Constraints on hyperon couplings from neutron star equations of state*

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Abstract

Based on the constituent quark picture of baryons and taking into account the contributions of isovector and strange mesons, we have developed the extended Zimanyi-Moszkowski model of dense baryon matter for studying neutron star (NS) equations of state (EOSs). Four sets of meson-hyperons coupling constants are investigated. The first is characterized by strong attractive $N\Sigma$ interaction while the others have repulsive $N\Sigma$ interactions. The second is characterized by strong attractive $\Lambda\Lambda$ interaction. The third has weak $\Lambda\Lambda$ but strong attractive $\Sigma\Sigma$ interactions. The last one has much weaker $\Sigma\Sigma$ interaction than the third one. By systematic analyses of the EOSs and mass sequences of NSs, it has been found that the strong attractive $N\Sigma$, $\Lambda\Lambda$ and $\Sigma\Sigma$ interactions are ruled out. The result is consistent to the most recent information on hyperon interactions from the experimental and theoretical investigations of hypernuclei.

1 Introduction

Through the equation of state (EOS) [1,2] the studying the properties of NSs becomes an important subject of nuclear physics. The recent refined nonrelativistic Brueckner-Hartree-Fock (NRBHF) calculations [3,4] however predicted the maximum masses of neutron stars which are lower than the famous canonical value $1.44\,M_\odot$ [5] obtained from the relativistic binary pulsar B1913+16. This has revealed the uncertainties in the models of baryonic matter at high densities and the models of baryon–baryon, especially hyperon-hyperon (YY) interactions. On the other hand, the renewed interests on the YY interactions [6-11] have been raised by the recent discovery of ${}^{6}_{\Lambda\Lambda}$He in the KEK-E373 experiment [12]. Those works however provided the results being inconsistent to each other. This has also revealed the theoretical uncertainties in the YY interactions and the model describing 6-body system.

As seen above the problem of the fundamental baryon-baryon interactions is inseparable from the problem of the models of many-body nuclear system or dense baryon matter. It is necessary to treat both the problems simultaneously in a realistic way as

*This paper is the revised version of CDS ext-2004-134 in which there has been found a bug in numerical code. Although the conclusion is not altered, the EOSs become stiffer. As the results the analyses in section 3 have been refined to some extent. I have also corrected several unclear contexts and mistypes.

possible. Inversely speaking, the physically realistic description of nuclear system is able to constrain the YY interactions to their realistic values. In this respect the relativistic theories are more suitable for dense medium than nonrelativistic ones. However the relativistic Brueckner-Hartree-Fock (RBHF) [13] calculations including all the baryon octets are not feasible at present. Therefore many studies of NS matter [2] have been done by the relativistic mean-field (RMF) theories based or inspired by the Walecka model [14].

The most widely used nonlinear extensions [15] of the Walecka model however are not realistic because they do not take into account the medium dependences of meson-baryon coupling constants, which are realized in the RBHF theory and are essentially important to dense baryon matter. In the RMF models including hyperons, the (hidden) strange mesons σ^* and ϕ [16] are necessary to realistic description of the YY interactions. In this respect the so-called DDRH model [17] is not suitable because of the difficulty to include the explicitly density-dependent $YY\sigma^*(\phi)$ coupling constants. Consequently, there remain only two possibilities of the realistic RMF models for NS matter. They are the Zimanyi-Moszkowski (ZM) model [18] and the quark-meson coupling (QMC) model [19]. In contrast to the DDRH model, these take into account the density dependences of the meson-baryon coupling constants implicitly through the effective masses of baryons in the medium.

The investigations of NSs by the ZM and QMC models including the strange mesons were performed in Refs. [20] and [21]. Although they predicted reasonable maximum NS masses $M_{ns} \approx 1.5\, M_{\odot}$ that are close to the canonical value $M_{ns} = 1.44\, M_{\odot}$, they are never physically realistic. The QMC model [21] adopted the baryon-baryon interactions deduced purely from SU(6) symmetry of baryon octet. It is believed that the realistic values of the interactions deviate from those values. On the other hand, the original ZM model is not a realistic model of nuclear matter because it cannot reproduce strong spin-orbit potentials of normal nuclei. The straightforward extension of the ZM model in Ref. [20] is therefore not physically reasonable.

