Symmetry-Energy Dependence of the Bulk Viscosity of Nuclear Matter

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We clarify how the weak-interaction-driven bulk viscosity ζ and the bulk relaxation time τ_Π of neutrino-transparent npe matter depend on the nuclear symmetry energy. We show that, at saturation density, the equation-of-state dependence of these transport quantities is fully determined by the experimentally constrained nuclear symmetry energy S and its slope L. Variations of L can change the bulk viscosity by orders of magnitude, which can affect both the dissipative and the conservative tidal response of neutron stars. This suggests that both conservative and dissipative effects encoded in the gravitational-wave signatures of binary neutron star inspirals may help constrain nuclear symmetry-energy properties.

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Introduction—Because of electric charge neutrality, neutron star matter exhibits isospin asymmetries much larger than those typically found in nuclei, with proton fractions around $\lesssim 10\%$ [1]. Thus, the symmetry energy E_{sym}, which in nuclear physics quantifies the energy cost of an isospin asymmetry, is crucial for the structure of neutron stars [2]. As a function of the baryon density n_B, the symmetry energy is often parametrized by its value S and its slope L at nuclear saturation density n_{sat}, both of which play an important role in the neutron star equation of state [3]. For an in-depth, comprehensive review on nuclear symmetry energy, see Ref. [4].

Constraints on the nuclear symmetry energy have been derived from both experimental and theoretical efforts. Experimental constraints are obtained, for example, through nuclear mass fitting [5–8] and neutron skin thickness measurements [9–14], while theoretical constraints can be obtained through chiral effective field theory (EFT) [15,16]. While S remains relatively well constrained, there is ongoing discussion on the value of L. Of particular interest are the recent parity-violating electron scattering neutron skin experiments of ^{208}Pb (PREX-I and PREX-II) [17] and ^{48}Ca (CREX) [4]. While most estimates lie around $L \approx 50$ MeV [18–20], PREX-I and II and CREX measurements suggest $L = 121 \pm 47$ and $L = -5 \pm 40$ MeV, respectively, to 68% confidence [4]. There have been efforts in accounting for the results from PREX [21,22] and reconciling both PREX and CREX results, see Refs. [4,23,24]. In particular, Ref. [4] shows that nuclear interactions optimally satisfying both measurements imply $L = 53 \pm 13$ MeV. Studies on the importance of L in neutron stars, showing, for instance, that it has a large impact on the stellar radius and tidal deformability, can be found in Refs. [2,25–28]. Furthermore, in the postmerger phase of binary neutron star collisions, Ref. [29] found that the amount of dynamically ejected mass increases with L, while gravitational-wave emission is mostly insensitive to variations of this quantity.

In a previous work [30], we proposed that independent information on the symmetry energy may be obtained by investigating bulk-viscous transport coefficients of neutron star matter in binary neutron star coalescence. There, bulk viscosity is expected to emerge as a consequence of weak-interaction-driven flavor equilibration dynamics [31]. Reference [30] showed that, at least for equations of state (EOS) derived using relativistic mean-field (RMF) models, the bulk viscosity transport coefficient ζ of matter composed of protons p, neutrons n, and electrons e (npe matter) can be sensitive to changes in S and L.

In this Letter, we work out in detail how ζ and the bulk relaxation time τ_Π of npe matter (in the low temperature regime) depend on the symmetry energy and its experimentally measured parameters. Using a range of values for S and L, we show that slight changes in L can change the zero frequency limit of ζ by orders of magnitude. We also argue that the out-of-β-equilibrium correction to the pressure of npe matter at saturation density present in the early inspiral phase of binary neutron star mergers dynamically responds to volume deformations in an approximate bulk elastic manner [32], with a bulk modulus determined by the

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values of S and L. We discuss the potential consequences of our findings for the detection of both conservative and dissipative effects encoded in the gravitational waves emitted by inspiraling neutron stars.

Bulk viscosity from chemical imbalance—Bulk-viscous effects can be relevant to the damping of r modes in neutron stars [33–39] and possibly to the gravitational-wave emission in binary neutron star mergers [40–48]. Furthermore, dissipative effects [49–51] contribute to the so-called dissipative tidal deformability $\tilde{\Xi}$ [49], which modifies the phase of gravitational waves emitted during the inspiral of a neutron star binary. An upper bound of $\tilde{\Xi} \lesssim 1200$ has been placed using data from the gravitational-wave event GW170817 in Refs. [50,51]. Finally, we note that tidal heating has been proposed as a way to probe strangeness degrees of freedom in neutron stars [52,53].

