

Modeling matter in SEOBNRV5THM: Generating fast and accurate effective-one-body waveforms for spin-aligned binary neutron stars

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We present SEOBNRV5THM, an accurate and fast gravitational-waveform model for quasicircular, spinning, nonprecessing binary neutron stars (BNSs) within the effective-one-body (EOB) formalism. It builds on the binary–black-hole approximant SEOBNRV5HM and, compared to its predecessor SEOBNRV4T, it (i) incorporates recent high-order post-Newtonian results in the inspiral, including higher-order adiabatic tidal contributions, spin-induced multipoles and dynamical tides for spin-aligned neutron stars, (ii) includes the gravitational modes $(\ell, |m|) = (2, 2), (3, 3), (2, 1), (4, 4), (5, 5), (3, 2)$, and $(4, 3)$, (iii) has a time of merger calibrated to BNS numerical-relativity (NR) simulations, (iv) accurately models the premerger $(2, 2)$ mode through a novel phenomenological ansatz, and (v) is 100 to 1000 times faster than its predecessor model for BNS systems with total mass $M \geq 2M_\odot$. Thus, SEOBNRV5THM can be used in Bayesian parameter estimation, which we perform for two BNS events observed by the LIGO-Virgo collaboration, GW170817 and GW190425. The model accurately reproduces BAM and SACRA NR waveforms with errors comparable to or lower than the intrinsic NR uncertainty. We validate the model against the other state-of-the-art BNS waveform models NRTIDALV3 and TEOBRESUMS and find differences only for highly spinning and highly tidally deformable BNSs, where there is no NR coverage and the models employ different spin prescriptions. Our model serves as a foundation for the development of subsequent SEOBNR waveform models with matter that incorporate further effects, such as spin precession and eccentricity, to be employed for upcoming observing runs of the LIGO-Virgo-KAGRA collaboration and future facilities on the ground.

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I. INTRODUCTION

The field of gravitational-wave (GW) astronomy has undergone rapid expansion since the initial detection of GWs from a binary–black-hole (BBH) merger in 2015 [1], and the first detection of a binary–neutron-star (BNS) merger only two years later, in the landmark event GW170817 [2]. Among the 90 GW detections of the LIGO-Virgo-KAGRA (LVK) collaboration up to and including the third observing run (O3) [3,4], GW170817 was the first to be observable in both GWs and electromagnetically across the entire spectrum from radio to gamma rays, yielding novel insights into astrophysics, tests of general relativity and cosmology

[2,5,6], as well as establishing the era of multimessenger astronomy with GWs. The fourth observing run (O4) of the LVK collaboration is currently ongoing, during which one binary system containing a neutron star (NS) has been observed so far [7].

The Bayesian analysis framework [8,9] facilitates the extraction of physical parameters, such as masses, spins, and the macroscopic description of matter, like the tidal deformability of a NS, from GW signals. To this end, observed data is compared with theoretical predictions derived from the Einstein field equations.

Over the next decades, the observed merger rates of BNSs will significantly increase, up to multiple thousands per year (see, e.g., Refs. [10,11]), as existing detector facilities become more sensitive, and as new facilities on the ground, such as the Einstein Telescope [12], and Cosmic Explorer [13,14], are constructed.

Two challenges have to be overcome to extract reliable and exhaustive knowledge. As detector sensitivities increase, higher signal-to-noise ratios demand more accurate waveform models to avoid parameter-estimation biases [15–18]. With a new plethora of observational data, we will additionally be challenged to drastically reduce waveform

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generation speed for the parameter estimation (PE), as the current Bayesian analysis approach will become computationally infeasible. PE codes based on machine-learning methods, notably neural posterior estimation, are also under continuous development to speed up inference studies [19–21].

Importantly, by enhancing the speed and precision of existing waveform models, significant insights into NSs can be obtained, going beyond the progress already obtained with GW170817 [2,22]. The structure of the NS is primarily described by a macroscopic equation of state (EOS) of the neutron-rich supranuclear-dense matter [23,24]. This EOS serves to supplement the Tolman-Oppenheimer-Volkoff equations, in order to make predictions about the NS's interior. Given that the matter in the core of a NS is a few times denser than the nuclear saturation density, it cannot be constrained by experiments in earth-based accelerators and reactors due to its extreme pressure and internal energy. Consequently, there are various realistic theoretical and phenomenological models in nuclear physics (see, e.g., Refs. [23,25–31]), which predict the state of such matter and would constrain the high-energy limit of nuclear physics if confirmed or rejected.

It appears that the inspiral and merger of BNS systems offer a viable opportunity to uncover the precise physics of matter at this most extreme state [2,32–35], in particular in tandem with future observations of, e.g., core-collapse supernovae [36], NSBH systems [37], and individual pulsars [38], to list a few. In BNS coalescence, the interior structure of NSs, particular its EOS, dictates how the NS reacts to a time-varying tidal field of a compact binary partner. Consequently, one can define the NS's EOS-dependent tidal deformability Λ_l , which quantifies the star's l th order multipolar deformation $\mathcal{Q}$ in response to an external tidal field $\mathcal{E}$ as, e.g., $\mathcal{Q}_{ij} = \Lambda_2 \mathcal{E}_{ij}$ for the quadrupole ($l = 2$) [32,39]. An important goal of BNS detection is then to determine this deformability as well as the individual star's mass, spin, and potentially other NS quantities, to measure the correct EOSs.

Elaborating on one of these NS parameters, it has also been predicted that, during the inspiral, the frequency of the time-varying tidal field can approach or cross the fundamental f -mode frequency of a NS. This can result in resonance phenomena that further augment the star's tidal response. These frequency-dependent *dynamical tides* have been studied in Refs. [32,40–51] and can be effectively described through separation-dependent tidal enhancement $\Lambda_l(r) = k_l^{\text{eff}}(r)\Lambda_l$ and depend through the resonance frequency $\omega_{f,l}$ on the EOS as well as the NS's spin. In addition to the change of the f -mode frequency and dynamical response in the inertial frame of the star, the spin also induces a multipolar structure in BHs and NSs, which can be captured by higher-order spin effects [52,53], and affects the conservative and dissipative dynamics of BNS systems.

It is possible to analytically solve the approximate Einstein field equation of two point particles with some associated multipolar structure to include the aforementioned matter effects. In post-Newtonian (PN) theory (see, e.g., Refs. [54–60] for reviews), Einstein's equations are expanded in powers of $(v/c)^2$, where v is the binary characteristic velocity and c is the speed of light. Within this framework, one can derive the effects of tidal interactions which start to affect the GW phase of a coalescing binary at the fifth PN order beyond the leading term [$\propto (v/c)^{10}$] [61] and are currently known beyond the next-to-next-to leading order or equivalently up to 7.5PN order [62–66]. The perturbative series resulting from these calculations has been shown to be accurate during the early inspiral phase [67]. However, as the merger phase is approached, the series becomes inaccurate due to the significant relativistic velocities of the objects ($v \sim c$), the nonlinear deformation of NSs, and the potential for matter exchange [68,69].

Consequently, numerical-relativity (NR) simulations have emerged as a valuable tool for modeling diverse matter phenomena when analytic methods become impractical, particularly in the vicinity of, and following, the merger. NR predicts the GW emission of BNS systems by numerically solving the Einstein equations augmented with prescriptions for hydrodynamic and particle physics effects on supercomputers. However, as these simulations are computationally intensive, they generally encompass only the late stage of the inspiral and merger [70–75].

Although these simulations are therefore our best effort at modeling the late stages of the coalescence of BNS systems, they are still prone to intrinsic errors, which must be paid careful attention to and which partially remain as open research issues. One typical error source for finite-difference codes are resolution effects, stemming from the numerical discretization of the Einstein equations and associated spacetime. Other error sources are, for example, the different phenomenological prescriptions to treat the hydrodynamics on these discretized grids; the omission of some physical effects like magnetic fields and their back-reaction on the matter, which is the case for older simulations like the ones used in this work [76,77]; or the numerical heating of the NSs, which can distinguish numerical NSs from astrophysical ones [78,79].

A successful approach for combining PN results, and other analytical approximation methods, with information from high-accuracy NR simulations and strong-gravity effects has been the effective-one-body (EOB) formalism, see Refs. [80–85]. It maps the dynamics of a compact binary to that of a test mass (or test spin) in a deformed Schwarzschild (or Kerr) background, with the deformation parameter being the symmetric mass ratio ν , supplemented with an associated energy loss due to GW emission. The resulting dissipative system of ordinary differential equations (ODEs) is then integrated numerically. EOB waveform

models of BBHs have been constructed for nonspinning [82,86–93], spinning [84,85,94–115], and eccentric binaries [116–123], and have recently been extended to scattering and dynamical captures [124–127]. They are developed in two flavors, namely SEOBNR in Refs. [106,107,110,121,128,129] and TEOBRESUMS in Refs. [93,112,114,118,130,131]. These models have also been extended to BNS systems, as SEOBNRV4T in Refs. [41,42] for the former and in Refs. [112,132–135] for the latter. TEOBRESUMS is furthermore developed as TEOBRESUMS-GIOTTO for quasicircular systems, and as TEOBRESUMS-DALÍ for eccentric systems. For simplicity, and as we only study quasicircular systems in this work, we refer to TEOBRESUMS-GIOTTO as the shorter TEOBRESUMS in the following.

As BNS waveforms are however in the sensitive frequency band of the detectors for thousands of wave cycles, the pure ODE integration and subsequent Fourier transformation of the waveform can become too computationally intensive for PE. To overcome this issue, there exist reduced-order models of EOB approximants, see Refs. [136,137], or one can use the postadiabatic approximations as in Refs. [129,138], machine learning techniques as in Refs. [139], and the stationary phase approximation instead of a fast Fourier transform of the waveform, see Ref. [140].

There also exist phenomenological GW models in Refs. [141,142], which combine and extrapolate PN, EOB, and NR waveforms across the parameter space of interest. The most important phenomenological BNS waveform approximant is the NRTIDAL family, developed in Refs. [70,143–145]. It aims to model the tidal phase contribution during the inspiral through the use of closed-form analytical expressions for the tidal contribution and an NR calibrated merger model. The matter contribution is then used to augment a given BBH waveform baseline to model BNS systems.

In this work, we introduce SEOBNRV5THM as the next step in modeling spin-aligned BNS systems accurately and fast in the SEOBNR framework. We have improved the model in comparison to SEOBNRV4T through the following:

- (i) the development of SEOBNRV5THM in the highly optimized PYSEOBNR Python package and with the modeling improvements of the point-particle case as developed in SEOBNRV5HM [115], which yields more accurate BNS waveforms that are faster by a factor 100 to 1000 when comparing to SEOBNRV4T for systems with total mass $M \geq 2M_\odot$ ¹;
- (ii) analytical improvements in the inclusion of up to relative 2.5PN order adiabatic tidal effects [63,65,66],

the inclusion of 3.5PN and 4PN order spin-induced multipolar effects for spin-squared, spin-cubed and spin-to-the-fourth contributions in the conservative sector [53] and in the dissipative sector [52], as well as dynamical tides for spinning NSs [32,41–43,46];

- (iii) development of an NR informed phenomenological premerger model; and in particular
- (iv) the calibration to a large set of NR BNS waveforms in the merger time.

In total, we utilize 64 NR waveforms from the BAM catalogs [73–75] and the SACRA catalogs [70,146], covering ten different EOSs, with mass ratios $q \in [1, 2]$, tidal deformabilities $\Lambda \in [43, 4090]$, and dimensionless spins $\chi \in [-0.18, 0.27]$.

The paper is organized as follows. In Sec. II, we summarize the EOB framework and present how we include tidal effects in the Hamiltonian, in the radiation-reaction force, and the waveform of SEOBNRV5HM to accurately model BNS systems and explain our premerger modeling approach. We also discuss the applicability of our model to highly spinning, antialigned NSs, which can exhibit f -mode resonances already during the early inspiral. In Sec. III we lay down our approach for the calibration against NR waveforms, where we pay close attention to the resolution of different NR runs and extrapolate our results to a common resolution across all employed simulations. In Sec. IV we discuss our implementation and perform tests to validate our model through comparisons to NR, and to the other state-of-the-art waveform models across the relevant parameter space for observations. We further perform PE for the two BNS detections GW170817 and GW190425 in Sec. V and compare our results to the literature. Finally, our main conclusion, directions for future improvements of the model, and recommendations for the region of validity of SEOBNRV5THM are discussed in Sec. VI.

A. Notation

In the following, we use natural units in which $c = G = 1$ unless stated otherwise, and we use overhead dots to signal time derivatives. We largely follow the notation as outlined in Ref. [115] and consider binaries with masses m_1 and m_2 , where $m_1 \geq m_2$. We also define the following derived quantities:

$$\begin{aligned} M &\equiv m_1 + m_2, & q &\equiv \frac{m_1}{m_2} \geq 1, & \nu &\equiv \frac{m_1 m_2}{M^2}, \\ \delta &\equiv \frac{m_1 - m_2}{M}, & \mu &\equiv \frac{m_1 m_2}{M}, & \mathcal{M} &= \nu^{3/5} M \\ X_i &\equiv \frac{m_i}{M}, & X &= 1 - 4\nu. \end{aligned} \quad (1)$$

¹Note that this applies to the comparison between the EOB models SEOBNRV4T and SEOBNRV5THM. There also exists the frequency-domain surrogate SEOBNRV4T_SURROGATE [137] for applications in PE, which is even faster than SEOBNRV5THM, as discussed in Sec. IID.

For NSs we define the adiabatic tidal deformability $\Lambda_i^{(i)}$ in accordance with the masses $i \in \{1, 2\}$. It is related to the dimensionless tidal Love number $k_i^{(i)}$ and the compactness $C_i = m_i/R_i$ defined through the NS radius R_i as

$$\Lambda_l^{(i)} = \frac{2}{(2l-1)!!} \frac{k_l^{(i)}}{\mathcal{C}_i^{2l+1}}, \quad (2a)$$

$$= \frac{2}{(2l-1)!!} k_l^{(i)} \frac{(R_i/M)^{2l+1}}{X_i^{2l+1}}. \quad (2b)$$

It should be highlighted that the tidal deformability as normalized through the compactness $\mathcal{C}_i$ in Eq. (2a) is given in terms of the individual masses m_i . If we want to work in units of total mass M , we therefore need to use $\Lambda_l^{(i)} X_i^{2l+1}$ as the dimensionless adiabatic tidal deformability in units of total mass, as can be seen from Eq. (2b).

We also define the quadrupolar tidal coupling constant,

$$\kappa_2^T = 3\nu(X_1^3\Lambda_2^{(1)} + X_2^3\Lambda_2^{(2)}), \quad (3)$$

which quantifies the leading order contribution of tidal effects.

For binaries with nonprecessing spins of magnitude S_1 and S_2 , we define the dimensionless spins,

$$\chi_i \equiv \frac{a_i}{m_i} = \frac{S_i}{m_i^2} \in [-1, 1], \quad (4)$$

where $i \in \{1, 2\}$, and define the following spin variables:

$$\begin{aligned} \chi_{\text{eff}} &\equiv \frac{(m_1\chi_1 + m_2\chi_2)}{m_1 + m_2}, \\ a_{\pm} &\equiv a_1 \pm a_2 = m_1\chi_1 \pm m_2\chi_2. \end{aligned} \quad (5)$$

We denote the spin-induced quadrupole, octupole, and hexadecupole constants as C_2 , C_3 , and C_4 , respectively. These constants equal 1 for BHs but are larger than 1 for NSs. We occasionally use $u \equiv M/r$ instead of the separation r , and introduce the Keplerian orbital frequency $\Omega_0 \equiv (r/M)^{-3/2}$.

Our conventions for the Fresnel integrals are

$$\begin{aligned} F_S(z) &= \int_0^z \sin\left(\frac{\pi t^2}{2}\right) dt, \\ F_C(z) &= \int_0^z \cos\left(\frac{\pi t^2}{2}\right) dt. \end{aligned} \quad (6)$$

As in the BBH case, we define the *merger time* of an aligned-spin BNS waveform to be the time where the (2,2) mode amplitude first peaks.

II. THE SEOBNRV5THM MODEL

In the following section, we describe our EOB model, highlighting the changes we make to the BBH baseline SEOBNRV5HM from Ref. [115]. More specifically, in Sec. II A we explain how we include spin-induced multipole moments and dynamical tides into the Hamiltonian of our model, and

the changes we make to the dynamical tide computation in comparison to SEOBNRV4T. In Sec. II B we then describe how the inspiral waveform is computed, and how we include relative 2.5PN adiabatic tidal and 3.5PN spin-induced multipole information into the inspiral waveform and fluxes. In Sec. II C we explain our novel approach to the termination and premerger waveform for BNS systems.

A. EOB dynamics with matter effects

The EOB formalism [80,81,83–85] maps the two-body dynamics onto the effective dynamics of a test body in a deformed Schwarzschild or Kerr background, with the deformation parametrized by the symmetric mass-ratio ν . The energy map relating the effective Hamiltonian H_{eff} and the two-body EOB Hamiltonian H_{EOB} is given by

$$H_{\text{EOB}} = M \sqrt{1 + 2\nu \left(\frac{H_{\text{eff}}}{\mu} - 1 \right)}. \quad (7)$$

We use the aligned-spin SEOBNRV5HM waveform model as our BBH baseline and refer interested readers to Ref. [115] for a more in-depth description of the model. In the following, we summarize the structure of the aligned-spin Hamiltonian, and its zero-spin limit, highlighting where matter effects enter into the model.

For aligned spins, the effective Hamiltonian reduces to the equatorial Kerr Hamiltonian in the test-body and BBH limit, with the Kerr spin $a = a_+$ mapped to the binary's total spin. It explicitly reads [53,115]

$$\begin{aligned} H_{\text{eff}}^{\text{align}} &\equiv \frac{M p_\phi (g_{a_+} a_+ + g_{a_-} \delta a_-) + \text{SO}_{\text{calib}} + G_{a^3}^{\text{align}}}{r^3 + a_+^2(r + 2M)} \\ &\quad + \sqrt{A^{\text{align}} \left(\mu^2 + p^2 + B_{np} p_r^2 + B_{npa} \frac{p_\phi^2 a_+^2}{r^2} + Q^{\text{align}} \right)}, \end{aligned} \quad (8)$$

where r is the binary's separation and ϕ the orbital phase, p_r , p_ϕ are their canonically conjugated momenta that satisfy $p^2 = p_r^2 + p_\phi^2/r^2$, $g_{a_{\pm}}$ are gyrogravitomagnetic factors, A^{align} , $B_{np(a)}$, $G_{a^3}^{\text{align}}$, and Q^{align} are EOB potentials containing resummed PN information, defined in Appendix A of Ref. [115], and SO_{calib} is a calibration parameter tuned to NR. The first term on the right-hand side includes the odd-in-spin contributions (in the numerator), while the second term (square root) includes the even-in-spin contributions. Some parts of the matter information influence the A^{align} potential, which reads more explicitly

$$A^{\text{align}} = \frac{a_+^2/r^2 + A_{\text{nos}} + A_{\text{ss}}^{\text{align}}}{1 + (1 + 2M/r)a_+^2/r^2}, \quad (9)$$

where A_{noS} and $A_{\text{SS}}^{\text{align}}$ are further potentials that contain PN information about the nonspinning and spin-spin (SS) sector.

