

Tidal Resonance in Binary Neutron Star Inspirals: A High-Precision Study in Numerical Relativity

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We investigate the tidal resonance of the fundamental (f) mode in spinning neutron stars, robustly tracing the onset of the excitation to its saturation, using numerical relativity for the first time. We performed long-term (≈ 15 orbits) fully relativistic simulations of a merger of two highly and retrogradely spinning neutron stars. The resonance window of the f mode is extended by self-interaction, and the nonlinear resonance continues up to the final plunging phase. We observe that the quasicircular orbit is maintained throughout since the dissipation of orbit motion due to the resonance is coherent with that due to gravitational waves. The f -mode resonance causes a variation in the stellar spin of $\gtrsim 6.3\%$ in the linear regime and much more as $\sim 33\%$ during the later nonlinear regime. At the merger, a phase shift of $\lesssim 40$ rad is rendered in the gravitational waveform as a consequence of the angular momentum and energy transfers into the neutron star oscillations.

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Introduction—Binary neutron star (BNS) mergers are among the most important systems to learn about the equation of state (EOS) of nuclear matter through measuring and analyzing the delicate imprint of tidal effects on gravitational waves (GWs); see, e.g., the reviews [1–5] and the references therein. How a neutron star (NS) responds to the tidal field sourced by its companion is sensitive to its internal structure that is determined by the EOS [6]. The response function, in general, depends on the tidal forcing frequency (about twice the orbital frequency) with the zeroth-order term being

$$\tilde{\Lambda} = \frac{16}{13M^5} [(M_1 + 11M_2)M_1^4\Lambda_1 + (M_2 + 11M_1)M_2^4\Lambda_2], \quad (1)$$

which is linked with the Love numbers of the participating NSs [7–11]. Here Λ_1 and Λ_2 are the tidal deformabilities of the two NSs, while M_1 and M_2 are their respective (Arnowitt-Deser-Misner) masses with the total mass of the binary given by $M = M_1 + M_2$. This tidal quantity has been constrained (though not significantly) by the first BNS event GW170817 [12–18]. As the binary approaches the

merger, the frequency-dependent response becomes enhanced to manifest dynamical tides [19–22]. This aspect has been incorporated in waveform modeling by interpreting dynamical tides as the excitation of the fundamental (f) mode of stellar oscillations [19–21, 23–26] or the somehow resummed contributions of higher post-Newtonian (PN) orders [27–30].

Developing accurate gravitational waveform models is essential to extract the tidal effects from the detected GWs, which requires one to study dynamical tides in detail [31–33]. To this end, a variety of analytic efforts has been devoted to advancing the knowledge about dynamically perturbed NSs [34–39], incorporating the high PN effects [40–43], improving on the modeling of tidal resonance [24–26, 44–47], and devising an adequate effective-one-body treatment for tidal effects [19, 23, 29, 48, 49]. However, analytic attempts have limited power in handling the nonlinear tidal response of late-inspiral BNSs, where numerical relativity (NR) is the unique tool to resolve the dynamics involved [49–57].

The state-of-the-art waveform models for BNSs include: (I) tidally tuned effective-one-body (EOB) models TEOBResumS [27, 29, 48, 58, 59] and SEOBNRv*T [19, 23, 49] (see Refs. [56, 60] for some comparison surveys); and (II) point-particle EOB or phenomenological phase models (IMPheom; [61, 62]) + close form of tide-related phase shift [63–70]. The construction of these models appeals to calibration with NR results in one way or another. Therefore, their validity could be limited to the parameter space covered by the NR waveforms in the literature. In this spirit, numerical studies of unexplored

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BNS parameters are essential to guide future waveform modeling.