In this Letter we consider neutrino-transparent npe matter in a macroscopic quasiequilibrium state described by $\{s, n_B, Y, u^\mu\}$, where s is the entropy density, $Y = n_e/n_B$ is the electron fraction (charge neutrality implies that the electron density n_e equals the proton density n_p), and u^μ is the local 4-velocity (normalized such $u_\mu u^\mu = 1$). At sufficiently low temperatures, the β-equilibrium condition is $\mu_n = \mu_p + \mu_e$, where μ_n, μ_p, and μ_e are the chemical potentials for neutrons, protons, and electrons, respectively. Therefore, in this regime, deviations from β equilibrium can be quantified by the nonzero value of the reaction affinity $\delta\mu \equiv \mu_n - \mu_p - \mu_e$ (at high temperatures, β equilibrium does not imply $\delta\mu = 0$, see Refs. [54–56]).

The dynamical equations that describe this system are [57,58]

$$\nabla_\mu (s u^\mu) = \frac{Q + \delta\mu \Gamma_e}{T}, \tag{1a}$$

$$(\varepsilon + \mathcal{P}) u^\mu \nabla_\mu u^\nu - \Delta^{\nu\alpha} \nabla_\alpha \mathcal{P} = 0, \tag{1b}$$

$$\nabla_\mu (n_B u^\mu) = 0, \tag{1c}$$

$$u^\mu \nabla_\mu Y = \frac{\Gamma_e}{n_B}, \tag{1d}$$

where Q stems from the (assumed isotropic) radiative energy loss due to neutrinos and Γ_e describes flavor equilibration rates associated with direct and modified Urca processes (see Supplemental Material for details [59]). Above, T is the temperature, $\varepsilon = \varepsilon(s, n_B, Y)$ is the system's energy density, $\mathcal{P} = \mathcal{P}(s, n_B, Y)$ is the pressure, $\Delta_{\mu\nu} = g_{\mu\nu} - u_\mu u_\nu$, and $g_{\mu\nu}$ is the (mostly minus) spacetime metric.

For small deviations around β equilibrium, one may approximate $\Gamma_e \sim \lambda \delta\mu$, where $\lambda \sim T^4$ for direct Urca and $\lambda \sim T^6$ for modified Urca processes (assuming, for simplicity, the so-called Fermi surface approximation, see Refs. [54,55,60] and Supplemental Material [59]). The out-of-β-equilibrium correction to the system's pressure, described by the bulk scalar $\Pi \equiv \mathcal{P} - \mathcal{P}|_{\delta\mu=0}$ [57], obeys in this limit a dynamical equation [57,61] *à la* Israel and Stewart [63],

$$u^\mu \nabla_\mu \Pi = -\frac{\Pi}{\tau_\Pi} - \frac{\zeta_0}{\tau_\Pi} \nabla_\mu u^\mu, \tag{2}$$

where the transport coefficients are given by

$$\tau_\Pi = -\frac{n_B}{\lambda} \left(\frac{\partial \delta\mu}{\partial Y}\bigg|_{n_B} \right)^{-1}, \tag{3a}$$

$$\zeta_0 = -\frac{P_1 n_B^2}{\lambda} \left(\frac{\partial \delta\mu}{\partial Y}\bigg|_{n_B} \right)^{-1} \frac{\partial \delta\mu}{\partial n_B}\bigg|_Y, \tag{3b}$$

and $P_1 = (\partial \mathcal{P}/\partial \delta\mu)_{n_B, \delta\mu=0}$. Above, we neglected T-dependent effects on the EOS. The contribution from Q to (2) is also negligible, see Supplemental Material [59]. Finally, we note that while here we focus on the small $\delta\mu$ case, the equivalence between bulk-viscous theories and chemical equilibration dynamics holds in the full nonlinear regime, see Ref. [64].