We remind that such a Hamiltonian includes most of the 5PN nonspinning contributions, together with spin-orbit information up to the next-to-next-to-leading order (NNLO), SS information to NNLO, as well as cubic- and quartic-in-spin terms at LO, corresponding to all PN information up to 4PN order for spin-aligned point particles. More details about the derivation of the generic-spin Hamiltonian, together with the full expressions of all terms and potentials, are given in Ref. [53].

Regarding matter effects, the SEOBNRV5THM Hamiltonian includes the spin-induced multipole moments C_l up to $l = 4$, to the same order as the BBH contributions, which in Ref. [53] are defined as $C_{\text{IES}2}$, $C_{\text{IBS}3}$, and $C_{\text{IES}4}$, respectively. These terms have been set to unity for the BBH baseline model SEOBNRV5HM, where spin-induced multipoles are absent, but enter now in the A^{align} , B_{np} , and Q^{align} expressions for $l = 2, 4$ and the $G_{a^3}^{\text{align}}$ expression for $l = 3$ as given explicitly in Eqs. (31)–(34) of Ref. [53].

In the nonspinning (noS) limit $a_1 = a_2 = 0$, the effective Hamiltonian reduces to [53]

$$H_{\text{eff}}^{\text{noS}} = \sqrt{A_{\text{noS}} \left[\mu^2 + A_{\text{noS}} \bar{D}_{\text{noS}} p_r^2 + \frac{p_\phi^2}{r^2} + Q_{\text{noS}} \right]}, \quad (10)$$

and the 5PN potentials $A_{\text{noS}}(r)$ and $\bar{D}_{\text{noS}}(r)$ are complete except for two quadratic-in- ν coefficients. They are respectively (1,5) and (2,3) Padé resummed. A_{noS} , $\bar{D}_{\text{noS}}(r)$, and Q_{noS} are furthermore respectively contained in A^{align} , B_{np} , and Q^{align} from Eq. (8).

As has been done for SEOBNRV4T in Ref. [137], we add NNLO dynamical tidal effects for the quadrupole ($l = 2$) and octupole ($l = 3$) into the Hamiltonian through the modification of the BBH A_{noS} potential, see Ref. [147],

$$A_{\text{noS}}(r) \mapsto A_{\text{noS}}^{\text{BBH}}(r) + A_{\text{noS}}^{\text{tidal}}(r), \quad (11)$$

where $A_{\text{noS}}^{\text{BBH}}(r) \equiv A_{\text{noS}}(r)$ of the BBH baseline and

$$\begin{aligned} A_{\text{noS}}^{\text{tidal}}(r) = & -3\Lambda_2^{(1)} k_2^{\text{eff},(1)}(r) X_1^4 X_2 u^6 \left[1 + \frac{5}{2} X_1 u + \left(3 + \frac{1}{8} X_1 + \frac{337}{28} X_1^2 \right) u^2 \right] \\ & - 15\Lambda_3^{(1)} k_3^{\text{eff},(1)}(r) X_1^6 X_2 u^8 \left[1 + \left(-2 + \frac{15}{2} X_1 \right) u + \left(\frac{8}{3} - \frac{311}{24} X_1 + \frac{110}{3} X_1^2 \right) u^2 \right] + (1 \leftrightarrow 2). \end{aligned} \quad (12)$$

In the above expression, $k_l^{\text{eff},(i)}(r)$ are the effective tidal enhancement factors due to dynamical tides, defined below in Eq. (13).

Note that we are only including relative 2PN adiabatic tidal information (up to u^2) in the Hamiltonian in Eq. (12). Our effective description of dynamical tides in $k_l^{\text{eff},(i)}(r)$ on the other hand is based on a leading-order two-timescale approximation [42,46]. Currently, the relative 3PN dynamical and adiabatic tidal correction to the Hamiltonian is known [50,51], but the calculation requires a counterterm whose finite contribution has yet to be determined for specific EOSs, and the Hamiltonian is not yet EOB resummed. These results can therefore not be immediately used in EOB models.

We note that although one could, as is done in TEOBRESUMS in Refs. [134,135], include the LO adiabatic tidal effects for $l \geq 4$, gravitomagnetic tidal effects, and the respective tidal adjustment of the $\bar{D}_{\text{noS}}$ potential, we leave the inclusion of these effects as future work. Our reasoning for this is that those effects were partially known but not included in SEOBNRV4T either, are outside the scope of this work, and have little impact on the waveform as one can see later in Sec. IV. We furthermore believe it to be more consistent to include only adiabatic information where the dynamical tidal extension is currently understood, which

is not the case for the aforementioned adiabatic tidal information.

Our computation of the effective tidal enhancement factor $k_l^{\text{eff}}(r)$ to be used in the $A_{\text{noS}}^{\text{tidal}}$ potential of Eq. (12) for spinning NSs is based on a two-timescale expansion between the orbital and oscillation timescales and taken from Ref. [46]. We however rewrite it here in a more condensed form inspired by the nonspinning $k_l^{\text{eff}}(r)$ as reported in Ref. [137]. It reads

$$k_l^{\text{eff}}(r) = a_l + b_l \left[f(x_f, \Omega_0) + \left(1 - \frac{\Delta \omega_{0l}^2}{\omega_{0l}^2} \right) \sqrt{\frac{\pi}{3\varepsilon}} Q(\hat{t}) \right], \quad (13)$$

where the coefficients for the quadrupole and octupole multipole moments are $(a_2, a_3) = (1/4, 3/8)$, and $(b_2, b_3) = (3/4, 5/8)$; the resonant term f and the trigonometric term Q are defined in Eqs. (14a)–(14c) below. We additionally omit the index (i) labeling the NS, but note that k_l^{eff} has to be computed for each of them individually before it can be used in Eq. (12).

As we employ the Keplerian orbital frequency $\Omega_0(r)$, it ultimately only depends on the separation r in addition to the NS parameters $(\Lambda_{l \geq 2}, \omega_{0l}, \Delta \omega_{0l})$, where ω_{0l} is the f -mode frequency of the 2^l multipole for a nonspinning

NS, and $\Delta\omega_{0l}$ is the shift in f -mode frequency due to the NS's spin. We denote the resulting f -mode frequency for spinning NSs by $\omega_{f,l}$; see Eq. (14c) below. All frequencies are given in units of inverse total mass M^{-1} in the equations below. The remaining equations for the computation of $k_l^{\text{eff}}(r)$ then read

$$f(x_f, \Omega_0) = \frac{(1 - \frac{\Delta\omega_{0l}^2}{\omega_{0l}^2})}{1 - \frac{(l\Omega_0 - \Delta\omega_{0l})^2}{\omega_{0l}^2}} + \frac{5}{6} \left(1 - \frac{\Delta\omega_{0l}}{\omega_{0l}}\right) \frac{x_f^2}{(1 - x_f^{5/3})}, \quad (14a)$$

$$Q(\hat{t}) = \cos\left(\frac{3}{8}\hat{t}^2\right) \left[1 + 2F_s\left(\frac{\sqrt{3}}{2\sqrt{\pi}}\hat{t}\right)\right] - \sin\left(\frac{3}{8}\hat{t}^2\right) \left[1 + 2F_c\left(\frac{\sqrt{3}}{2\sqrt{\pi}}\hat{t}\right)\right], \quad (14b)$$

$$\omega_{f,l} = \omega_{0l} + \Delta\omega_{0l}, \quad x_f = \frac{\omega_{f,l}}{l\Omega_0}, \quad \epsilon = \frac{256\nu}{5} \left(\frac{M\omega_{f,l}}{l}\right)^{5/3}, \quad \hat{t} = \frac{8}{5\sqrt{\epsilon}}[1 - x_f^{5/3}], \quad (14c)$$

where x_f is the ratio between the forced f mode and driving orbital frequency, ϵ is a small dimensionless parameter used in the two-timescale expansion [46], and $\hat{t} \equiv \sqrt{\epsilon}(\phi - \phi_{\text{res}})$ is a rescaled shifted dimensionless phase variable that describes the orbital phasing around resonance, denoted by “res.” We remind that the Fresnel integrals used in Eq. (14b) are defined in Eq. (6).

It has to be noted that the $f(x_f)$ of Eq. (14a) has a removable singularity for the resonance $x_f = 1 \Leftrightarrow l\Omega_0 = \omega_{f,l}$. Similarly to what has been done in SEOBNRV4T, we take care of this singularity by switching to the fourth order Taylor expansion of f in the vicinity of the pole $|x_f - 1| < 10^{-2}$, while keeping ω_{0l} and $\Delta\omega_{0l}$ fixed and replacing $\Omega_0 = \omega_{f,l}/(lx_f)$ prior to the expansion. This procedure only introduces a negligible discontinuity at the transition region of the order 10^{-10} .

By close inspection of Eq. (14c), we find that

$$\hat{t} \propto -(r/M)^{5/2} + \text{const} \quad (15)$$

grows exponentially for large orbital separations r . As a consequence, the trigonometric functions and Fresnel integrals used in the definition of $Q(\hat{t})$ of Eq. (14b) become numerically unstable for large separations r during the early inspiral in both the spinning and nonspinning case.

In contrast to what has been done in SEOBNRV4T, we therefore use the Taylor expansion of Q around $1/\hat{t} = 0$, or equivalently $r \rightarrow +\infty$,

$$Q(\hat{t}) = -\frac{4}{\sqrt{3}\pi} \frac{1}{\hat{t}} + \mathcal{O}(\hat{t}^{-5}), \quad (16)$$

for the early inspiral, where the tides behave nearly adiabatically. In the interval $\hat{t} \in [-45, -44]$, we then transition to the full Eq. (14b) for $Q(\hat{t})$. We have decided on this transition interval for $\hat{t}$ as through this approach all relative and absolute errors of $Q(\hat{t})$ and $Q'(\hat{t})$ are below 10^{-4} .

We want to highlight that this change alone increases the waveform generation speed for long waveforms starting at 20 Hz by close to 2 orders of magnitude compared to SEOBNRV4T, as the adaptive ODE solver is not picking up numerical noise. It should additionally increase the model's accuracy in the early inspiral.

We further employ quasiuniversal relations as given in Refs. [16,148–152] to relate the free NS parameters $(\Lambda_{l>2}, \omega_{0l}, \Delta\omega_{0l}, C_l)$ in the model to the tidal deformability Λ_2 . This way, we can compute accurate waveforms for different EOSs. One can however also specify those parameters freely for uses in PE or to study their impact on the waveform. A more precise description of the quasiuniversal relations used for SEOBNRV5THM can be found in Appendix A.

It is known that for antialigned (counterrotating) NSs the total f -mode frequency $\omega_{f,l}$ can become zero and even negative for high spins, which makes the f mode unstable. In particular, Eq. (14a) can develop a proper singularity in the denominator of the first term for $l\Omega_0 - \Delta\omega_{0l} = \omega_{0l}$ for large, negative $\Delta\omega_{0l}$ associated to counterrotating NSs. Such a singularity makes the ODE integration impossible, but has not been considered previously for SEOBNRV4T. This regime of large (antialigned) spins requires an extension of the model, possibly through the addition of non-linear dynamical tidal effects, which we leave for future investigations.

This effect is essentially the Chandrasekhar–Friedman–Schutz, or CFS, instability (see Refs. [153,154]), where counterrotating linear perturbations on a spinning NS can become unstable to the emission of GWs. This leads to a spin-down of the NS over secular timescales, much shorter than the inspiral timescale, until $\omega_{f,l} > 0$ is reached and the instability disappears. This effect, in tandem with other instabilities like the r -mode instability in nascent stars, therefore limits the maximum allowed spin rate of astrophysical NSs expected in BNS systems, see the extensive reviews of Refs. [155,156].

In order to still be able to generate waveforms across the full parameter space for this work, we apply a smooth

asymptotic cutoff on the spin shift to ensure $|\Delta\omega_{0l}/\omega_{0l}| < 0.9$ on the $\Delta\omega_{0l}$ as taken from Ref. [150]. While this generated waveform in the instability regime is unphysical, this regime is also astrophysically implausible for NSs in inspiraling binaries. The observed maximum NS spin rates furthermore suggest the validity of our procedure [157,158], given that they hint at a maximum spin rate below the onset of the CFS instability for the $l \in \{2, 3\}$ f modes [156].

Recently, the authors of Ref. [159] have studied anti-aligned systems for the dynamical tides in the formulation of Ref. [46] for highly spinning antialigned NSs and followed up on their work in Ref. [160]. They find that the crossing of the resonance can induce eccentricity in the originally quasicircular system. We find the same development of eccentricity when naively continuing the integration of the EOB system of ODEs. The resulting waveform of such a continued integration is however clearly unphysical due to the assumption of quasicircularity in the model and in particular in the quasicircular derivation of k_l^{eff} from Refs. [42,46]. It was also shown before in Ref. [161] that the current description of dynamical tides in terms of the effective (circular-orbit) tidal deformability k_l^{eff} is not accurate in modeling eccentric BNS systems. While generic-orbit EOB Hamiltonians not based on k_l^{eff} have been derived [42], a complete dynamical-tidal EOB waveform model for eccentric orbits is missing.

As our model assumes quasicircularity and because we do not introduce the tidal torque that is suggested in Ref. [159] or the tidal spin that is discussed in Ref. [160], we have opted to terminate the waveform prior to the development of eccentricity, as soon as $\dot{p}_{r_*} > 0$.

In contrast, we want to note that the authors of Ref. [162] have NR simulated one such BNS system with large antialigned spins with the SACRA code and found no signs of eccentricity developing. Instead, they report a change of the spins of the individual NSs as they go through resonance and gain spin angular momentum from the orbit due to tidal torques. Currently, there only exists the very recent work of Ref. [160] investigating the magnitude of the dynamical tidal contribution to the spin, which could lead to angular momentum transfer during f -mode resonance, but has not been verified against NR. We therefore note the open question of the waveform of high spinning NSs and leave the development and inclusion of a consistent framework for angular momentum exchange due to f -mode resonances as future work.

For now, we suggest to only use the model up to medium NS spins of $|\chi_{\text{NS}}| \lesssim 0.3$ for this reason.

The resulting Hamiltonian obtained after incorporating all relevant matter information can then be used to solve the equations of motion of the aligned-spin BNS system. They are given in Refs. [90,115] and read

$$\dot{r} = \xi \frac{\partial H_{\text{EOB}}}{\partial p_{r_*}}, \quad \dot{p}_{r_*} = -\xi \frac{\partial H_{\text{EOB}}}{\partial r} + \frac{p_{r_*}}{p_\phi} \mathcal{F}_\phi, \quad (17a)$$

$$\dot{\phi} = \frac{\partial H_{\text{EOB}}}{\partial p_\phi}, \quad \dot{p}_\phi = \mathcal{F}_\phi, \quad (17b)$$

where p_{r_*} and ξ relate to the tortoise coordinates defined below, and the RR force $\mathcal{F}_\phi$ is obtained by summing the GW modes in factorized form $h_{\ell m}^{\text{F}}$ as in Refs. [87,88,163,164], which we define in Sec. II B, that is

$$\mathcal{F}_\phi \equiv -\frac{M\Omega}{8\pi} \sum_{\ell=2}^8 \sum_{m=1}^{\ell} m^2 |d_L h_{\ell m}^{\text{F}}|^2, \quad (18)$$

where $\Omega = \dot{\phi}$ is the orbital frequency, and d_L is the luminosity distance of the binary to the observer.

Like in the BBH baseline, we use the tortoise-coordinate p_{r_*} instead of p_r , since it improves the stability of the equations of motion during the plunge and close to merger [96,115,163]. The tortoise-coordinate r_* is defined through

$$\frac{dr_*}{dr} = \frac{1}{\xi(r)} \quad (19)$$

and its conjugate momentum is given by

$$p_{r_*} = \xi(r) p_r. \quad (20)$$

We further choose $\xi(r)$ as

$$\xi(r) = \frac{\sqrt{\bar{D}_{\text{noS}}(A_{\text{noS}} + a_+^2/r^2)}}{1 + a_+^2/r^2}. \quad (21)$$

Note that ξ has been constructed to reduce to the Kerr value $(dr/dr_*) = (r^2 - 2Mr + a_+^2)/(r^2 + a_+^2)$ in the test particle limit $\nu \rightarrow 0$.

As in SEOBNRV5HM, we use quasicircular adiabatic initial conditions derived in Ref. [85], and then solve for the binary dynamics by numerically integrating Eqs. (17a) and (17b) using an eighth order Dormand-Prince integrator.

B. EOB inspiral waveforms with matter effects

Like the Hamiltonian, our waveforms are built on the BBH SEOBNRV5HM multipolar waveform [115]. In the following, we give an overview of the waveform computation highlighting our changes to include matter effects.

In general, the complex linear combination of GW polarizations, $h(t) \equiv h_+(t) - ih_\times(t)$, can be expanded in the basis of -2 spin-weighted spherical harmonics [90] as follows:

$$h(t; \lambda, \iota, \varphi_0) = \sum_{\ell \geq 2} \sum_{|m| \leq \ell} {}_{-2}Y_{\ell m}(\iota, \varphi_0) h_{\ell m}(t; \lambda), \quad (22)$$

where λ denotes the intrinsic parameters of the compact binary, such as mass ratio q , spins $\chi_{1,2}$, and the NS

parameters $(\Lambda_{l \geq 2}, \omega_l, \Delta\omega_l, C_l)$. The parameters (i, φ_0) describe the binary's inclination angle (computed with respect to the direction perpendicular to the orbital plane) and the azimuthal direction to the observer, respectively.

In SEOBNRV5THM, the GW modes defined in Eq. (22) are decomposed into inspiral modes and premerger modes close to merger. All modes are further tapered to zero after the merger. SEOBNRV5THM can compute and return the waveforms of the dominant (2,2) mode and the largest subdominant modes: (3,3), (2,1), (4,4) (3,2), (5,5), and (4,3).

In this work, we model the premerger of the dominant (2,2) mode phenomenologically up to merger, as we will describe in detail in Sec. II C. For the subdominant modes, however, our premerger waveform only consists of a tapering of the inspiral-waveform close to and after merger, which we discuss in Appendix B.

For aligned-spin binaries it holds that $h_{\ell m} = (-1)^\ell h_{\ell -m}^*$, therefore we restrict the discussion to (ℓ, m) modes with $m > 0$. We have

$$h_{\ell m}(t) = \begin{cases} h_{\ell m}^{\text{inspiral}}(t), & t < t_{\text{trans}} \\ h_{\ell m}^{\text{premerger}}(t), & t \geq t_{\text{trans}}, \end{cases} \quad (23)$$

where we define the transition time t_{trans} and $h_{\ell m}^{\text{premerger}}(t)$ later in Sec. II C.

