Strangeness and $\Delta$ resonance in compact stars with relativistic-mean-field models

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We explore the effects of strangeness and $\Delta$ resonance in baryonic matter and compact stars within the relativistic-mean-field (RMF) models. The covariant density functional PKDD is adopted for N-N interaction, parameters fixed based on finite hypernuclei and neutron stars are taken for the hyperon-meson couplings, and the universal baryon-meson coupling scheme is adopted for the $\Delta$-meson couplings. In light of the recent observations of GW170817 with the dimensionless combined tidal deformability $197 \leq \tilde{\Lambda} \leq 720$, we find it is essential to include the $\Delta$ resonances in compact stars, and small $\Delta$-$b$ coupling $g_{\Delta b}$ is favored if the mass $2.27_{-0.15}^{+0.17}M_\odot$ of PSR J2215+5135 is confirmed.

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I. INTRODUCTION

The recent observation of gravitational waves from the binary neutron star merger event GW170817 suggests that the merging objects are compact [1, 2]. Assuming low spin priors, the dimensionless combined tidal deformability $\tilde{\Lambda}$ is considered to be less than 720 at 90% confidence level [3], while a lower limit with $\tilde{\Lambda} \geq 197$ is obtained based on electromagnetic observations of the transient counterpart AT2017gfo [4]. Even though the observations of neutron stars’ radii are controversial and depend on specific assumptions, the recent measurements seem to be converging and lie at the lower end of 10-14 km range [2, 5–9]. The combined constraints on the tidal deformability and radii of neutron stars indicate a soft equation of states (EoS), where many covariant density functionals are in jeopardy [10, 11]. A possible solution to this problem is to introduce new degrees of freedom, e.g., $\Delta$ resonances, hyperons, and deconfined quarks [12]. As one increases the density of nuclear matter, the inevitable emergence of $\Delta$ isobars, hyperons, and quarks can soften the EoSs significantly and reduce the radius and tidal deformability of the corresponding compact stars, which can be consistent with these recent observations.

However, a soft EoS will result in compact stars with too small masses that cannot reach two solar mass as observed in pulsars PSR J1614-2230 (1.928±0.017$M_\odot$) [13, 14] and PSR J0348+0432 (2.01±0.04$M_\odot$) [15], i.e., the Hyperon Puzzle [16] or $\Delta$ Puzzle [17]. Extensive efforts were made to resolve the Hyperon Puzzle [18–26] and $\Delta$ Puzzle [37–40]. Nevertheless, with the constrained observable tidal deformability of GW170817 [1, 3, 4], these solutions may be challenged, especially for the latest observation of a more massive PSR J2215+5135 (2.27$^{+0.17}_{-0.15}$ $M_\odot$) [41].

To satisfy these stringent observational constraints, we seriously consider the possible existence of both $\Delta$ isobars and hyperons in neutron stars. Since relativistic-mean-field (RMF) models [42–49] have been successfully adopted to describe finite (hyper)nuclei [50–60] and baryonic matter [61–68], in this work the EoSs of baryonic matter are obtained based on RMF model. More specifically, we adopt the covariant density functional PKDD [69], while the hyperon-meson couplings are fixed based on our previous investigations on hypernuclei and neutron stars [36, 60, 70]. For the $\Delta$-meson couplings, as in Ref. [17], we adopt the universal baryon-meson coupling scheme, while a vanishing $\Delta$-$b$ coupling is considered as well. It is found that the observational tidal deformability and mass of PSR J2215+5135 can be reproduced only by including $\Delta$ isobars in neutron stars.

The paper is organized as follows. In Sec. II, we present the formalism of RMF model for baryonic matter, the choices of baryon-meson couplings, the conditions for obtaining the EoSs of neutron star matter, and the formalism to determine the structures of compact stars. Results and discussions are given in Sec. III. We make a summary in Sec. IV.

