The Possibility of the Secondary Object in GW190814 as a Neutron Star

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Received 2020 August 10; revised 2020 September 19; accepted 2020 September 21; published 2020 November 19

\textbf{Abstract}

A compact object was observed with a mass of 2.50–2.67 \(M_\odot\) by LIGO Scientific and Virgo collaborations (LVC) in GW190814, which provides a great challenge to investigations of supranuclear matter. To study this object, the properties of the neutron star are systematically calculated within the latest density-dependent relativistic mean-field (DDRMF) parameterizations, which are determined by the ground-state properties of spherical nuclei. The maximum masses of the neutron star calculated by DD-MEX and DD-LZ1 sets can be around 2.55 \(M_\odot\) with quite stiff equations of state generated by their strong repulsive contributions from vector potentials at high densities. Their maximum speeds of sound \(c_s/c\) are smaller than \(\sqrt{0.8}\) at the center of the neutron star, and the dimensionless tidal deformabilities at 1.4 \(M_\odot\) are less than 800. Furthermore, the radii of 1.4 \(M_\odot\) also satisfy the constraint from the observation of simultaneous mass–radius measurements (Neutron star Interior Composition Explorer). Therefore, we conclude that one cannot exclude the possibility of the secondary object in GW190814 as a neutron star composed of hadron matter from DDRMF models.

\textit{Unified Astronomy Thesaurus concepts:} Neutron stars (1108); Compact objects (288); Nuclear astrophysics (1129); Gravitational wave astronomy (675)

\textbf{1. Introduction}

The rapid progress of astronomical-observable techniques provides not only great challenges but also many opportunities for investigations of neutron stars. In the past decade, measurements of massive neutron stars have successively broken through our recognition of their maximum masses, from PSR J1614-2230 (1.928 \(\pm\) 0.017 \(M_\odot\)) (Demorest \textit{et al.} 2010; Fonseca \textit{et al.} 2016), to PSR J0348+0432 (2.01 \(\pm\) 0.04 \(M_\odot\)) (Antoniadis \textit{et al.} 2013), to PSR J0740+6620 (2.14\textsuperscript{+0.10}_{-0.09} \(M_\odot\)) (Cromartie \textit{et al.} 2020). For the first time, the mass and radius of PSR J0030+0451 were simultaneously measured by the Neutron star Interior Composition Explorer (NICER) collaboration, who drew the first-ever map of the neutron star (Raaijmakers \textit{et al.} 2019). PSR J0030+0451 was reported by two independent analysis groups to have a mass of 1.44\textsuperscript{+0.15}_{-0.14} \(M_\odot\) with a radius of 13.02\textsuperscript{+1.24}_{-1.06} km (Miller \textit{et al.} 2019) and a mass of 1.34\textsuperscript{+0.16}_{-0.16} \(M_\odot\) with a radius of 12.71\textsuperscript{+1.14}_{-1.10} km (Riley \textit{et al.} 2019).

At the same time, the multi-messenger astronomy era has begun with the successful operation of gravitational wave detectors, LIGO Scientific and Virgo Collaborations (LVC), which were the first to receive the gravitational waves generated by a binary neutron star (BNS) merger GW170817 event (Abbott \textit{et al.} 2017a, 2017b, 2018). The tidal deformability of the neutron star was estimated from the signals, which becomes a new constraint on the equation of state of neutron star matter. The total mass of this BNS system in GW170817 is around 2.7 \(M_\odot\) and the mass of the heavier component is around 1.16–1.60 \(M_\odot\) with lower-spin priors, while the maximum mass of the neutron star can approach 1.89 \(M_\odot\) with high-spin priors (Abbott \textit{et al.} 2019). After that, the second possible BNS merger was observed in 2019 April as GW190425, with the total mass 3.4\textsuperscript{+0.3}_{-0.2} \(M_\odot\). The mass range of components is from 1.12 to 2.52 \(M_\odot\) with high-spin priors (Abbott \textit{et al.} 2020a). Several months later, a new event of a compact binary merger with a 22.2–24.3 \(M_\odot\) black hole and a compact component with a mass of 2.50–2.67 \(M_\odot\) was reported by LVC as GW190814. The secondary object of GW190814 has attracted a lot of attention, since it may be either the heaviest neutron star or the lightest black hole ever discovered (Abbott \textit{et al.} 2020b).

Since then, many interesting works have been proposed to explain the secondary object of GW190814. Most \textit{et al.} (2020) have argued that it could be a rapidly rotating neutron star. Tan \textit{et al.} (2020) have considered that it may be a heavy neutron star, including the deconfined QCD matter in the core (see also Dexheimer \textit{et al.} 2020). The possibility of a super-fast pulsar was assumed by Zhang \& Li (2020). There were also many other models that supported the neutron star with 2.50–2.60 \(M_\odot\) under the constraints on the properties of the 1.4 \(M_\odot\) neutron star (Lim \textit{et al.} 2020; Roupas 2020; Toskaros \textit{et al.} 2020). On the other hand, it was also concluded that GW190814 may be a binary black hole merger by Fattoyev \textit{et al.} (2020) and Tews \textit{et al.} (2020).