The factorized inspiral modes are written as [115]

$$h_{\ell m}^F = h_{\ell m}^N \hat{S}_{\text{eff}} T_{\ell m} f_{\ell m} e^{i\delta_{\ell m}}, \quad (24)$$

where details regarding the individual factors can be found in Eqs. (26)–(33) of Ref. [115]. The different factors in the modes depend on the set of ODE dynamical quantities (r, ϕ, Ω) .

We include tidal information into the amplitude $f_{\ell m}$ and phase factors $\delta_{\ell m}$ of the factorized modes in Eq. (24), which are computed such that the expansion of $h_{\ell m}^F$ agrees with the PN-expanded modes. For nonspinning binaries, $f_{\ell m}$ is further resummed as $f_{\ell m} = (\rho_{\ell m})^\ell$ to reduce the magnitude of the 1PN coefficient, which grows linearly with ℓ [88]. Following Refs. [104, 105, 164], we separate the nonspinning, spin, and tidal contributions for the odd m modes, such that

$$f_{\ell m} = \begin{cases} (\rho_{\ell m}^{\text{noS}} + \rho_{\ell m}^S + \rho_{\ell m}^{\text{tidal}})^\ell, & m \text{ even}, \\ (\rho_{\ell m}^{\text{noS}})^\ell + f_{\ell m}^S + f_{\ell m}^{\text{tidal}}, & m \text{ odd}, \end{cases} \quad (25)$$

where $\rho_{\ell m}^{\text{noS}}$ ($\rho_{\ell m}^S$) is the nonspinning (spinning) part of $\rho_{\ell m}$, while $f_{\ell m}^S$ is the spin part of $f_{\ell m}$ and $f_{\ell m}^{\text{tidal}}$ as well as $\rho_{\ell m}^{\text{tidal}}$ contain the factorized form of the relative 2.5PN adiabatic gravitoelectric tidal modes from Ref. [63]. The latter were found to be partially incorrect and were corrected in Refs. [65, 66], and are given in their factorized form in Ref. [66] and its supplementary material.

For the other tidal EOB models TEOBRESUMS and SEOBNRV5T, the tidal information in the modes has instead been added to the factorized point-mass modes through an ansatz of the form

$$h_{\ell m} = h_{\ell m}^F + h_{\ell m}^{\text{tidal}}. \quad (26)$$

We highlight that SEOBNRV5THM is therefore currently the only EOB model where tidal effects are instead included in the modes in their factorized form.

We employ the factorized $f_{\ell m}^S$ modes for the spin-induced multipole moments from Ref. [52], Sec. IV B, Appendix B, as well as its supplementary material. We include the spin-induced multipole moments in the modes consistently to their inclusion in the Hamiltonian up to 3.5PN order in SS, S^3 and S^4 .

In comparison with SEOBNRV4T, however, we have decided to omit the effective tidal enhancement of the adiabatic tidal modes due to dynamical effects from our model ([137] and Eq. (16) of Ref. [165]). This is because the expression used in LALSIMULATION is different from its published counterparts in Refs. [22, 161, 165]. As this discrepancy on how the dynamical tides should affect the waveform's modes is not resolved, and as this effect only becomes relevant towards merger, which we model phenomenologically, we have decided to omit this for now. We plan to revisit the effect of dynamical tides on the waveform modes in later work.

For the odd m modes, both $f_{\ell m}^{\text{tidal}}$ as well as the spinning $f_{\ell m}^S$ expressions contain terms that scale as $f_{\ell m} \propto 1/\delta$. This is a removable divergence, as the Newtonian factor in the factorized waveform $h_{\ell m}^N \propto \delta$ in Eq. (24), and those terms cancel, to give a finite result. We take care of this by factorizing out the $1/\delta$ dependence in the equal-mass case in the $f_{\ell m}$ expression as well as the Newtonian prefactor.

We employ nonquasicircular (NQC) corrections N_{22} , but only for the amplitude of the (2,2) mode and not in its phase or for any of the higher modes. NQCs are complex functions multiplied by the factorized waveform,

$$h_{22} = h_{22}^F N_{22}(r, \Omega, p_r, \lambda), \quad (27)$$

which are used in SEOBNRV5HM (also for the higher modes), SEOBNRV4T as well as TEOBRESUMS. They are aimed to model relevant deviations from circularity during plunge and towards the merger. It has to be noted that

- (i) SEOBNRV4T uses the BBH NQCs one recovers from running the BBH model with the same parameters (q, χ_1, χ_2) and applies N_{22}^{BBH} to the BNS waveform [137], while
- (ii) TEOBRESUMS computes BNS tuned NQCs that enforce BNS quantities at merger as fitted from NR, like frequency and its first time derivative as well as amplitude at merger [135].

We denote as “input values” the quantities like amplitude or frequency at merger that are extracted from NR and fit

across the parameter space. In the next section we will give our reasoning for replacing the NQC corrections for the phase with a phenomenological premerger model tuned to NR.

C. EOB premerger waveforms with matter effects

In the SEOBNRV5 model, the input values and peak of the (2,2) mode are enforced at $t_{\text{peak}}^{22} = t_{\text{match}}$. We use a similar approach as is being used in the BBH model SEOBNRV5HM [115], where t_{match} is computed as

$$t_{\text{match}} \equiv t_{\text{ref}} - \Delta t_{22}, \quad (28)$$

with the reference time $t_{\text{ref}} = t_{\text{ISCO}}$ being defined as the time at which the separation crosses $r = r_{\text{ISCO}}$. The radius of the geodesic innermost stable circular orbit (ISCO) in Kerr spacetime r_{ISCO} is defined in Ref. [166] and taken to be of a remnant BH with the same mass and spin one would predict from the merger of a BBH system of the same masses and spins as our compact objects, computed with NR fitting formulas from Refs. [167,168]. $\Delta t_{22} = \Delta t_{22}(\lambda)$ is a calibration parameter, to be determined by comparing against NR simulations in Sec. III.

We note that the SEOBNRV5 family of waveform models consistently employs the remnant ISCO crossing as a reference time in Eq. (28) instead of, e.g., a peak in the orbital angular frequency as is being used in the SEOBNRV5THM and TEOBRESUMS families. The ISCO approach has been found to be more robust across the BBH parameter space for SEOBNRV5HM, as the orbital angular frequency does not necessarily peak for all systems if the resummation of the A potential of the Hamiltonian is not specifically designed to enforce it. It is furthermore beneficial to employ a similar approach to determine t_{match} across all SEOBNRV5 waveform models, as this facilitates a future combination of matter effects, spin precession, and eccentricity in one SEOBNR waveform model.

In contrast to the BBH baseline model, we employ Eq. (28) in SEOBNRV5THM with a modified reference time,

$$t_{\text{ref}}^{\text{BNS}} = t_{2\text{ISCO}}, \quad (29)$$

which is defined to be the time when the separation crosses

$$r(t_{\text{ref}}^{\text{BNS}}) = 2r_{\text{ISCO}}, \quad (30)$$

which is twice the ISCO radius of the remnant BH that one would expect if the coalescing NSs were BHs, i.e., employing the same fitting formulas as in SEOBNRV5HM. The reason for this choice is that the BBH ISCO is not necessarily crossed for BNS systems. Rather, the ODE integration will stop prior to the ISCO crossing in many regions across the BNS parameter space, typically because there will be a peak in p_{r^*} .

This is reasonable from a physical point of view, as the NSs are likely to touch, become disrupted, or strongly deformed instead of plunging into each other like BHs. The ISCO radius of a NS is, for example, expected to be comparable to its physical radius $r_{\text{ISCO}} \approx r_{\text{NS}}$, but can be larger or smaller dependent on the NS's EOS and resulting compactness. Consequently, for different BNS systems, the ISCO crossing of the EOB system and the NSs touching each other and merging could occur in any order. Changing the arbitrary reference point to the earlier time of the $r_{\text{ref}} = 2r_{\text{ISCO}}$ crossing that is reached across the relevant parameter space therefore gives us more control on t_{match} . This control is of importance for a more accurate modeling of the time of merger and the premerger waveform, as we discuss below, but also to increase the robustness of our EOB model under small perturbations of the intrinsic parameters.

Regarding the waveform prior to the peak emission, we opt to take a different approach to what has been presented in the literature so far to model the nonlinear matter and gravity effects close to merger. Instead of BNS-tuned NQCs as have been developed and used in TEOBRESUMS [135], we develop a phenomenological model for the phase and amplitude close to merger. Our reason for this is that, on some parts of the parameter space, BNS-tuned NQCs can lead to a dephasing and overestimated amplitude of the waveform towards merger in comparison to NR. This can be seen in the top panel of Fig. 1, where we show the frequency-time evolution around merger for the exemplary SACRA simulation of a $(1.25 + 1.25)M_{\odot}$ BNS for different premerger models. We remind that an incorrect frequency evolution translates into a dephasing over time in comparison to NR.

When one uses the factorized waveform from Eq. (24) (without application of NQCs or a premerger modeling and only up to t_{match}), the frequency evolution of the EOB model is too “soft” in comparison with the NR one towards merger, i.e., the frequency does not increase fast enough towards merger. We find this result across all NR simulations that we have investigated and one could argue that this might be due to friction when the NSs touch and exchange matter [146], as well as higher order PN effects. Notably, there also exists the possibility that the development of tidal spin, even for nonspinning NSs, might “harden” the frequency evolution [160]. All of these phenomena are however yet to be modeled analytically in an efficient way. For example, a full evolution of the NS mode amplitudes during f -mode resonance (in contrast to our effective description) would increase the waveform generation speed substantially [42,160].

Despite the fact that TEOBRESUMS has BNS-tuned NQCs, its time of merger is not calibrated to NR, and it therefore enforces the NR input values ~ 1 ms too late for this particular simulation. As a consequence, one can see that its frequency evolution is also too soft, highlighting the

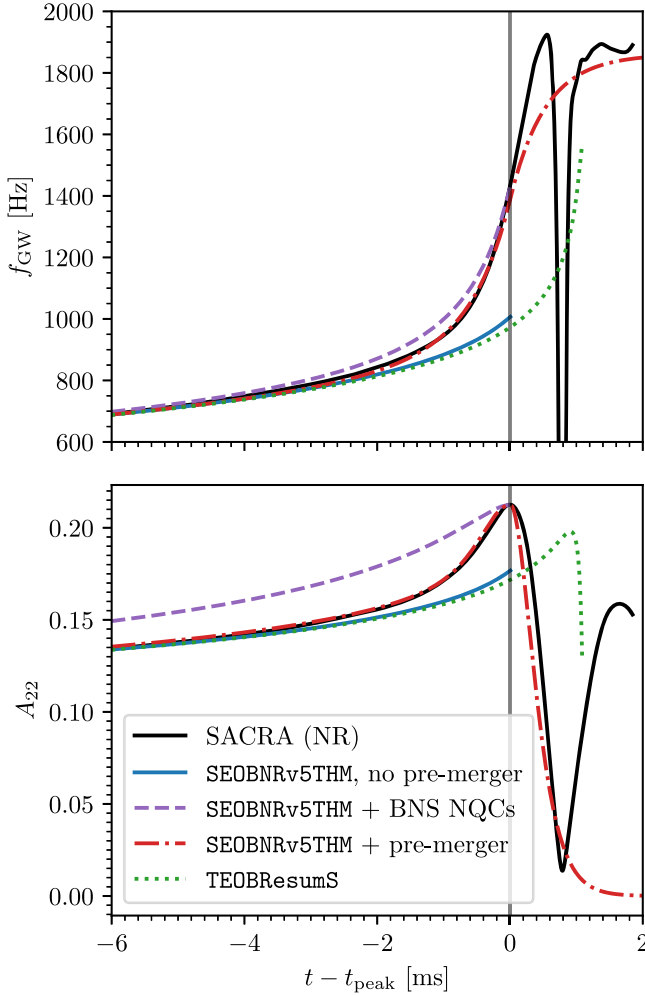


FIG. 1. The frequency-time (top) and amplitude-time evolution (bottom) of different premerger prescriptions compared to an exemplary SACRA simulation with EOS 15H for a $(1.25 + 1.25)M_\odot$ BNS. The NR waveform is shown as a solid, black line and the SEOBNRv5THM factorized waveform from Eq. (24) without application of any premerger model as a blue, solid line. The application of BNS-tuned NQCs on top of the factorized waveform is shown as a purple, dashed line and the phenomenological premerger and tapered model we use in SEOBNRv5THM as outlined in this section as a dot-dashed, red line. We further show TEOBRESUMS as a dotted, green line and indicate the time of merger of the NR simulation $t_{\text{peak}} = t_{\text{match}}$ with a vertical line.

importance of an accurate prediction of the merger time. This is also observable later in Sec. IV A, where TEOBRESUMS can accumulate dephasing towards merger for some NR simulations.

In contrast, note that the application of BNS-tuned NQCs to the SEOBNRv5THM waveform inspired by the TEOBRESUMS approach makes the frequency-evolution too hard, i.e., the frequency is too high during the premerger, in particular, as the NQCs are enforced at the correct time of merger and enforce the correct frequency and first frequency derivative at merger. For this work, we have tried out different

functional forms for the NQCs, by employing the same NQCs as are being used in SEOBNRv5HM [115],

$$N_{\ell m} = \left[1 + \frac{\hat{p}_{r_*}^2}{(r\Omega)^2} \left(a_1^{h_{\ell m}} + \frac{a_2^{h_{\ell m}}}{\hat{r}} \right) \right] \times \exp \left[i \left(b_1^{h_{\ell m}} \frac{\hat{p}_{r_*}}{r\Omega} + b_2^{h_{\ell m}} \frac{\hat{p}_{r_*}^3}{r\Omega} \right) \right]; \quad (31)$$

as well as the same form with modified powers of $\hat{r} = r/M$ and $\hat{p}_{r_*} = p_{r_*}/\mu$ to “harden” the amplitude and frequency evolution; and by testing the different NQC prescription employed by TEOBRESUMS, given in Eqs. (33)–(37) of Ref. [135] instead. We have found no satisfactory results across the parameter space covered by NR when employing NQCs with the input parameters enforced at the NR time of merger. Rather, we find that some of the NR waveforms will always be modeled more inaccurately due to the application of NQCs as described above.

We argue that this can be expected as the functional form of the NQCs is chosen to model deviations from quasi-circularity during the plunge. As a BNS waveform is largely driven by matter effects around merger instead, we suggest that they should not blindly be used to enforce BNS input values. Very similar observations also hold for the amplitude towards merger, as one can see in the bottom panel of Fig. 1. Note that for both panels we employ NQCs that enforce the correct frequency, first frequency derivative, amplitude, and a peak of the amplitude at t_{match} through Eq. (31).

Although it might be worthwhile to study how matter effects should modify the NQC functional form, starting from theoretical considerations in the future, in this work, we instead suggest a different, more agnostic approach. It is motivated by the idea of applying a simple exponential function to the frequency and amplitude of SEOBNRv5THM, which accurately models the frequency evolution as well as the amplitude evolution of a given BNS system up to merger across the entire parameter space covered by NR, as we show in detail in Sec. IV.

The (2,2) mode containing the premerger modeling and postmerger tapering has the form

$$h_{22}^{\text{premerge}}(t) = A_{22}^{\text{phen}}(t) W(t) e^{-i\phi_{22}^{\text{phen}}(t)}, \quad (32)$$

where the phenomenologically modeled phase $\phi_{22}^{\text{phen}}(t)$ and amplitude $A_{22}^{\text{phen}}(t)$ are defined below in Eqs. (38) and (43), respectively, and $W(t)$ is a Planck window taper, inspired by the tapering of SEOBNRv4T [137],

$$W(t) = \left[1 + \exp \left(\frac{t - t_{\text{match}}}{\tau_W} \right) \right]^{-1} \quad (33)$$

which in our case is centered at t_{match} and with width τ_W that is defined in Eq. (44) below.

Starting from the factorized waveform as given by Eq. (24),

$$h_{22}^F(t) = A_{22}(t)e^{-i\phi_{22}(t)}, \quad (34)$$

we want to boost the angular frequency $\omega_{22} = \dot{\phi}_{22}$ as

$$\omega_{22}^{\text{phen}}(t) = \omega_{22}(t)\beta(t|B_\omega, \tau_\omega, t_{\text{match}}), \quad (35)$$

by an exponential ansatz of the form

$$\beta(t|B, \tau, t_{\text{match}}) = 1 + B \exp\left(\frac{t - t_{\text{match}}}{\tau}\right). \quad (36)$$

This ansatz keeps the frequency evolution based on the EOB waveform, but exponentially hardens it towards merger as depicted in Fig. 1 and suggested by NR. Below we apply the same ansatz to the amplitude as well, but with different coefficients B and τ .

To obtain the phase of the premerger waveform $\phi_{22}^{\text{phen}}(t)$, we integrate the angular frequency for premerger times $t \leq t_{\text{match}}$,

$$\begin{aligned} \phi_{22}^{\text{phen}}(t) &= \phi_{22}(t_{\text{trans}}) \\ &+ \int_{t_{\text{trans}}}^t \omega_{22}(t')(1 + B_\omega e^{(t' - t_{\text{match}})/\tau_\omega}) dt', \end{aligned} \quad (37)$$

where t_{trans} is defined in Eq. (46) below and marks the start of the premerger waveform. This integral can be written in the form

$$\begin{aligned} \phi_{22}^{\text{phen}}(t) &= \phi_{22}(t) + B_\omega \phi_{22}(t') e^{(t' - t_{\text{match}})/\tau_\omega} \Big|_{t_{\text{trans}}}^t \\ &- \frac{B_\omega}{\tau_\omega} \int_{t_{\text{trans}}}^t \phi_{22}(t') e^{(t' - t_{\text{match}})/\tau_\omega} dt', \end{aligned} \quad (38)$$

where the integral in the last line is performed numerically through a cumulative sum of the integrand on a uniform sample grid of 10^3 points between t_{trans} and t_{match} , i.e., through the rectangle rule. Given the phenomenological nature of our premerger approach, we have decided against the usage of a more accurate, but computationally more intensive higher-order integration method.

We use fits for the BNS-tuned input values of the merger frequency ω_{22}^{NR} and merger amplitude A_{22}^{NR} from Ref. [169]. From Eqs. (35) and (36), one can see that

$$B_\omega \equiv \frac{\omega_{22}^{\text{NR}}}{\omega_{22}} - 1 \quad (39)$$

enforces the correct input value frequency at t_{match} . We further choose $\tau_\omega = 1.75\pi/\omega_{22}^{\text{NR}}$ as we find it produces a

good match with NR across all simulations and can be interpreted as the boost taking effect over the last 1.75 orbital cycles.