II. THEORETICAL FRAMEWORK

The Lagrangian density of RMF models is given as

\[
\mathcal{L} = \sum_b \bar{\psi}_b [i \gamma^\mu \partial_\mu - m_b - g_{b\sigma} \sigma - g_{b\omega} \gamma^\mu \omega_\mu - g_{b\rho} \gamma^\mu \rho_\mu - g_{b\omega} \gamma^\mu \omega_\mu + \frac{1}{2} m_b^2 \sigma^2 - \frac{1}{4} \omega_\mu \omega^{\mu \nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^{\mu \nu} - \frac{1}{4} \rho^{\mu \nu} \cdot \rho^{\nu \mu} + \frac{1}{2} m_\rho^2 \rho_\mu \cdot \rho^\mu - \frac{1}{4} A_\mu \cdot A^{\mu \nu} + \sum_{l=e,\mu} \bar{\psi}_l [i \gamma^\mu \partial_\mu - m_l + e \gamma^\mu A_\mu] \psi_l, \tag{1}
\]
with the field tensors
\[ \omega_{\mu \nu} = \partial_{\mu} \omega_{\nu} - \partial_{\nu} \omega_{\mu}, \]
\[ \rho_{\mu \nu} = \partial_{\mu} \rho_{\nu} - \partial_{\nu} \rho_{\mu}, \]
\[ A_{\mu \nu} = \partial_{\mu} A_{\nu} - \partial_{\nu} A_{\mu}. \]

The included baryons here are nucleons, hyperons (\( \Lambda^0 \), \( \Sigma^{+,0,-} \), and \( \Xi^{0,-} \)), and \( \Delta \) resonance. To describe the baryon-baryon interactions, the isoscalar-scalar channel (\( \sigma \)), isoscalar-vector channel (\( \omega \)) and isovector-vector channel (\( \rho \)) are considered.

Based on the Typel-Woter ansatz [56], the density dependence of coupling constants \( g_{\Omega b} (\xi = \sigma, \omega) \) are obtained with
\[ g_{\Omega b}(n) = g_{\Omega b}(n_0) a_\xi \frac{1 + b_\xi (n/n_0 + d_\xi)}{1 + c_\xi (n/n_0 + e_\xi)^2}, \]
where \( n \) is the density of nuclear matter with \( n_0 \) being the saturation density. Note that a different formula is adopted for the \( \rho \) meson, i.e.,
\[ g_{\rho b}(n) = g_{\rho b}(n_0) \exp[-a_\rho (n/n_0 - 1)]. \]

The space-like components of the vector fields \( \omega_{\mu}\rho \) and \( \rho_{\mu} \) vanish, leaving only the time components \( \omega_0 \) and \( \rho_0 \). Meanwhile, the charge conservation guarantees that only the 3rd component in the isospin space of \( \rho_0 \) survives. In the mean-field and no-sea approximations, the single particle (s.p.) Dirac equations for baryons and the Klein-Gordon equations for mesons and photon can be obtained from the variational procedure.

For the \( N-N \) interactions, we adopt the covariant density functional PKDD [69], which gives the saturation density \( n_0 = 0.149552 \text{ fm}^{-3} \), saturation energy \( E_0 = -16.267 \text{ MeV} \), incompressibility \( K = 262.181 \text{ MeV} \) and symmetry energy \( E_{\text{sym}} = 36.790 \text{ MeV} \).