The mass, radius, and tidal deformabilities of neutron stars are mainly determined by the equation of state (EOS) of neutron star matter, i.e., the relation between energy density and pressure. Many attempts have been made to obtain the EOS of supranuclear matter in neutron stars from different models. It can be assumed as a simple polynomial in terms of pressure and energy density (Annala \textit{et al.} 2018). It also can be generated by the nuclear density functional theories (DFT; Vautherin \& Brink 1972; Shen \textit{et al.} 1998; Douchin \& Haensel 2001; Shen 2002; Long \textit{et al.} 2006, 2007; Sun \textit{et al.} 2008; Dutra \textit{et al.} 2012; Bao \textit{et al.} 2014a; Bao \& Shen 2014b), where the nucleon–nucleon (NN) interaction was effectively determined by fitting the ground-state properties of finite nuclei or the empirical saturation properties of infinite nuclear matter. Moreover, ab initio methods with realistic nuclear potentials extracted from NN scattering are available to study the neutron star (Akmal \textit{et al.} 1998; Li \textit{et al.} 2006; Carlson \textit{et al.} 2015;
are the antisymmetry tensor represents the wave function of the nucleon denote the and

were not carefully discussed due to the de

extrapolated to the supranuclear matter (\(\sim 5\rho_B\)), most of them can reasonably describe the properties of massive neutron stars around \(2.0\ M_\odot\). There are very few EOSs from the covariant density functional theory (CDFT) that can generate the mass of neutron stars above \(2.5\ M_\odot\), such as the NL3 parameter set (Lalazissis et al. 1997). However, the radius of \(1.4\ M_\odot\) from NL3 is too large to satisfy the recent constraints from the observations of GW170817 and NICER. Therefore, a new parameter set, BigApple, was proposed to generate a \(2.6\ M_\odot\) neutron star and reproduce the observables of finite nuclei and NICER (Fattoyev et al. 2020).

The CDFT has achieved great successes in the fields of nuclear physics and astrophysics. The first available version of CDFT was proposed by Walecka with the Hartree approximation, i.e., the \(\sigma – \omega\) model (Walecka 1974), which is also called the relativistic mean-field (RMF) model. Then, the \(\rho\) meson, the nonlinear terms of \(\sigma\) and \(\omega\) mesons, and the coupling terms with the \(\rho\) meson to the \(\sigma\) or \(\omega\) meson were introduced step by step in the RMF model (Boguta & Bodmer 1977; Serot 1979; Sugahara & Toki 1994; Horowitz & Piekarewicz 2001). These nonlinear RMF models can precisely describe the ground-state properties of most nuclei in the nuclide chart. Meanwhile, the contributions of the exchange terms in the mean-field approximation were considered at the end of the 1970s in the framework of the relativistic Hartree–Fock (RHF) model, where the pion effect can be taken into account (Brockmann & Bodmer 1977; Serot 1979; Sugahara & Toki 1994; Long et al. 2006, 2007). The picture of meson exchange can be simplified as a point contact interaction when the meson masses are assumed to have an infinite value, which avoids solving the equation of motion for mesons (Nikolaus et al. 1992). This point coupling RMF model is also widely applied to study the nuclear mass table (Zhao et al. 2010). Furthermore, the nonlinear terms of various mesons could be replaced by the density-dependent meson-nucleon coupling constants in the density-dependent RMF (DDRMF) and DDRHF models, which consider that the nuclear medium effect originated from the relativistic Brueckner–Hartree–Fock model (Brockmann & Toki 1992).

Ten years ago, it was showed by Sun et al. (Sun et al. 2008) that some parameterizations of DDRMF and DDRHF models generated massive neutron stars around \(2.33–2.48\ M_\odot\), such as PKDD (Long et al. 2004), DD-ME1 (Nikšić et al. 2002), DD-ME2 (Lalazissis et al. 2005), PKO1, PKO2, and PKO3 (Long et al. 2006) sets, whereas properties of the neutron star at \(1.4\ M_\odot\) were not carefully discussed due to the deficiencies of astronomical observables. In 2020, several DDRMF parameters, DD-MEX (Taninah et al. 2020), DD-LZ1 (Wei et al. 2020), and DDV, DDTV, DDVTD (Typel & Terrero 2020) were proposed by different groups by fitting ground-state properties of spherical finite nuclei, which considered the parametric correlations, shell evaluations, and tensor couplings of the vector mesons to nucleons, respectively. Therefore, it is necessary to systematically calculate the properties of the neutron star with these latest DDRMF parameterizations and discuss the possibility of the secondary object of GW190817 being a neutron star.

This paper is arranged as follows. The theoretical descriptions of the DDRMF model and neutron star matter are shown in Section 2. In Section 3, properties of nuclear matter and neutron star are presented and discussed with various DDRMF models. A summary and discussion are given in Section 4.

2. The Density-dependent Relativistic Mean-field Model in the Neutron Star

In the DDRMF model, the nucleons usually interact with each other in the nuclear system by exchanging scalar-isoscalar (\(\sigma\)), vector-isoscalar (\(\omega\)), and vector–isoscalar (\(\rho\)) mesons. In some models, the scalar-isoscalar (\(\delta\)) meson is also taken into account to consider the isovector effect on the scalar potential of the nucleon. The DDRMF Lagrangian density can be written as:

\[
\mathcal{L}_{DD} = \sum_{i=p, n} \bar{\psi}_i \left[ -i\gamma^\mu \left( \partial_\mu - \Gamma_\mu\rho_B - \frac{\Gamma_\mu\rho_B}{2}\rho_\mu\right) \psi_i \right] - (M - \Gamma_\sigma\rho_B^2)\sigma_i - \Gamma_\delta\rho_B\delta_i + \frac{\Gamma_\delta\rho_B}{2}\delta_i \rho_\mu \delta_j + \frac{1}{2}(\partial^\mu\sigma_i\partial_\mu - m_\sigma^2\sigma_i) + \frac{1}{2}(\partial^\mu\delta_i\partial_\mu - m_\delta^2\delta_i) + \frac{W_{\mu\nu}W_{\mu\nu} + \frac{1}{2}m_\omega^2\omega_i\omega_j + \frac{1}{4}R_{\mu\nu\rho\sigma}R_{\mu\nu\rho\sigma}}{2}\rho_\mu, \tag{1}
\]