We simply continue the waveform frequency from the merger onwards to an asymptotic value of $\omega_{22}^{\text{asympt}} = 1.2\omega_{22}^{\text{NR}}$ by setting

$$\omega_{22}^{\text{phen}}(t) = \omega_{22}^{\text{asympt}} - \Delta\omega e^{-(t - t_{\text{match}})/\tau_{\text{taper}}} \quad (40)$$

for $t > t_{\text{match}}$ and by tuning $\Delta\omega$ and τ_{taper} such that the first frequency derivative is kept continuous. One finds the tapered phase by analytically integrating this equation. We highlight that a continuous first frequency derivative was not enforced in the tapering of SEOBNRV4T [137], which has a discontinuous $\dot{\omega}_{22}$ at merger.

Regarding the premerger model for the amplitude $A_{22}^{\text{phen}}(t)$, we want

- (i) to enforce the correct amplitude at merger,
- (ii) the peak to occur at t_{match} , and
- (iii) to taper the waveform to zero in a continuously differentiable way

while recovering the amplitude evolution as given by NR faithfully across all simulations used in this work. Our approach also has to produce a robust model without waveform generation errors, unintended artifacts or pathologies across the parameter space.

To this end, we have decided to employ one NQC that enforces the correct amplitude at merger and then using the boost function from Eq. (35) to enforce the NR calibrated merger amplitude and time of merger in the presence of the tapering $W(t)$, equalling one function per constraint.

We employ an amplitude NQC inherited from the NQCs as used in SEOBNRV5HM, see Eq. (31),

$$N_{22}(t) = 1 + \frac{p_{r_*}^2}{(r\Omega)^2} a_1^{h_{22}} \quad (41)$$

with one free parameter $a_1^{h_{22}}$ tuned such that

$$A_{22}(t_{\text{match}})N_{22}(t_{\text{match}}) = A_{22}^{\text{NR}} \quad (42)$$

with A_{22}^{NR} from Ref. [169].

We then boost the amplitude by applying the ansatz from Eq. (35) up to t_{match} . After the merger, we linearly continue the resulting amplitude based on the amplitude after application of the NQC and boost. Prior to windowing, we therefore arrive at a continuously differentiable, phenomenological amplitude,

$$A_{22}^{\text{phen}}(t) = \begin{cases} A_{22}(t)N_{22}(t)\beta(t|B_A, \tau_A, t_{\text{match}}) & \text{for } t \leq t_{\text{match}}, \\ A_{22}^{\text{phen}}(t_{\text{match}}) + \dot{A}_{22}^{\text{phen}}(t_{\text{match}})(t - t_{\text{match}}) & \text{for } t > t_{\text{match}}, \end{cases} \quad (43)$$

with two free parameters B_A and τ_A .

TABLE I. A summary of the free parameters and functions used in the phenomenological premerger model presented in this section. The free parameters are heuristically set to recover agreement with NR, while some parameters are constrained to recover input values, continuity in the first derivative throughout the merger, or consistency of the assumption of positive decay times $\tau > 0$.

Function	Definition	Parameter	Constraint
Phase boost $\phi_{22}^{\text{phen}}(t)$	Eq. (38)	B_ω	Input value
		τ_ω	Free
		$\omega_{22}^{\text{asympt}}$	Free
Postmerger phase	Eq. (40)	$(\Delta\omega, \tau_{\text{taper}})$	Continuity
Tapering $W(t)$	Eq. (33)	τ_W	Merger time
NQC N_{22}	Eq. (41)	$d_1^{h_{22}}$	Input value
Amplitude boost $A^{\text{phen}}(t)$	Eq. (43)	$B_A = 1$	Input value
		τ_A	Consistency
Residual freedom	Eq. (44)	τ_{free}	Free

By inserting Eqs. (32)–(43) into each other, one can check that the correct amplitude at merger t_{match} requires $B_A = 1$. By differencing those equations at $t = t_{\text{match}}$, one can also see that defining

$$\begin{aligned}\tau_A &\equiv \frac{\tau_{\text{free}}}{1 + [1 + \text{sgn}(\mathcal{A})]\tau_A}, \\ \tau_W &\equiv \frac{\tau_{\text{free}}}{1 + [-1 + \text{sgn}(\mathcal{A})]\tau_A}\end{aligned}\quad (44)$$

yields positive decay times $(\tau_A, \tau_W) > 0$ independently on the behavior of the original amplitude A_{22} as well as a correct time of merger. The free parameter $\tau_{\text{free}} = 1.5\tau_\omega$ is set heuristically to faithfully model the premerger amplitude across all NR simulations we have studied. We have further used the quantity

$$\mathcal{A} = \left. \frac{d \ln(A_{22}^{\text{phen}})}{dt} \right|_{t_{\text{match}}} = \frac{\dot{A}_{22}^{\text{phen}}(t_{\text{match}})}{A_{22}^{\text{phen}}(t_{\text{match}})}. \quad (45)$$

Through this procedure we get a phenomenological premerger model, which proves to be accurate and robust across the whole parameter space as we show in Sec. IV. We summarize all parameters which we have introduced for this purpose in Table I.

As a last step, we set the beginning of our premerger-tapered waveform t_{trans} from Eq. (23) as

$$t_{\text{trans}} = t_{\text{match}} - 15 \max(\tau_\omega, \tau_A, \tau_W), \quad (46)$$

the maximum of all of the exponential decay times.

For low tidal deformabilities $\kappa_2^T < 20$, we smoothly transition the matching time $t_{\text{match}}^{\text{BNS}}$ of Eq. (28) and the input values ω_{22}^{NR} and A_{22}^{NR} to the BBH quantities as used in SEOBNRV5HM through the smooth transition function,

$$\sigma(x) = \begin{cases} 0 & \text{for } x \leq 0, \\ \frac{\exp(\frac{-1}{x})}{\exp(\frac{-1}{x}) + \exp(\frac{-1}{1-x})} & \text{for } 0 < x < 1, \\ 1 & \text{for } x \geq 1, \end{cases} \quad (47)$$

rescaled to be applied on the interval $\kappa_2^T \in [10, 20]$. This way, we recover the BBH baseline to the 10^{-4} to 10^{-3} mismatch level in the absence of matter as we discuss further in Sec. IV C. There is, however, no exact match to BBH waveforms, due to the absence of a merger-ringdown model as well as the different premerger modeling.

As it was not the scope of this work to model the merger waveform for the higher modes, we simply taper them similarly to the SEOBNRV4T approach in an unmodeled way. We describe the tapering for the subdominant modes in Appendix B and leave the development of a similar but faster approach for the higher modes as future work.

The reasons for this choice are that higher order modes are suppressed for binaries of similar masses $q \lesssim 3$ like BNS systems [107], and that the mode frequencies scale approximately as $f_{lm} \simeq (m/2)f_{22}$ [74]. High m modes are therefore currently outside of the detector bands close to merger, and a correct premerger model of the (2,2) mode is of greater importance. Additionally, only very few BNS simulations have a decomposition of the NR waveform into spherical harmonics for us to compare or calibrate against. We therefore leave the extension of our model for the merger of all higher modes as future work.

We however perform a validation study of the higher modes compared to NR in Appendix C, which shows that the current model is nonetheless able to generate accurate higher modes until close to merger.

D. Waveform generation speed

We use the same procedure as is being used in SEOBNRV4HM for the PA approximation, see Eqs. (17)–(20) of Ref. [115], which allows us to speed up the computation of very long waveforms, as derived in Refs. [91, 138]. This technique has previously been used extensively in the TEOBRESUMS family of models [114, 130, 131, 170].

The idea is to obtain and solve explicit algebraic equations for the momenta over a grid of radii covering the inspiral, by perturbatively solving for the deviations from circularity $p_{r_*} = 0$. This way one can circumvent the explicit ODE integration to get an accurate and fast waveform. Like SEOBNRV5HM, we employ the PA approximation at the eighth order.

It has to be noted that the usage of the PA approximation, our changes in the k_l^{eff} explained around Eq. (16), and our code optimizations have increased the waveform generation speed of SEOBNRV5THM by 2 to 4 orders of magnitude in comparison to SEOBNRV4T for systems with $M \geq 1M_\odot$, which can be seen in Fig. 2. This allows us to use the current model for the PE of BNS events, as we show in Sec. V.

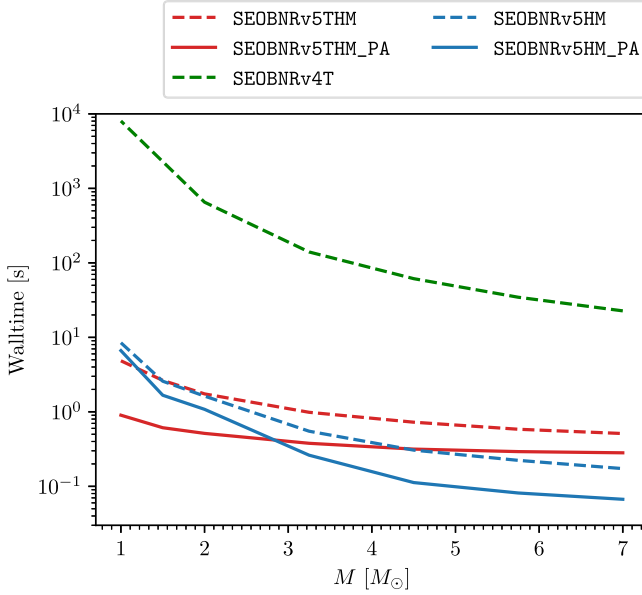


FIG. 2. Benchmarks for SEOBNRV5THM, our baseline SEOBNRV5HM, and the predecessor approximant SEOBNRV4T for a BNS system ($q = 1.5, \Lambda_1 = 700, \Lambda_2 = 800, \chi_1 = 0.1, \chi_2 = 0.2$) starting from 20 Hz where we only compute the (2,2) mode. For the SEOBNRV5 models, we also show the benchmarks if the PA approximation is applied as solid lines. We call both tidal approximants at a sampling rate of $f_{\text{sample}} = 8192$ Hz and note that for low total-mass systems, SEOBNRV5HM, which is run with $\Lambda_1 = \Lambda_2 = 0$, needs a higher sampling rate to resolve the ringdown, which is why SEOBNRV5THM can become faster than its baseline. These benchmarks were run on an Apple M2 processor.

We highlight however that the speed of the SEOBNRV4T model has been greatly increased through the development of the frequency-domain SEOBNRV4T_SURROGATE in Ref. [137]. This has made the SEOBNRV4T model even faster than SEOBNRV5THM, and further allows the application of SEOBNRV4T in PE. As the surrogate has been built in frequency domain, its waveform generation speed is however not directly comparable to time-domain models, in particular for very long waveforms. This is also why we omit the timing comparison to the surrogate and NRTIDALV3 models from Fig. 2.

For reference, the generation of a single frequency-domain SEOBNRV4T_SURROGATE waveform from 20 to 4096 Hz with frequency sampling rate $\Delta f = 1/1024$ Hz takes on the order of 0.1 s, largely independent of total mass. The same holds approximately true for an IMRPHEMOMXAS_NRTIDALV3 waveform. In contrast, the generation of an EOB waveform in frequency domain is slowed down by an additional discrete Fourier transform. Compared to a surrogate, an EOB model has however the advantage of extensibility, i.e., new analytical or numerical information can be easily included as it becomes available. Furthermore, PE with SEOBNRV5THM is only twice as expensive as PE with IMRPHEMOMXAS_NRTIDALV3 when

using PARALLEL BILBY, even though the waveform generation cost is a few times higher, as we discuss in Sec. V. Note that this does not necessarily apply to all PE codes, in particular PE codes employing multibanding [21,171,172].

III. CALIBRATION TO NUMERICAL RELATIVITY WAVEFORMS

An advantage of the time-domain EOB models over frequency-domain models is the access to the binary's separation $r(t)$, enabling improved estimates of the merger time via Eq. (28).

Accurately predicting the merger time is crucial as an observable, potentially detectable by future facilities like Einstein Telescope [169]. Of more importance is, however, its use in refining the premerger waveform modeling as described in Sec. II B, as the matching time t_{match} significantly affects premerger frequency and amplitude evolution.

For the calibration and validation of SEOBNRV5THM, 45 NR waveforms simulated by the SACRA code [70,71,146], and 19 waveforms simulated with the BAM code [73–75] were employed. We provide more detailed information about the individual simulations in Appendix D. We use NR input values from Ref. [169] for the peak amplitude and frequency. The only quantity of SEOBNRV5THM calibrated in this work is thus the time of merger Δt_{22} given in Eq. (28). We do, however, also employ NR fits for the BBH ISCO of the BBH remnant used in the computation of t_{ref} , as explained below Eq. (28), as well as for the BNS input values as used in our phenomenological premerger model in Sec. II C.

For all simulations, we align our model at the beginning of the respective NR simulation by minimizing the (2,2) mode phase difference between a given NR simulation and our model,

$$\min_{(t_0, \phi_0)} \int_{t_{\text{start}}}^{t_{\text{end}}} |\phi_{22}^{\text{NR}}(t) - \phi_{22}^{\text{EOB}}(t - t_0) - \phi_0| dt, \quad (48)$$

where $\phi_{22}(t) = \arg[h_{22}(t)]$. We align SACRA waveforms, which tend to be longer than the BAM counterparts, over seven waveform cycles and BAM waveforms over five cycles, setting t_{start} at a point where the junk radiation has decayed. We then extract the Δt_{22} of Eq. (28) that matches the merger of the NR simulation to t_{match} of SEOBNRV5THM.

We find that the time of merger strongly depends on the employed resolution of the NR simulation, where for the convergent BAM simulations, which have lower total resolution than the SACRA simulations, a higher resolution typically corresponds to a later time of merger.

To have data for Δt_{22} that behaves continuously across the parameter space and can be fitted, we therefore have to extrapolate the resolution dependent Δt_{22} to a common resolution across the NR simulations used in this work. We illustrate this in Fig. 3. We observe that the extrapolation of

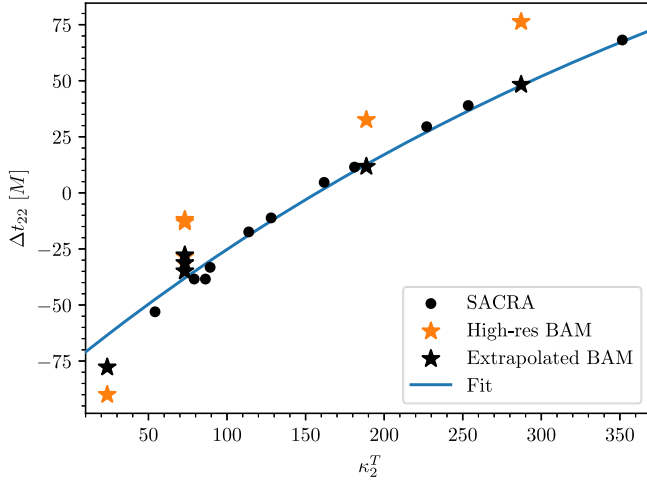


FIG. 3. The resulting Δt_{22} for the equal-mass NR simulations prior to and after the extrapolation to a common resolution. We show the recovered Δt_{22} that matches the merger of the NR simulation to t_{attach} from Eq. (28) for all used equal-mass simulations, given their respective quadrupolar tidal coupling constant κ_2^T . We indicate with black dots the employed SACRA simulations and with stars the BAM simulations. The originally recovered Δt_{22} values for the highest-resolution BAM simulation are shown in orange, and the Δt_{22} extrapolated to SACRA resolution in black. We also indicate the resulting fit from Eq. (52) as a blue line.

the merger times for BAM, as described below, makes Δt_{22} continuous and fittable across the parameter space. Note that we fit Δt_{22} across tidal parameters and mass ratio simultaneously, which yields the slight mismatch for the low tides, equal-mass case in Fig. 3.

As all SACRA simulations employ a uniform highest grid spacing of $n_{\text{grid}} = 182$, while the BAM simulations have lower resolutions,² but are convergent at a clear convergence order, we have opted to extrapolate the BAM Δt_{22} to a value that corresponds to the SACRA resolution as depicted in Fig. 4.

For this matter we extract the time of merger among the four highest resolution BAM simulations,³ and fit the resolution-dependent merger time as

$$t_{\text{peak}}(1/n_{\text{grid}}) = a/n_{\text{grid}}^2 + b. \quad (49)$$

We extract the value of $t_{\text{peak}}(1/n_{\text{SACRA}})$ and then calibrate as an extension of Eq. (28):

$$t_{\text{peak}}^{22} + \Delta t_{\text{peak}}^{(\text{extrap})} = t_{\text{ref}} - \Delta t_{22}, \quad (50)$$

²The way in which the SACRA and BAM simulations label their grid is different in such a way that the SACRA grid number is twice the one of BAM.

³If only three runs are available or the fourth run has a too low resolution, we opt to only use the three highest resolution runs for the fit.

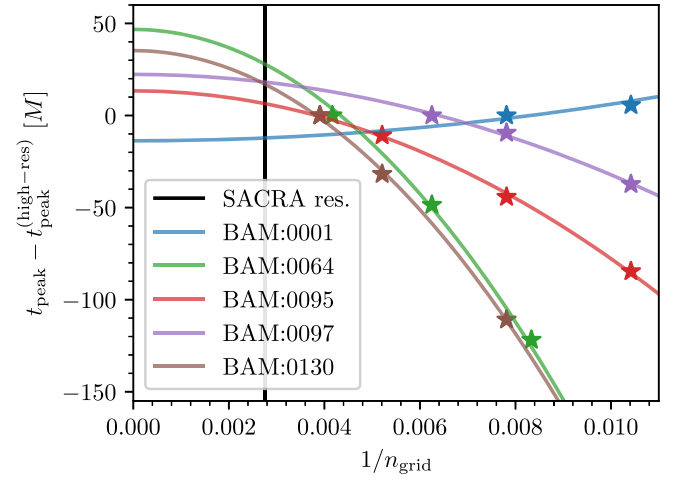


FIG. 4. Extrapolation of the merger time for some BAM simulations of different resolutions $1/n_{\text{grid}}$. We show the difference between the time of merger t_{peak} for runs at different resolutions and the time of merger of the respective highest resolution run $t_{\text{peak}}^{(\text{high-res})}$ of a given system as stars of the same color. We also include the respective quadratic fit of the data as a solid line. The resolution employed by SACRA simulations is indicated by a black vertical line, from where we extract $\Delta t_{\text{peak}}^{(\text{extrap})}$ as the value of the quadratic fit at the SACRA resolution, see Eq. (51).

where

$$\Delta t_{\text{peak}}^{(\text{extrap})} = t_{\text{peak}}(1/n_{\text{SACRA}}) - t_{\text{peak}}^{(\text{high-res})}. \quad (51)$$

We have also explored the possibility of extrapolating both SACRA and BAM to infinite resolution, but as the SACRA waveforms are not clearly convergent, we have found their merger times to not be clearly convergent. By trying to extrapolate the SACRA merger times, one therefore introduces noise and the resulting Δt_{22} becomes less continuous and therefore fittable. Additionally, the agreement with NR in Sec. IV worsens through this procedure, as we are not calibrating to the finite-resolution NR waveforms anymore but their extrapolations instead.