Equation (2) shows that small deviations in pressure away from β equilibrium in npe matter are governed by the equations of a relativistic viscoelastic material in which Π displays a delay in response to time-dependent deformations [64]. Assuming, for simplicity, small amplitude baryon density oscillations $\sim e^{-i\omega t}$ around a spatially uniform β-equilibrated state, linear response gives

$$\delta\Pi = -\frac{\zeta_0}{1 - i\tau_\Pi \omega} \delta(\nabla_\mu u^\mu) = \frac{i}{\omega} G_R(\omega) \delta(\nabla_\mu u^\mu), \tag{4}$$

where we defined the retarded Green's function G_R that relates the stress and the expansion rate (see Supplemental Material for details [59]). The ω-dependent bulk viscosity can be extracted from the imaginary part of the Green's function,

$$\zeta(\omega) = \frac{1}{\omega} \operatorname{Im} G_R = \frac{\zeta_0}{1 + \tau_\Pi^2 \omega^2}, \tag{5}$$

defined in [31] (for a review, see Ref. [65]). Different phenomena [66] are encoded in (2). When $\tau_\Pi \omega \ll 1$, one may find $\Pi \sim -\zeta_0 \nabla_\mu u^\mu$ and the system is in the Navier-Stokes viscous fluid regime [57]. When $\tau_\Pi \omega \sim 1$, the system is in the resonant regime where Π responds on the same timescale associated with macroscopic motion, leading to maximal dissipation. This regime may be realized in neutron star mergers, see Refs. [42–45,54,55]. On the other hand, when $\tau_\Pi \omega \gg 1$, the system is in the "frozen" regime [57,67,68] where Y stays fixed, macroscopic motion becomes approximately reversible, and there is basically no dissipation [57].

Because of the strong temperature dependence of the Urca rates, τ_Π and ζ_0 are strongly enhanced at low temperatures, and neutrino-transparent npe matter is expected to be in the frozen Y regime [57,67,68]. In this case, because the equilibrium charge fraction $Y_{\rm eq}$ depends on density while Y remains constant, the system still drops out of β equilibrium as matter compresses or expands. As in equilibrium, the full pressure becomes a function of the baryon density, but this function is given by the EOS at fixed proton fraction, which is not in β equilibrium. This is relevant for early-to-late binary neutron star inspirals, where $T\sim 10^5$ K, tidal forces deform the stars with frequencies $\omega\sim$ kHz [51], and $\delta\mu\neq 0$.

We here note that, formally, in this frozen regime $\zeta_0, \tau_\Pi \to \infty$, $\zeta(\omega)\to \zeta_0/(\tau_\Pi\omega)^2$, and the out-of-$\beta$-equilibrium pressure deviation Π is approximately described by the following dynamical equation:

$$u^\mu \nabla_\mu \Pi = -\frac{\zeta_0}{\tau_\Pi}\nabla_\mu u^\mu, \tag{6}$$

which corresponds to a relativistic generalization of Hooke's law of elasticity [32,64], even though, of course, there is no underlying lattice in this fluid. This approximate bulk elastic response is characterized by a bulk modulus coefficient given by ζ_0/τ_Π, which does not depend on the rates, see Eq. (3), and is fully determined by the cold matter EOS. We also note that, in this limit, the negative bulk modulus is exactly the real part of G_R defined in Eq. (4) and represents local reactive conservative processes.

Symmetry-energy dependence of transport coefficients—The energy per baryon $E(n_B,\delta)$ may be expanded in powers of the isospin asymmetry $\delta\equiv 1\text{–}2Y$ as follows:

$$E(n_B,\delta)\approx E(n_B,\delta=0)+E_{\rm sym,2}(n_B)\delta^2 + E_{\rm sym,4}(n_B)\delta^4+\mathcal{O}(\delta^6), \tag{7}$$

where $E_{\rm sym,2}(n_B)$, also denoted as $E_{\rm sym}(n_B)$, is usually referred to as the nuclear symmetry energy [18,69–72]. Neglecting $\mathcal{O}(\delta^4)$ terms, (7) is reduced to the empirical parabolic approximation [73,74], and $E_{\rm sym}(n_B)$ can be approximated by the difference between symmetric nuclear matter and pure neutron matter,