where \(\psi_i\) represents the wave function of the nucleon (proton or neutron). \(\sigma, \omega_i, \rho_i, \) and \(\delta\) denote the \(\sigma, \omega, \rho, \) and \(\delta\) mesons, respectively. \(W_{\mu\nu}\) and \(R_{\mu\nu\rho\sigma}\) are the antisymmetry tensor fields of \(\omega\) and \(\rho\) mesons. In nuclear matter, the tensor coupling between the vector meson and nucleon does not provide any contributions. Therefore, it is neglected in the present Lagrangian. The coupling constants between mesons and nucleons are density-dependent in the DDRMF model, which was first proposed by Brockmann & Toki (1992). It takes into account that the NN interaction in dense matter is affected by the nuclear medium. The density-dependent behaviors of the coupling constants have many styles. In CDFT, there are two types of densities, i.e., the scalar density (\(\rho_\sigma\)) and the vector density (\(\rho_\omega\)). In principle, the coupling constants in DDRMF can be dependent on scalar density or vector density. In this work, we focus on the parameterizations of DDRMF depending on the vector density, which only influences the self-energy instead of total energy. Coupling constants of \(\sigma\) and \(\omega\) mesons are usually expressed as a fraction function of the vector density. In DD2 (Nikšić et al. 2002), DD-ME1, DD-ME2, DDME-X, DDV, DDVT, and DDVTD parameterizations, they are given as:

\[
\Gamma_i(\rho_B) = \Gamma_i(\rho_{B0})f_i(x), \quad \text{with} \quad f_i(x) = \frac{a_i + b_i(x + d_i)^2}{1 + c_i(x + d_i)^2}, \quad x = \rho_B/\rho_{B0}, \tag{2}
\]

for \(i = \sigma, \omega, \rho_{B0}\) is the saturation density of symmetric nuclear matter. Five constraints on the coupling constants \(f_i(1) = 1, f_i(0) = 0, f_i''(1) = f_i''(0) = 1\) can reduce the numbers of independent parameters to three in Equation (2). The first
two constraints lead to
\[ a_i = 1 + c_i (1 + d_i)^2, \quad 3c_i d_i^2 = 1. \]

For the isovector mesons \( \rho \) and \( \delta \), their coupling constants are,
\[ \Gamma_i(\rho_B) = \Gamma_i(\rho_B) \exp[-a_i(x - 1)]. \]

While in DD-LZ1 parameterization, the coefficient in front of the fraction function, \( \Gamma_i \), is fixed at \( \rho_B = 0 \) for \( i = \sigma, \omega, \rightarrow \)
\[ \Gamma_i(\rho_B) = \Gamma_i(0) f_i(x). \]

There are only four constraint conditions as \( f_i(0) = 1 \) and \( f''_i(0) = 0 \) for \( \sigma \) and \( \omega \) coupling constants in DD-LZ1. The constraint \( f''_i(1) = f''_i(1) \) is removed in DD-LZ1 parameterization, which can give more precise shell evaluations of finite nuclei around \( Z = 58 \) and 92 (Wei et al. 2020). For the \( \rho \) meson, its coupling constant is also changed accordingly as
\[ \Gamma_i(\rho_B) = \Gamma_i(0) \exp(-a_i \rho_B x). \]

To solve the nuclear many-body system in the DDRMF model, the mean-field approximation must be adopted following the nonlinear RMF models, in which various mesons are treated as classical fields as
\[ \sigma \rightarrow \langle \sigma \rangle = \sigma, \quad \omega \rightarrow \langle \omega \rangle = \omega, \quad \rho \rightarrow \langle \rho \rangle = \rho, \quad \delta \rightarrow \langle \delta \rangle = \delta, \quad (\psi) \rightarrow \psi. \]

The space components of the vector meson are removed in the parity conservation system. In addition, the spatial derivatives about nucleons and mesons are neglected in the infinite nuclear matter due to its transformation invariance. Finally, using the Euler–Lagrange equation, the equations of motion for nucleons and mesons are obtained:
\[
\sum_{i=\text{pro,n}} \left[ i \gamma_i \partial_i - \gamma \frac{\partial}{\partial \gamma} \left( \Gamma_i(\rho_B) \omega + \frac{\Gamma_i(\rho_B)}{2} \rho \gamma + \Sigma_{\text{Re}}(\rho_B) - M_B^{*} \right) \right] \psi_i = 0.
\]

The isospin third components of nucleons are defined as \( \gamma_3 = 1 \) and \( \gamma_3 = -1 \) for protons and neutrons, respectively. A rearrangement term \( \Sigma_{\text{Re}} \) will be introduced into Equation (8) due to the density dependence of coupling constants,
\[
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\]

where the scalar, vector densities, and their isospin components are generated by the expectation value of nucleon fields,
\[
\rho_s = \langle \psi \overline{\psi} \rangle = \rho_{0s} + \rho_{ms},
\]
\[
\rho_{3s} = \langle \psi \gamma_3 \overline{\psi} \rangle = \rho_{0s} - \rho_{ms},
\]
\[
\rho_B = \langle \psi \overline{\psi} \rangle = \rho_{0B} + \rho_{mB},
\]
\[
\rho_3 = \langle \psi \gamma_3 \overline{\psi} \rangle = \rho_{0B} - \rho_{mB}.
\]

The effective masses of nucleons in Equation (8) are dependent on the scalar mesons \( \sigma \) and \( \delta \),
\[
M_B^{*} = M - \Gamma_\sigma(\rho_B) \sigma - \Gamma_\delta(\rho_B) \delta,
\]
\[
M_B^{*} = M - \Gamma_\sigma(\rho_B) \sigma + \Gamma_\delta(\rho_B) \delta
\]
and the corresponding effective energies of nucleons have the following form because of the mass–energy relation,
\[
E_B^{*} = \sqrt{k_{F_B}^{2} + (M_B^{*})^2},
\]
where \( k_{F_B} \) is the Fermi momentum of nucleons.