This is also the reason why in Sec. IV C our mismatches compared to the SACRA simulations are better than our mismatches to BAM simulations, as we are not calibrating to the highest resolution BAM simulations directly.

The resulting values of Δt_{22} at SACRA resolution are then indeed continuous across the parameter space, as one can see in Fig. 5, where our fit is shown in the background together with the parameters used for training and validation. We employ the fitting function,

$$\Delta t_{22}(X, \kappa_2^T) = a - b e^{-c \kappa_2^T (1 + dX + eX^2) + f} - gX, \quad (52)$$

where $X = 1 - 4\nu \in [0, 1]$ and use 80% of the simulations for the fit while leaving 20% of the simulations for validation.

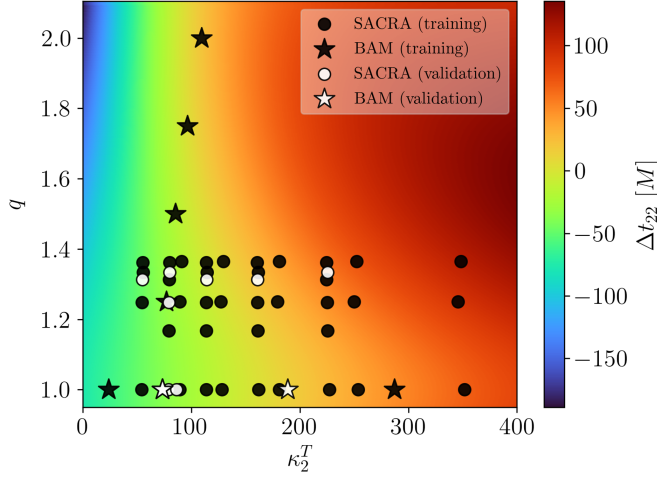


FIG. 5. Δt_{22} fit across the parameter space of the employed simulations. In the background of the image we show the result of our fit, while the data points show the respective mass-ratio q and quadrupolar tidal coupling constant κ_2^T of the employed non-spinning simulations. We indicate SACRA simulations as circles and BAM simulations as stars. Points with a black (white) interior have been used for fitting (validation).

The resulting fitting parameters are shown in Table II. We choose the above functional form to ensure an asymptote $a - gX$ for tidal deformabilities outside of the calibrated region $\kappa_2^T \gg 400$, which is suggested by the data and furthermore keeps our t_{match} bound across the whole parameter space as $\Delta t_{22} \in [-970, 186]M$ for any system, which prevents waveform generation failures. The quality of the fit is expressed in Table III.

TABLE II. The fitting parameters for Δt_{22} in Eq. (52).

a	b	c	d
215.5214	177.0887	1.932128×10^{-3}	31.37043
e	f	g	...
59.45997	0.500273	892.9965	...

TABLE III. The quality of the SEOBNRV5THM Δt_{22} calibration and parameters for the fit. The table lists the mean absolute error and 90% percentile of the absolute error of the (extrapolated) t_{match} in units of total mass and for a $(1.4 + 1.4M_\odot)$ BNS system in ms. We also compare the absolute error of the SEOBNRV4T merger time after alignment as a comparison of a model without calibration.

	Mean absolute error		90% absolute error	
	[M]	[ms]	[M]	[ms]
Fitting	4.87	0.0672	10.03	0.138
Validation	3.74	0.0515	6.74	0.0929
SEOBNRV4T	10.33	0.1425	13.83	0.191

We did not calibrate Δt_{22} to spinning BNS systems, as there are not enough NR simulations to cover a large portion of the parameter space, but have instead opted to use $\Delta t_{22} + \Delta t_{22}^{\text{spin}}$ from the baseline model SEOBNRV5HM; see Eqs. (73) and (80) of Ref. [115] to include the effects of nonprecessing, spinning systems on the merger time. As t_{match} is computed in a linear way in both cases, this should include all point-mass spin effects that augment the merger time for BNS systems.

IV. VALIDATION OF TIDAL WAVEFORM APPROXIMANTS

To assess the impact of the changes and improvements introduced in the SEOBNRV5THM waveform model, we compare it to the set of NR simulations described in Sec. III, and to other state-of-the-art BNS approximants, namely the predecessor model SEOBNRV4T, as well as SEOBNRV5_ROM_NRTIDALV3 and TEOBRESUMS-GIOTTO. We also assess the accuracy of SEOBNRV5THM when modeling BBH systems by comparing against SEOBNRV5HM.

We do so by showing a select few time-domain phasings of BNS simulations with the available waveform approximants in Sec. IV A, by performing unfaithfulness computations to NR simulations in Sec. IV B, as well as between different waveform approximants across the parameter space of interest in Sec. IV C.

A. Phasing comparison

To compare the different approximants in time domain to NR simulations, we use the same alignment procedure described in Eq. (48). As the BAM waveforms are convergent at a clear convergence order, we Richardson extrapolate the phase as in Ref. [145] to infinite resolution. To this end, we redefine the phase of the highest resolution NR waveform as

$$\phi_{\text{Richardson}} = \frac{s\phi_{\text{high}} - \phi_{\text{low}}}{s - 1}, \quad (53)$$

where

$$s = \left(\frac{n_{\text{high}}}{n_{\text{low}}} \right)^k > 1, \quad (54)$$

k is the convergence order found through the standard procedure (see, e.g., Refs. [68,75]), n are the number of grid points, and ϕ is the phase of the highest and second highest (low) resolution waveforms.

As described in Sec. III, the merger time depends on the employed resolution of the NR simulation. When Richardson extrapolating the waveform through the above equation, one therefore runs into the situation where one of the simulations has merged and the other is still in its

premerger phase, which creates unphysical artifacts in the frequency and therefore phase.

This effect is amplified when one instead tries to Richardson extrapolate the amplitude time series, due to the peak at merger of the individual runs. If a lower-resolution run merges earlier than the highest-resolution run, a Richardson extrapolated amplitude will unphysically increase following the lower-resolution merger, which would distract from our results. As this scenario applies to nearly all of the NR simulations we study in this work, and as a correctly modeled phase is of more importance than a correctly modeled amplitude with regards to PE, we do not extrapolate the amplitude in our comparisons. We instead combine the highest-resolution amplitude with the Richardson extrapolated phases to arrive in our strain h_{22}^{NR} against which we then compare.

Similarly to what has been done for NRTIDALV3 in Ref. [145], we also indicate the numerical uncertainty as shaded error bands in the phase difference $\Delta\phi = \phi_{22}^{\text{approx}} - \phi_{22}^{\text{NR}}$. In particular, as the SACRA simulations are not convergent at a clear convergence order, we calculate the phase difference between the highest and third-highest resolutions of a given system. The errors may be underestimated in this case. For NR simulations where we find a clear convergence order, we use the phase difference between the Richardson-extrapolated waveform and the highest resolution simulation as an error estimate.

The results of the time-domain comparison can be seen in Fig. 6. We want to highlight the effects of our premerger modeling as described in Sec. II B and of our calibration of the merger time as described in Sec. III. Although most of the approximants stay inside the error band of the phase difference $\Delta\phi$, it is evident that SEOBNRV5THM consistently stays in and close to the NR phase all the way to merger for both SACRA and BAM.

SEOBNRV5THM additionally predicts the merger time best for the SACRA simulations and models the amplitude correctly all the way up to merger. For BAM waveforms this is not necessarily the case, but we want to remind the reader that we do not Richardson extrapolate the amplitude and that the merger time is therefore shown for the highest resolution BAM simulation, not the prediction for SACRA-level resolutions that we calibrated against, as described in Sec. III.

It is therefore to be expected that our approach models the immediate amplitude of the highest resolution SACRA simulations correctly all the way up to merger, while predicting a later merger and, hence, a different amplitude evolution from the highest resolution BAM simulations. We hope that, in the future, higher resolution BAM simulations will support our methodology of predicting the extrapolated merger time.

Note that, for some simulations, the phase difference of all approximants may show jumps in Fig. 6 during the postmerger. This is due to sudden changes in the NR

frequency ω_{22}^{NR} soon after merger and prior to collapse of the remnant. This abrupt frequency modulation is associated to violent radial bounces of the remnant's core prior to collapse, see, e.g., Refs. [169, 173], and is not modeled by any approximant so far. As a consequence, all waveform models may show a discontinuous phase difference during the postmerger, if a simulation exhibits this feature.

It is furthermore interesting to note that the inclusion of the spin-shifted dynamical tides lowers the agreement with the highest resolution aligned-spin BAM simulations in comparison to SEOBNRV4T or TEOBRESUMS, but that it does perform comparable to these models when one looks at the Richardson-extrapolated waveform. As all approximants lie within the NR error band, however, it is not yet possible to distinguish the accuracy of including the spin-shifted dynamical tides in BNS simulations (but see also the investigation in Ref. [46], including a comparison to high-accuracy simulations of NS-BH binaries).

B. Faithfulness against numerical relativity

We want to compare all available BNS approximants against all of the NR simulations used in this work by computing mismatches between the approximants and simulations. The mismatch (or unfaithfulness) between two waveforms $h_1(t)$ and $h_2(t)$ is given by

$$1 - \mathcal{M} = 1 - \max_{(\phi_c, t_c)} \frac{(h_1(\phi_c, t_c)|h_2)}{\sqrt{(h_1|h_1)(h_2|h_2)}}, \quad (55)$$

where the overlap is defined in the frequency domain as

$$(h_1|h_2) = 4\text{Re} \int_{f_{\min}}^{f_{\max}} \frac{\tilde{h}_1^*(f)\tilde{h}_2(f)}{S_n(f)} df. \quad (56)$$

The maximization of the overlap over some arbitrary phase ϕ_c and time shift t_c in the time domain corresponds to a phase and frequency shift in frequency domain and ensures the alignment between the two waveforms similar to Eq. (48). In line with the procedure for SEOBNRV5THM in Ref. [115], we use the design zero-detuned high-power noise power-spectral-density (PSD) of Advanced LIGO from Ref. [174] for our mismatch computations.

We taper the beginning of time-domain waveforms with a Planck window, and to ensure no mismatch contributions from the tapering we choose the lower bound of the mismatch integral as $f_{\min} = 1.25f_{\text{start}}$, with f_{start} the point of the NR waveform where the junk radiation has decayed.

To estimate the NR error, we compute mismatches between the waveforms that were also used to create the error bands in the previous subsection, i.e., the third-highest resolution SACRA waveforms or the highest resolution waveforms for BAM, where we compare the approximants against the Richardson-extrapolated waveforms.

For SACRA we then set the upper bound of the mismatch integral as $f_{\max} = f_{\text{mrg}}$ of the highest resolution simulation.

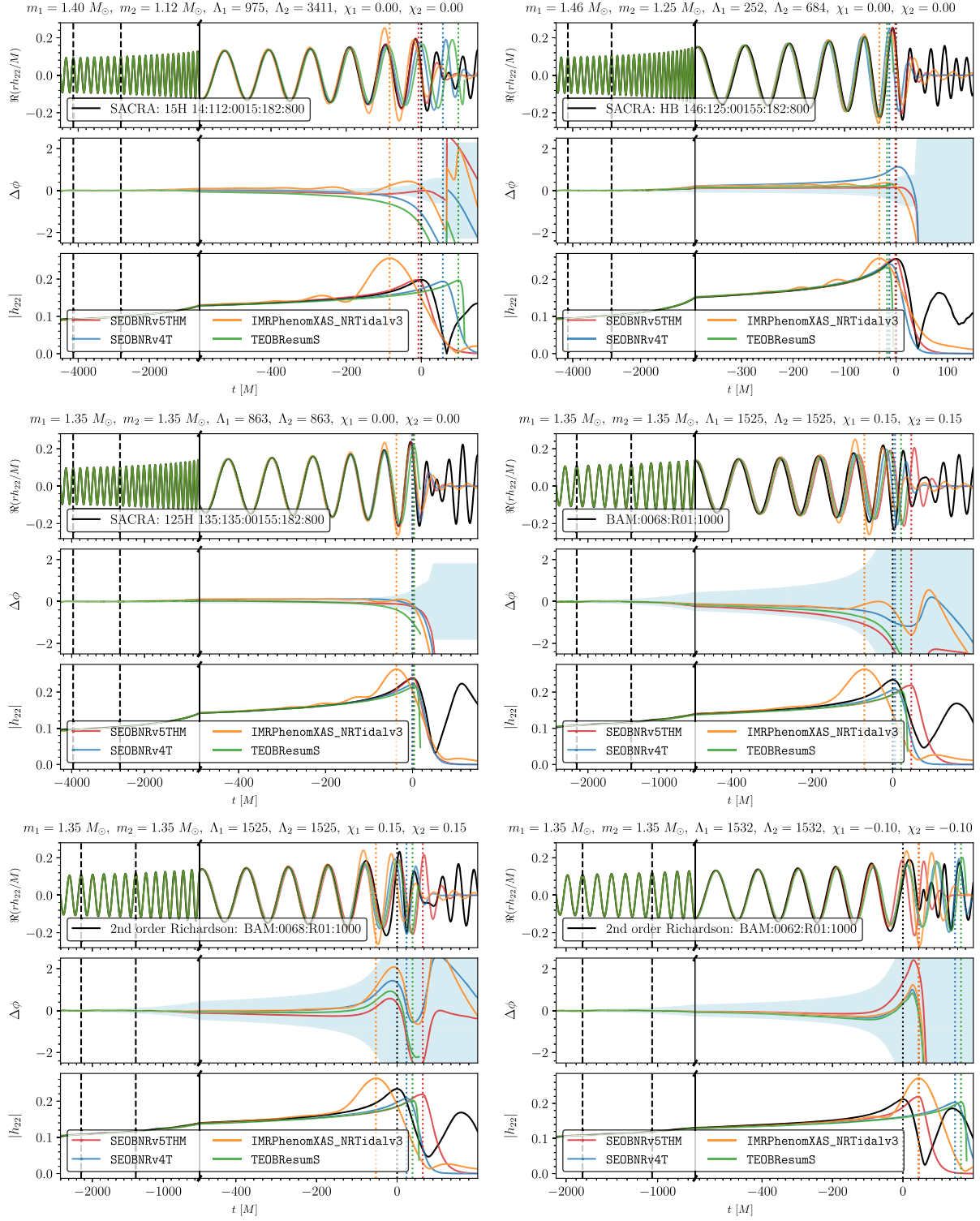


FIG. 6. Time-domain comparisons for three SACRA and two BAM simulations against four approximants as different colors. For each NR simulation, we state the BNS parameters above the figure. For each figure, the upper panel shows the real part of the (2,2) mode, while the middle panel shows the phase difference between the waveform model and the NR waveform, and the lowest panel shows the amplitude of the (2,2) mode. We show the inspiral and merger with two different timescales to focus on the merger. We further indicate the merger times of the approximants with vertical dotted lines. We compare against the BAM:0068 simulation 2 times, once against the highest resolution simulation and once against the Richardson-extrapolated waveform at convergence order 2, which we indicate in the legend.

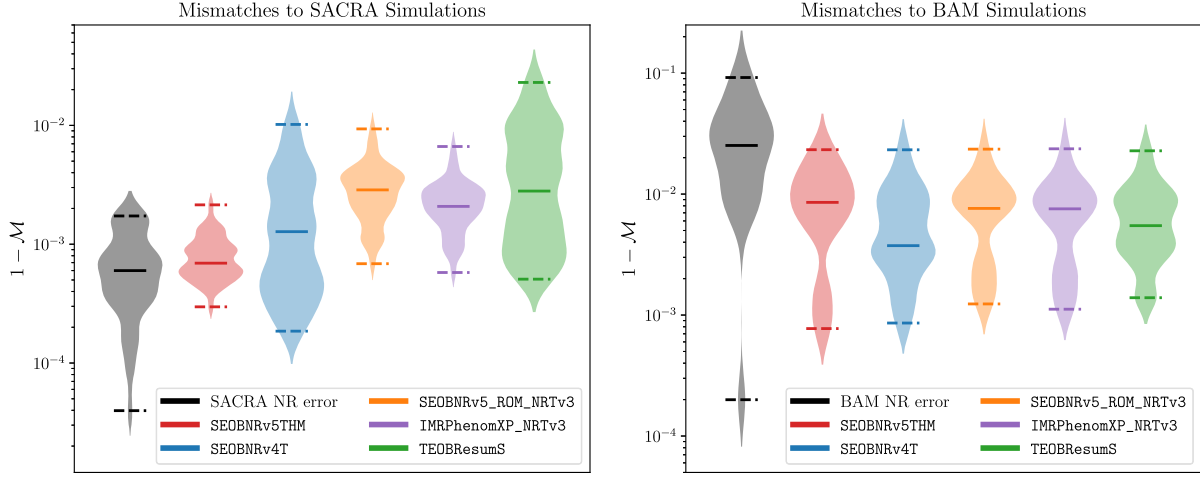


FIG. 7. Mismatch comparison against SACRA waveforms on the left and BAM waveforms on the right, starting at 1.25 times the starting frequency of the simulation, and ending at the merger time for SACRA simulations, or the minimum of the merger times of the simulations used for the Richardson extrapolation for BAM simulations. We use the advanced LIGO PSD and indicate with a solid line the median mismatch across all simulations and with dashed lines the minimum and maximum mismatch across all simulations. We highlight the different scales of the y axis between the left and the right figure.

For BAM, on the other hand, we choose the frequency of the Richardson-extrapolated waveform that corresponds to the minimum merger time of the two simulations used in the Richardson extrapolation,

$$f_{\max} = f_{\text{Richardson}}(\min(t_{\text{peak}}^{\text{high}}, t_{\text{peak}}^{\text{low}})). \quad (57)$$

We employ this procedure to reduce the NR error mismatches in comparison to BAM, which are otherwise larger due to artifacts in the Richardson-extrapolated waveforms beyond this time.

As the TEOBRESUMS waveform terminates at merger, see Ref. [135], we taper the waveform with a similar approach as in SEOBNRv4T prior to Fourier transforming. We choose a sampling rate of $2^{13} = 8192$ Hz for all of the approximants, only compare the (2,2) mode of all waveforms, and show our results as violin plots of the smoothed mismatch histograms in Fig. 7.

We note that, when looking at the 45 SACRA simulations, SEOBNRv5THM has the lowest maximum and median mismatch of all BNS approximants and remains comparable to the SACRA NR error. This shows the overall accuracy of the model due to the calibration and premerger modeling. When comparing against BAM, the model performs slightly worse, but comparable to the other approximants. We, however, want to raise awareness to the higher NR error mismatches for these simulations. It is therefore hard to decide on a best approximant for the BAM simulations, as all approximants lie within the NR error.

