Strangeness and Δ resonance in compact stars with relativistic-mean-field models

Ting-Ting Sun (孙亭亭),[1] Shi-Sheng Zhang (张时声),[2] Qiu-Lan Zhang (张秋兰),[3] and Cheng-Jun Xia (夏铖君)[3,*]

[1]*School of Physics and Engineering, Zhengzhou University, Zhengzhou 450001, China*
[2]*School of Physics and Nuclear Energy Engineering, Beihang University, Beijing 100191, China*
[3]*School of Information Science and Engineering, Ningbo Institute of Technology, Zhejiang University, Ningbo 315100, China*

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We explore the effects of strangeness and Δ resonance in baryonic matter and compact stars within the relativistic-mean-field models. The parameter set 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 Δ-meson couplings. In light of the recent observations of GW170817 with the dimensionless combined tidal deformability $197 \leq \bar{\Lambda} \leq 720$, we find it is essential to include the Δ resonances in compact stars, and small Δ-ρ coupling $g_{\rho\Delta}$ is favored if the mass $2.27^{+0.17}_{-0.15}\ M_\odot$ 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 $\bar{\Lambda}$ is considered to be less than 720 at 90% confidence level [3], while a lower limit with $\bar{\Lambda} \geq 197$ 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 the 10–14 km range [2,5–10]. The combined constraints on the tidal deformability and radii of neutron stars indicate a soft equation of state (EOS), where many covariant density functionals are in jeopardy [11,12]. A possible solution to this problem is to introduce new degrees of freedom, e.g., Δ resonances, hyperons, and deconfined quarks [13]. As one increases the density of nuclear matter, the inevitable emergence of Δ isobars, hyperons, and quarks can soften the EOS 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 ($1.928 \pm 0.017\ M_\odot$) [14,15] and PSR J0348 + 0432 ($2.01 \pm 0.04\ M_\odot$) [16], i.e., the Hyperon Puzzle [17] or Δ Puzzle [18]. Extensive efforts were made to resolve the Hyperon Puzzle [19–37] and Δ Puzzle [38–41]. Nevertheless, with the constrained observable tidal deformability of GW170817 [1,3,4], those solutions may be challenged, especially for the latest observation of a more massive PSR J2215 + 5135 ($2.27^{+0.17}_{-0.15}\ M_\odot$) [42].

To satisfy these stringent observational constraints, we consider the possible existence of both Δ isobars and hyperons in neutron stars. Since relativistic-mean-field (RMF) models [43–50] have been successfully adopted to describe finite (hyper)nuclei [51–61] and baryonic matter [62–69], in this work the EOS of baryonic matter are obtained based on the RMF model. More specifically, we adopt the covariant density functional PKDD [70], while the hyperon-meson couplings are fixed based on our previous investigations on hypernuclei and neutron stars [37,61,71]. For the Δ-meson couplings, as in Ref. [18], we adopt the universal baryon-meson coupling scheme, while a vanishing Δ-ρ 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 Δ isobars in neutron stars.

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

*cjxia@itp.ac.cn

II. THEORETICAL FRAMEWORK

The Lagrangian density of RMF models is given as

$$\begin{aligned}\mathcal{L} = &\sum_{\mathrm{b}}\bar{\psi}_{\mathrm{b}}[i\gamma^{\mu}\partial_{\mu} - m_{\mathrm{b}} - g_{\sigma\mathrm{b}}\sigma - g_{\omega\mathrm{b}}\gamma^{\mu}\omega_{\mu} \\ &- g_{\rho\mathrm{b}}\gamma^{\mu}\boldsymbol{\tau}_{\mathrm{b}}\cdot\boldsymbol{\rho}_{\mu} - \gamma^{\mu}A_{\mu}q_{\mathrm{b}}]\psi_{\mathrm{b}} + \frac{1}{2}\partial_{\mu}\sigma\partial^{\mu}\sigma \\ &- \frac{1}{2}m_{\sigma}^{2}\sigma^{2} - \frac{1}{4}\omega_{\mu\nu}\omega^{\mu\nu} + \frac{1}{2}m_{\omega}^{2}\omega_{\mu}\omega^{\mu} \\ &- \frac{1}{4}\boldsymbol{\rho}_{\mu\nu}\cdot\boldsymbol{\rho}^{\mu\nu} + \frac{1}{2}m_{\rho}^{2}\boldsymbol{\rho}_{\mu}\cdot\boldsymbol{\rho}^{\mu} - \frac{1}{4}A_{\mu\nu}A^{\mu\nu} \\ &+ \sum_{l=e,\mu}\bar{\psi}_{l}[i\gamma^{\mu}\partial_{\mu} - m_{l} + e\gamma^{\mu}A_{\mu}]\psi_{l}, \end{aligned} \tag{1}$$

with the field tensors

$$\begin{aligned}\omega_{\mu\nu} &= \partial_{\mu}\omega_{\nu} - \partial_{\nu}\omega_{\mu}, \\ \boldsymbol{\rho}_{\mu\nu} &= \partial_{\mu}\boldsymbol{\rho}_{\nu} - \partial_{\nu}\boldsymbol{\rho}_{\mu}, \\ A_{\mu\nu} &= \partial_{\mu}A_{\nu} - \partial_{\nu}A_{\mu}. \end{aligned} \tag{2}$$