The defect of the ZM model has been eliminated in Refs. [22,23] by modifying it based on the constituent quark picture of nucleons. This modified ZM (MZM) model can be extended to hyperons unambiguously in contrast to the original ZM model [24,25]. Then the MZM model was applied to NS matter in Refs. [26,27], but the strange mesons were not considered. The extended ZM (EZM) model taking into account strange mesons has been first developed in Ref. [28] to investigate the isoscalar strange hadronic matter.

In the present work we apply the EZM model to NS matter by further extending it to introduce the isovector mesons δ and ρ. In the next section we derive the effective renormalized meson-baryon coupling constants in the medium, which are the essential ingredients of the EZM model. Then the RMF theory of the NS matter is developed. In section 3 the properties of NS matter are calculated. We investigate several possibilities of the YY interactions and compare the EZM model with the other RMF models. We summarize our investigation and draw conclusions in section 4.

2 Formalism

In this work we consider the contributions of the isoscalar mesons σ and ω, their strange counterparts σ^* and ϕ mesons and the isovector mesons δ and ρ. Their masses are taken to be $m_\sigma = 550\,\text{MeV}$, $m_\omega = 783\,\text{MeV}$, $m_{\sigma^*} = 975\,\text{MeV}$, $m_\phi = 1020\,\text{MeV}$, $m_\rho = 769\,\text{MeV}$ and $m_\delta = 983\,\text{MeV}$.The masses of baryons are assumed to be $M_N = 938.9\,\text{MeV}$, $M_\Lambda = 1115.6\,\text{MeV}$, $M_\Sigma = 1193.05\,\text{MeV}$ and $M_\Xi = 1318.1\,\text{MeV}$. The essential concept of the present version of the EZM model is the same as the previous works [22,23,26-28].

2.1 Effective renormalized coupling constants

Because the nucleons have no strange contents, the renormalized meson (Π)-nucleon (N) coupling constants $g^*_{NN\Pi}$ of the EZM model in isospin-asymmetric medium are the same as those of the MZM model derived in Ref. [23]. Here we show only the results for convenience:

$$g^*_{pp\,\sigma(\omega)} = \left[(1-\lambda_N) + \lambda_N m^*_p\right] g_{NN\sigma(\omega)}, \tag{1}$$

$$g^*_{nn\sigma(\omega)} = \left[(1-\lambda_N) + \lambda_N m^*_n\right] g_{NN\sigma(\omega)}, \tag{2}$$

$$g^*_{pp\,\delta(\rho)} = \left[(1-\lambda_N) + \lambda_N \left(2\,m^*_n - m^*_p\right)\right] g_{NN\delta(\rho)}, \tag{3}$$

$$g^*_{nn\,\delta(\rho)} = \left[(1-\lambda_N) + \lambda_N \left(2\,m^*_p - m^*_n\right)\right] g_{NN\delta(\rho)}, \tag{4}$$

where $g_{NN\Pi}$ is the free coupling constant and the renormalization constant λ_N is

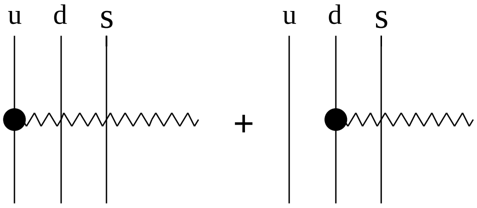
$$\lambda_N = 1/3. \tag{5}$$

The quantities $m^*_{p\,(n)}$ are the ratios of the effective nucleon masses $M^*_{p\,(n)}$ in medium to the free mass:

$$m^*_{p\,(n)} = M^*_{p\,(n)}/M_N = \left(M_N + S_{p\,(n)}\right)/M_N, \tag{6}$$

where $S_{p(n)}$ are the scalar potentials of protons (neutrons). M^*_p and M^*_n are different from each other owing to the isovector scalar mean-field by δ meson. Because the renormalized coupling constants depend on the effective masses, they are determined self-consistently in nuclear medium.