With the energy-momentum tensor in a uniform system, the energy density, \( \varepsilon \), and pressure, \( P \), of infinite nuclear matter can be obtained, respectively, as
\[
\varepsilon_{\text{DD}} = \frac{1}{2} \rho_0^2 \omega^2 - \frac{1}{2} \rho_0^2 \sigma^2 + \frac{1}{2} \rho_0^2 \delta^2 + P_{\text{kin}}^{\text{pro,n}} + P_{\text{kin}}^{\text{M}},
\]
\[
P_{\text{DD}} = \rho_{0s} \Sigma_{\text{Re}}(\rho_B) - \frac{1}{2} \rho_0^2 \sigma^2 + \frac{1}{2} \rho_0^2 \delta^2 + 2P_{\text{kin}}^{\text{pro,n}} + P_{\text{kin}}^{\text{M}},
\]
where the contributions from kinetic energy are
\[
\varepsilon_{\text{kin}}^{\text{pro,n}} = \frac{\gamma}{2\pi^2} \int_{k_{F_B}}^{k_{0}} k^2 \frac{1}{\sqrt{k^2 + M_{*}^{2}}} dk = \frac{\gamma}{16\pi^2} \int_{k_{F_B}}^{k_{0}} k^2 \log \left( \frac{k^2 + M_{*}^{2}}{k_{F_B}^2 + M_{*}^{2}} \right),
\]
\[
P_{\text{kin}}^{\text{pro,n}} = \frac{\gamma}{6\pi^2} \int_{k_{F_B}}^{k_{0}} k^2 \frac{1}{\sqrt{k^2 + M_{*}^{2}}} dk = \frac{\gamma}{48\pi^2} \int_{k_{F_B}}^{k_{0}} k^2 \log \left( \frac{k^2 + M_{*}^{2}}{k_{F_B}^2 + M_{*}^{2}} \right),
\]
\[
\gamma = 2 \text{ is the spin degeneracy factor. The binding energy per nucleon can be defined by}
\]
\[
E = \frac{\varepsilon}{A} - M.
\]
\[
The symmetry energy is calculated by
\]
\[
E_{\text{symDD}} = \frac{1}{2} \frac{\partial^2 E}{\partial B^2}.
\]
where $\beta$ is the asymmetry factor, defined as $\beta = (\rho_B - \rho_p) / (\rho_B + \rho_p)$ and its slope, $L_{DD}$, is given by

$$L_{DD} = 3 \rho_B \frac{\partial E_{\text{sym}}}{\partial \rho_B} \bigg|_{\rho_B = \rho_0}. \quad (17)$$

Actually, both of them can be derived analytically in the RMF model (Dutra et al. 2014).

The scalar potential and vector potential of nucleons are expressed as,

$$U_S = \Gamma_{\omega}(\rho_B) \sigma + \Gamma_3(\rho_B) \delta \tau_3, \quad (18)$$

$$U_V = \Gamma_{\omega}(\rho_B) \omega + \frac{1}{2} \Gamma_3(\rho_B) \rho \tau_3$$

$$\quad + \left[ \frac{- \partial \Gamma_{\omega}(\rho_B)}{\partial \rho_B} \sigma - \frac{\partial \Gamma_3(\rho_B)}{\partial \rho_B} \delta \rho \tau_3 \right]$$

$$\quad + \frac{\partial \Gamma_{\omega}(\rho_B)}{\partial \rho_B} \omega \rho B + \frac{1}{2} \frac{\partial \Gamma_3(\rho_B)}{\partial \rho_B} \rho \delta \rho \tau_3 \right], \quad (19)$$

where the derivative terms in the vector potential originate from the density dependence of coupling constants.

The outer core part of a neutron star, which almost dominates its mass and radius, is usually treated as the uniform matter composed of neutrons, protons, and leptons. They are stably existing with the conditions of beta equilibrium and charge neutrality. Therefore the chemical potentials of nucleons and leptons are very important; these can be derived from the thermodynamics equations at zero temperature,

$$\mu_{\text{e}} = \sqrt{k_{\text{B}} T + m_{\text{e}}^2} + \left[ \Gamma_{\omega}(\rho_B) \omega + \frac{1}{2} \Gamma_3(\rho_B) \rho \tau_3 + \Sigma_{\text{g}}(\rho_B) \right],$$

$$\mu_{\text{l}} = \sqrt{k_{\text{B}} T + m_{\text{l}}^2}. \quad (20)$$

In neutron star matter, with the density increasing, the muon will be onset when the electron chemical potential $\mu_{\text{e}}$ is larger than the muon rest mass, i.e., $\mu_{\text{e}} > m_{\text{e}} = 106.55$ MeV. Hence, the beta equilibrium condition now can be expressed by

$$\mu_{\text{e}} = \mu_{\text{e}} = \mu_{\text{n}} - \mu_{\text{p}}. \quad (21)$$

The charge neutrality condition has the following form:

$$\rho_{\text{B}0} = \rho_{\text{e}} + \rho_{\text{p}}. \quad (22)$$

The pressure and energy density will be obtained as a function of nucleon density within the constraints of Equations (21) and (22). The Tolman–Oppenheimer–Volkoff (TOV) equation (Oppenheimer & Volkoff 1939; Tolman 1939) describes a spherically symmetric star in the gravitational equilibrium from general relativity,

$$\frac{dP}{dr} = -\frac{GM(r)E(r)}{r^2} \left[ 1 + \frac{P(r)}{E(r)} \right] \left[ 1 + \frac{4\pi \tau^2 P(r)}{M(r)} \right],$$

$$\frac{dM(r)}{dr} = 4\pi r^2 E(r), \quad (23)$$

where $P$ and $M$ are the pressure and mass of the neutron star at $r$. Furthermore, the tidal deformability becomes a typical property of the neutron star after the observation of the gravitational wave from BNS merger, which characterizes the deformation of a compact object in an external field generated by another star. The tidal deformability of a neutron star can be reduced as a dimensionless form,

$$\Lambda = \frac{2}{3} k_2 C^{-5}, \quad (24)$$

where $C = GM/R$ is the compactness parameter. The second-order Love number $k_2$ (Hinderer 2008; Hinderer et al. 2010) is given by

$$k_2 = \frac{8C^5}{5} (1 - 2C)^2 [2 + 2C(\Upsilon_R - 1) - \Upsilon_R] \frac{\Upsilon_R}{\Upsilon_R^2}$$

$$\times [2C(6 - 3\Upsilon_R + 3C(5\Upsilon_R - 8))]$$

$$+ 4C^3[13 - 11\Upsilon_R + C(3\Upsilon_R - 2) + 2C^2(1 + \Upsilon_R)]$$

$$+ 3(1 - 2C)^2[2 - \Upsilon_R + 2C(\Upsilon_R - 1)\ln(1 - 2C)]^{-1}. \quad (25)$$