$$E_{\rm sym}(n_B)\approx E(n_B,\delta=1)-E(n_B,\delta=0). \tag{8}$$

This approximation has been shown to work reasonably well in the context of microscopic asymmetric matter calculations of chiral nucleon-nucleon and three-nucleon interactions at $n_B\lesssim n_{\rm sat}$ [75–79] and in phenomenological models, even at higher densities, as shown for a chiral mean-field model in a recent work by two of the authors and collaborators [80]. On the other hand, the expansion of $E_{\rm sym}(n_B)$ around $n_{\rm sat}$ gives

$$E_{\rm sym}(n_B)=S+L\left(\frac{n_B-n_{\rm sat}}{3n_{\rm sat}}\right) + \frac{K_{\rm sym}}{2}\left(\frac{n_B-n_{\rm sat}}{3n_{\rm sat}}\right)^2+\ldots, \tag{9}$$

which defines S, the symmetry slope L, the curvature $K_{\rm sym}$, and other higher-order coefficients computed at saturation density.

We note here that the bulk viscosity can be connected to the symmetry-energy expansion by noting that $\delta\mu$ is the derivative of the energy per baryon $E+E_\ell$ with respect to $n_I/n_B\equiv\delta/2$, where E_ℓ is the contribution from the leptons (treated here as a free electron gas). This gives

$$\delta\mu = 4E_{\rm sym}(n_B)(1-2Y)-\left(\frac{\partial E_\ell}{\partial Y}\right)_{n_B}. \tag{10}$$

Equation (10) shows that $\delta\mu$ follows from the symmetry energy and the energy of the electrons. As a consequence, the EOS part of the transport coefficients also depends solely on the symmetry energy and electron contributions to the energy per baryon. In fact, modulo the contribution from the rates and leptons, one finds schematically

$$\tau_\Pi\sim\frac{1}{E_{\rm sym}(n_B)},\quad \zeta_0\sim\left(\frac{\partial E_{\rm sym}(n_B)}{\partial n_B}\Big/E_{\rm sym}(n_B)\right)^2. \tag{11}$$

Thus, as long as the parabolic approximation in the asymmetry parameter δ holds, the transport coefficients can be reliably calculated using experimentally observed properties of the symmetry energy [81]. We remark that Eq. (11) was derived without making an expansion in the baryon density. Nevertheless, we note that the quantities are only fully determined by S and L when they are evaluated at $n_{\rm sat}$. This shows that models for cold npe matter can only give different values for ζ_0 and τ_Π if they do not have the same basic nuclear symmetry properties (assuming they use similar rates).

To assess the validity of the parabolic approximation in the calculation of bulk-viscous transport coefficients beyond RMF modeling, we calculate ζ_0 and τ_Π using as input chiral EFT calculations from Ref. [16], which provided uncertainty bands for three different nuclear Hamiltonians. Following [83], we cover the entire band of energy per nucleon E by linearly interpolating between the upper and lower bounds of chiral EFT, $E_{\rm up}$ and $E_{\rm low}$,

$$E_\sigma=(1-\sigma)E_{\rm up}+\sigma E_{\rm low}. \tag{12}$$

To perform the interpolation, we use the following parametrization [84–86]:

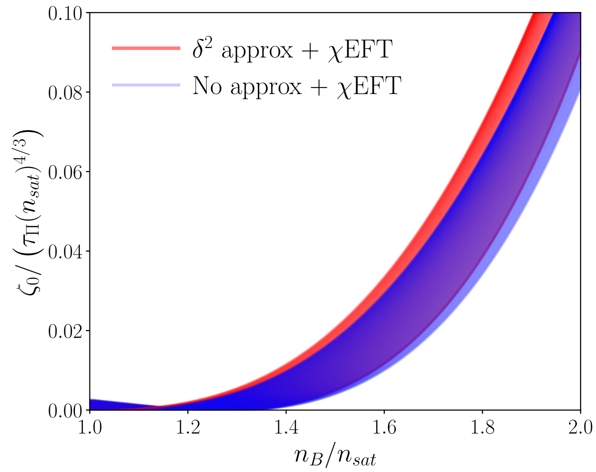

FIG. 1. The dimensionless bulk modulus as a function of the baryon density. The blue curves are calculated directly from the chiral EFT parametrization [16,83–86]. The red curves are calculated using the parabolic symmetry-energy approximation of the chiral EFT parametrization. The purple region denotes the overlap between blue and red curves (darker colors reflect a higher density of curves).