In particular, mergers of retrogradely spinning NSs are valuable for accurately studying how the f -mode resonance develops to a nonlinear phase and its saturation, as the reduced mode frequency [71–74] can ensure resonance before the plunge. Some literature, such as [75–80], has simulated BNSs with antialigned spins. However, the spin parameters considered in these studies only allowed for resonance to occur in the merger phase, when the system becomes essentially one body, thus obscuring the f -mode resonance effects. In this Letter, we aim to provide a detailed analysis of f -mode resonance, self-consistently demonstrating its nonlinear excitation and backreaction to the NS. For this purpose, we performed high-precision long-term general relativity (GR) simulations (covering the last ≈ 15 orbits), enabled by cutting-edge numerical techniques, to robustly resolve this relativistic hydrodynamical process for the first time. The geometrical units are assumed throughout this Letter, unless explicitly stated otherwise.

f-mode effects—Dynamical tides can be effectively modeled as a set of spherical harmonic oscillators in linear order, each representing a star’s quasinormal mode (though see Ref. [81]). When the star is immersed in a tidal field exerted by a companion, these oscillators are enforced, with a mode resonating when the tidal frequency (Ω_{tid}) matches its characteristic frequency [82,83]. For the quasicircular inspiral and focusing on the dominant tidal effect, the forcing rate approximates the instantaneous frequency (ω_{gw}) of emitted GWs (i.e., $\Omega_{\text{tid}} \simeq \omega_{\text{gw}}$). Among the spectra of stellar oscillations, the f mode couples most strongly to the tidal field, and thus its (perhaps resonant; see, e.g., Ref. [84] for discussions) excitation will adjust the binary motion by efficiently tapping off the binary’s orbital energy to fuel its kinetic energy [85–89]. As a result, the merger will occur noticeably earlier, and the resulting waveform dephasing is important for data analysis of future GW observatories [20,25,32,33,90].

If the amplitude of a mode grows enough, the nonlinear self-interaction will become influential [44] (see also Ref. [91]). To encode such an effect in the evolution of f mode’s amplitude (c_f), we phenomenologically introduce self-interacting terms to the Hamiltonian,

$$\mathcal{H} = \frac{1}{2} \dot{c}_f^2 + \frac{\omega_f^2 c_f^2}{2} + \frac{\eta}{4} c_f^4, \quad (2)$$

while leaving out the contributions of $\mathcal{O}(c_f^5)$. Here, the overhead dot denotes the time derivative, ω_f is the f mode’s frequency, and η is a constant setting the degree of the self-interaction. The associated evolution equation is written as [92]

$$\ddot{c}_f + \omega_{f,\text{eff}}^2 c_f = \mathcal{F}_{\text{tid}} = A \cos(\Omega_{\text{tid}} t) \quad (3)$$

with $\omega_{f,\text{eff}}^2 = \omega_f^2 + \eta c_f^2$ and the magnitude of tidal force A . If the f mode gets excited to a nonlinear regime such that $\omega_f^2 \gg \eta c_f^2$ no longer holds, the Hamiltonian (2) manifests an anharmonic oscillator, with the resonance condition given by $\Delta_{\text{eff}}^2 = \omega_{f,\text{eff}}^2 - \Omega_{\text{tid}}^2 \approx 0$ [93]. This nonlinear tidal effect, combined with another nonlinear influence from spin variation in the star [94,95], can, in some cases, lock the mode in resonance with the tidal frequency.

In this Letter, we focus on a representative binary in which the involved NSs spin retrogradely with respect to the rotation axis of the system (referred to as the z axis and the equatorial plane of the orbit is assumed to lie on the x - y plane). The z component of the angular momentum of one of the NSs (say NS 1) is computed via [cf. Eq. (59) in [96]]

$$J_1 = \int q_b [(q_x - h\bar{u}_x)(y - \bar{y}_1) - (q_y - h\bar{u}_y)(x - \bar{x}_1)] d\mathcal{V}, \quad (4)$$

where the volume integral is taken over the interior of the NS. In the above expression, we define $q_b = \sqrt{\gamma} \rho w$ and $q_{x,y} = h u_{x,y}$ with γ the determinant of the spatial metric, ρ the rest mass density, w the Lorentz factor, $u_{x(y)}$ the $x(y)$ component of the covariant four-velocity, and h the specific enthalpy, respectively. The overhead bar denotes the volume-average value for the quantities; for example, \bar{u}_x is the approximate velocity at the mass center of the NS which is located at (\bar{x}_1, \bar{y}_1) . In GR, the spin of each member of a binary cannot be gauge invariantly defined. In this sense, the analysis based on definition (4) should be viewed as an approximate indicator.