The included baryons here are nucleons, hyperons (Λ^0, $\Sigma^{+,0,-}$, and $\Xi^{0,-}$), and Δ resonance. To describe the baryon-baryon interactions, the isoscalar-scalar channel (σ), isoscalar-vector channel (ω) and isovector-vector channel ($\boldsymbol{\rho}$) are considered.

Based on the Typel-Wolter ansatz [57], the density dependence of coupling constants $g_{\xi\mathrm{b}}$ ($\xi = \sigma, \omega$) are obtained with

$$g_{\xi\mathrm{b}}(n) = g_{\xi\mathrm{b}}(n_0)a_{\xi}\frac{1 + b_{\xi}(n/n_0 + d_{\xi})^2}{1 + c_{\xi}(n/n_0 + e_{\xi})^2}, \tag{3}$$

where n is the density of nuclear matter with n_0 being the saturation density. Note that a different formula is adopted for the $\boldsymbol{\rho}$ meson, i.e.,

$$g_{\rho\mathrm{b}}(n) = g_{\rho\mathrm{b}}(n_0)\exp[-a_{\rho}(n/n_0 - 1)]. \tag{4}$$

For a system with time-reversal symmetry, the spacelike components of the vector fields ω_{μ} and $\boldsymbol{\rho}_{\mu}$ vanish, leaving only the time components ω_0 and $\boldsymbol{\rho}_0$. Meanwhile, the charge conservation guarantees that only the third component in the isospin space of $\boldsymbol{\rho}_0$ 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 photons can be obtained from the variational procedure.

For the N-N interactions, we adopt the covariant density functional PKDD [70], which gives the saturation density $n_0 = 0.149552\ \mathrm{fm}^{-3}$, saturation energy $E_0 = -16.267$ MeV, incompressibility $K = 262.181$ MeV, and symmetry energy $E_{\mathrm{sym}} = 36.790$ MeV.

Beside nucleons, we also consider the effects of strangeness and Δ resonance, i.e., Λ, Ξ, Σ, and Δ baryons. For the Λ-ω coupling, according to our previous investigations [37], the mass of PSR J0348 + 0432 can only be attained with large values of $g_{\omega\Lambda}$ at fixed Λ potential well depth ($V_{\Lambda} = -29.786$ MeV) in symmetric nuclear matter ($n_p = n_n = n_0/2$). Thus, in this work we suppose $g_{\omega\Lambda} = g_{\omega N}$, which gives $g_{\sigma\Lambda} = 0.878 g_{\sigma N}$. Similarly, we fix the Ξ-meson and Σ-meson couplings with $g_{\omega\Xi} = g_{\omega\Sigma} = g_{\omega N}$, $g_{\sigma\Xi} = 0.844 g_{\sigma N}$, and $g_{\sigma\Sigma} = 0.878 g_{\sigma N}$, which corresponds to the potential well depths $V_{\Xi} = -16.276$ MeV and $V_{\Sigma} = -29.957$ MeV [71]. Note that there is some ambiguity on the potential well depth V_{Σ}, where the (π^-, K^+) reactions on medium-to-heavy nuclei indicate a repulsive potential [72–75] while the observation of a ${}^4_{\Sigma}$He bound state in the (K^-, π^-) reaction favors an attractive potential [76]. For the hyperon-ρ couplings, we take $g_{\rho\Lambda} = 0$ and $g_{\rho\Xi} = g_{\rho\Sigma} = g_{\rho N}$ according to their isospin characters [61,71]. In principle, in consideration of the hyperon-hyperon interactions such as the weakly attractive Λ-Λ interaction, the exchange of σ^* and ϕ mesons between hyperons should also be taken into account. However, according to the recipe of various baryon-meson couplings inspired by the symmetries of the baryon octet [19,21,77,78], taking $g_{\omega\Lambda} = g_{\omega\Xi} = g_{\omega\Sigma} = g_{\omega N}$ and $g_{\phi N} = 0$ indicates vanishing hyperon-ϕ couplings. In such cases, the contributions from σ^* and ϕ mesons are neglected in our Lagrange density (1).