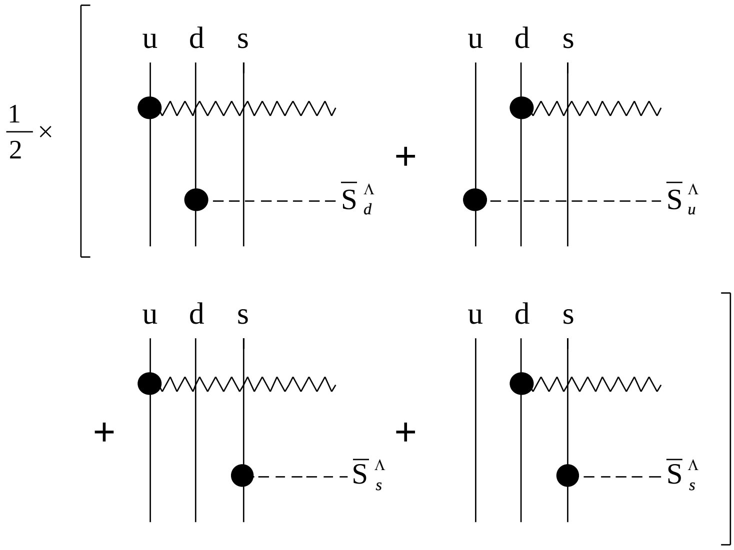
Next, using the intuitive schematic methods, we consider the renormalized meson-hyperons coupling constants of the EZM model that are the extensions of those derived in Ref. [28] to take into account the isovector mesons. First the Λ hyperon is considered. In the constituent quark model (QCM) of baryons, the free $\Lambda\Lambda\sigma$ or $\Lambda\Lambda\omega$ coupling is schematically written by



or expressed by

$$g_{\Lambda\Lambda\sigma(\omega)} = 2\, g^{\Lambda}_{qq\sigma(\omega)}. \tag{7}$$

The wavy lines denote σ or ω mesons and q is u or d quark. Although the other two quarks disconnected from the wavy lines are the spectators, they are also embedded in mean-fields if the Λ hyperons compose baryon matter. In the RMF model of baryon matter, the mass of Λ hyperon in the medium must be reduced by the scalar mean-fields as the effective masses of nucleons in Eq. (6). In the QCM this means that the masses of u, d and s quarks are also reduced by their scalar potentials. Therefore we can consider the following medium correction to $g_{\Lambda\Lambda\sigma(\omega)}$:



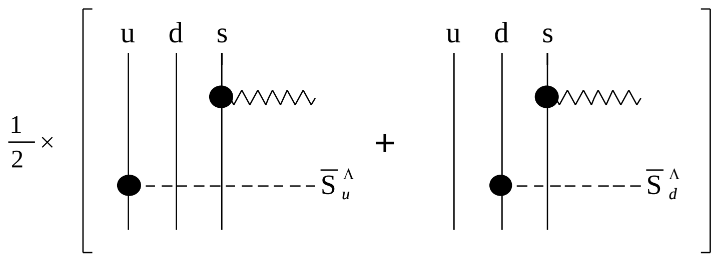
The dashed lines are the effects of the mean-fields on quarks defined by

$$\bar{S}^{Y}_{u,d,s} = S^{Y}_{u,d,s}/M_Y, \tag{8}$$

where S^Y_u, S^Y_d and S^Y_s are the scalar potentials of u, d and s quarks in the hyperon Y. Adding the above correction to Eq. (7), we have the renormalized (or effective) $\Lambda\Lambda\sigma(\omega)$ coupling constant $g^*_{\Lambda\Lambda\sigma(\omega)}$ in the medium,

$$g^*_{\Lambda\Lambda\sigma(\omega)} = g_{\Lambda\Lambda\sigma(\omega)} + \frac{1}{2}\left(\bar{S}^{\Lambda}_u + \bar{S}^{\Lambda}_d + 2\bar{S}^{\Lambda}_s\right) g^{\Lambda}_{qq\sigma(\omega)} = \left[1 + \frac{1}{4}\left(\bar{S}^{\Lambda}_u + \bar{S}^{\Lambda}_d\right) + \frac{1}{2}\bar{S}^{\Lambda}_s\right] g_{\Lambda\Lambda\sigma(\omega)}. \tag{9}$$