The low outlier mismatch of order 10^{-4} corresponds to BAM:0001, which has a very low mismatch but also possesses the lowest tidal deformabilities of all of the BNS simulations at hand. Note, however, that the highest-resolution run of BAM:0001 is rather low compared to the

resolutions of the other BAM simulations, as can be seen in Fig. 4. The NR error computed with the Richardson extrapolation might therefore be underestimated in this particular case.

C. Mismatch comparison to other tidal waveform approximants

Finally, we want to compare our model with other tidal models for nonspinning and aligned-spin configurations across the relevant BNS parameter space. For BNS systems we therefore compute the mismatch between SEOBNRv5THM and TEOBRESUMS, SEOBNRv4T, as well as SEOBNRv5_ROM_NRTIDALV3 and IMRPHENOMXP_NRTIDALV3.

All mismatches computed in this section are performed as in Eq. (55) for the (2,2) mode of the approximants between $f_{\min} = 20$ Hz, and $f_{\max} = 2048$ Hz. For computational efficiency we employ the SEOBNRv4T-surrogate from Ref. [137] instead of SEOBNRv4T directly. We call SEOBNRv5THM consistently in frequency domain or time-domain dependent on the native domain of the approximant it is compared to, to avoid artifacts from applying different Fourier transform prescriptions to the individual waveforms. When comparing against the time-domain waveform TEOBRESUMS, we use sampling rates of $2^{13} = 8192$ Hz, while we use $\Delta f = 1/1024$ Hz when comparing to the frequency-domain waveforms SEOBNRv4T-surrogate, NRTIDALV3, and SEOBNRv5THM, which we call in frequency domain.

In the nonspinning case, the mismatches are computed for 2000 random configurations of masses $m_{A,B} \in [0.5, 3]M_{\odot}$ and tidal deformabilities $\Lambda_2^{(1,2)} \in [0, 5000]$, for comparisons done with TEOBRESUMS and NRTIDALV3. We reduce the mass range to $m_{A,B} \in [1, 3]M_{\odot}$ for the SEOBNRv4T-surrogate, which is only available up to mass ratio $q \leq 3$. We find

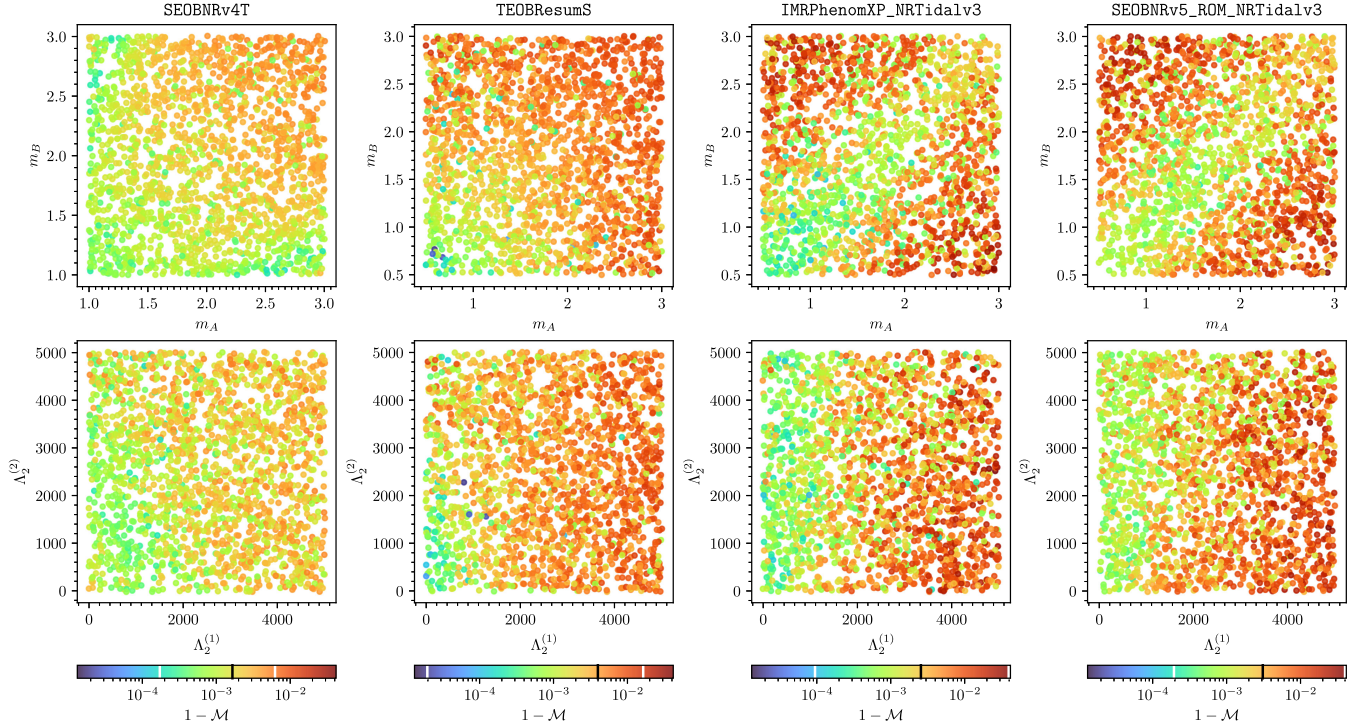


FIG. 8. Mismatch comparisons of SEOBNRV5THM with other tidal waveform models for nonspinning configurations in each column. Each mismatch comparison contains 2000 random configurations with tidal deformability $\Lambda_2^{(1,2)} \in [0, 5000]$ and $m_{A,B} \in [0.5, 3]M_\odot$, except for SEOBNRV4T, where we set the mass range to $m_{A,B} \in [1, 3]M_\odot$. For each subfigure we include a color bar indicating the values of the mismatches and state the name of the approximant in the title. We keep the minimum and maximum value in the color bar consistent across the different figures, but indicate the minimum (white), maximum (white), and median (black) mismatch against each approximant with solid vertical lines on the color bar.

that the mismatch increases with higher total mass, as the effect of the differing premerger prescriptions of the waveform approximants reaches more of the sensitive detector band. The reduction of the mass range for lower mass systems should therefore not affect our results for SEOBNRV4T drastically in comparison to the other waveform approximants.

The changes that we have introduced into SEOBNRV5THM in comparison to its baseline model SEOBNRV5HM, namely the omission of the phase NQCs and our novel premerger modeling approach, additionally yield differences in the waveform compared to SEOBNRV5HM even when the tidal deformabilities are set to zero. We also check that these changes have a negligible impact on PE by computing the mismatches of spin-aligned BBH systems between the two approximants.

For the BBH model comparison with SEOBNRV5THM, we reduce the mass range to $m_{A,B} \in [1, 3]M_\odot$, set $\Lambda_2^{(1,2)} = 0$ for SEOBNRV5THM, and consider an additional spin range $\chi_{1,2} \in [-0.5, 0.5]$. SEOBNRV5HM models the ringdown of the resulting BH and requires the QNMs to be resolved, which means that we need a high sampling frequency of the low total-mass BBH for this waveform model. For consistency, we thus call both SEOBNRV5 approximants in frequency domain up to a maximum frequency of $2^{15} = 32768$ Hz

(which translates into a high sampling frequency in time domain) when comparing them. The mismatches to SEOBNRV5HM stay below the 1.2×10^{-3} limit across the described BBH parameter space and are therefore not shown in this work. We find that the uncertainty in BNS waveform modeling introduced by tidal effects dominates the mismatch and therefore PE, and our changes to the BBH waveform will have a negligible systematic effect in the study of BNS signals.

The resulting mismatches for the BNS approximants can be seen in Fig. 8. Note that, while we assume the mass ordering of the primary and secondary $m_1 \geq m_2$, we will use randomly chosen m_A and m_B in the top row of this figure to ignore this mass ordering. Otherwise, only half of the figure would be filled with data points. In the lower rows, however, we show the tidal deformability $\Lambda_2^{(1)}$ ($\Lambda_2^{(2)}$) associated to the primary (secondary). This results in the individual subfigures possessing different symmetries.

Our findings are in line with the results of a similar study performed for NRTIDALV3 in Ref. [145], and our mismatches with SEOBNRV4T and TEOBRESUMS are lower than the reported mismatches of NRTIDALV3 with SEOBNRV4T and TEOBRESUMS. This is in line with our finding that the highest mismatches are in comparison to the NRTIDALV3 models, which goes up to a few 10^{-2} . Note also that, for

near equal-mass systems, the NRTIDALV3 models have lower mismatches to SEOBNRV5THM than the other EOB BNS models have to SEOBNRV5THM. For unequal-mass systems, however, the mismatches between NRTIDALV3 models and SEOBNRV5THM increase quite strongly. This mismatch dependence on mass ratio for NRTIDALV2 and NRTIDALV3 has been previously reported in Refs. [144,145] and appears to be a systematic effect for this class of approximants. It might be insightful to study its origin in the future. A possible hypothesis to be tested would be, for example, that this is caused by one of the analytical approximations on which the NRTIDAL-family is built being more valid for equal mass binaries than for unequal mass ones; namely, the approximation that the GW phase can be

linearly decomposed into different contributions, one of which being the tidal contribution [143].

The mismatches between SEOBNRV5THM and the other EOB approximants SEOBNRV4T and TEOBRESUMS are below or around the 10^{-2} level across the nonspinning BNS parameter space. This serves as a consistency check that there are no pathologies in the model and that SEOBNRV5THM stays in line with past results. The largest mismatches with the other EOB BNS models are found for high primary tidal deformability $\Lambda_2^{(1)}$ and high total mass M . Another noteworthy result is that the change of BBH baseline between SEOBNRV5_ROM_NRTIDALV3 and IMRPHENOMXP_NRTIDALV3 only plays a minor role in the mismatches of nonspinning BNS systems.

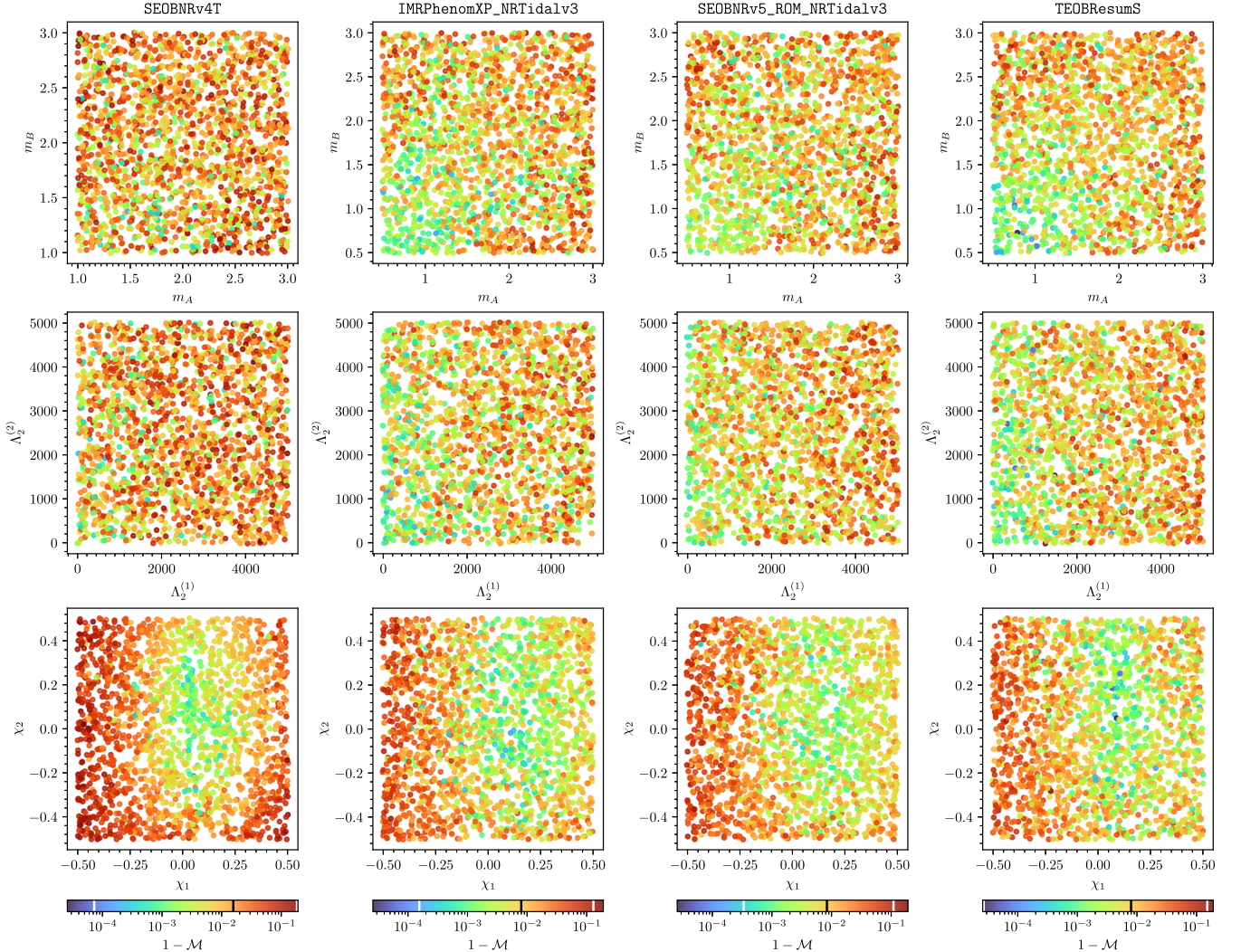


FIG. 9. Mismatch comparisons of SEOBNRV5THM with other tidal waveform models for aligned-spin configurations in each column. Each subfigure has 2000 random configurations of mass $m_{A,B} = [0.5, 3]$, tidal deformability $\Lambda_2^{(1,2)} = [0, 5000]$, and spin $\chi_{1,2} = [-0.5, 0.5]$, except for SEOBNRV4T, where the mass range is $m_{A,B} = [1, 3]$. For each subfigure we include a color bar indicating the values of the mismatches and state the name of the approximant in the title. We keep the minimum and maximum value in the color bar consistent across the different figures, but indicate the minimum (white), maximum (white), and median (black) mismatch against each approximant with solid vertical lines on the color bar.

In addition to the 2000 random configurations of the masses and tidal deformabilities that we generated for the nonspinning case, we also generate 2000 random aligned-spin BNS waveforms $\chi_{1,2} = [-0.5, 0.5]$ for the comparison of SEOBNRV5THM against SEOBNRV4T, SEOBNRV5_ROM_NRTIDALV3, IMRPHENOMXP_NRTIDALV3, and TEOBRESUMS, which can be seen in Fig. 9. For context, the maximum dimensionless spin an NS can have before it breaks up is about 0.67 [175], or about a spin rate of 1000 Hz (depending on the equation of state). The fastest-spinning pulsar known, J1748-2446ad, has a spin rate of 716 Hz [157], which translates into a dimensionless spin of about 0.4 [158]. It furthermore resides in the globular cluster Terzan 5 and has a light companion, which makes it a candidate for the production of GW sources due to dynamical capture in its future [159,176].

The mismatches here are much higher than for nonspinning systems and are mainly driven by highly spinning BNS configurations, with the highest mismatches being found for nonzero primary's spin $|\chi_1| \gtrsim 0.2$ and in particular for antialigned BNS systems $\chi_1 \lesssim -0.25$. The mismatches are the highest with SEOBNRV4T, which does not include the effect of spin-induced multipole moments above the LO quadrupole, in contrast to all other approximants studied in this work. For highly spinning systems, the mismatches can approach and exceed the 10^{-1} level, which is due to the spin-shifted dynamical tides [46] that have not been included in another waveform model that has been reviewed by the LVK collaboration so far.

Since we terminate antialigned BNS systems with large spin magnitudes prior to the development of eccentricity as described in Sec. II A, these waveforms are much shorter than the waveforms from approximants that do not include spin-shifted dynamical tides. Note that there is evidence from NR that such systems indeed have a much shorter inspiral phase [162], although further studies are needed.

For now, this early merger that we find for SEOBNRV5THM should explain the high mismatches up to the order of 10^{-1} for binary systems containing highly spinning, antialigned NSs in Fig. 9.

We leave the correct modeling of postresonance BNS waveforms as future work and note that the approximant should not be blindly trusted for NSs with large negative spins $|\chi_i| \lesssim 0.3$. We have still included the spin-shift effect in our model, as it is a theoretical prediction valid for low spins and serves as a good starting point to incorporate future analytical developments.

V. PARAMETER ESTIMATION

As a last validation to test the applicability of the SEOBNRV5THM model for data analysis, we reanalyze the two BNS detections GW170817 [2] and GW190425 [35] of the LVK collaboration. For this purpose, we utilize PARALLEL BILBY [177], which performs GW parameter estimation using a parallelized nested sampler named DYNesty [178]. In PARALLEL BILBY, the Bayes' theorem is employed:

$$P(\theta|d, \mathcal{H}) = \frac{\mathcal{L}(d|\theta, \mathcal{H})p(\theta|\mathcal{H})}{E(d|\mathcal{H})}, \quad (58)$$

where $P(\theta|d, \mathcal{H})$ is the posterior probability distribution of the waveform parameters θ given some data d , which in our case consists of the observed strain data, and hypothesis $\mathcal{H}$. $p(\theta|\mathcal{H})$ is the prior probability distribution, and $E(d|\mathcal{H})$ is the evidence, serving as normalization constant to $P(\theta|d, \mathcal{H})$. Meanwhile, $\mathcal{L}(d|\theta, \mathcal{H})$ is the likelihood of obtaining the data d containing a waveform with parameters θ given the hypothesis $\mathcal{H}$, derived from a Gaussian noise model with power spectral density extracted from the data around the GW event. For further details, we refer to

TABLE IV. Settings and evaluation times for the different PE runs on two real GW events GW170817 and GW190425 performed in this work. We employ the PARALLEL BILBY [177] sampler with the waveform models SEOBNRV5THM and IMRPHENOMXAS_NRTIDALV3. We list the maximum frequency ($f_{\max}$) starting from $f_{\min} = 20$ Hz in the data settings, while the number of live points (nlive), and the number of autocorrelation times (nact) are specified in the sampler settings. We use a duration of 128s for all of the runs in this table. The time reported is walltime, while the total computational cost in CPU hours can be obtained by multiplying this time by the reported number of CPU cores employed. In general, we find that an analysis using SEOBNRV5THM takes on the order of twice as long as an analysis using IMRPHENOMXAS_NRTIDALV3.