Table I. Strangeness number \( S \), mass \( M \), third component of isospin \( \tau_3 \), total angular momentum and parity \( J^P \), charge \( q \), and coupling constants \( \alpha_\xi = g_{\Omega b}/g_{\Omega N} (\xi = \sigma, \omega, \rho) \) for \( \Lambda^0, \Xi^{0,-}, \Sigma^{+,0,-} \), and \( \Delta \) baryons.

| \( S \) | \( M \) (MeV) | \( \tau_3 \) | \( J^P \) | \( q \) (e) | \( \alpha_\sigma \) | \( \alpha_\omega \) | \( \alpha_\rho \) |
|----|---------|--------|--------|--------|-------------|-------------|-------------|
| \( \Lambda^0 \) | -1 | 1115.6 | 0 (1/2)+ | 0 | 0.878 | 1 | 0 |
| \( \Xi^{0,-} \) | -2 | 1314.9 | +1 (1/2)+ | 0 | 0.844 | 1 | 1 |
| \( \Sigma^{+,0,-} \) | -3 | 1321.3 | -1 (1/2)+ | -1 | 0.844 | 1 | 1 |
| \( \Delta \) | -0 | 1197.4 | -1 (1/2)+ | -1 | 0.878 | 1 | 1 |
| \( \Delta^+ \) | 0 | 1232 \pm 120 | +3 (3/2)+ | +2 | 1 | 1 | 0.1 |
| \( \Delta^- \) | 0 | 1232 \pm 120 | +1 (3/2)+ | +1 | 1 | 1 | 0.1 |
| \( \Delta^0 \) | 0 | 1232 \pm 120 | 0 (3/2)+ | 0 | 1 | 1 | 0.1 |
| \( \Delta^- \) | 0 | 1232 \pm 120 | -3 (3/2)+ | -1 | 1 | 1 | 0.1 |

Based on the Lagrangian density in Eq. (1), the meson fields are obtained by solving
\[ m_\sigma^2 \sigma = - \sum_b g_{\sigma b} n_b, \]
\[ m_\omega^2 \omega_0 = \sum_b g_{\omega b} n_b, \]
\[ m_\rho^2 \rho_{0,3} = \sum_b g_{\rho b} n_b, \]
with the number density \( n_b = \langle \bar{b} \psi_b^\dagger \psi_b \rangle \) and scalar density \( n_\sigma^s = \langle \bar{b} \sigma \psi_b \rangle \) of baryon type b, which are given in Eqs. (8) and (9). Here we take \( \sigma, \omega_0 \) and \( \rho_{0,3} \) as their mean values.
At zero temperature, with no sea approximation, the energy density can be determined by
\[ E = \sum_{i=b,l} \varepsilon_i(\nu_i, m_i^2) + \sum_{\xi = \sigma, \omega, \rho} \frac{1}{2} m_\xi^2 \xi^2, \tag{6} \]
in which the kinetic energy density of fermion \( i \) is
\[ \varepsilon_i(\nu_i, m_i) = \int_0^{\nu_i} \frac{f_i p^2}{2 \pi^2} \sqrt{p^2 + m_i^2} dp \]
\[ = \frac{f_i m_i^4}{16 \pi^2} \left[ \nu_i (2 x_i^2 + 1) \sqrt{x_i^2 + 1 - \text{arcsh}(x_i)} \right]. \tag{7} \]

Here we have defined \( x_i = \nu_i / m_i \) with \( \nu_i \) being the Fermi momentum and \( f_i = 2 J_i + 1 \) the degeneracy factor of particle type \( i \). Note that in Eq. (6), the baryon effective mass is defined as \( m_b^0 = m_b + g_{\sigma b} \sigma \), while the mass of leptons remain constants with \( m_l^0 = m_l \). The source currents of fermion \( i \) are given by
\[ n_i = \langle \bar{\psi}_i \gamma^0 \psi_i \rangle = \frac{f_i \nu_i^3}{6 \pi^2}, \tag{8} \]
\[ n_i^e = \langle \bar{\psi}_i \psi_i \rangle = \frac{f_i m_i^4}{4 \pi^2} \left[ x_i \sqrt{x_i^2 + 1 - \text{arcsh}(x_i)} \right]. \tag{9} \]