Similarly, the medium correction to $\Lambda\Lambda\sigma^*(\phi)$ coupling constant in the EZM model is depicted by



Thus the renormalized coupling constant is given by

$$g^*_{\Lambda\Lambda\sigma^*(\phi)} = g_{\Lambda\Lambda\sigma^*(\phi)} + \frac{1}{2}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right) g^\Lambda_{ss\,\sigma^*(\phi)} = \left[1 + \frac{1}{2}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right)\right] g_{\Lambda\Lambda\sigma^*(\phi)}, \tag{10}$$

where $g_{\Lambda\Lambda\sigma^*(\phi)} = g^{\ \Lambda}_{ss\sigma^*(\phi)}$ is used. Because the Λ is charge neutral,

$$g^*_{\Lambda\Lambda\delta(\rho)} = 0. \tag{11}$$

The renormalized meson-Σ^0 coupling constants have the same form as those of Λ:

$$g^*_{\Sigma^0\Sigma^0\sigma(\omega)} = \left[1 + \frac{1}{4}\left(\bar{S}_u^{\Sigma^0} + \bar{S}_d^{\Sigma^0}\right) + \frac{1}{2}S_s^{\Sigma^0}\right] g_{\Sigma\Sigma\sigma(\omega)}, \tag{12}$$

$$g^*_{\Sigma^0\Sigma^0\sigma^*(\phi)} = \left[1 + \frac{1}{2}\left(\bar{S}_u^{\Sigma^0} + \bar{S}_d^{\Sigma^0}\right)\right] g_{\Sigma\Sigma\sigma^*(\phi)}, \tag{13}$$

$$g^*_{\Sigma^0\Sigma^0\delta(\rho)} = 0. \tag{14}$$

Clearly, for the charged Σ's we have

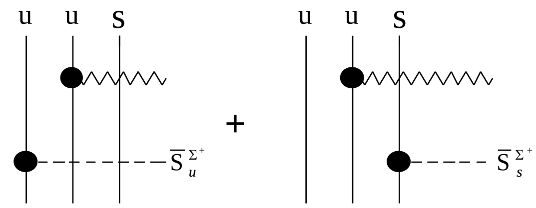
$$g^*_{\Sigma^+\Sigma^+\sigma(\omega)} = \left[1 + \frac{1}{2}\left(\bar{S}_u^{\Sigma^+} + \bar{S}_s^{\Sigma^+}\right)\right] g_{\Sigma\Sigma\sigma(\omega)}, \tag{15}$$

$$g^*_{\Sigma^+\Sigma^+\sigma^*(\phi)} = \left(1 + \bar{S}_u^{\,\Sigma^+}\right) g_{\Sigma\Sigma\sigma^*(\phi)}, \tag{16}$$

$$g^*_{\Sigma^-\Sigma^-\sigma(\omega)} = \left[1 + \frac{1}{2}\left(\bar{S}_d^{\Sigma^-} + \bar{S}_s^{\Sigma^-}\right)\right] g_{\Sigma\Sigma\sigma(\omega)}, \tag{17}$$

$$g^*_{\Sigma^-\Sigma^-\sigma^*(\phi)} = \left(1 + \bar{S}_d^{\Sigma^-}\right) g_{\Sigma\Sigma\sigma^*(\phi)}. \tag{18}$$

Next, the medium correction to $\Sigma^+\Sigma^+\delta(\rho)$ coupling constant in the EZM model is depicted by