For the Δ-ω and Δ-σ couplings, they are often chosen to be close to the N-ω and N-σ couplings, i.e., $g_{\omega\Delta} \approx g_{\omega N}$ and $g_{\sigma\Delta} \approx g_{\sigma N}$ [18,38,40,79,80], which can be attributed to the similar potential depths of Δ's and nucleons in a nuclear medium according to the data analyses of photoabsorption, electron-nucleus, and pion-nucleus scattering [18]. Slight deviations from those values were also explored in Ref. [81]. However, little is known for the Δ-ρ coupling, while the linear dependence of the onset density $n^{\mathrm{crit}}_{\Delta^-}$ with $g_{\rho\Delta}$ was reported in Refs. [39,40]. Therefore, in this work we adopt the universal baryon-meson coupling scheme with $g_{\omega\Delta} = g_{\omega N}$, $g_{\sigma\Delta} = g_{\sigma N}$, and $g_{\rho\Delta} = g_{\rho N}$. To see the possibility of smaller $g_{\rho\Delta}$, we also study the cases with $g_{\rho\Delta} = 0$. Since the Δ baryons have a Breit-Wigner mass distribution around the centroid mass 1232 MeV with a width of about 120 MeV, the variation of m_{Δ} has sizable effects on baryonic matter and structures of compact stars [39]. In this work, we adopt various Δ masses with $m_{\Delta} = 1112$, 1232, and 1352 MeV.

In Table I, we list properties and coupling constants for baryons other than nucleons in Eq. (1). Meanwhile, it is worth mentioning that the covariant density functional PKDD adopted here is phenomenological, where the nucleon-meson coupling constants are fixed according to the masses of spherical nuclei, the incompressibility, saturation density, and symmetry energy of nuclear matter [70]. In light of the recent developments of microscopic

TABLE I. Strangeness number S, mass M, third component of isospin τ_3, total angular momentum and parity J^P, charge q, and coupling constants $\alpha_\xi = g_{\xi b}/g_{\xi N}$ ($\xi = \sigma$, ω, and ρ) for Λ^0, $\Xi^{0,-}$, and $\Sigma^{+,0,-}$ hyperons and Δ baryons.

	S	M (MeV)	τ_3	J^P	q (e)	α_σ	α_ω	α_ρ
Λ^0	−1	1115.6	0	$(1/2)^+$	0	0.878	1	0
Ξ^0	−2	1314.9	+1	$(1/2)^+$	0	0.844	1	1
Ξ^-	−2	1321.3	−1	$(1/2)^+$	−1	0.844	1	1
Σ^+	−1	1189.4	+1	$(1/2)^+$	+1	0.878	1	1
Σ^0	−1	1192.5	0	$(1/2)^+$	0	0.878	1	1
Σ^-	−1	1197.4	−1	$(1/2)^+$	−1	0.878	1	1
Δ^{++}	0	1232 ± 120	+3	$(3/2)^+$	+2	1	1	0,1
Δ^+	0	1232 ± 120	+1	$(3/2)^+$	+1	1	1	0,1
Δ^0	0	1232 ± 120	0	$(3/2)^+$	0	1	1	0,1
Δ^-	0	1232 ± 120	−3	$(3/2)^+$	−1	1	1	0,1

many-body calculations in describing finite nuclei and nuclear matter starting from realistic nucleon-nucleon interactions [82–88], a more refined adjustment of parameters incorporating those results are necessary. A possible way to reach this in the RMF model is to introduce density-dependent coupling constants derived from self-energies of Dirac-Brueckner calculations of nuclear matter [57,89], which are found decreasing with density and can be reproduced with Eqs. (3) and (4).

Based on the Lagrangian density in Eq. (1), the meson fields are obtained by solving

$$\begin{aligned} m_\sigma^2 \sigma &= -\sum_{\mathrm{b}} g_{\sigma \mathrm{b}} n_{\mathrm{b}}^{\mathrm{s}}, \\ m_\omega^2 \omega_0 &= \sum_{\mathrm{b}} g_{\omega \mathrm{b}} n_{\mathrm{b}}, \\ m_\rho^2 \rho_{0,3} &= \sum_{\mathrm{b}} g_{\rho \mathrm{b}} \tau_{\mathrm{b},3} n_{\mathrm{b}}, \end{aligned} \tag{5}$$

with the number density $n_{\mathrm{b}} = \langle \bar\psi_{\mathrm{b}} \gamma^0 \psi_{\mathrm{b}} \rangle$ and scalar density $n_{\mathrm{b}}^{\mathrm{s}} = \langle \bar\psi_{\mathrm{b}} \psi_{\mathrm{b}} \rangle$ of baryon type b, which are given in Eqs. (8) and (9). Here we take σ, ω_0, and $\rho_{0,3}$ as their mean values.