Here, $\Upsilon_R = y(R)$. $y(r)$ satisfies the following first-order differential equation,

$$r \frac{dy(r)}{dr} + y^2(r) + y(r)F(r) + r^2Q(r) = 0,$$

where $F(r)$ and $Q(r)$ are functions related to the pressure and energy density

$$F(r) = \left[ 1 - \frac{2M(r)}{r} \right]^{-1} [1 - 4\pi r^2 \{E(r) - P(r)\}],$$

$$r^2Q(r) = \left\{ \frac{4\pi r^2}{E(r)} \right\} \left[ 5E(r) + 9P(r) + \frac{\frac{E(r)}{\delta E(r)}}{P(r)} \right] - 6 \right\} \times \left[ 1 - \frac{2M(r)}{r} \right]^{-1} - \left[ \frac{2M(r)}{r} + 2 \times 4\pi r^2 P(r) \right] \times \left[ 1 - \frac{2M(r)}{r} \right]^{-2}. \quad (27)$$

The second Love number corresponds to the initial condition $y(0) = 2$. It is also related to the speed of sound in compact matter, $c_s$,

$$c_s^2 = \frac{\partial P(\varepsilon)}{\partial E}. \quad (28)$$

### 3. Results and Discussion

First, masses of nucleons and mesons, coupling constants between nucleons and mesons, and saturation densities of symmetric nuclear matter, $\rho_{\text{B}0}$, in DD2 (Typel et al. 2010), DD-ME1 (Nikšić et al. 2002), DD-ME2 (Lalazissis et al. 2005), DDME-X (Taninah et al. 2020), DDV, DDVT, DDVTD (Typel & Terrero 2020), and DD-LZ1 (Wei et al. 2020) sets are all listed in Table 1.

The mass of the $\sigma$ meson is fitted as a free parameter in the DDRMF model. The coefficients of meson coupling constants, $\Gamma_{\text{e}}$, in DD-LZ1 are the values at zero density, while other parameter sets adopted the values at nuclear saturation densities. The magnitudes of $\Gamma_{\text{e}}(\rho_{\text{B}0})$, $\Gamma_{\text{e}}(\rho_{\text{B}0})$, and $\Gamma_{\text{e}}(\rho_{\text{B}0})$ in DD2, DDME-1, DD-ME2, DD-MEX, and DDV are consistent with each other. The tensor couplings between vector mesons and nucleons were considered in DDVT and DDVTD, where $\Gamma_{\text{e}}(\rho_{\text{B}0})$ and $\Gamma_{\text{e}}(\rho_{\text{B}0})$ have significant
differences compared to other parameter sets. In addition, the δ meson is included in the DDVTD set.

To show the density-dependent behaviors of these coupling constants more clearly, they are plotted as functions of the vector density in Figure 1. It can be found that all of these coupling constants decrease when the nuclear density becomes larger due to the nuclear medium effect. For the ρ meson coupling constants in panel (c), all parameter sets have very similar density-dependent behaviors in the whole density region. In DDVT and DDVTD, the tensor coupling constants play obvious roles in finite nuclei due to their derivative forms; however, they do not provide any contribution in nuclear matter. Their coupling constants of ω and ρ mesons in panel (a) and panel (b) are dramatically smaller than other sets. Furthermore, the coupling constants from several typical nonlinear RMF models, NL3 (Lalazissis et al. 1997), TM1 (Sugahara & Toki 1994), IUFSU (Horowitz & Piekarewicz 2001), and BigApple (Fattoyev et al. 2020) are also shown to compare their differences with those in the DDRMF model.

At the low density region, the coupling constants in DDRMF models are usually stronger than those in nonlinear RMF modes, while weaker at higher density.

With these DDRMF parameter sets, the saturation properties of nuclear matter can be calculated, such as the saturation density, the binding energy, the incompressibility, the symmetry energy, the slope of symmetry energy, and the effective nucleon mass. In Table 2, these properties calculated by various DDRMF models are listed, whose uncertainties of different parameter sets are very small in saturation density, binding energy, incompressibility, and symmetry energy. The slopes of symmetry energy from different models, \( L \), are around 40–70 MeV, which also satisfies the recent constraints, \( L = 59.57 \pm 10.06 \) MeV (Zhang et al. 2020). On the other hand, the effective nucleon masses in DDVT and DDVTD are relatively larger, since their scalar coupling strengths are much smaller compared to other sets.