The chemical potentials for baryons \( \mu_b \), and leptons \( \mu_l \) are
\[ \mu_b = g_{n b} \omega_0 + g_{\sigma b} \sigma_0 n_0 + \Sigma_b^R + \sqrt{v_b^2 + m_b^2}, \tag{10} \]
\[ \mu_l = \sqrt{v_l^2 + m_l^2}, \tag{11} \]
with the “rearrangement” term
\[ \Sigma_b^R = \sum_b \left( \frac{dg_{\sigma b}}{dn} \sigma n^b + \frac{dg_{n b}}{dn} \omega n_b + \frac{dg_{\sigma b}}{dn} \rho_{\sigma b} n_{\sigma b} \right). \tag{12} \]

Then the pressure is expressed by
\[ P = \sum_i \mu_i n_i - E. \tag{13} \]

For neutron star matter, it should fulfill the charge neutrality condition
\[ \sum_i q_i n_i = 0, \tag{14} \]
with \( q_i \) being the charge of particle type \( i \). To reach the lowest energy, particles will undergo weak reactions until the \( \beta \)-equilibrium condition is satisfied, i.e.,
\[ \mu_b = \mu_n - q_b \mu_e, \quad \mu_l = \mu_e. \tag{15} \]

The EoS of neutron star matter can be obtained from Eqs. (6) and (13), which is the input of the Tolman-Oppenheimer-Volkov (TOV) equation
\[ \frac{dP}{dr} = -\frac{G M E (1 + P/E)(1 + 4\pi^3 P/M)}{1 - 2GM/r}. \tag{16} \]

By solving the TOV equation with the subsidiary condition
\[ \frac{dM(r)}{dr} = 4\pi Er^2, \tag{17} \]
we get the relation of mass \( M \) and radius \( R \) of a neutron star. Here, the gravity constant \( G = 6.707 \times 10^{-45} \text{ MeV}^{-2} \). The tidal deformability of a compact star is extracted from
\[ \Lambda = \frac{2k_2}{3} \left( \frac{R}{G M} \right)^5, \tag{18} \]
where \( k_2 \) is the second Love number and can be fixed simultaneously with the structures of compact stars [81–83].

III. RESULTS AND DISCUSSIONS

At given total baryon number density \( n \), the properties of neutron star matter can be obtained by fulfilling the conditions of baryon number conservation with \( n = \sum_b n_b \), charge neutrality in Eq. (14), and chemical equilibrium in Eq. (15) simultaneously. The particle number density for each species is determined by Eq. (8), where the corresponding values are presented as functions of the total baryon number density \( n \) in Figs. 1 and 2. By including \( \Lambda^0 \) in nuclear matter, as indicated by the dashed curves in Fig. 1, the densities of protons and neutrons are slightly reduced on the emergence of \( \Lambda^0 \). If we also include other hyperons (\( \Xi^{-0,-} \) and \( \Sigma^{+, 0,-} \)) (dash-dotted curves), since similar potential well depths are adopted for \( \Lambda^0 \)’s and \( \Sigma^0 \)’s, the \( \Sigma^- \) firstly appears at \( n = 0.27 \text{ fm}^{-3} \) due to the negative charge it carries. In
such cases, the number densities of leptons are decreased while protons are increased. Meanwhile, the onset density of $\Lambda^0$ is increased from $n = 0.39$ fm$^{-3}$ to 0.46 fm$^{-3}$ due to the inclusion of the negatively charged $\Sigma^-$. Since $\Xi^0_-$ possess the largest masses, their onset densities are much larger with $n^\text{crit}_{\Xi^-} = 1.2$ fm$^{-3}$ and $n^\text{crit}_{\Xi^-} > n^\text{crit}_{\Sigma^-}$, which exceed the density limit of Fig. 1.