At zero temperature, with the no-sea approximation, the energy density can be determined by

$$E = \sum_{i=\mathrm{b},l} \varepsilon_i(\nu_i, m_i^*) + \sum_{\xi=\sigma,\omega,\rho} \frac{1}{2} m_\xi^2 \xi^2, \tag{6}$$

in which the kinetic energy density of fermion i is

$$\begin{aligned} \varepsilon_i(\nu_i, m_i) &= \int_0^{\nu_i} \frac{f_i p^2}{2\pi^2} \sqrt{p^2 + m_i^2}\, \mathrm{d}p \\ &= \frac{f_i m_i^4}{16\pi^2} \left[x_i (2x_i^2 + 1) \sqrt{x_i^2 + 1} - \mathrm{arcsh}(x_i) \right]. \end{aligned} \tag{7}$$

Here we have defined $x_i \equiv \nu_i / m_i$ with ν_i being the Fermi momentum and $f_i = 2J_i + 1$ the degeneracy factor of particle type i. Note that in Eq. (6), the baryon effective mass is defined as $m_{\mathrm{b}}^* \equiv m_{\mathrm{b}} + g_{\sigma \mathrm{b}} \sigma$, while the mass of leptons remains constants with $m_l^* \equiv m_l$. The source currents of fermion i are given by

$$n_i = \langle \bar\psi_i \gamma^0 \psi_i \rangle = \frac{f_i \nu_i^3}{6\pi^2}, \tag{8}$$

$$n_i^{\mathrm{s}} = \langle \bar\psi_i \psi_i \rangle = \frac{f_i m_i^3}{4\pi^2} \left[x_i \sqrt{x_i^2 + 1} - \mathrm{arcsh}(x_i) \right]. \tag{9}$$

The chemical potentials for baryons μ_{b} and leptons μ_l are

$$\mu_{\mathrm{b}} = g_{\omega \mathrm{b}} \omega_0 + g_{\rho \mathrm{b}} \tau_{\mathrm{b},3} \rho_{0,3} + \Sigma_{\mathrm{b}}^{\mathrm{R}} + \sqrt{\nu_{\mathrm{b}}^2 + m_{\mathrm{b}}^{*2}}, \tag{10}$$

$$\mu_l = \sqrt{\nu_l^2 + m_l^2}, \tag{11}$$

with the "rearrangement" term

$$\Sigma_{\mathrm{b}}^{\mathrm{R}} = \sum_{\mathrm{b}} \left(\frac{\mathrm{d} g_{\sigma \mathrm{b}}}{\mathrm{d}n} \sigma n_{\mathrm{b}}^{\mathrm{s}} + \frac{\mathrm{d} g_{\omega \mathrm{b}}}{\mathrm{d}n} \omega_0 n_{\mathrm{b}} + \frac{\mathrm{d} g_{\rho \mathrm{b}}}{\mathrm{d}n} \rho_{0,3} \tau_{\mathrm{b},3} n_{\mathrm{b}} \right). \tag{12}$$

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 β-equilibrium condition is satisfied, i.e.,

$$\mu_{\mathrm{b}} = \mu_n - q_{\mathrm{b}} \mu_e, \qquad \mu_\mu = \mu_e. \tag{15}$$

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

$$\frac{\mathrm{d}P}{\mathrm{d}r} = -\frac{GME}{r^2} \frac{(1 + P/E)(1 + 4\pi r^3 P/M)}{1 - 2GM/r}. \tag{16}$$

By solving the TOV equation with the subsidiary condition

$$\frac{\mathrm{d}M(r)}{\mathrm{d}r} = 4\pi E r^2, \tag{17}$$

we get the relation of mass M and radius R of a neutron star. Here, the gravity constant $G = 6.707 \times 10^{-45}$ MeV^{-2}. The tidal deformability of a compact star is extracted from

$$\Lambda = \frac{2k_2}{3}\left(\frac{R}{GM}\right)^5, \tag{18}$$

where k_2 is the second Love number and can be fixed simultaneously with the structures of compact stars [90–92].