The binding energies per nucleon for symmetric nuclear matter in Figure 2(a) and pure neutron matter in Figure 2(b) as functions of vector density are plotted with the present DDRMF parameterizations. These equations of state (EOSs)

### Table 1

| Parameter | Value |
|-----------|-------|
| \( m_n \) (MeV) | 938.900000 |
| \( m_p \) (MeV) | 938.900000 |
| \( m_\omega \) (MeV) | 538.619216 |
| \( m_{\rho} \) (MeV) | 769.0000 |
| \( m_{\delta} \) (MeV) | 0.158100 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.149 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.152 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.152 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.153 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.153 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.153 |
| \( \rho_0 \) (fm\(^{-1}\)) | 0.153 |

Figure 1. Coupling constants of \( \omega, \sigma, \rho \) and \( \rho \) mesons as functions of vector density in various DDRMF models and several nonlinear RMF models.
of nuclear matter below 0.2 fm$^{-3}$ are almost identical since all the parameters were determined by properties of finite nuclei, whose central density is around nuclear saturation density $r_0 \sim 0.15$ fm$^{-3}$. Their differences increase from 0.30 fm$^{-3}$. In symmetric nuclear matter, they are separated into the softer group with DDV, DDVT, and DDVTD, and the stiffer group with DD2, DD-ME1, DD-ME2, DD-MEX, and DD-LZ1. The scalar and vector coupling strengths in the softer group sets are obviously weaker than those in the stiffer group sets. The binding energy of pure neutron matter from DDV is larger than the ones from DDVT and DDVTD. The DDV set has the largest slope of symmetry energy in the present DDRMF parameterizations. This slope will determine the density-dependent behaviors of symmetry energy and the binding energy of pure neutron matter, due to $E/A(\beta = 1) \approx E/A(\beta = 0) + E_{\text{sym}}$ at a fixed density.

In general, it is very difficult to measure properties of nuclear matter more than twice the nuclear saturation density from finite nuclei. Recently, the experiments about heavy-ion collisions have provided us with some useful information to constrain the EOS of nuclear matter at high density. In Figure 3, the pressures in symmetric nuclear matter as functions of density from various DDRMF models are shown and compared to the constraints from heavy-ion collisions. We can find that the EOSs from the softer group sets are completely consistent with the experiment data, while the other group is indeed stiffer than the heavy-ion collision constraints. We also note that the BigApple and NL3 sets also have similar situations in the work by Fattoyev et al. (2020). However, we want to emphasize here that the constraints from the heavy-ion collisions are strongly model-dependent, which is determined by many inputs, such as the NN interaction. To our knowledge, there were few investigations about heavy-ion collisions that adopted the RMF model as the NN interaction. Therefore, it cannot be claimed with certainty that the EOSs generated by DD2, DD-ME1, DD-ME2, DD-MEX, and DD-LZ1

Table 2: Nuclear Matter Properties at Saturation Density Generated by Present DDRMF Parameterizations

| Model       | $\rho_0$ [fm$^{-3}$] | $E/A$ [MeV] | $K_0$ [MeV] | $E_{\text{sym}}$ [MeV] | $L$ [MeV] | $M_0^s/M$ | $M_0^v/M$ |
|-------------|----------------------|--------------|-------------|----------------------|-----------|-----------|-----------|
| DD-LZ1      | 0.1585               | -16.126      | 32.016      | 42.467               | 0.558     | 0.558     |
| DD2         | 0.149                | -16.916      | 31.635      | 54.933               | 0.563     | 0.562     |
| DD-ME1      | 0.152                | -16.668      | 33.060      | 55.428               | 0.578     | 0.578     |
| DD-ME2      | 0.152                | -16.233      | 32.31       | 51.265               | 0.572     | 0.572     |
| DD-MEX      | 0.1518               | -16.14       | 32.269      | 49.692               | 0.556     | 0.556     |
| DDV         | 0.1511               | -16.097      | 33.589      | 69.646               | 0.586     | 0.585     |
| DDVT        | 0.1536               | -16.924      | 31.558      | 42.348               | 0.667     | 0.667     |
| DDVTD       | 0.1536               | -16.915      | 31.817      | 42.583               | 0.666     | 0.666     |

Figure 2. EOSs of symmetric nuclear matter ($\beta = 0$) in panel (a) and pure neutron matter ($\beta = 1$) in panel (b) from various DDRMF models.

Figure 3. Pressures as a function of vector density of symmetric nuclear matter with various DDRMF parameter sets and the constraints from the heavy-ion collisions.
parameterizations are clearly excluded by the constraints of heavy-ion collisions.

To explain the stiff EOSs at high density of DD2, DD-ME1, DD-ME2, DD-MEX, and DD-LZ1 sets, the vector potentials in Figure 4(a) and scalar potentials in Figure 4(b) for symmetric nuclear matter from the present DDRMF parameterizations are shown. The scalar potentials in these sets are very similar, while the vector potentials from different parameter sets have significant differences. The softer group sets provide the weakest vector potentials, which have analogous magnitudes to the scalar potentials of nucleons. On the other hand, DD-ME2, DD-MEX, and DD-LZ1 generate the strongest vector components, which are almost twice those from DDV, DDVT, and DDVTD, because their ω coupling constants are largest at high density regions. Therefore, they provide very stiff EOSs.

Together with the conditions of beta equilibrium and charge neutrality, the EOSs of neutron star matter with the DDRMF model can be obtained in Figure 5, which shows the pressures of neutron star matter as a function of energy density. In the crust region of a neutron star, the EOS in the nonuniform matter generated by TM1 parameterization with the Thomas–Fermi approximation is adopted. In the core of the neutron star, EOSs of the uniform matter are calculated with various DDRMF sets. Their density-dependent behaviors are very similar to those in symmetric nuclear matter. At the high density region, the stiffer group sets provide higher pressures due to the stronger vector potentials. The joint constraints on EOS extracted from the GW170817 and GW190814 are shown as a shaded band here. When the energy density is smaller than 600 MeV fm$^{-3}$, the EOSs from stiffer group sets satisfy the constraints from the gravitational wave detection, while the pressures obtained from softer group sets start to become lower than the constraint band from $\varepsilon = 300$ MeV fm$^{-3}$. Furthermore, the EOS from the TM1 set is also given, which is stiffer than those from DDRMF at the intermediate density region and becomes softer when density is higher. Since the slope of symmetry energy in TM1 is around 110 MeV. It is much larger than those derived from present DDRMF models, whose $L$ are around 40–70 MeV. The slope of symmetry energy mainly influences the pressures in intermediate density. At higher density, the vector potentials in DDRMF from the stiffer group sets are stronger than the one in TM1.