The energy (ΔM_{res}) and angular momentum (ΔJ_{res}) transferred during a tidal resonance in NS 1 will build up a differential rotation and yield a change in its spin as

$$\begin{aligned} \chi_1 + \Delta\chi_1 &= \frac{J_1 + \Delta J_{\text{res}}}{(M_1 + \Delta M_{\text{res}})^2} \\ &\approx \frac{J_1}{M_1^2} \left(1 + \frac{\Delta J_{\text{res}}}{J_1} - \frac{2\Delta M_{\text{res}}}{M_1} \right), \end{aligned} \quad (5)$$

where $\chi_1 := J_1/M_1^2$ is a dimensionless spin. As indicated by the numerical results below, the binary orbit will remain quasicircular after a resonance (i.e., no sizable eccentricity is induced). The variations in the energy and angular momentum can thus be related through the orbital angular velocity Ω_{orb} as (see, e.g., Refs. [85,97])

$$\Delta J_{\text{res}} \simeq \frac{\Delta M_{\text{res}}}{\Omega_{\text{orb}}} = \frac{2\Delta M_{\text{res}}}{\omega_f}, \quad (6)$$

which in turn yields the fractional change in spin as

$$\begin{aligned} \frac{\Delta\chi_1}{\chi_1} &= 2 \frac{\Delta M_{\text{res}}}{M_1} \left[(M_1 \omega_f)^{-1} \frac{M_1^2}{J_1} - 1 \right] \\ &\simeq -7\% \left(\frac{\Delta M_{\text{res}}/M_\odot}{5 \times 10^{-4}} \right) \left[\left(\frac{0.03}{M_1 \omega_f} \right) \left(\frac{-0.48}{\chi_1} \right) - 1 \right]. \end{aligned} \quad (7)$$

Note that, for $\chi_1 < 0$, the negative fractional change indicates a positive value of $\Delta\chi_1$. In addition to the tidally induced oscillations inside the NS, the tidal force torques the NS in an effect to align the tidal bulge on the NS with the external tidal potential. Consequently, the NS gains angular momentum in the orbital direction [88,98]. However, this effect is subdominant to the spin-up caused by f -mode resonance.

The above Newtonian picture can be approximately brought into relativistic tidal response [19,22,23,49], while it is worth noting that this interpretation is an approximation to the response in reality since the quasinormal modes are not complete in GR [81,99,100]. In addition, the quantities defined above depend on gauge and lose strict meaning in GR. That said, we will use them as indicators of what is happening in the GR simulation.

Numerical scheme—We adopt the code SACRA-MPI [101–103], which employs the Baumgarte-Shapiro-Shibata-Nakamura-puncture formalism [104–107] with the Z4c constraint propagation prescription [108] to integrate Einstein’s equation with an adaptive mesh (2:1) refinement algorithm. Implementation details are in the cited articles. For the simulations presented here, the cell-centered grid is configured with four comoving, concentric finer boxes per NS and six coarser domains encompassing both stacks of finer domains. All domains use a grid of $(2N, 2N, N)$ points in the (x, y, z) direction for an even number N , where an equatorial mirror symmetry on the $z = 0$ orbital plane is imposed. We choose $N = 118, 126, 158, 174$, and 190, and the finest domain’s size is set to 14.77 km with a grid spacing denoted as $\Delta x = 14.77 \text{ km}/(N + 1/2)$. The key improvements, implemented in the latest version [109], for obtaining an accurate simulation of f -mode resonance (see Supplemental Material [110] for the evaluation) include the Harten-Lax–Van Leer contact (HLLC) Riemann solver [116,117], refluxing treatment, and a high-order (sixth) interpolation at the mesh boundary.