$$
\begin{aligned}
E(n_B, Y) = m_B &+ T_0\left[\frac{3}{5}\left(Y^{3/5} + (1-Y)^{3/5}\right)\left(\frac{2n_B}{n_{\rm sat}}\right)^{2/3}\right.\\
&- [(2\alpha - 4\alpha_L)Y(1-Y) + \alpha_L]\frac{n_B}{n_{\rm sat}}\\
&\left.+ [(2\eta - 4\eta_L)Y(1-Y) + \alpha_L]\left(\frac{2n_B}{n_{\rm sat}}\right)^{\gamma}\right], \quad (13)
\end{aligned}
$$

where α, α_L, η, η_L, and γ are fitting parameters [87]. After fitting (13), we calculate S, L, τ_Π, and ζ_0 for each obtained EOS (assuming $n_{\rm sat} = 0.16\ {\rm fm}^{-3}$). In Fig. 1, we plot the dimensionless bulk modulus $\zeta_0/(\tau_\Pi n_{\rm sat}^{4/3})$ as a function of n_B for all chiral EFT parametrizations. In Fig. 2, we plot ζ_0 at $n_{\rm sat}$ as a function of S or L, which also shows the correlation between S and L for our chiral EFT parametrizations. The blue curves are calculated directly from the chiral EFT parametrization, while the red curves are calculated using the respective parabolic symmetry-energy approximations of the chiral EFT parametrizations.

One can see that the parabolic approximation in the asymmetry parameter δ is excellent, especially at low densities and relatively higher L. This indicates that (11) very accurately captures the dependence of τ_Π and ζ_0 with $E_{\rm sym}$ in neutrino-transparent npe matter. Also, even with the tight constraints on L imposed from chiral EFT [16], ζ_0 remains strongly sensitive and changes by orders of magnitude when L changes.

Transport coefficients at saturation—The EOS-dependent contributions to the transport coefficients significantly simplify at saturation density, being solely determined by the experimentally measured quantities S and L. In fact, direct calculation [89] gives

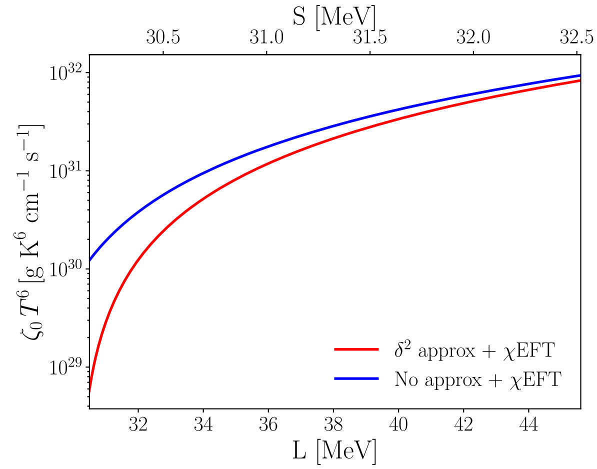

FIG. 2. Bulk viscosity at $n_{\rm sat}$ as a function of S or L. The blue curve is calculated directly from a chiral EFT parametrization [16,83–86]. The red curve is calculated using the parabolic symmetry-energy approximation of the chiral EFT parametrization.

$$
\tau_\Pi(n_{\rm sat}) = \frac{n_{\rm sat}}{\lambda}\left[8S + \frac{\partial^2 E_\ell}{\partial Y^2}\right]^{-1}, \quad (14a)
$$

$$
\zeta_0(n_{\rm sat}) = \frac{n_{\rm sat}^4}{\lambda}\left[8S + \frac{\partial^2 E_\ell}{\partial Y^2}\right]^{-2} \times \left[\frac{4L}{3n_{\rm sat}}(2Y-1) + \frac{\partial^2 E_\ell}{\partial n_B \partial Y} - \frac{1}{n_{\rm sat}}\frac{\partial E_\ell}{\partial Y}\right]^2. \quad (14b)
$$

.

In Fig. 3, we take $S \in [30, 40]$ and $L \in [30, 150]$ MeV, and we plot ζ_0 vs τ_Π using all combinations of S and L, without taking into account their correlations. As the figure shows, $\zeta_0(n_{\rm sat})$ is strongly sensitive to L.