III. RESULTS AND DISCUSSIONS

At a 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. Similar to Ref. [18], by varying the parameter a_ρ in Eq. (4), we examine the dependence of onset densities of Δ s and hyperons n_b^{crit} on the symmetry energy slope L, which is fixed by fulfilling $\mu_b|_{\nu_b=0} = \mu_n - q_b\mu_e$. A linear dependence of L (in MeV) on a_ρ is obtained, i.e., $L = 110.3 - 109.5a_\rho$. The variation of $n_{\Lambda^0}^{\mathrm{crit}}$, $n_{\Sigma^-}^{\mathrm{crit}}$, and $n_{\Delta^-}^{\mathrm{crit}}$ are presented in Fig. 1, while the onset densities for other Δ s and hyperons are much larger. For β-stable nuclear matter, the values of μ_e and $\rho_{0,3}$ are increasing with L. Consequently, the obtained $n_{\Lambda^0}^{\mathrm{crit}}$ and $n_{\Sigma^-}^{\mathrm{crit}}$ are decreasing with L while $n_{\Delta^-}^{\mathrm{crit}}$ is increasing, which is consistent with the trends in ([18], Fig. 1). If we take $g_{\rho\Delta} = 0$, the obtained $n_{\Delta^-}^{\mathrm{crit}}$ for $m_\Delta = 1232$ MeV and 1352 MeV are decreasing with L since the contribution of $\rho_{0,3}$ becomes irrelevant. Meanwhile, for the cases with $m_\Delta = 1112$ MeV, $n_{\Delta^-}^{\mathrm{crit}}$ is even smaller than the saturation density. Since μ_e is decreasing with L at subsaturation densities, the corresponding $n_{\Delta^-}^{\mathrm{crit}}$ (black solid curve) increases with L. Finally, it is worth mentioning that the variation of $n_{\Delta^-}^{\mathrm{crit}}$ with respect to L is insignificant comparing with m_Δ due to its relatively larger uncertainty.

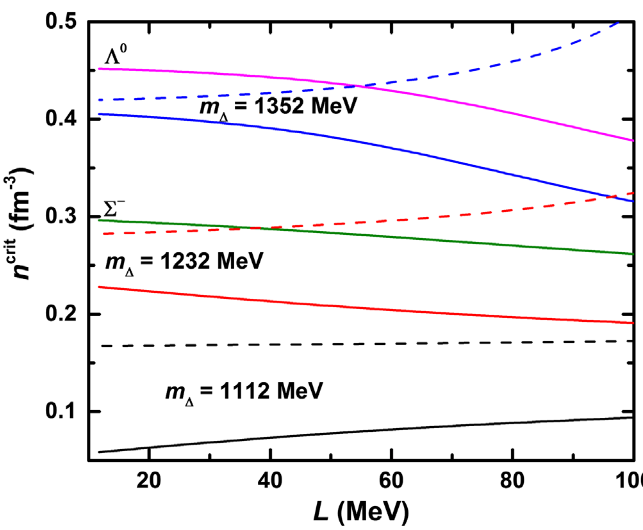

FIG. 1. The onset densities of hyperons and Δ's in β-stable nuclear matter as functions of the symmetry energy slope L. The black ($m_\Delta = 1112$ MeV), red ($m_\Delta = 1232$ MeV), and blue ($m_\Delta = 1352$ MeV) curves correspond to the onset densities of Δ^-, where the solid ones and dashed ones are obtained with $g_{\rho\Delta} = 0$ and $g_{\rho N}$, respectively.

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. 2 and 3. By including Λ^0 in nuclear matter, as indicated by the dashed curves in Fig. 2, the densities of protons and neutrons are slightly reduced on the emergence of Λ^0. If we also include other hyperons such as $\Xi^{0,-}$ and $\Sigma^{+,0,-}$ (dash-dotted curves), since similar potential well depths are adopted for Λ's and Σ's, the Σ^- first appears at $n = 0.27$ fm^{-3} due to the negative charge it carries. In such cases, the number densities of leptons decrease while those of protons increase. Meanwhile, the onset density of Λ^0 is increased from $n = 0.39$ fm^{-3} to 0.46 fm^{-3} due to the inclusion of the negatively charged Σ^-. Since $\Xi^{0,-}$ possess the largest masses, their onset densities are much larger with $n_{\Xi^-}^{\mathrm{crit}} = 1.2$ fm^{-3} and $n_{\Xi^0}^{\mathrm{crit}} > n_{\Xi^-}^{\mathrm{crit}}$, which exceed the density limit of Fig. 2.