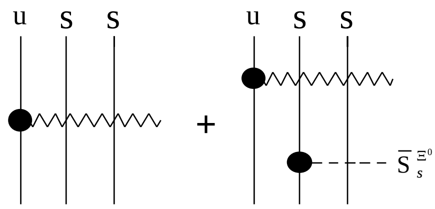
Thus the renormalized coupling constant becomes

$$g^*_{\Sigma^+\Sigma^+\delta(\rho)} = g_{\Sigma\Sigma\delta(\rho)} + \left(\bar{S}_u^{\Sigma^+} + \bar{S}_s^{\Sigma^+}\right) g^{\Sigma}_{qq\delta(\rho)} = \left[1 + \frac{1}{2}\left(\bar{S}_u^{\Sigma^+} + \bar{S}_s^{\Sigma^+}\right)\right] g_{\Sigma\Sigma\delta(\rho)}, \quad (19)$$

where $g_{\Sigma\Sigma\delta(\rho)} = 2\, g^{\Sigma}_{qq\delta(\rho)}$ is used. Clearly, for Σ^- we have

$$g^*_{\Sigma^-\Sigma^-\delta(\rho)} = g_{\Sigma\Sigma\delta(\rho)} + \left(\bar{S}_d^{\Sigma^-} + \bar{S}_s^{\Sigma^-}\right) g^{\Sigma}_{qq\delta(\rho)} = \left[1 + \frac{1}{2}\left(\bar{S}_d^{\Sigma^-} + \bar{S}_s^{\Sigma^-}\right)\right] g_{\Sigma\Sigma\delta(\rho)}. \quad (20)$$

Then we consider the Ξ hyperons. In the EZM model the renormalized $\Xi^0\Xi^0\sigma(\omega, \delta, \rho)$ coupling in the mean-fields is depicted by

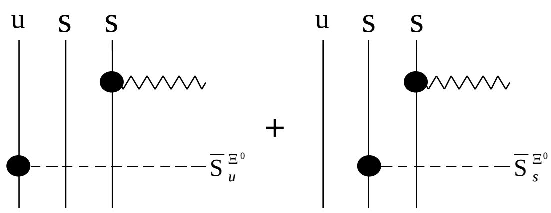


or expressed by

$$g^*_{\Xi^0\Xi^0\sigma,\omega,\delta,\rho} = g_{\Xi\Xi\sigma,\omega,\delta,\rho} + \bar{S}_s^{\Xi^0} g^{\Xi}_{uu\sigma,\omega,\delta,\rho} = \left(1 + \bar{S}_s^{\Xi^0}\right) g_{\Xi\Xi\sigma,\omega,\delta,\rho}, \quad (21)$$

where $g_{\Xi\Xi\sigma,\omega,\delta,\rho} = g^{\Xi}_{qq\sigma,\omega,\delta,\rho}$ is used.

The medium correction to $\Xi^0\Xi^0\sigma^*(\phi)$ coupling in the EZM model is depicted by



Thus the renormalized coupling constant is

$$g^*_{\Xi^0\Xi^0\sigma^*(\phi)} = g_{\Xi\Xi\sigma^*(\phi)} + \left(\bar{S}_u^{\Xi^0} + \bar{S}_s^{\Xi^0}\right) g^{\Xi}_{ss\sigma^*(\phi)} = \left[1 + \frac{1}{2}\left(\bar{S}_u^{\Xi^0} + \bar{S}_s^{\Xi^0}\right)\right] g_{\Xi\Xi\sigma^*(\phi)}, \quad (22)$$

where $g_{\Xi\Xi\sigma^*(\phi)} = 2\, g^{\Xi}_{ss\sigma^*(\phi)}$ is used.

Clearly, for Ξ^- we have

$$g^*_{\Xi^-\Xi^-\sigma,\omega,\delta,\rho} = \left(1 + \bar{S}_s^{\Xi^-}\right) g_{\Xi\Xi\sigma,\omega,\delta,\rho}, \tag{23}$$

$$g^*_{\Xi^-\Xi^-\sigma^*(\phi)} = \left[1 + \frac{1}{2}\left(\bar{S}_d^{\Xi^-} + \bar{S}_s^{\Xi^-}\right)\right] g_{\Xi\Xi\sigma^*(\phi)}. \tag{24}$$

In Appendix the effective meson-baryon couplings will be derived again by more stringent method using the relativistic SU(6) model of baryon octet.

