between the two cores and the GW amplitude.

We find that for fixed total binary masses, the frequencies of the secondary GW peaks show a tight relation with the compactness of nonrotating, individual NSs. These relations follow similar trends as the relation between the dominant peak frequency and the NS compactness (or, equivalently, the NS radius, as shown in Refs. [28,29]). For the dominant peak and for the secondary peak associated with the orbital motion of the antipodal bulges, we find the relations between frequency and compactness to be similarly tight. We rule out the existence of a universal, mass-independent relation for secondary peaks as proposed in Refs. [36,37].

Identifying the type of a merger event and especially the determination of the quasiradial frequency by $f_0 = f_{\text{peak}} - f_{2-0}$ can lead to an estimate of the threshold

mass for black hole collapse and thus to an estimate of the maximum mass of nonrotating NSs (see Ref. [32]). The insights from our classification scheme and from the frequency dependencies found here can provide a basis for constructing future GW template waveforms, increasing the changes for the observation of the postmerger GW emission [40].

We will further investigate whether the frequency differences between the dominant and secondary peaks can clarify the nature of a detected secondary peak or if detailed comparisons between spectra of different binary masses are needed. We will also explore unequal-mass binaries (anticipating a strong impact of the mass ratio on f_{spiral}) and analyze the relevance of our classification scheme for the mass ejection and torus formation of NS mergers and for accompanying phenomena, such as r-process nucleosynthesis [60,61], electromagnetic counterparts [62–64] and short gamma-ray bursts [61,65].

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- [1] G. M. Harry and LIGO Scientific Collaboration, *Classical Quantum Gravity* **27**, 084006 (2010).
 - [2] F. Acernese, P. Amico, M. Alshourbagy, F. Antonucci, S. Aoudia, S. Avino, D. Babusci, G. Ballardin, F. Barone, L. Barsotti *et al.*, *Classical Quantum Gravity* **23**, S635 (2006).
 - [3] J. Abadie *et al.*, *Classical Quantum Gravity* **27**, 173001 (2010).
 - [4] M. D. Duez, *Classical Quantum Gravity* **27**, 114002 (2010).
 - [5] N. Andersson, V. Ferrari, D. I. Jones, K. D. Kokkotas, B. Krishnan, J. S. Read, L. Rezzolla, and B. Zink, *Gen. Relativ. Gravit.* **43**, 409 (2011).
 - [6] T. W. Baumgarte and S. L. Shapiro, *Numerical Relativity: Solving Einstein’s Equations on the Computer* (Cambridge University Press, Cambridge, 2010).
 - [7] J. A. Faber and F. A. Rasio, *Living Rev. Relativity* **15**, 8 (2012).
 - [8] L. Rezzolla and O. Zanotti, *Relativistic Hydrodynamics* (Oxford University Press, New York, 2013).
 - [9] X. Zhuge, J. M. Centrella, and S. L. W. McMillan, *Phys. Rev. D* **50**, 6247 (1994).
 - [10] M. Ruffert, H.-T. Janka, and G. Schaefer, *Astron. Astrophys.* **311**, 532 (1996).
 - [11] S. Rosswog, M. B. Davies, F.-K. Thielemann, and T. Piran, *Astron. Astrophys.* **360**, 171 (2000).
 - [12] R. Oechslin, S. Rosswog, and F.-K. Thielemann, *Phys. Rev. D* **65**, 103005 (2002).
 - [13] M. Shibata and K. Uryū, *Prog. Theor. Phys.* **107**, 265 (2002).
 - [14] J. A. Faber, P. Grandclément, and F. A. Rasio, *Phys. Rev. D* **69**, 124036 (2004).
 - [15] M. Shibata, *Phys. Rev. Lett.* **94**, 201101 (2005).
 - [16] M. Shibata, K. Taniguchi, and K. Uryū, *Phys. Rev. D* **71**, 084021 (2005).
 - [17] M. Shibata and K. Taniguchi, *Phys. Rev. D* **73**, 064027 (2006).
 - [18] R. Oechslin, H.-T. Janka, and A. Marek, *Astron. Astrophys.* **467**, 395 (2007).