The effects of Δ resonances are also studied and the results are shown in Fig. 3. To consider the Breit-Wigner mass distribution of the Δ baryons and the possible in-medium mass shift [39], three masses $m_\Delta = 1112$, 1232, and 1352 MeV are adopted in our calculation. Note that the nucleon effective mass $m_N^* \equiv m_N + g_{\sigma N}\sigma$ 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. 3 we do not present the results with $m_N^* < 0$ at the higher densities. For all Δ baryons, the negatively charged Δ^- appears first as we increase the density. The onset density of Δ^- is found to increase both

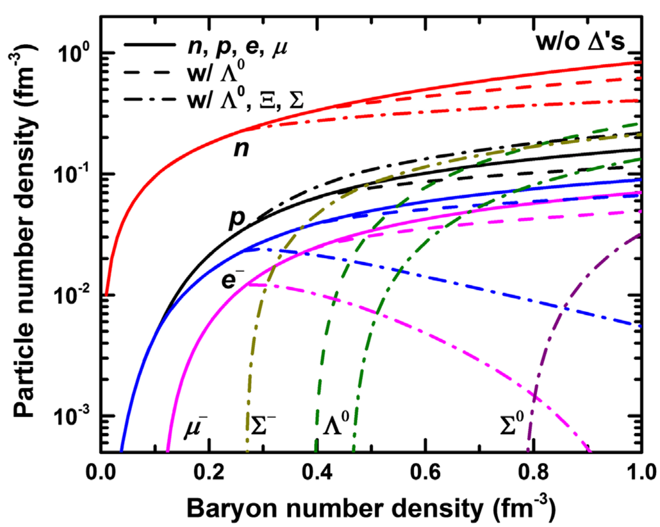

FIG. 2. Particle number densities for baryons and leptons in neutron star matter as functions of the total baryon number density n without Δ resonances.

with m_Δ and $g_{\rho\Delta}$, which is consistent with previous findings [39,40]. For massive Δ's ($m_\Delta = 1352$ MeV), the effects of Δ resonance are insignificant and only Δ^- appears. In the comparison with hyperons, the massive Δ^- appears at larger densities than Σ^-, where the densities of hyperons are similar as the cases in Fig. 2. If we adopt smaller values of m_Δ and $g_{\rho\Delta}$, the effects of Δ resonances become important, where Δ^-, Δ^0, Δ^+, and Δ^{++} appear sequentially as the density increases. 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 hyperon left is Λ^0, which appears at a much larger density $n^{\rm crit}_{\Lambda^0} = 0.74\ \mathrm{fm}^{-3}$. Note that a first-order phase transition from nuclear matter to Δ matter takes place in the density range $n = 0.083$–$0.17\ \mathrm{fm}^{-3}$, where we have shown the corresponding densities in the lower left panel of Fig. 3.

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. 4 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 Δ resonances and adopt the largest mass, i.e., $m_\Delta = 1352$ MeV, the EOS is modified slightly at high density regions since only Δ^- appears at insignificant densities n_{Δ^-}. Moreover, adopting smaller values of m_Δ and $g_{\rho\Delta}$ would result in softer EOS, where in the extreme case of $m_\Delta = 1112$ MeV and $g_{\rho\Delta} = 0$, a softest EOS is obtained for neutron star matter.

Based on the EOS displayed in Fig. 4, the structure of a neutron star can be determined by solving the TOV equation in Eq. (16). For neutron star matter at subsaturation densities ($n \le 0.08\ \mathrm{fm}^{-3}$), we adopt the EOS presented in Refs. [93–95], where the properties of crystalized matter that form the neutron star crust can be well described. In Fig. 5 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 Δ resonances are considered. The obtained results are compared with the observational masses of PSR