The effects of $\Delta$ resonances are also studied and the results are shown in Fig. 2. To consider the Breit-Wigner mass distribution of the $\Delta$ baryons and the possible in-medium mass shift [38], three masses $m_\Delta = 1112$ MeV, 1232 MeV and 1352 MeV are adopted in our calculation. Note that the nucleon effective mass $m^*_N = m_N + g_{\rho N} \sigma$ may become negative at higher densities. This is out of the scope of our current study and we do not consider such cases. Thus, when we adopt $m_\Delta = 1112$ MeV, 1232 MeV and $g_{\rho \Delta} = g_{\rho N}$, in Fig. 2 we do not present the results with $m^*_N < 0$ at the higher densities. For all $\Delta$ baryons, the negatively charged $\Delta^-$ appears firstly as we increase the density. The onset density of $\Delta^-$ is found to increase both with $m_\Delta$ and $g_{\rho \Delta}$, which is consistent with previous findings [38, 39]. For massive $\Delta$’s ($m_\Delta = 1352$ MeV), the effects of $\Delta$ resonance are insignificant and only $\Delta^- \Sigma^+$ appears. In the comparison with hyperons, the massive $\Delta^-$ appears at larger densities than $\Sigma^-$, where the densities of hyperons are similar as the cases in Fig. 1. If we adopt smaller values of $m_\Delta$ and $g_{\rho \Delta}$, the effects of $\Delta$ resonances become important, where $\Delta^-$, $\Delta^0$, $\Delta^+$, and $\Delta^{++}$ appear sequentially as increasing the density. Consequently, hyperons are hindered and appear only at larger densities. In the extreme case of $m_\Delta = 1112$ MeV and $g_{\rho \Delta} = 0$, the only left hyperon is $\Lambda^0$, which appears at a much larger density $n = 0.74$ fm$^{-3}$. Note that a first-order phase transition from nuclear matter to $\Delta$ matter.

FIG. 2. (Color online) Same as Fig. 1 but including $\Delta$ resonances with $m_\Delta = 1112$ MeV, 1232 MeV, and 1352 MeV. Left: $g_{\rho \Delta} = 0$; Right: $g_{\rho \Delta} = g_{\rho N}$.
takes place in the density range $n = 0.083 - 0.17 \text{ fm}^{-3}$, where we have shown the corresponding densities in the lower left panel of Fig. 2.

Based on the number density of each species, the energy density $E$ and pressure $P$ of neutron star matter can be obtained from Eqs. (6) and (13). In Fig. 3 we present the energy per baryon of neutron star matter as a function of the baryon number density. As expected, the EoS becomes soft once we include new degrees of freedom. For hyperonic matter (dash-dotted curve), if we consider $\Delta$ resonances and adopt the largest mass, i.e., $m_\Delta = 1352 \text{ MeV}$, the EoS is modified slightly at high density regions since only $\Delta^-$ appears at insignificant densities $n_\Delta^-$. Moreover, adopting smaller values of $m_\Delta$ and $g_\rho\Delta$ would result in softer EoSs, where in the extreme case of $m_\Delta = 1112 \text{ MeV}$ and $g_\rho\Delta = 0$, a softest EoS is obtained for neutron star matter.

Based on the EoSs displayed in Fig. 3, the structure of a neutron star can be determined by solving the TOV equation in Eq. (16). For neutron star matter at sub-saturation densities ($n \leq 0.08 \text{ fm}^{-3}$), we adopt the EoS presented in Refs. [84–86], where the properties of crystallized matter that forms the neutron star crust can be well described. In Fig. 4 we show the masses of compact stars as functions of radius (Left panel) and central baryon number density (Right panel), where the possible existence of hyperons and $\Delta$ resonances are considered. The obtained results are compared with the observational masses of PSR J0348+0432 ($2.01 \pm 0.04 \text{ M}_\odot$) [15] and PSR J2215+5135 ($2.27^{+0.17}_{-0.15} \text{ M}_\odot$) [41]. As we include more degrees of freedom, the maximum mass and radii of compact stars become smaller. For compact stars including $\Delta$ resonances, if we adopt $m_\Delta = 1112 \text{ MeV}$ and $g_\rho\Delta = g_\rho N$, the maximum mass does not reach the lower limit of PSR J2215+5135. This can be fixed by using smaller values of $\rho^2$ couplings, e.g., $g_\rho\Delta = 0$. Due to the occurrence of a first-order phase transition at small densities ($n = 0.083-0.17 \text{ fm}^{-3}$), a smallest radius with $R = 11.3 \text{ km}$ for $1.4 \text{ M}_\odot$ compact star is obtained, which is consistent with the recent measurements of neutron star radii [2, 5–9].