Event	Model	Data settings	Sampler settings		Computing resources	
		$f_{\max}$ [Hz]	nact	nlive	nodes $\times$ cores	Run-time
GW170817	IMRPHENOMXAS_	2048	65	1024	12×32	15 h
	NRTIDALV3					
	SEOBNRV5THM	2048	65	1024	16×32	1 d 6 h
	SEOBNRV5THM	1024	20	1024	12×32	18 h
GW190425	IMRPHENOMXAS_	2048	65	2048	16×32	1 d 3 h
	NRTIDALV3					
	SEOBNRV5THM	2048	65	2048	16×32	1 d 20 h

Refs. [179,180]. In our study, we use GW170817 and GW190425 using the open data from Ref. [181], and analyze these events with SEOBNRV5THM, IMRPHENOMXP_NRTIDALV3 (where we use the posteriors from the study performed in Ref. [145]), and IMRPHENOMXAS_NRTIDALV3, and also compare their results with each other. For the analysis of both events, the LVK collaboration employed a low-spin prior $|\chi| \leq 0.05$ and a high-spin prior analysis $|\chi| \leq 0.89$ [2,35]. For this work we have opted to only repeat the low-spin prior analysis. We remind of our finding in Sec. IV C

that the tidal approximants should not be blindly trusted for highly spinning NSs $|\chi| \gtrsim 0.3$. We follow the standard LVK collaboration analyses [2,22,34,35,179,180].

The data settings, sampler settings, and computing resources for the individual analyses performed in this work are listed in Table IV. We highlight that, when studying the total cost in terms of the CPU hours, SEOBNRV5THM appears to be $\mathcal{O}(\text{twice})$ as costly as the same analysis when employing IMRPHENOMXAS_NRTIDALV3. It appears that the reason for this is the order 1 s overhead per likelihood

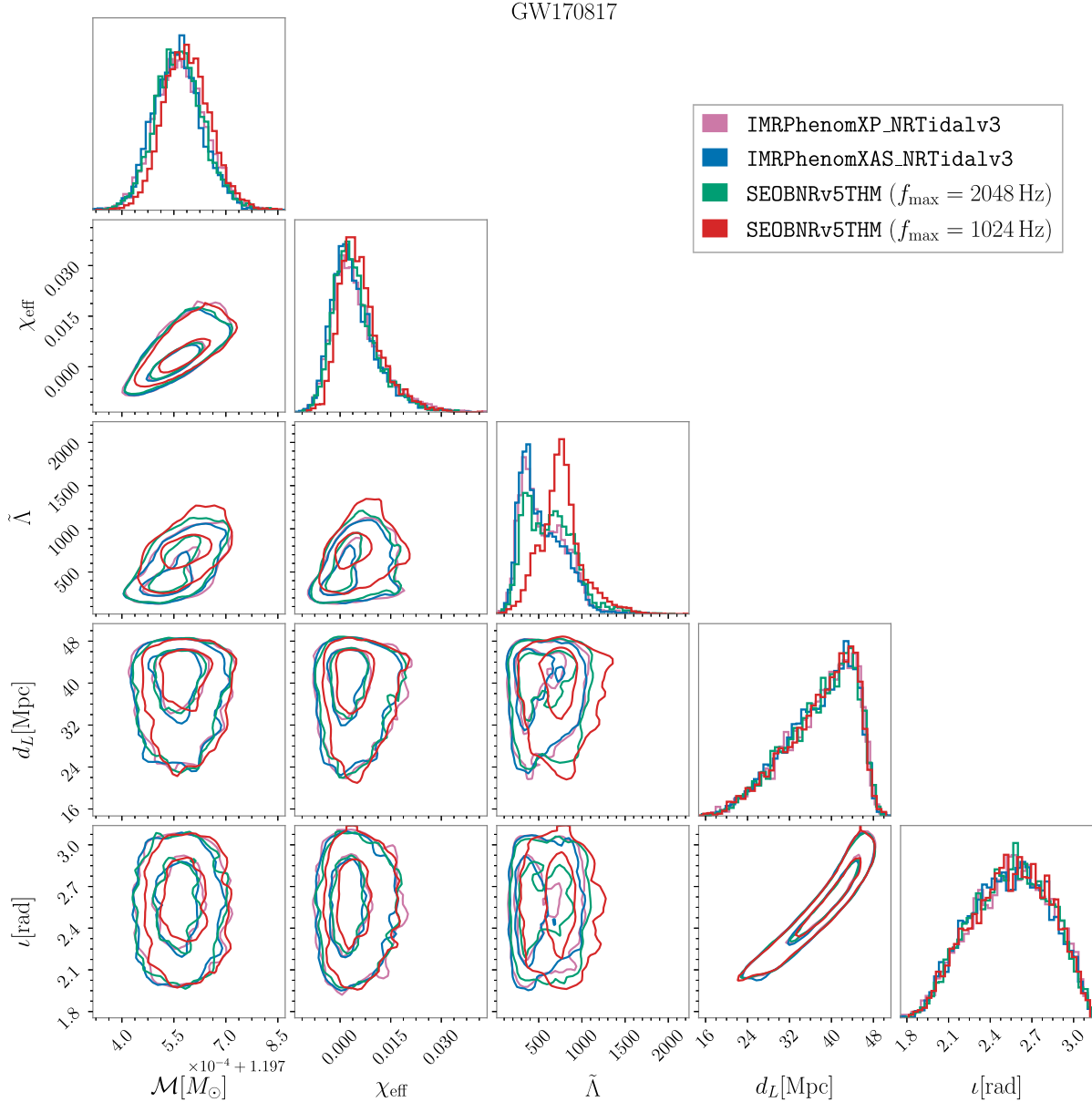


FIG. 10. The marginalized 1D and 2D posterior probability distributions for selected parameters of GW170817, obtained with IMRPHENOMXP_NRTIDALV3 (blue), IMRPHENOMXAS_NRTIDALV3 (orange), and SEOBNRV5THM (green/red). We run SEOBNRV5THM twice with different high-frequency cutoffs, at either 2048 Hz, as used for the NRTIDALV3 approximants (green), and at 1024 Hz in line with the analysis in Ref. [16]. The parameters shown here are the detector-frame chirp mass of the system $\mathcal{M}$ in solar masses, the effective spin χ_{eff} , the binary tidal deformability $\tilde{\Lambda}$, the luminosity distance d_L in Mpc, and the inclination angle i_0 . The 68% and 90% confidence intervals are indicated by contours for the 2D posterior plots.

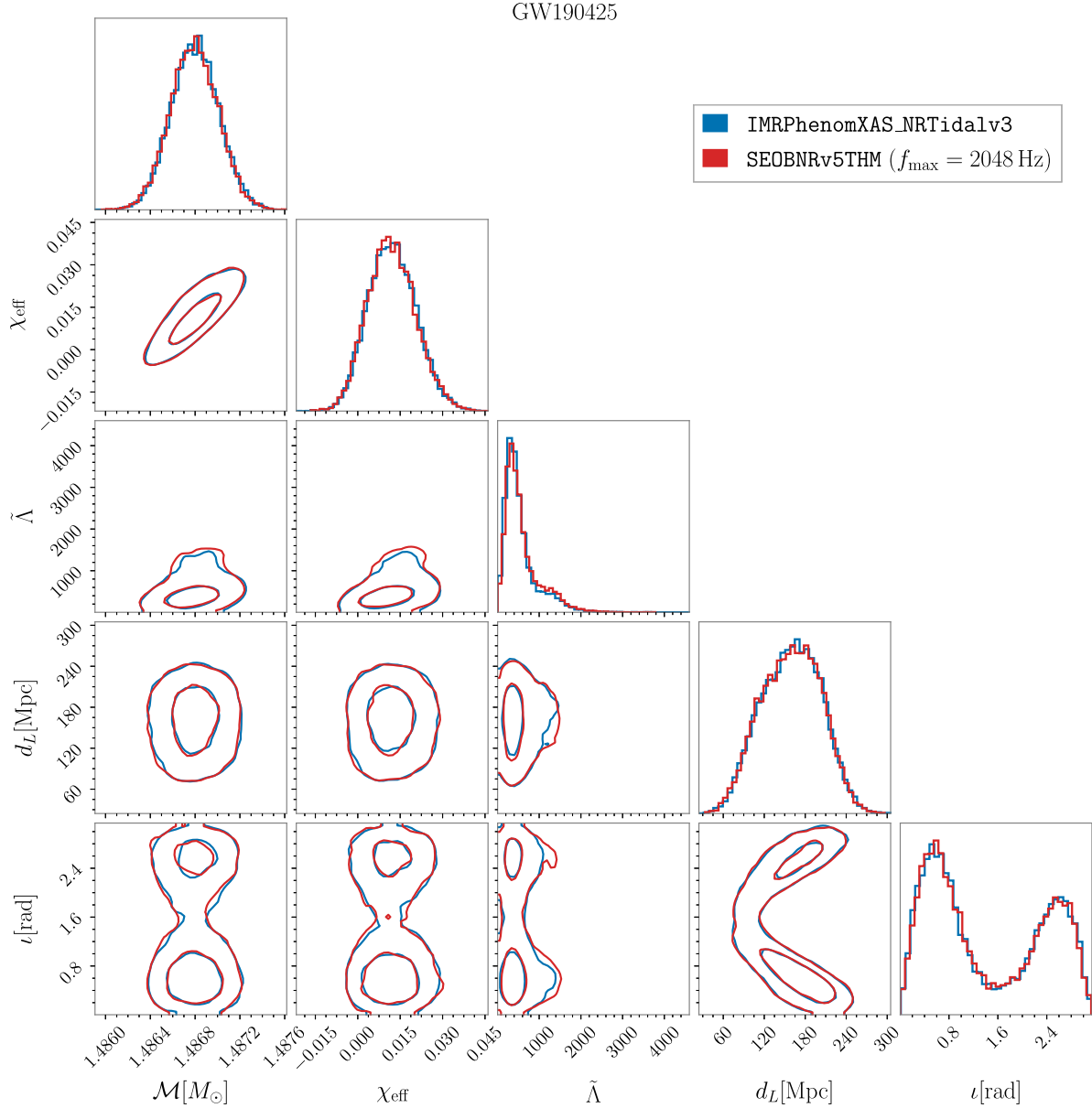


FIG. 11. The marginalized 1D and 2D posterior probability distributions for selected parameters of GW190425, obtained with IMRPHENOMXAS_NRTIDALV3 (blue), and SEOBNRV5THM (red) up to a high-frequency cutoff of 2048 Hz. The parameters shown here are the detector-frame chirp mass of the system $\mathcal{M}$ in solar masses, the effective spin χ_{eff} , the binary tidal deformability $\tilde{\Lambda}$, the luminosity distance d_L in Mpc, and the inclination angle ι . The 68% and 90% confidence intervals are indicated by contours for the 2D posterior plots.

evaluation in PARALLEL BILBY,⁴ which is shared between SEOBNRV5THM and IMRPHENOMXAS_NRTIDALV3. The effect of the much higher waveform generation speed

⁴We calculate the overhead by accessing the number of likelihood evaluations N and the total sampling time T from the result file of the respective PARALLEL BILBY runs. The time for one likelihood evaluation is then estimated as $T n_{\text{nodes}} n_{\text{cores}} / N$ with n_{nodes} (n_{cores}) the number of employed nodes (cores). The likelihood evaluation time is on average 2.6 s for SEOBNRV5THM and 1.4 s for IMRPHENOMXAS_NRTIDALV3 across both events studied in this work at $f_{\text{max}} = 2048$ Hz.

of IMRPHENOMXAS_NRTIDALV3 is therefore reduced during PE.

We show the results of the inferred marginalized 1D and 2D posterior probability distribution of a selection of source parameters for GW170817 in Fig. 10, for a low-spin prior with $\chi \leq 0.05$. We show the detector-frame chirp mass of the system $\mathcal{M}$ in solar masses, the effective spin χ_{eff} from Eq. (5), the luminosity distance d_L in Mpc, and the inclination angle ι . We also include here the dimensionless tidal deformability defined in terms of the individual masses and tidal deformabilities of the stars, see Ref. [22], which is related to κ_2^T :

$$\tilde{\Lambda} = \frac{16(m_1 + 12m_2)m_1^4\Lambda_2^{(1)} + (m_2 + 12m_1)m_2^4\Lambda_2^{(2)}}{13M^5}. \quad (59)$$

We further fix the sky position as recovered from the electromagnetic counterpart for GW170817.

From Fig. 10, we note that the performance of SEOBNRV5THM in terms of the inferred parameters is consistent with the results of IMRPHEMOPV2_NRTIDALV2 [34] when analyzing the signal up to 2048 Hz.

The main difference is the slightly larger secondary peak of SEOBNRV5THM for $\tilde{\Lambda}$ compared to the other models. Note that all pf the posterior weight goes towards this secondary peak when analyzing the event only up to 1024 Hz.

It has been shown in Refs. [16,182,183] that the first peak might be related to the glitch mitigation in Livingston. As discussed in Ref. [16], above 1 kHz only a signal to noise of $\lesssim 1$ should have been deposited for GW170817-like events (which roughly corresponds to 3% of the total SNR), further hinting at the origin of the first peak being related to data quality issues. When only performing PE up to a high-frequency cutoff of 1 kHz, our results are furthermore fully consistent with the results for TEOBRESUMS in Ref. [16].

The results for the marginalized posterior distributions for GW190425 can be found in Fig. 11, where all results are fully consistent between IMRPHEMOPV2_NRTIDALV3 and SEOBNRV5THM.

VI. CONCLUSION

In this work, we introduce the new waveform approximant SEOBNRV5THM for BNS systems. This model represents a significant advancement over the previous model SEOBNRV4T [42,137]. The key improvements include

- (i) new PN information, in particular of the spin-shifted dynamical tides [46], 7.5PN order adiabatic effects in the waveform modes [63,65,66], and 3.5PN spin-induced multipole moments [52,53],
- (ii) the development of a premerger model and calibration to NR waveforms across a wide range of tidal deformabilities with mass ratios up to $q = 2$,
- (iii) the higher modes,
- (iv) drastically improved speed for $M \geq 2M_\odot$ systems by a factor 100 to 1000, making it fast enough to be used in PE for current detectors, and
- (v) the state-of-the-art BBH model SEOBNRV5HM [115] as the BBH baseline.

In summary, the model exhibits enhancements in both speed and accuracy when compared to the prior generation model SEOBNRV4T.

To validate SEOBNRV5THM, we study its dephasing and mismatch relative to existing NR waveforms, its predecessor SEOBNRV4T, as well as the other state-of-the-art BNS waveform models TEOBRESUMS-GIOTTO, IMRPHEMOPV2_NRTIDALV3, and SEOBNRV5_ROM_NRTIDALV3. Our findings reveal that, while all waveform models examined in this

study perform similarly in predicting waveforms from BAM simulations, only SEOBNRV5THM generates waveforms that are of the same order of error as the intrinsic NR uncertainty for both BAM and SACRA simulations. A noteworthy caveat of this result is, however, that although the SACRA simulations were generated at a higher resolutions than the BAM simulations, they, in contrast to the BAM simulations, do not converge at a clear convergence order. Our error estimate for the BAM simulations is therefore more robust than our error estimate for the SACRA simulations, which might be underestimated.

A comparison of the current-generation BNS waveform approximants across the relevant BNS parameter space reveals noticeable differences amongst the waveforms for spin magnitudes $|\chi| \gtrsim 0.3$, high primary tidal deformabilities $\Lambda_2^{(1)} \gtrsim 4000$, and high total masses $M \gtrsim 6M_\odot$. This is to be expected due to the limited number of NR simulations in these parts of the parameter space. Furthermore, while the EOB BNS models appear to agree for mass ratios up to $q = 6$, noticeable differences between SEOBNRV5THM and NRTIDAL-variants appear for large mass ratios $q \gtrsim 2$.

As SEOBNRV5THM is, however, the only current model that includes the shift in the f -mode frequency due to spin, we believe that it should be the most accurate spin-aligned BNS model for small to medium spins $|\chi| \lesssim 0.3$ in comparison to the other available approximants, and that it can be employed for masses $m_{A,B} \in [0.5, 3]M_\odot$, and for dimensionless tidal deformabilities $\Lambda_{A,B} \in [0, 5000]$, although the model is robust enough to generate waveforms also for higher mass ratios and highly spinning BNS systems, if one starts the evolution from a low enough initial frequency.

Suitable future developments for SEOBNRV5THM would be in order of presumed feasibility:

- (i) the inclusion of the LO adiabatic tidal effects for $l \geq 4$, gravitomagnetic tidal effects, and the respective tidal adjustment of the $\bar{D}_{\text{nos}}$ potential as is done in TEOBRESUMS [133,184];
- (ii) the further calibration and validation to new BNS and black-hole-neutron-star simulations [162,185,186];
- (iii) the accurate modeling of the premerger signal of the higher-order modes, as SEOBNRV5THM only calibrates the (2,2) mode during the premerger;
- (iv) the application of the stationary-phase approximation when Fourier transforming [140];
- (v) the continued study of resonance phenomena for highly spinning, antialigned NSs through analytical and numerical advances [159,162], including the tidal spin [42,46,160];
- (vi) the addition of precession and eccentric effects on top of the aligned-spin model; and
- (vii) the development of a full theoretical understanding of dynamical tides [51], in particular for precessing and eccentric BNS systems.

We believe that SEOBNRV5THM is an excellent starting point for such future research.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

The data supporting this study's findings are available within the article.

APPENDIX A: QUASIUNIVERSAL RELATIONS

In this Appendix, we collect all of the universal relations used in the waveform model. We use the spin-induced multipole moments C_l , as a function of Λ_2 , from Refs. [145,149], the octupolar tidal deformability Λ_3 from Ref. [148], and ω_{0l} from Refs. [148,151].

To find a good estimate for $\omega_{f,l}(\Lambda, \chi)$ we compute the nonspinning moment of inertia I [149] and the NS compactness and radius [16]. We then use Ref. [152] to obtain the spin-adjusted moment of inertia $I(\Lambda, \Omega)$ and invert this numerically to get the spin-angular-frequency $\Omega = I/\chi$. Finally, we use Ω in Ref. [150] to arrive at $\Delta\omega_{0l}$.

APPENDIX B: TAPERING OF THE SUBDOMINANT MODES

The tapering of the subdominant $(\ell, |m|) = (3, 3)$, $(2, 1)$, $(4, 4)$, $(5, 5)$, $(3, 2)$, and $(4, 3)$ modes builds heavily on the tapering of SEOBNRV4T [137], but we change it to satisfy a continuously differentiable frequency evolution as described in Sec. II B.