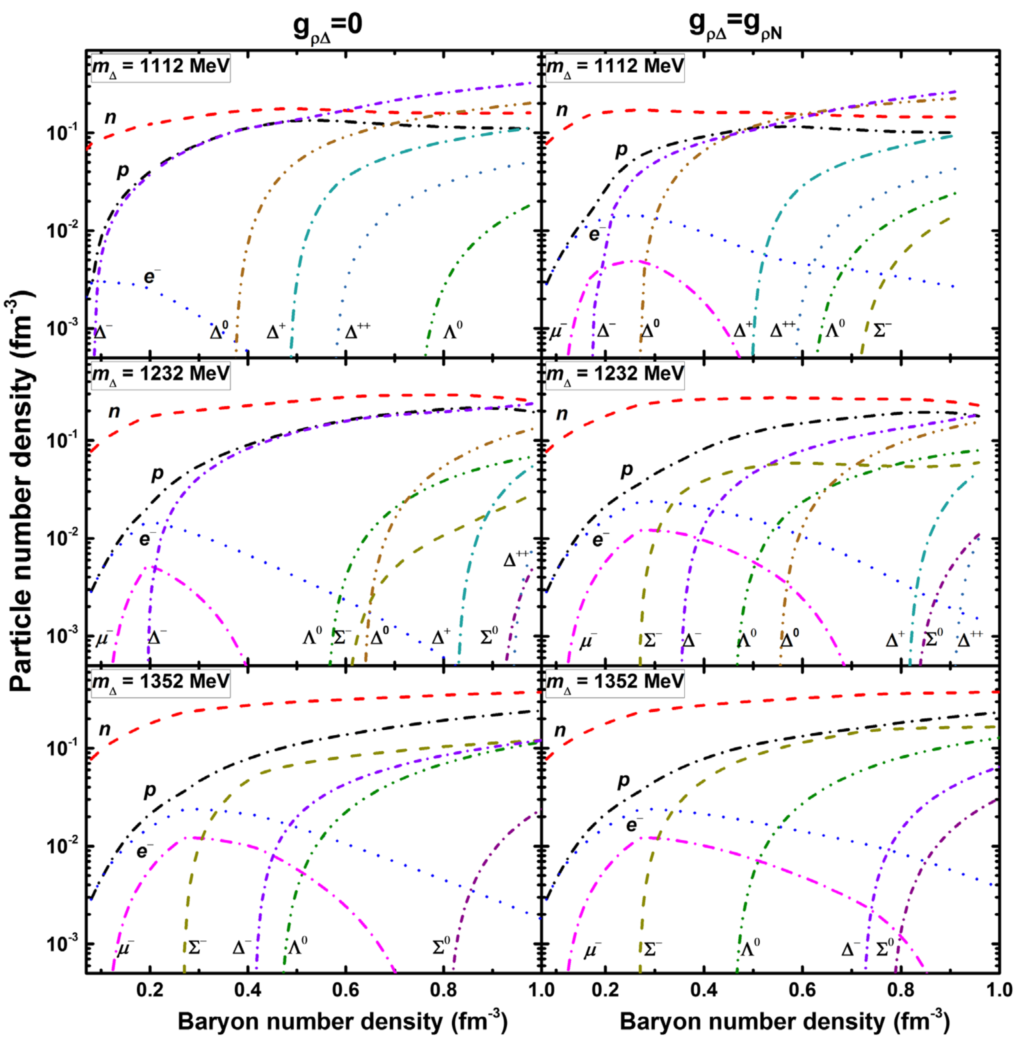

FIG. 3. Same as Fig. 2 but including Δ resonances with $m_\Delta = 1112$, 1232, and 1352 MeV.

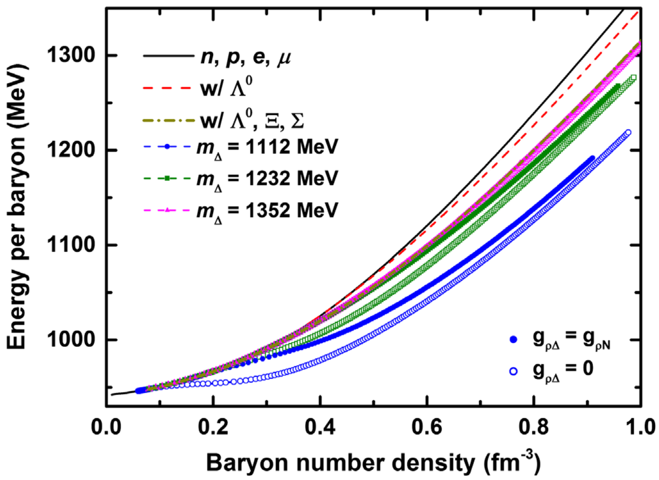

FIG. 4. The energy per baryon of neutron star matter as functions of the baryon number density n. The solid and open symbols are results obtained with $g_{\rho\Delta} = g_{\rho N}$ and 0, respectively. The same convention is adopted for the following figures.

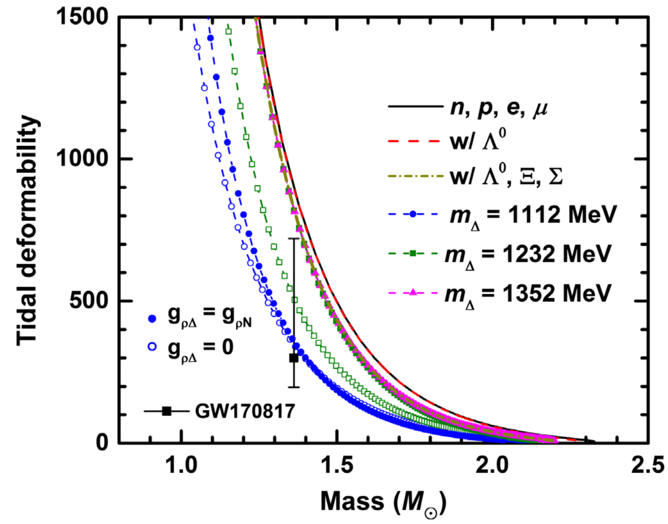

FIG. 6. The tidal deformabilities of compact stars as functions of their masses. The recent constraint obtained with the binary neutron star merger event GW170817 is indicated with the black solid box [1,3,4].