Another important constraint is the tidal deformability of the compact stars, which can be obtained based on Eq. (18). In Fig. 5 we present the tidal deformabilities of compact stars corresponding to those in Fig. 4. The ob-
servation of binary neutron star merger event GW170817 have set the dimensionless combined tidal deformability $197 \leq \bar{\Lambda} \leq 720$ \cite{3, 4}, which is a mass-weighted linear combination of tidal deformabilities \cite{88}

$$\bar{\Lambda} = \frac{16}{13} \left( m_2 + 12m_2 \right) m_1 \Lambda_1 + \left( m_2 + 12m_1 \right) m_2 \Lambda_2. \quad (19)$$

Since $\bar{\Lambda}$ is insensitive to the mass ratio $m_2/m_1$ \cite{88}, combined with the best measured chirp mass $M = (m_1 m_2)^{3/5} (m_1 + m_2)^{-1/5} = 1.186 \pm 0.001 M_{\odot}$ \cite{3}, in Fig. 5 we show the corresponding constraint on the tidal deformability $\Lambda = \Lambda_1 + \Lambda_2$ at $m_1 = m_2 = 1.362 M_{\odot}$. It is found that the observational tidal deformability has put a strong constraint on the compositions of compact stars, so that the $\Delta$ resonances have to be included. Meanwhile, as discussed before, a small enough $\Delta$-coupling $g_{\rho \Delta}$ should also be adopted for compact stars to reach the mass of PSR J2215+5135.

IV. CONCLUSION

We explore the possible existence of hyperons and $\Delta$ resonances in compact stars. The properties of baryonic matter is obtained based on the RMF models. For the $N$-$N$ interactions, we adopt the covariant density functional PKDD \cite{69}, while the hyperon-meson couplings are fixed based on our previous investigations on hypernuclei and neutron stars \cite{36, 70}. For the $\Delta$-meson couplings, we adopt the universal baryon-meson coupling scheme. Meanwhile, to consider the possibility of smaller $g_{\rho \Delta}$ and mass variations, we also study the cases with $g_{\rho \Delta} = 0$ and various $\Delta$ masses with $m_\Delta = 1112$ MeV, 1232 MeV, and 1352 MeV. The EoSs of neutron star matter become softer once we include new degrees of freedom. By solving the TOV equation with these EoSs, we obtained the masses, radii, and tidal deformabilities of the corresponding compact stars. Comparing with the dimensionless combined tidal deformability $197 \leq \bar{\Lambda} \leq 720$ constrained according to the recent observations of GW170817 \cite{3, 4}, we find it is essential to include the $\Delta$ resonances in compact stars, and the $\Delta$-$\rho$ coupling $g_{\rho \Delta}$ should be small enough if the mass of PSR J2215+5135 $(2.7^{+0.17}_{-0.15} M_{\odot})$ \cite{41} is confirmed.