To ensure that the f -mode resonance is well resolved, we need to consider an inspiraling BNS for which the resonance window is well separated from the merger phase. For this purpose, one is restricted to study BNSs of sufficiently large spin retrograde to the orbit. In fact, the f mode may not even reach a resonance before the merger if the spin magnitude is lower than a certain threshold, depending on the masses of the binary and the EOS (e.g., Ref. [84]). We consider an equal-mass binary with two NSs spinning with $\chi = -0.48$, where the negative sign denotes the retrograde direction. To prepare the quasiequilibrium states of highly spinning quasicircular configurations, we use the public spectral code FUKA [118,119] and follow

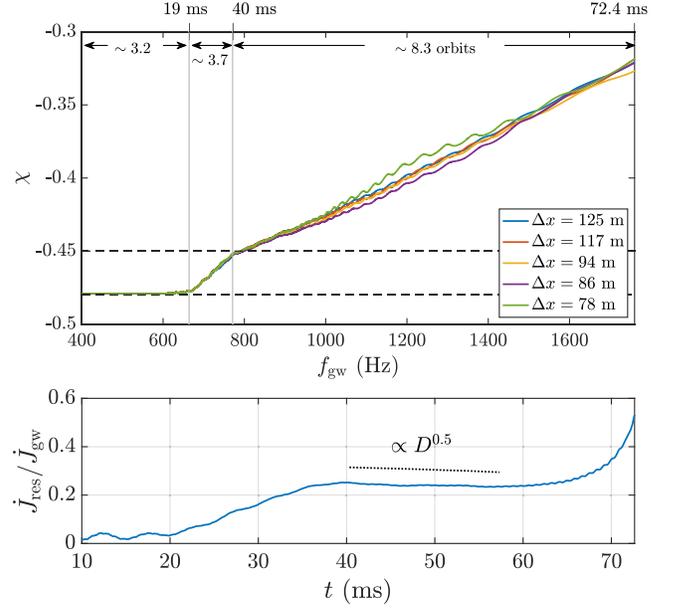


FIG. 1. Top: local measure of the dimensionless spin [Eq. (4)] of an NS in the considered equal-mass BNS as a function of f_{gw} for the grid resolutions adopted. The dashed horizontal lines indicate the dimensionless spin at the onset (lower) and offset (upper) of the linear resonance window, which is independent of the resolutions. Time duration and orbits covered before, during, and after this window are also shown. Bottom: the ratio between rates of the angular momentum transfer into an NS via f -mode resonance and the angular momentum loss via GWs as a function of time. In the bottom panel, we only show the result with the highest resolution adopted.

the method sketched in [120] to remove the residual eccentricity until $e < 10^{-3}$. The initial state is prepared at $M\Omega_{\text{orb}} = 0.0129$. We adopt the piecewise-polytropic approximation of the cold EOS SLy4 [121] and augment a Γ -law thermal component to it in the manner detailed in, e.g., Ref. [50]. We have chosen the index for the thermal component as $\Gamma_{\text{th}} = 1.8$. Varying this value will not significantly impact the resonance phenomena studies here since heating effects essentially play no role before the merger phase [50] while the resonance occurs before that phase.

Simulation results—The binary evolves ≈ 15 orbits before the merger and the convergence order for the resulted GW phase at the merger is estimated as $p_{\text{conv}} \simeq 3.851$ ($\lesssim 4$; see Supplemental Material [110] for the details). We trace the evolution of the local measure of spin [Eq. (4)] for one of the NSs in the top panel of Fig. 1. The spin does not vary before the resonance in the first 21 ms (~ 3.3 orbits), implying that the tidal torque is weak even when the orbital separation is $D \lesssim 24M$ ($= 65 \text{ km}$). For the specific spin of $\chi = -0.48$, we estimate the f mode’s frequency via the universal relation (5) of [72] as $\omega_f \approx 775 \text{ Hz}$. Its resonance begins when $f_{\text{gw}} = \omega_{\text{gw}}/2\pi \gtrsim 670 \text{ Hz}$ and can be divided into two regimes as indicated by the two slopes of spin-up shown in the evolution