- [19] R. Oechslin and H.-T. Janka, *Phys. Rev. Lett.* **99**, 121102 (2007).
- [20] M. Anderson, E. W. Hirschmann, L. Lehner, S. L. Liebling, P. M. Motl, D. Neilsen, C. Palenzuela, and J. E. Tohline, *Phys. Rev. D* **77**, 024006 (2008).
- [21] Y. T. Liu, S. L. Shapiro, Z. B. Etienne, and K. Taniguchi, *Phys. Rev. D* **78**, 024012 (2008).
- [22] L. Baiotti, B. Giacomazzo, and L. Rezzolla, *Phys. Rev. D* **78**, 084033 (2008).
- [23] K. Kiuchi, Y. Sekiguchi, M. Shibata, and K. Taniguchi, *Phys. Rev. D* **80**, 064037 (2009).
- [24] B. Giacomazzo, L. Rezzolla, and L. Baiotti, *Phys. Rev. D* **83**, 044014 (2011).
- [25] N. Stergioulas, A. Bauswein, K. Zagkouris, and H.-T. Janka, *Mon. Not. R. Astron. Soc.* **418**, 427 (2011).
- [26] K. Hotokezaka, K. Kyutoku, H. Okawa, M. Shibata, and K. Kiuchi, *Phys. Rev. D* **83**, 124008 (2011).
- [27] Y. Sekiguchi, K. Kiuchi, K. Kyutoku, and M. Shibata, *Phys. Rev. Lett.* **107**, 051102 (2011).
- [28] A. Bauswein and H.-T. Janka, *Phys. Rev. Lett.* **108**, 011101 (2012).
- [29] A. Bauswein, H.-T. Janka, K. Hebeler, and A. Schwenk, *Phys. Rev. D* **86**, 063001 (2012).
- [30] V. Paschalidis, Z. B. Etienne, and S. L. Shapiro, *Phys. Rev. D* **86**, 064032 (2012).
- [31] 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).
- [32] A. Bauswein, T. W. Baumgarte, and H.-T. Janka, *Phys. Rev. Lett.* **111**, 131101 (2013).
- [33] K. Hotokezaka, K. Kiuchi, K. Kyutoku, T. Muranushi, Y. Sekiguchi, M. Shibata, and K. Taniguchi, *Phys. Rev. D* **88**, 044026 (2013).
- [34] D. Neilsen, S. L. Liebling, M. Anderson, L. Lehner, E. O'Connor, and C. Palenzuela, *Phys. Rev. D* **89**, 104029 (2014).
- [35] S. Bernuzzi, T. Dietrich, W. Tichy, and B. Brügmann, *Phys. Rev. D* **89**, 104021 (2014).
- [36] K. Takami, L. Rezzolla, and L. Baiotti, *Phys. Rev. Lett.* **113**, 091104 (2014).
- [37] K. Takami, L. Rezzolla, and L. Baiotti, *Phys. Rev. D* **91**, 064001 (2015).
- [38] W. Kastaun and F. Galeazzi, *Phys. Rev. D* **91**, 064027 (2015).
- [39] A. Bauswein, N. Stergioulas, and H.-T. Janka, *Phys. Rev. D* **90**, 023002 (2014).
- [40] J. Clark, A. Bauswein, L. Cadonati, H.-T. Janka, C. Pankow, and N. Stergioulas, *Phys. Rev. D* **90**, 062004 (2014).
- [41] C. Messenger, K. Takami, S. Gossan, L. Rezzolla, and B. S. Sathyaprakash, *Phys. Rev. X* **4**, 041004 (2014).
- [42] S. Hild, *Classical Quantum Gravity* **29**, 124006 (2012).
- [43] R. X. Adhikari, *Rev. Mod. Phys.* **86**, 121 (2014).
- [44] J. Miller, L. Barsotti, S. Vitale, P. Fritschel, M. Evans, and D. Sigg, *Phys. Rev. D* **91**, 062005 (2015).
- [45] S. Hild, S. Chelkowski, A. Freise, J. Franc, N. Morgado, R. Flaminio, and R. DeSalvo, *Classical Quantum Gravity* **27**, 0815003 (2010).
- [46] J. Isenberg and J. Nester, in *General Relativity and Gravitation* (Plenum Press, New York, 1980), p. 23.
- [47] J. R. Wilson, G. J. Mathews, and P. Marronetti, *Phys. Rev. D* **54**, 1317 (1996).
- [48] G. Faye and G. Schäfer, *Phys. Rev. D* **68**, 084001 (2003).
- [49] A. Bauswein, H.-T. Janka, and R. Oechslin, *Phys. Rev. D* **82**, 084043 (2010).
- [50] D. R. Lorimer, *Living Rev. Relativity* **11**, 8 (2008).
- [51] M. Hempel and J. Schaffner-Bielich, *Nucl. Phys.* **A837**, 210 (2010).
- [52] S. Typel, G. Röpke, T. Klähn, D. Blaschke, and H. H. Wolter, *Phys. Rev. C* **81**, 015803 (2010).
- [53] Other references in the literature call the dominant frequency f_2 .
- [54] G. A. Lalazissis, J. König, and P. Ring, *Phys. Rev. C* **55**, 540 (1997).
- [55] J. M. Lattimer and F. D. Swesty, *Nucl. Phys.* **A535**, 331 (1991).
- [56] A. Bauswein, S. Goriely, and H.-T. Janka, *Astrophys. J.* **773**, 78 (2013).
- [57] K. Kiuchi, Y. Sekiguchi, M. Shibata, and K. Taniguchi, *Phys. Rev. Lett.* **104**, 141101 (2010).
- [58] C. L. Rodriguez, B. Farr, V. Raymond, W. M. Farr, T. B. Littenberg, D. Fazi, and V. Kalogera, *Astrophys. J.* **784**, 119 (2014).
- [59] J. M. Lattimer, *Annu. Rev. Nucl. Part. Sci.* **62**, 485 (2012).
- [60] J. M. Lattimer, F. Mackie, D. G. Ravenhall, and D. N. Schramm, *Astrophys. J.* **213**, 225 (1977).
- [61] D. Eichler, M. Livio, T. Piran, and D. N. Schramm, *Nature (London)* **340**, 126 (1989).
- [62] L.-X. Li and B. Paczyński, *Astrophys. J. Lett.* **507**, L59 (1998).
- [63] S. R. Kulkarni, [arXiv:astro-ph/0510256](https://arxiv.org/abs/astro-ph/0510256).
- [64] B. D. Metzger, G. Martínez-Pinedo, S. Darbha, E. Quataert, A. Arcones, D. Kasen, R. Thomas, P. Nugent, I. V. Panov, and N. T. Zinner, *Mon. Not. R. Astron. Soc.* **406**, 2650 (2010).
- [65] B. Paczynski, *Astrophys. J. Lett.* **308**, L43 (1986).