Using (14a) and (14b), one can see that the nuclear physics contribution to the bulk modulus ζ_0/τ_Π is fully determined by $n_{\rm sat}$, S, and L (above $n_{\rm sat}$, the bulk modulus depends on higher-order coefficients such as $K_{\rm sym}$). We find $\zeta_0/(\tau_\Pi n_{\rm sat}^{4/3}) = 0.004$ for $L \sim 50$ MeV and $\zeta_0/(\tau_\Pi n_{\rm sat}^{4/3}) = 0.056$ for $L \sim 100$ MeV. Tighter constraints on L are needed to nail down the value of the bulk modulus of npe matter.

Final remarks—In this Letter, we have shown that the bulk-viscous transport properties ζ_0 and τ_Π of npe matter at saturation density are directly constrained by knowledge on the coefficients S and L of the symmetry-energy expansion. In particular, ζ_0 can vary by orders of magnitude by varying L. The strong dependence of ζ_0 with L is particularly relevant given the upper bound on the dissipative tidal deformability found in Ref. [49], which was translated into an upper bound on a density-averaged bulk viscosity $\langle\zeta\rangle_A \lesssim 10^{31}\ {\rm g\ cm^{-1}\ s^{-1}}$ in the inspiraling neutron stars [49]. One can

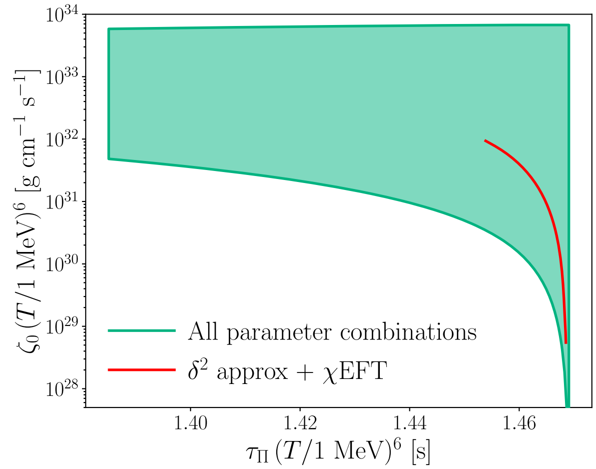

FIG. 3. Bulk viscosity vs relaxation time at n_{sat}. The green box is calculated using $S \in [30, 40]$ and $L \in [30, 150]$ MeV. The red curve is calculated using the chiral EFT parametrizations [16,83–86].

see that ζ_0 at saturation can become much larger than this bound on $\langle\zeta\rangle_A$ if L increases. In addition, as illustrated by Eq. (6), the Israel-Stewart equation also captures conservative dynamics and should, thus, contribute to the conservative dynamical Love number $k_l^{(2)}$ [90]. As Fig. 3 shows, the bulk modulus ζ_0/τ_Π can increase by orders of magnitude as L increases and is not suppressed by temperature. Therefore, we also expect a potentially detectable effect induced by bulk viscosity through the conservative dynamical tides, though the precise relation requires a new dedicated analysis that is left for future work. These findings suggest the interesting possibility that one may use both the conservative and dissipative tidal response to place an upper bound on L. While our results are limited to nuclear saturation and slightly larger densities, they should be informative of the averaged bulk viscosity computed taking into account the different layers of the star.

For the dissipative processes, it is unclear how the upper bound on the bulk viscosity from [49] is affected by the very large relaxation times found at low temperatures, which induce approximate elastic response of the bulk scalar Π. Furthermore, it is important to remark that our results are only valid for the transport coefficients of npe matter at low temperatures and at saturation density. The average $\langle\zeta\rangle_A$ should also receive contributions from the stellar crust, below n_{sat}, and from the inner core, where n_B can reach a few times n_{sat} and exotic degrees of freedom, such as hyperons and quarks, may be found and contribute to bulk viscosity. Using results from Ref. [91], we estimate that densities below saturation at the stellar crust should contribute less than 1% to $\langle\zeta\rangle_A$. Nevertheless, the precise contributions to $\langle\zeta\rangle_A$ from the inner core remain to be determined. Further studies, combining general-relativistic calculations with a more complete description of neutron star matter from crust to core, are needed to determine how gravitational waves emitted during the early-to-late inspiral may encode information on the nuclear symmetry energy.

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Data availability—The data that support the findings of this article are openly available [88].

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