In Figure 6, the pressures as functions of density in neutron star matter from DDRMF models are given. The pressures from

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**Figure 4.** Scalar and vector potentials as a function of the vector density from various DDRMF models.

**Figure 5.** Pressure $P$ vs. energy density $\varepsilon$ of neutron star matter from DDRMF models and joint constraints from GW170817 and GW190814.

**Figure 6.** EOSs of neutron star matter with different DDRMF models. The corresponding speeds of sound in units of the speed of light are shown in the insert.
Figure 7. Neutron star masses as functions of radius and the central baryon density. Constraints from astronomical observables for the massive neutron star, NICER, and GW170817 are also shown.

Figure 8. Tidal deformabilities from various DDRMF models as functions of neutron star mass. The mass regions of massive neutron stars are also plotted.

the stiffer group sets are obviously larger than those generated by the softer group sets. The speeds of sound in neutron star matter, \( c_s \), with the unit of light speed are plotted in the insert. The \( c_s^2 \) from softer group sets are much lower than those from other parameterizations, which are around 0.6 at \( \rho_B = 1.0 \) fm\(^{-3}\). They are consistent with the results from nonlinear RMF models (Hu et al. 2020). The speed of sound from stiffer group EOSs rapidly increase from \( \rho_B = 0.2 \) fm\(^{-3}\) and \( c_s^2 \) reach around 0.8 at high density. They will be constants less than one as the density continues growing. In fact, the EOS and speed of light of the BigApple set in the nonlinear RMF model are very similar to the present work, where a \( \omega - \rho \) coupling term was included to reduce the slope of symmetry energy, and its vector–isovector coupling constant is very strong, as we show in Figure 1 (Fattoyev et al. 2020).

The mass–radius relation of a static neutron star can be solved by TOV Equation (23), where the EOS of neutron star matter is used. In Figure 7, the mass–radius \( (M–R) \) relation in panel (a) and mass-central density \( (M–\rho_B) \) relation in panel (b) from various DDRMF models are shown. From panel (a), it can be found that the maximum masses of neutron stars in softer group sets are around 1.85–1.93 \( M_\odot \), and the corresponding radii are 9.85–10.34 km. These results only can explain the existence of PSR J1614-2230 (1.928 ± 0.017 \( M_\odot \)). As we discussed before, the EOSs from these three parameter sets are relatively soft due to their small vector potentials. The maximum masses calculated by DD2, DD-ME1, and DD-ME2 sets are about 2.42–2.48 \( M_\odot \), which are consistent with the available investigations (Sun et al. 2008). DD-MEX and DD-LZ1 can support the neutron star above 2.5 \( M_\odot \) because of their strongest repulsive contributions from the \( \omega \) meson. Their maximum masses can approach 2.56 \( M_\odot \), which are in accord with the observed mass of the secondary compact object in GW190814, 2.50–2.67 \( M_\odot \). In addition to the constraints from the observables of massive neutron stars, PSR J1614-2230, PSR J034+0432, and PSR J0740+6620, recently the mass and radius of the neutron star at the intermediate mass region were measured simultaneously for PSR J0030+0451 by NICER. Its mass and radius were reported around 1.4 \( M_\odot \) and 13 km (Miller et al. 2019; Riley et al. 2019). These constraints from NICER are plotted in panel (a). It can be found that the \( M–R \) relations from stiffer group parameterizations around 1.4 \( M_\odot \) completely satisfy the observables from NICER, while the radii of the neutron star at 1.4 \( M_\odot \) from DDVT, and DDVTD are around 11.4 km, which are smaller than the possible radii of J0030+0451. The \( R_{1.4} \) of DDV is 12.2 km since its slope of symmetry energy is obviously larger than those of DDVT and DDVTD. When the isoscalar properties of RMF models are the same, the slope of symmetry energy can influence the radii of the neutron star at 1.4 \( M_\odot \). By the recent investigations (Ji et al. 2019; Hu et al. 2020). Meanwhile, the resulting \( R_{1.4} \) from the softer group are completely comparable to the constraints extracted from GW170817, \( R_{1.4} = 11.9 \) ± 1.4 km (Abbott et al. 2018). However, the radii, 12.9–13.1 km from stiffer EOSs are obviously larger.

The \( M–\rho_B \) relations in panel (b) from present DDRMF models can be separated into two groups. The first group can only generate the neutron star around 1.9 \( M_\odot \) at the central densities \( \sim 8 \rho_{B0} \) from softer group EOSs. The second group can support neutron stars around 2.5 \( M_\odot \), where the central densities locate at \( 5 \rho_{B0} \). The corresponding speeds of sound are less than \( \sqrt{0.8} c \) from Figure 6.
With the rapid development of gravitational wave detectors, the tidal deformability of the neutron star can be extracted from the BNS merger. It can be calculated with the Love number by solving a first-order differential equation, Equation (25). In Figure 8, the dimensionless tidal deformabilities, $\Lambda$, of the neutron star as a function of their masses from DDRMF models are shown. These dimensionless tidal deformabilities decrease with the neutron star mass and become very small at the maximum masses. Their values in softer group sets are significantly lower than those from other parameterizations, since $\Lambda \propto R^2$ approximately from Equation (24). The radii of the neutron star from the softer group EOSs are smaller compared to the stiffer EOSs. The corresponding $\Lambda$ at $1.4 \, M_\odot$, $\Lambda_{1.4}$ are from 274.91 to 390.01, while a recent analysis by LVC gives $\Lambda_{1.4} = 190^{+390}_{-120}$ from GW170817 (Abbott et al. 2018). Due to the larger radii and speeds of sound of neutron stars in stiffer group EOSs, the $\Lambda$ are relatively higher and $\Lambda_{1.4}$ are between 639.03 and 790.01. Furthermore, the tidal deformabilities at $2.0 \, M_\odot$ from these two types of EOSs have obvious differences. For the softer EOSs, $\Lambda$ almost approach zero, while they are around 100 for the stiffer EOSs at $2.0 \, M_\odot$. In fact, there are none of the stiffer EOSs, whose tidal deformabilities are compatible with the present constraints from GW170817. Once the BNS merger, whose components are around $2 \, M_\odot$, is more precisely measured by the advanced gravitational wave detectors in the future, the EOSs of the neutron star can hopefully be determined well.