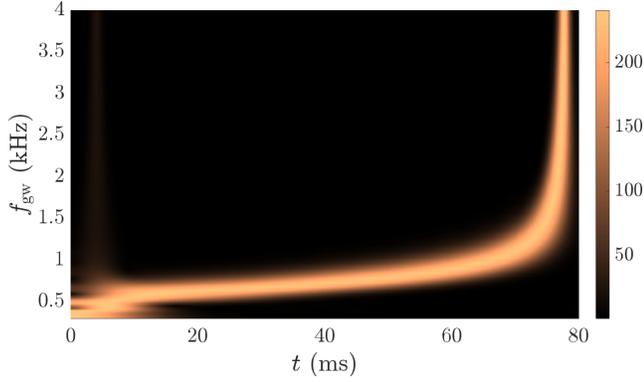


FIG. 2. Time-frequency map of numerically computed Ψ_4 with the highest resolution adopted. The amplitude scale is arbitrary. In the first 10 ms, low-frequency noise due to initial gauge self-adjustment and junk radiation is observed. After that, the frequency evolution follows a chirp pattern.

of J . As will be detailed below, these two phases correspond to linear and nonlinear regimes of f -mode evolution. The first epoch lasts for 21 ms and undergoes ~ 3.7 orbits, during which the NS's spin increases by $\sim 6.3\%$ as indicated by the two horizontal lines in the plot. According to Eq. (7), the star gains the energy of $\Delta M_{\text{res}} \sim 4.5 \times 10^{-4} M_{\odot}$ in this phase. This energy can also be estimated by the difference in binding energy between the binary at the end of this phase (i.e., at $\Omega_{\text{orb}} \simeq 390$ Hz) and a quasicircular, quasiequilibrium BNS at the same orbital frequency. In the latter case, orbital decay is due to GW emission, whereas it arises from both GW and tidal excitation in simulation. The alternative, gauge-invariant estimate gives $\Delta M_{\text{res}} \sim 8 \times 10^{-4} M_{\odot}$, agreeing in order of magnitude with the gauge-dependent method. The absorbed energy drives the f -mode oscillation and sustains a degree of differential rotation (see Supplemental Material [110] for the profile).

An angular momentum transfer rate accompanies the f -mode excitation (\dot{J}_{res}), which is derived from the numerical data and compared with the angular momentum emission rate via GW [\dot{J}_{gw} ; Eq. (2.11) of [103]] in the bottom panel of Fig. 1. We see that \dot{J}_{res} kicks in at $t \simeq 20$ ms and grows to $\lesssim 0.25 \dot{J}_{\text{gw}}$ at $t \simeq 40$ ms. After this point, the ratio between them is broadly “locked” until the merger begins at $\gtrsim 70$ ms. In particular, the tidal force, scaling as $\propto D^{-3}$ (e.g., Ref. [122]), continues to efficiently torque the NS through the nonlinear evolution of f mode. The ratio $\dot{J}_{\text{res}}/\dot{J}_{\text{gw}}$ can then be estimated as $\propto D^{-3}/D^{-3.5} = D^{0.5}$ during the locking (dotted line). Summing the contributions of the f -mode resonance in two NSs, we see about a 50% increase in the dissipation rate of the binary's angular momentum throughout the nonlinear regime of resonance. The loss rate of orbital energy is also enhanced by this fraction since the binary remains quasicircular as we will see below.

After the merger, a black hole promptly forms, even though the total mass of the binary is only $M = 2.7 M_{\odot}$.