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\begin{thebibliography}{99}
\bibitem{1} LIGO Scientific Collaboration and Virgo Collaboration (LIGO Scientific Collaboration and Virgo Collaboration), Phys. Rev. Lett. \textbf{119}, 161101 (2017).
\bibitem{2} The LIGO Scientific Collaboration and the Virgo Collaboration, arXiv \textbf{1805.11581} (2018).
\bibitem{3} The LIGO Scientific Collaboration and the Virgo Collaboration, arXiv \textbf{1805.11579} (2018).
\bibitem{4} M. W. Coughlin, T. Dietrich, Z. Doctor, D. Kasen, S. Coughlin, A. Jerkstrand, G. Leloudas, O. McBrien, B. D. Metzger, R. O’Shaughnessy, and S. J. Smartt, arXiv \textbf{1805.09371} (2018).
\bibitem{5} S. Guillot, M. Servillat, N. A. Webb, and R. E. Rutledge, Astrophys. J. \textbf{772}, 7 (2013).
\bibitem{6} J. Lattimer and A. Steiner, Eur. Phys. J. A \textbf{50}, 40 (2014).
\bibitem{7} F. Özel and P. Freire, Annu. Rev. Astron. Astrophys. \textbf{54}, 401 (2016).
\bibitem{8} Z.-S. Li, Z.-J. Qu, L. Chen, Y.-J. Guo, J.-L. Qu, and R.-X. Xu, Astrophys. J. \textbf{795}, 56 (2015).
\bibitem{9} A. W. Steiner, C. O. Heinke, S. Bogdanov, C. K. Li, W. C. G. Ho, A. Bahramian, and S. Han, Mon. Not. R. Astron. Soc. \textbf{476}, 421 (2018).
\bibitem{10} Z.-Y. Zhu, E.-P. Zhou, and A. Li, arXiv \textbf{1802.05510} (2018).
\bibitem{11} T. Malik, N. Alam, M. Fortin, C. Providência, B. K. Agrawal, T. K. Jha, B. Kumar, and S. K. Patra, arXiv \textbf{1805.11963} (2018).
\bibitem{12} R. Gomes, P. Char, and S. Schramm, arXiv \textbf{1806.04763} (2018).
\bibitem{13} P. B. Demorest, T. Pennucci, S. M. Ransom, M. S. E. Roberts, and J. W. T. Hessels, Nature \textbf{467}, 1081 (2010).
\bibitem{14} E. Fonseca, T. T. Pennucci, J. A. Ellis, I. H. Stairs, D. J. Nice, S. M. Ransom, P. B. Demorest, Z. Arzoumanian, K. Crowter, T. Dolch, R. D. Ferdman, M. E. Gonzalez, G. Jones, M. L. Jones, M. T. Lam, L. Levin, M. A. McLaughlin, K. Stovall, J. K. Swiggum, and W. Zhu, Astrophys. J. \textbf{832}, 167 (2016).
\bibitem{15} J. Antoniadis, P. C. C. Freire, N. Wex, T. M. Tauris, R. S. Lynch, M. H. van Kerkwijk, M. Kramer, C. Bassa, V. S. Dhillon, T. Driebe, J. W. T. Hessels, V. M. Kaspi, V. I. Kondratiev, N. Langer, T. R. Marsh, M. A. McLaughlin, T. T. Pennucci, S. M. Ransom, I. H. Stairs, J. van Leeuwen, J. P. W. Verbiest, and D. G. Whelan, Science \textbf{340}, 6131 (2013).
\bibitem{16} I. Vidaña, AIP Conf. Proc. \textbf{1645}, 79 (2015).
\bibitem{17} A. Drago, A. Lavagno, G. Pagliara, and D. Pigato, Phys. Rev. C \textbf{90}, 065809 (2014).
\bibitem{18} S. Weissenborn, D. Chatterjee, and J. Schaffner-Bielich, Phys. Rev. C \textbf{85}, 065802 (2012).
\bibitem{19} Bednarek, I., Haensel, P., Zdunik, J. L., Bejger, M., and Maïka, R., Astron. Astrophys. \textbf{543}, A157 (2012).
\bibitem{20} M. Oertel, C. Providência, F. Guinilini, and A. R. Raduta, J. Phys. G: Nucl. Part. Phys. \textbf{42}, 075202 (2015).
\bibitem{21} K. Maslov, E. Kolomeitsev, and D. Voskresensky, Phys. Lett. B \textbf{748}, 369 (2015).
\bibitem{22} K. Maslov, E. Kolomeitsev, and D. Voskresensky, Nucl. Phys. A \textbf{950}, 64 (2016).
\end{thebibliography}
| Author(s) | Title and Reference |
|----------|---------------------|
| T. Takatsuka, S. Nishizaki, and Y. Yamamoto | Eur. Phys. J. A 13, 213 (2002) |
| I. Vidaña, D. Logoteta, C. Providência, A. Pols, and I. Bombaci | Europhys. Lett. 94, 11002 (2011) |
| J. Meng and S. G. Zhou | J. Phys. G: Nucl. Part. Phys. |
| N. Paar, D. Vretenar, E. Khan, and G. Colò | Rep. Prog. Phys. 88, 022801 (2015) |
| D. Lonardoni, A. Lovato, S. Gandolfi, and F. Pedervera | Phys. Rev. Lett. 114, 092301 (2015) |
| H. Togashi, E. Hiya, Y. Yamamoto, and M. Takano | Phys. Rev. C 93, 035808 (2016) |
| S. Weissenborn, I. Sagert, G. Pagliara, M. Hempel, and J. Schaffner-Bielicke | Astrophys. J. 740, L14 (2011) |
| J. Mareš and B. K. Jennings | Phys. Rev. C 49, 2472 (1994) |
| C. Y. Song, J. M. Yao, H. F. Lv, and J. Meng | Int. J. Mod. Phys. E 19, 2538 (2010) |
| Y. Tanimura and K. Hagino | Phys. Rev. C 85, 014306 (2012) |
| X.-S. Wang, H.-Y. Sang, J.-H. Wang, and H.-F. Lv | Commun. Theor. Phys. 60, 479 (2013) |
| P.-G. Reinhard | Rep. Prog. Phys. 52, 439 (1989) |
| R. Ring, J. Meng, H. Toki, S. Zhou, S. Zhang, W. Long, and L. Geng | Phys. Rev. Lett. 114, 052501 (2015) |
| J. Meng, S. G. Zhou, and J. Meng | J. Phys. G: Nucl. Part. Phys. 42, 093101 (2015) |
| J. Meng | ed., Relativistic Density Functional for Nuclear Structure: International Review of Nuclear Physics, Vol. 10 (World Scientific Pub Co Pte Lt, 2016) |