J0348 + 0432 ($2.01 \pm 0.04\ M_\odot$) [16] and PSR J2215 + 5135 ($2.27^{+0.17}_{-0.15}\ M_\odot$) [42]. As we include more degrees of freedom, the maximum mass and radii of compact stars become smaller. For compact stars including Δ resonances, if we adopt $m_\Delta = 1112$ 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 $\rho - \Delta$ couplings, e.g., $g_{\rho\Delta} = 0$. Because of the occurrence of a first-order phase transition at small densities ($n = 0.083$–$0.17\ \mathrm{fm}^{-3}$), the smallest radius with $R = 11.3$ km for the 1.4 $M_\odot$ compact star is obtained, which is consistent with the recent measurements of neutron star radii [2,5–9].

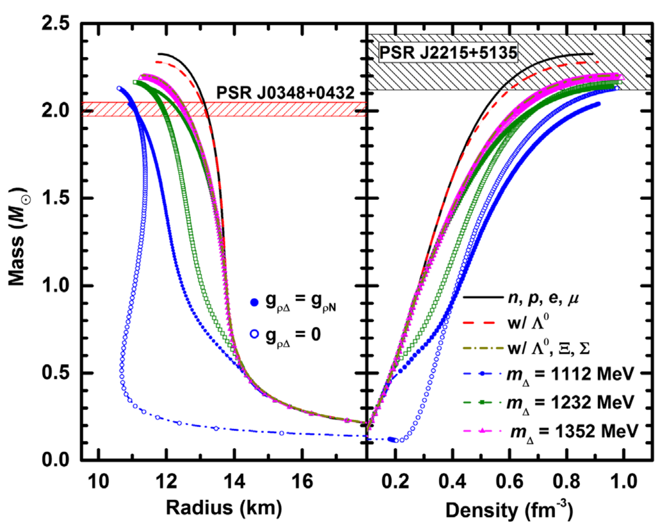

FIG. 5. The obtained mass-radius relations of compact stars including the possible existence of hyperons and Δ resonances. The masses of pulsars PSR J0348 + 0432 ($2.01 \pm 0.04\ M_\odot$) [16] and PSR J2215 + 5135 ($2.27^{+0.17}_{-0.15}\ M_\odot$) [42] are indicated with horizonal bands.

Another important constraint is the tidal deformability of the compact stars, which can be obtained based on Eq. (18). In Fig. 6 we present the tidal deformabilities of compact stars corresponding to those in Fig. 5. The observation of binary neutron star merger event GW170817 has set the dimensionless combined tidal deformability $197 \le \bar{\Lambda} \le 720$ [3,4], which is a mass-weighted linear combination of tidal deformabilities [96]

$$\bar{\Lambda} = \frac{16}{13} \frac{(m_1 + 12 m_2) m_1^4 \Lambda_1 + (m_2 + 12 m_1) m_2^4 \Lambda_2}{(m_1 + m_2)^5}. \quad (19)$$

Since $\bar{\Lambda}$ is insensitive to the mass ratio m_2/m_1 [97], combined with the best measured chirp mass $\mathcal{M} = (m_1 m_2)^{3/5} (m_1 + m_2)^{-1/5} = 1.186 \pm 0.001\ M_\odot$ [3], in Fig. 6 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 Δ resonances have to be included. Meanwhile, as discussed before, a small enough Δ-ρ 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 Δ 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 [70], while the hyperon-meson couplings are fixed based on our previous investigations on hypernuclei and neutron stars [37,71]. For the Δ-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 Δ masses with $m_\Delta = 1112$, 1232, and 1352 MeV. The EOS of neutron star matter become softer once we include new degrees of freedom. By solving the TOV equation with these EOS, we obtained the masses, radii, and tidal deformabilities of the corresponding compact stars. Comparing with the dimensionless combined tidal deformability $197 \leq \tilde{\Lambda} \leq 720$ constrained according to the recent observations of GW170817 [3,4], we find it is essential to include the Δ resonances in compact stars, and the Δ-ρ coupling $g_{\rho\Delta}$ should be small enough if the mass of PSR J2215 + 5135 ($2.27^{+0.17}_{-0.15}\ M_\odot$) [42] is confirmed.

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