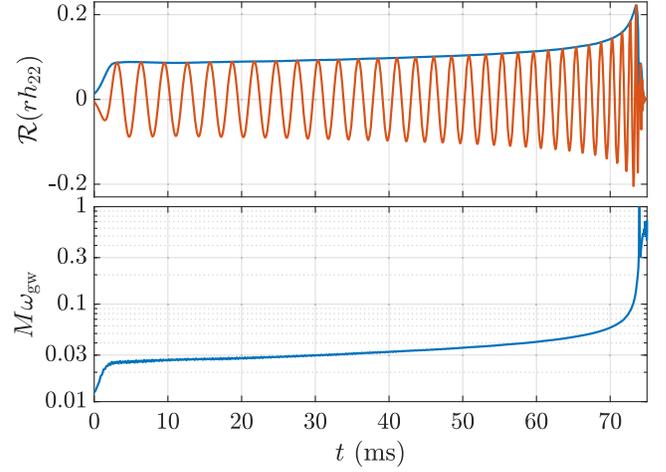


FIG. 3. Top: waveform covering the last ≈ 30 cycles of inspiral phase (red) and its amplitude over time (blue). Bottom: evolution of GW angular frequency. The shown results are obtained from the highest resolution run.

This outcome is due to the binary's much less angular momentum compared to an irrotational binary with the same NS masses [50]. The remnant black hole has a dimensionless spin of $\chi \simeq 0.74$ and the irreducible mass of $M_{\text{irr}} \lesssim 2.41 M_{\odot}$.

Effect of resonance on GW signal—The fixed-frequency method is often adopted to derive waveforms from the numerically computed Ψ_4 functions [123–125]. This method applies a frequency cutoff to filter out low-frequency noise and spurious contributions below a chosen threshold. Recently, Yu *et al.* [47] raised concerns about its accuracy for GWs generated after the onset of f -mode resonance, suggesting that a non-negligible GW component could result from the excited f mode. To ascertain that the method is legitimate, i.e., the contribution of f -mode emission is negligible, we show in Fig. 2 the spectrogram of Ψ_4 function that is extracted at $r = 600 M_{\odot}$. Apart from the blurring in the first 10 ms, which results from the initial data, the signal shows a chirp pattern of the instantaneous frequency up to the merger. The absent clue of a f -mode-related spectral density validates the use of the fix-frequency method for deriving waveform from Ψ_4 . Accordingly, we apply this method to derive waveforms throughout (see Refs. [102,103] for the detailed formulas).

The strain of the (2,2) component of GWs (top) and its time-frequency representation (bottom) are shown in Fig. 3. We can see that the f -mode resonance induces essentially no eccentricity regardless of the linear or nonlinear regimes. The orbit largely remains quasicircular up to the merger phase, as no oscillatory behavior is observed for either the amplitude or the frequency of the waveform. The quasicircularity after the resonance is further evidenced by the acceleration rate of GW's phasing, viz. $Q_w \equiv \omega_{\text{gw}}^2/\dot{\omega}_{\text{gw}}$ [28,51,52,126]. We diagnose Q_w of the numerical waveform and compare it to the EOB waveform model

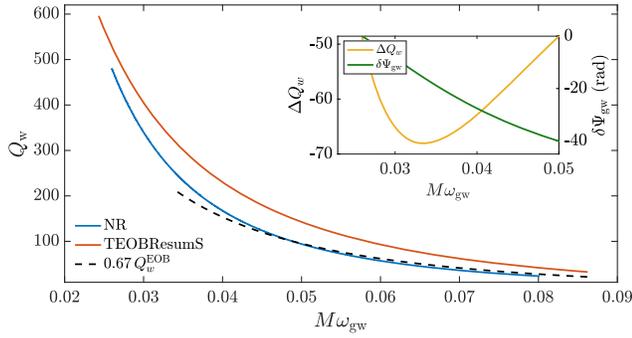


FIG. 4. Gauge-invariant and dimensionless measure Q_w of the phase acceleration for the numerical (blue) and the TEOBResumS (red) waveforms. The prediction on Q_w modified from the EOB model by the enhanced dissipation due to f -mode resonance is overplotted (dashed). The inset window draws the difference between the EOB and NR data (yellow) and the associated dephasing estimated from Eq. (8) (green).