Finally, properties of neutron stars, i.e., the maximum mass ($M_{\text{max}}/M_\odot$), the corresponding radius ($R_{\text{max}}$), the central density ($\rho_c$), the radius ($R_{1.4}$), and dimensionless tidal deformability ($\Lambda_{1.4}$) at $1.4 \, M_\odot$ from present DDRMF models are listed in Table 3, respectively.

| Table 3 | Neutron Star Properties from Various DDRMF Models |
|---------|-----------------------------------------------|
|         | DD-LZ1 | DD2 | DD-ME1 | DD-ME2 | DD-MEX | DDV | DDVT | DDVTD |
| $M_{\text{max}}/M_\odot$ | 2.5545 | 2.4168 | 2.4426 | 2.4829 | 2.5566 | 1.9317 | 1.9251 | 1.8507 |
| $R_{\text{max}}$ [km] | 12.178 | 11.826 | 11.885 | 12.012 | 12.274 | 10.336 | 10.023 | 9.850 |
| $\rho_{\text{max}}$ [fm$^{-3}$] | 0.786 | 0.845 | 0.832 | 0.813 | 0.777 | 1.188 | 1.237 | 1.306 |
| $R_{1.4}$ [km] | 12.864 | 12.938 | 12.931 | 12.961 | 13.118 | 12.195 | 11.511 | 11.396 |
| $\Lambda_{1.4}$ | 727.071 | 639.032 | 686.786 | 730.737 | 790.051 | 390.005 | 301.388 | 274.908 |

4. Summary and Perspectives

The latest density-dependent relativistic mean-field (DDRMF) parameterizations were systematically applied to investigate the properties of neutron stars, i.e., DD2, DD-ME1, DD-ME2, DD-MEX, DD-LZ1, DDV, DDVT, and DDVTD sets. All of them were determined by fitting properties of spherical finite nuclei and have the same density-dependent function forms for meson coupling constants. Their densities, binding energies, incompressibilities, and symmetry energies at saturation points of symmetric nuclear matter are almost identical.

The EOSs of symmetric nuclear matter and pure neutron matter from present sets were separated into the softer type and the stiffer one at the high density region. The softer EOSs are generated by the DDV, DDVT, and DDVTD, whose coupling strengths of $\sigma$ and $\omega$ mesons are weaker compared to other sets. Their vector and scalar potentials have comparable magnitudes, while the vector potentials are much larger than the scalar ones in stiffer EOSs given by DD2, DD-ME1, DD-ME2, DD-MEX, and DD-LZ1. Their pressures in symmetric nuclear matter at $2 \sim 4\rho_0$ were a little bit higher than the present constraints from heavy-ion collisions, while the softer EOSs satisfied these constraints.

The TOV equation was solved using the EOSs of neutron star matter, where the nucleons and leptons are in conditions of beta equilibrium and charge neutrality, generated by present DDRMF models. The softer EOSs from DDV, DDVT, and DDVTD can only support the neutron stars with maximum masses around $1.0 \, M_\odot$ at $10 \, \text{km}$ and tidal deformabilities at $1.4 \, M_\odot$, $\Lambda_{1.4} = 274$–390. The stiffer EOSs can generate very massive neutron stars around $2.5 \, M_\odot$. In particular, the DD-MEX and DD-LZ1 parameter sets can even produce neutron stars with masses of $2.55 \, M_\odot$, which can explain the secondary object in GW190814 with a mass of $2.50$–$2.67 \, M_\odot$. Furthermore, their radii at $1.4 \, M_\odot$ are also consistent with the constraints from NICER including the simultaneous mass and radius measurements, while their $\Lambda_{1.4}$ were around 639–790, which are not compatible with the values extracted from GW170718 $\Lambda_{1.4} = 190^{+390}_{-120}$.

In this work, we found that several parameterizations in DDRMF can provide very massive neutron stars. The DDRMF theory properly takes into account the nuclear medium effect in the effective NN potentials through the density-dependent coupling constants. Therefore, it gives a reasonable description for the finite nuclei at low density and neutron stars at a very high density at the same time. In particular, the short-range correlations between nucleons at high density will become stronger. As a result, it provides very massive neutron stars due to the strong repulsive contribution from vector mesons at high density. The stiffer EOSs may slightly exceed the constraints of EOS from heavy-ion collisions and tidal deformability from GW170817. However, due to the strong model dependence of these two constraints deduced from the present measurements, we cannot exclude the possibility of the secondary object of GW190814 as a neutron star consisting of nucleons and leptons. We have shown that the stiffer EOSs give a dimensionless tidal deformability around 100 for the $2 \, M_\odot$ massive neutron star, while the softer ones give less than 10. Therefore, the more precise measurement of dimensionless tidal deformability by the gravitational wave detectors will help us to determine the proper EOSs in the future.

We acknowledge the works that first proposed the idea discussed within this paper, such as Most et al. (2020), Tan et al. (2020), Toskaros et al. (2020). This work was supported in part by the National Natural Science Foundation of China (grants No. 11775119, No. 11675083, and No. 11405116), the Natural Science Foundation of Tianjin, and China Scholarship Council (grant No. 201906205013 and No. 201906255002).
The Astrophysical Journal, 904:39 (10pp), 2020 November 20

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