**Notes:**
- The references are cited in the text to indicate sources for the extracted content.
[76] J. Schaffner and I. N. Mishustin, Phys. Rev. C 53, 1416 (1996).
[77] T. Miyatsu, M.-K. Cheoun, and K. Saito, Phys. Rev. C 88, 015802 (2013).
[78] Z. Li, G. Mao, Y. Zhuo, and W. Greiner, Phys. Rev. C 56, 1570 (1997).
[79] D. Kosov, C. Fuchs, B. Martemyanov, and A. Faessler, Phys. Lett. B 421, 37 (1998).
[80] T. Schürhoff, S. Schramm, and V. Dexheimer, ApJ 724, L74 (2010).
[81] T. Damour and A. Nagar, Phys. Rev. D 80, 084035 (2009).
[82] T. Hinderer, B. D. Lackey, R. N. Lang, and J. S. Read, Phys. Rev. D 81, 123016 (2010).
[83] S. Postnikov, M. Prakash, and J. M. Lattimer, Phys. Rev. D 82, 024016 (2010).
[84] R. P. Feynman, N. Metropolis, and E. Teller, Phys. Rev. 75, 1561 (1949).
[85] G. Baym, C. Pethick, and P. Sutherland, Astrophys. J. 170, 299 (1971).
[86] J. W. Negele and D. Vautherin, Nucl. Phys. A 207, 298 (1973).
[87] M. Favata, Phys. Rev. Lett. 112, 101101 (2014).
[88] S. A. Bhat and D. Bandyopadhyay, arXiv 1807.06437 (2018).