TEOBResumS in Fig. 4 while noting that the difference between TEOBResumS and SEOBNRv*T models is much less than the deviation of the NR result from either of them for this BNS configuration. We use the Love number of a static NS with $M = 1.35M_\odot$, which is $\Lambda = 389.4$, to generate this EOB waveform. Their overall agreement is poor, and the inset window shows the difference $\Delta Q_w = Q_w^{\text{NR}} - Q_w^{\text{EOB}}$ (yellow) between the NR result and the EOB data.

To identify the primary source of deviation, we compare Q_w of this EOB model with an openly available numerical binary black hole waveform (serial number 1500) from the database of the SXS Collaboration [127,128] with similar parameters. In particular, we generate a TEOBResumS waveform and compare it with the SXS data for an equal-mass binary with both black holes spinning at $\chi = -0.4842$. While we do not present these results here, we found that the deviation ΔQ_w oscillates around zero with a peak value $\simeq 4$ all the way up to the merger. The deviation shown in Fig. 4 suggests a more pronounced discrepancy when compared to BNS mergers with the given spin parameter. We now turn to argue that f -mode resonance can result in this deviation.

During the resonance, a portion of the binary’s orbital energy is transferred into the NSs. This energy transfer leads to an increased amount of $\dot{\omega}_{\text{gw}}$, and thus Q_w is lower than the situation where the resonance is absent. Given the sustained quasicircularity of the orbit, the $\dot{J}_{\text{res}}/\dot{J}_{\text{gw}}$ ratio (cf. Fig. 1) suggests an approximately 50% increase in the energy loss rate of the orbit (\mathcal{F}). This effectively enhanced \mathcal{F} can be translated to a decrease in Q_w by a factor of ≈ 1.5 using the energy balance equation $\dot{\omega}_{\text{gw}} = -\mathcal{F}/(dE/d\omega_{\text{gw}})$, where E is the binding energy of the binary. In Fig. 4, we overplot a curve of $0.67Q_w$, which Q_w^{NR} indeed asymptotically approaches [129]. Although the local measurement of \dot{J}_{res} is a gauge-dependent quantity, the observed impact on Q_w can be explained by the orbital dissipation enhanced by

it. We also estimate an overall phase shift in the waveform due to the resonance through

$$\delta\Psi_{\text{gw}} = \int \Delta Q_w d(\ln \omega_{\text{gw}}), \quad (8)$$

which is depicted in the inset of Fig. 4 (green). The phase shift due to the f -mode resonance is enormous and amounts to $\lesssim 6.4$ GW cycles (or $\lesssim 3.2$ orbits).

Outlook—We reported the first GR simulation for f -mode resonance that is performed by a new implementation presented in [109]. Our results underscore the importance of f -mode effects for accurate GW data analysis and offer valuable insights into refining current waveform models in the tidal resonance sector. In particular, we show that the resonance window can be robustly resolved, which reflects an intrinsic feature of tidal interactions in GR. Information about its backreaction to the NS and how it develops from the linear regime and saturates to the nonlinear stage can thus provide hints for the future modeling of other systems, such as eccentric BNS (e.g., Ref. [130]) or close encounter resonance. However, the locking phenomenon between nonlinear resonance and GW dissipation is specific for quasicircular binaries. Assuming that the influence of the resonance in the two NSs could be linearly added to render the gross effects, our result can also apply to black hole–NS binaries.

This Letter demonstrates that the numerical scheme has now achieved the precision levels that allow us to explore the dynamical tides with high fidelity. This advancement opens a new path for systematically investigating the f -mode resonance. The Letter also presents an important initiative to complete the NR dataset in the untouched BNS parameter space, which will be key to next-generation waveform modeling of BNS mergers.

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