The steps for the tapering of the inspiral-plunge waveform $h_{\ell m}^{\text{premerge}} = A_{\ell m} \exp(-i\phi_{\ell m})$ are

- (i) linearly continue the intermediate amplitude from the inspiral-plunge amplitude

$$\hat{A}_{\ell m}(t) = \begin{cases} A_{\ell m}(t) & \text{for } t \leq t_{\text{match}}, \\ A_{\ell m}(t_{\text{match}}) + \dot{A}_{\ell m}(t_{\text{match}})(t - t_{\text{match}}) & \text{for } t > t_{\text{match}}; \end{cases} \quad (\text{B1})$$

- (ii) extract the final waveform phase $\phi_{\text{match}}^{\ell m} = \phi_{\ell m}(t_{\text{match}})$ and frequency $\omega_{\text{match}}^{\ell m} = \omega_{\ell m}(t_{\text{match}})$ and the respective phase and frequency $12M$ prior to t_{match} : $\phi_{\text{freq}}^{\ell m} = \phi_{\ell m}(t_{\text{match}} - 12M)$ and $\omega_{\text{freq}}^{\ell m} = \omega_{\ell m}(t_{\text{match}} - 12M)$;
- (iii) let the phase transition from $t_{\text{match}} - 12M$ onwards to the asymptotic frequency ω_{match} through a tapering of the form

$$\hat{\omega}_{\ell m}(t) = \begin{cases} \omega_{\ell m}(t) & \text{for } t \leq t_{\text{match}} - 12M, \\ \omega_{\text{match}}^{\ell m} - \Delta\omega_{\ell m} \exp\left(-\frac{t - t_{\text{match}}}{\tau_{\text{freq}}^{\ell m}}\right) & \text{for } t > t_{\text{match}} - 12M, \end{cases} \quad (\text{B2})$$

where $\Delta\omega_{\ell m} = \omega_{\text{match}}^{\ell m} - \omega_{\text{freq}}^{\ell m}$ and $\tau_{\text{freq}}^{\ell m}$ is chosen to ensure continuity of the first frequency derivative across the whole waveform;

- (iv) use the Planck taper from Eq. (33) with decay time tuned to be the same as for the dominant (2,2) mode $\tau_W = 0.5\pi/\omega_{22}(t_{\text{match}})$ centered $15M$ prior to t_{match} ,

$$W(t) = \left\{ 1 + \exp\left[\frac{t - (t_{\text{match}} - 15M)}{\tau_W}\right] \right\}^{-1}; \quad (\text{B3})$$

- (v) the tapered waveform then has the form

$$h_{\ell m}(t) = \hat{A}_{\ell m}(t)W(t) \exp\left(-i \int^t \hat{\omega}_{\ell m}(t') dt'\right), \quad (\text{B4})$$

with

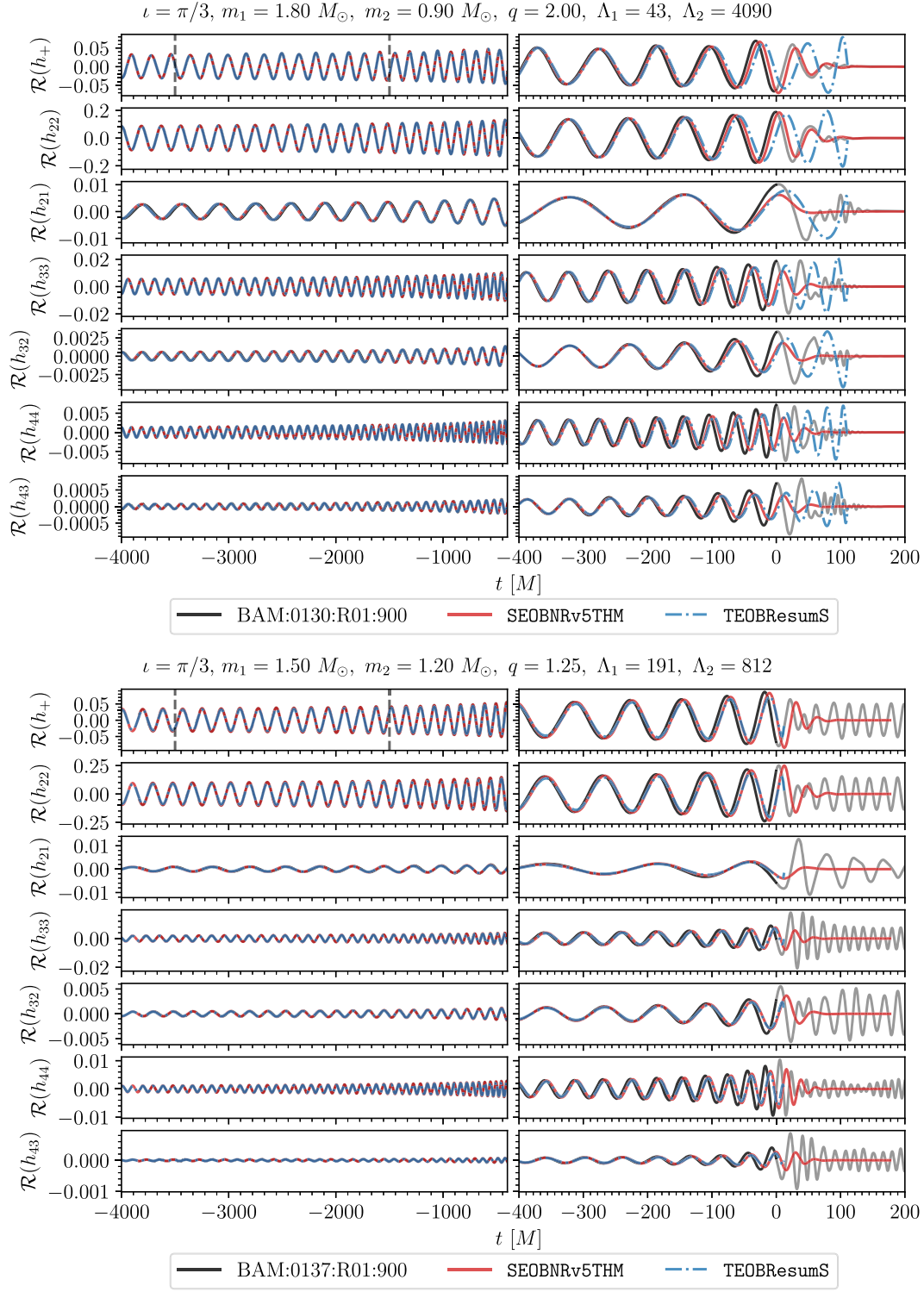


FIG. 12. Comparison between the GW polarization and modes of SEOBNRV5THM and TEOBRESUMS against NR simulations (black), where the postmerger NR waveform is shown in gray. The system parameters are stated above each figure and we depict the real part of the $+$ polarization h_+ in the first row, and the higher modes as aligned through the polarization in the lower rows. The vertical dashed lines mark the time window where the alignment against h_+^{NR} is done. Note that TEOBRESUMS has no tapering prescription in time domain and ends abruptly.

$$\int^t \hat{\omega}_{\ell m}(t') dt' = \phi_{\text{match}}^{\ell m} + \omega_{\text{match}}^{\ell m}(t - t_{\text{freq}}) + \tau_{\text{freq}}^{\ell m} \Delta \omega_{\ell m} \left[\exp\left(-\frac{t - t_{\text{match}}}{\tau_{\text{freq}}^{\ell m}}\right) - 1 \right]. \quad (\text{B5})$$

$$\phi_{\ell m}(t) \mapsto \phi_{\ell m}(t - t_c) e^{im\phi_c}. \quad (\text{C5})$$

APPENDIX C: VALIDATION OF THE HIGHER MODES

In Fig. 12 we show the gravitational polarization $h_+ = \text{Re}(h)$ from Eq. (22) as well as the higher modes

$$(\ell, m) \in \{(2, 2), (2, 1), (3, 3), (3, 2), (4, 4), (4, 3)\} \quad (\text{C1})$$

for two of the four publicly available BNS NR simulations with available higher-modes decomposition [74] at an inclination $\iota = \pi/3$.

To align the modes, we minimize the difference in the frequency-time series of the NR and EOB waveforms,

$$\min_{(t_c, \phi_c)} \int_{t_{\text{start}}}^{t_{\text{end}}} |\omega_+^{(\text{NR})}(t, \iota, \phi_c) - \omega_+^{(\text{EOB})}(t - t_c, \iota, \phi_c)|^2 dt, \quad (\text{C2})$$

where

$$\omega_+(t, \iota, \phi_c) = \frac{d\phi_+(t, \iota, \phi_c)}{dt} \quad (\text{C3})$$

is the angular frequency of the $+$ polarization,

$$h_+ = A_+(t, \iota, \phi_c) e^{i\phi_+(t, \iota, \phi_c)}. \quad (\text{C4})$$

We then transform each GW mode by the obtained azimuthal angle,

We note that we only show comparisons to the highest-resolution simulations. We find that we introduce artifacts into the waveform phase when performing a Richardson interpolation towards merger, and that the alignment against the polarization becomes harder and tends to not align the higher modes well, which we leave to be studied more in depth for the future.

We make the observation that, for near equal-mass systems $q \sim 1$, both TEOBRESUMS and SEOBNRV5THM agree very well with each other, while for higher mass ratios $q \sim 2$ SEOBNRV5THM performs better than TEOBRESUMS due to the more accurate time of merger across all modes, which remains true also for the two simulations not shown in this work. We remind the reader that the time-domain waveform of TEOBRESUMS has no tapering prescription and therefore ends abruptly at the merger.

APPENDIX D: NR SIMULATIONS USED IN THIS WORK

In Tables V and VI we show the NS parameters of all of the 64 NR waveforms from SACRA [70,71,146] and CoRe (using the BAM code) [73–75] that were used in this work partially in the calibration of the merger time and fully in the validation of SEOBNRV5THM in Sec. IV. We list the individual masses of the stars $m_{1,2}$, the mass ratio q , dimensionless tidal deformabilities $\Lambda_2^{(1,2)}$, and radii $R_{1,2}$. The radii data for the SACRA waveforms were adapted from Refs. [70,71,146], while we compute the radii for the CORE (BAM) waveforms using the adiabatic Love number and tidal deformability, i.e., from $\Lambda_2^{(i)} = (2/3)k_2^{(i)} R_i^5/m_i^5$.

TABLE V. Properties of the 64 NR waveforms used for SEOBNRV5THM. For each waveform, we indicate the EOS, the individual masses $m_{1,2}$ [$M_\odot$] of the individual bodies, the mass ratio q , the dimensionless tidal deformabilities $\Lambda_2^{(1,2)}$, the radii $R_{1,2}$ [km], and dimensionless spins $\chi_{1,2}$.

NR waveform name	EOS	m_1	m_2	q	$\Lambda_2^{(1)}$	$\Lambda_2^{(2)}$	R_1	R_2	χ_1	χ_2
CoRe:BAM:0001	2B	1.35	1.35	1.00	127	127	9.72	9.72	0	0
CoRe:BAM:0037	H4	1.37	1.37	1.00	1006	1006	13.54	13.54	0	0
CoRe:BAM:0039	H4	1.37	1.37	1.00	1002	1002	13.54	13.54	0.27	0.27
CoRe:BAM:0039	H4	1.37	1.37	1.00	1002	1002	13.54	13.54	0.27	0.27
CoRe:BAM:0062	MS1b	1.35	1.35	1.00	1532	1532	14.02	14.02	−0.18	−0.18
CoRe:BAM:0064	MS1b	1.35	1.35	1.00	1532	1532	14.01	14.01	0	0
CoRe:BAM:0066	MS1b	1.35	1.35	1.00	1532	1532	14.02	14.02	0.18	0.18
CoRe:BAM:0068	MS1b	1.35	1.35	1.00	1525	1525	14.01	14.01	0.27	0.27
CoRe:BAM:0095	SLy	1.35	1.35	1.00	390	390	11.46	11.46	0	0
CoRe:BAM:0096	SLy	1.35	1.35	1.00	390	390	11.46	11.46	0	0
CoRe:BAM:0097	SLy	1.35	1.35	1.00	390	390	11.46	11.46	0	0
CoRe:BAM:0101	SLy	1.35	1.35	1.00	390	390	11.46	11.46	0.10	0.10
CoRe:BAM:0104	SLy	1.35	1.35	1.00	388	388	11.46	11.46	0.19	0.19

(Table continued)

TABLE V. (*Continued*)

NR waveform name	EOS	m_1	m_2	q	$\Lambda_2^{(1)}$	$\Lambda_2^{(2)}$	R_1	R_2	χ_1	χ_2
CoRe:BAM:0129	SLy	1.35	1.35	1.00	390	390	11.47	11.47	0.17	0.17
CoRe:BAM:0130	SLy	1.80	0.90	2.00	43	4090	11.14	11.36	0	0
CoRe:BAM:0131	SLy	1.72	0.98	1.75	66	2535	11.27	11.39	0	0
CoRe:BAM:0132	SLy	1.35	1.35	1.00	390	390	11.46	11.46	0	0
CoRe:BAM:0136	SLy	1.62	1.08	1.50	108	1489	11.35	11.42	0	0
CoRe:BAM:0137	SLy	1.50	1.20	1.25	191	812	11.42	11.45	0	0
CoRe:BAM:0141	MS1b	1.40	1.37	1.02	1260	1416	14.06	14.03	0.21	0.08
CoRe:BAM:0144	SLy	1.35	1.35	1.00	390	390	11.47	11.47	-0.17	-0.17

TABLE VI. Properties of the 64 NR waveforms used for SEOBNRV5THM. For each waveform, we indicate the EOS, the individual masses $m_{1,2}$ of the individual bodies, the mass ratio q , the dimensionless tidal deformabilities $\Lambda_2^{(1,2)}$, the radii $R_{1,2}$, and dimensionless spins $\chi_{1,2}$.

NR waveform name	EOS	$m_1[M_\odot]$	$m_2[M_\odot]$	q	$\Lambda_2^{(1)}$	$\Lambda_2^{(2)}$	R_1 [km]	R_2 [km]	χ_1	χ_2
SACRA:15H_135_135_00155_182_135	15H	1.35	1.35	1.00	1211	1211	13.69	13.69	0	0
SACRA:125H_135_135_00155_182_135	125H	1.35	1.35	1.00	863	863	12.97	12.97	0	0
SACRA:H_135_135_00155_182_135	H	1.35	1.35	1.00	607	607	12.27	12.27	0	0
SACRA:HB_135_135_00155_182_135	HB	1.35	1.35	1.00	422	422	11.61	11.61	0	0
SACRA:B_135_135_00155_182_135	B	1.35	1.35	1.00	289	289	10.96	10.96	0	0
SACRA:15H_125_146_00155_182_135	15H	1.46	1.25	1.17	760	1871	13.72	13.65	0	0
SACRA:125H_125_146_00155_182_135	125H	1.46	1.25	1.17	535	1351	12.99	12.94	0	0
SACRA:H_125_146_00155_182_135	H	1.46	1.25	1.17	369	966	12.18	12.26	0	0
SACRA:HB_125_146_00155_182_135	HB	1.46	1.25	1.17	252	684	11.59	11.61	0	0
SACRA:15H_121_151_00155_182_135	15H	1.51	1.21	1.25	625	2238	13.73	13.63	0	0
SACRA:125H_121_151_00155_182_135	125H	1.51	1.21	1.25	435	1621	12.98	12.93	0	0
SACRA:H_121_151_00155_182_135	H	1.51	1.21	1.25	298	1163	12.26	12.25	0	0
SACRA:HB_121_151_00155_182_135	HB	1.51	1.21	1.25	200	827	11.57	11.60	0	0
SACRA:B_121_151_00155_182_135	B	1.51	1.21	1.25	131	581	10.89	10.98	0	0
SACRA:15H_118_155_00155_182_135	15H	1.55	1.18	1.31	530	2575	13.74	13.62	0	0
SACRA:125H_118_155_00155_182_135	125H	1.55	1.18	1.31	366	1875	12.98	12.92	0	0
SACRA:H_118_155_00155_182_135	H	1.55	1.18	1.31	249	1354	12.26	12.24	0	0
SACRA:HB_118_155_00155_182_135	HB	1.55	1.18	1.31	165	966	11.55	11.60	0	0
SACRA:B_118_155_00155_182_135	B	1.55	1.18	1.31	107	681	10.87	10.98	0	0
SACRA:15H_117_156_00155_182_135	15H	1.56	1.17	1.33	509	2692	13.74	13.61	0	0
SACRA:125H_117_156_00155_182_135	125H	1.56	1.17	1.33	350	1963	12.98	12.91	0	0
SACRA:H_117_156_00155_182_135	H	1.56	1.17	1.33	238	1415	12.25	12.24	0	0
SACRA:HB_117_156_00155_182_135	HB	1.56	1.17	1.33	157	1013	11.55	11.60	0	0
SACRA:B_117_156_00155_182_135	B	1.56	1.17	1.33	101	719	10.86	10.98	0	0
SACRA:15H_116_158_00155_182_135	15H	1.58	1.16	1.36	465	2863	13.73	13.60	0	0
SACRA:125H_116_158_00155_182_135	125H	1.58	1.16	1.36	319	2085	12.98	12.90	0	0
SACRA:H_116_158_00155_182_135	H	1.58	1.16	1.36	215	1506	12.25	12.23	0	0
SACRA:HB_116_158_00155_182_135	HB	1.58	1.16	1.36	142	1079	11.53	11.59	0	0
SACRA:B_116_158_00155_182_135	B	1.58	1.16	1.36	91	765	10.84	11.98	0	0
SACRA:15H_125_125_0015_182_135	15H	1.25	1.25	1.00	1875	1875	13.65	13.65	0	0
SACRA:125H_125_125_0015_182_135	125H	1.25	1.25	1.00	1352	1352	12.94	12.94	0	0
SACRA:H_125_125_0015_182_135	H	1.25	1.25	1.00	966	966	12.26	12.26	0	0
SACRA:HB_125_125_0015_182_135	HB	1.25	1.25	1.00	683	683	11.61	11.61	0	0
SACRA:B_125_125_0015_182_135	B	1.25	1.25	1.00	476	476	10.98	10.98	0	0
SACRA:15H_112_140_0015_182_135	15H	1.40	1.12	1.25	975	3411	13.71	13.58	0	0
SACRA:125H_112_140_0015_182_135	125H	1.40	1.12	1.25	693	2490	12.98	12.89	0	0
SACRA:H_112_140_0015_182_135	H	1.40	1.12	1.25	484	1812	12.28	12.23	0	0
SACRA:HB_112_140_0015_182_135	HB	1.40	1.12	1.25	333	1304	11.60	11.59	0	0
SACRA:B_112_140_0015_182_135	B	1.40	1.12	1.25	225	933	10.95	10.97	0	0

(Table continued)

TABLE VI. (Continued)

NR waveform name	EOS	$m_1[M_\odot]$	$m_2[M_\odot]$	q	$\Lambda_2^{(1)}$	$\Lambda_2^{(2)}$	R_1 [km]	R_2 [km]	χ_1	χ_2
SACRA:15H_107_146_0015_182_135	15H	1.46	1.07	1.36	760	4361	13.72	13.54	0	0
SACRA:125H_107_146_0015_182_135	125H	1.46	1.07	1.36	535	3196	12.99	12.86	0	0
SACRA:H_107_146_0015_182_135	H	1.46	1.07	1.36	369	2329	12.18	12.22	0	0
SACRA:HB_107_146_0015_182_135	HB	1.46	1.07	1.36	252	1695	11.59	11.60	0	0
SACRA:B_107_146_0015_182_135	B	1.46	1.07	1.36	168	1216	10.92	10.97	0	0
SACRA:SFHo_135_135_00155_182_135	SFHo	1.35	1.35	1.00	460	460	11.91	11.91	0	0

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