2.2 The RMF model of NS matter

The Lagrangian of NS matter in the RMF model has the common form:

$$\begin{aligned}\mathcal{L} \; = & \sum_{\substack{B=p,n,\Lambda,\Sigma^+,\\ \Sigma^0,\Sigma^-,\Xi^0,\Xi^-}} \bar{\psi}_B \left(\not{p} - M_B^* - \gamma^0 V_B\right) \psi_B + \sum_{l=e^-,\mu^-} \bar{\psi}_l \left(\not{p} - m_l\right) \psi_l \\ & - \frac{1}{2} m_\sigma^2 \langle\sigma\rangle^2 + \frac{1}{2} m_\omega^2 \langle\omega_0\rangle^2 - \frac{1}{2} m_\delta^2 \langle\delta_3\rangle^2 + \frac{1}{2} m_\rho^2 \langle\rho_{03}\rangle^2 - \frac{1}{2} m_{\sigma^*}^2 \langle\sigma^*\rangle^2 + \frac{1}{2} m_\phi^2 \langle\phi_0\rangle^2 ,\end{aligned} \tag{25}$$

where ψ_B and ψ_l are the Dirac fields of baryons and leptons, $\langle\sigma\rangle$, $\langle\omega_0\rangle$, $\langle\delta_3\rangle$, $\langle\rho_{03}\rangle$, $\langle\sigma^*\rangle$ and $\langle\phi_0\rangle$ are the mean-fields, m_l is the mass of each lepton. The energy density is

$$\begin{aligned}\mathcal{E} \; = & \sum_B \left(\langle E_k^*\rangle_B + V_B\right) \rho_B + \sum_l \langle E_k\rangle_l \rho_l + \frac{1}{2} m_\sigma^2 \langle\sigma\rangle^2 + \frac{1}{2} m_\delta^2 \langle\delta_3\rangle^2 \\ & + \frac{1}{2} m_{\sigma^*}^2 \langle\sigma^*\rangle^2 - \frac{1}{2} m_\omega^2 \langle\omega_0\rangle^2 - \frac{1}{2} m_\rho^2 \langle\rho_{03}\rangle^2 - \frac{1}{2} m_\phi^2 \langle\phi_0\rangle^2 ,\end{aligned} \tag{26}$$

where $\langle E_k^*\rangle_B$ and $\langle E_k\rangle_l$ are the average kinetic energies of baryons and leptons, and ρ_B and ρ_l are their vector densities.

The effective masses M_B^* of the baryons are

$$M_B^* = M_B + S_B. \tag{27}$$

Using the renormalized meson-baryon coupling constants, the scalar potentials S_B are given by

$$S_B = -g^*_{BB\sigma} \langle\sigma\rangle - g^*_{BB\delta} \langle\delta_3\rangle I_{3B} - g^*_{BB\sigma^*} \langle\sigma^*\rangle , \tag{28}$$

where $I_{3B} = \{1, -1, 0, 1, 0, -1, 1, -1\}$ for $B = \{p, n, \Lambda, \Sigma^+, \Sigma^0, \Sigma^-, \Xi^0, \Xi^-\}$. On the other hand, the vector potentials V_B are given by

$$V_B = g^*_{BB\omega} \langle\omega_0\rangle + g^*_{BB\rho} \langle\rho_{03}\rangle I_{3B} + g^*_{BB\phi} \langle\phi_0\rangle . \tag{29}$$

Etremizing Eq. (26) by the vector mean-fields, we have

$$\langle \omega_0 \rangle = \sum_B \frac{g^*_{BB\omega}}{m^2_\omega}\, \rho_B, \tag{30}$$

$$\langle \rho_{03} \rangle = \sum_B I_{3B}\, \frac{g^*_{BB\rho}}{m^2_\rho}\, \rho_B, \tag{31}$$

$$\langle \phi_0 \rangle = \sum_Y \frac{g^*_{YY\phi}}{m^2_\phi}\, \rho_Y. \tag{32}$$

Thus the energy density becomes

$$\begin{aligned} \mathcal{E} \;=\; & \frac{1}{4}\sum_B \left(3E^*_{BF}\rho_B + M^*_B \rho_{BS}\right) + \frac{1}{4}\sum_l \left(3E_{lF}\rho_l + m_l \rho_{lS}\right) + \frac{1}{2}\, m^2_\sigma \left\langle \sigma \right\rangle^2 + \frac{1}{2}\, m^2_\omega \left\langle \omega_0 \right\rangle^2 \\ & + \frac{1}{2}\, m^2_\delta \left\langle \delta_3 \right\rangle^2 + \frac{1}{2}\, m^2_\rho \left\langle \rho_{03} \right\rangle^2 + \frac{1}{2}\, m^2_{\sigma^*} \left\langle \sigma^* \right\rangle^2 + \frac{1}{2}\, m^2_\phi \left\langle \phi_0 \right\rangle^2, \end{aligned} \tag{33}$$

where ρ_{BS} and ρ_{lS} are the scalar densities of baryons and leptons. The pressure becomes

$$\begin{aligned} P \;=\; & \frac{1}{4}\sum_B \left(E^*_{BF}\rho_B - M^*_B \rho_{BS}\right) + \frac{1}{4}\sum_l \left(E_{lF}\rho_l - m_l \rho_{lS}\right) - \frac{1}{2}\, m^2_\sigma \left\langle \sigma \right\rangle^2 + \frac{1}{2}\, m^2_\omega \left\langle \omega_0 \right\rangle^2 \\ & - \frac{1}{2}\, m^2_\delta \left\langle \delta_3 \right\rangle^2 + \frac{1}{2}\, m^2_\rho \left\langle \rho_{03} \right\rangle^2 - \frac{1}{2}\, m^2_{\sigma^*} \left\langle \sigma^* \right\rangle^2 + \frac{1}{2}\, m^2_\phi \left\langle \phi_0 \right\rangle^2. \end{aligned} \tag{34}$$

Equation (33) is a function of three independent fundamental quantities m^*_p, m^*_n and m^*_Λ, the effective masses of proton, neutron and Λ. The scalar mean-fields, the scalar potentials and the effective masses of the other hyperons are expressed by them. The expressions of $\langle \sigma \rangle$ and $\langle \delta_3 \rangle$ have been derived in Ref. [23]. The results are

$$\langle \sigma \rangle = \frac{\left(1 - m^*_p\right) g^*_{nn\delta} + \left(1 - m^*_n\right) g^*_{pp\delta}}{C_N} M_N, \tag{35}$$

$$\langle \delta_3 \rangle = \frac{\left(m^*_n - 1\right) g^*_{pp\sigma} - \left(m^*_p - 1\right) g^*_{nn\sigma}}{C_N} M_N, \tag{36}$$

where

$$C_N = \; g^*_{pp\sigma}\, g^*_{nn\delta} + \; g^*_{nn\sigma}\, g^*_{pp\delta}. \tag{37}$$

The effective mass of Λ is defined by

$$m^*_\Lambda = 1 + \bar{S}_\Lambda, \tag{38}$$

$$\bar{S}_\Lambda = \bar{S}^\Lambda_u + \bar{S}^\Lambda_d + \bar{S}^\Lambda_s, \tag{39}$$

where from Eqs. (9) and (10)

$$\bar{S}_u^\Lambda + \bar{S}_d^\Lambda = -\frac{g^*_{\Lambda\Lambda\sigma}\langle\sigma\rangle}{M_\Lambda} = -\left[1 + \frac{1}{4}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right) + \frac{1}{2}\bar{S}_s^\Lambda\right]\bar{\sigma}_\Lambda, \tag{40}$$

$$\bar{S}_s^\Lambda = -\frac{g^*_{\Lambda\Lambda\sigma^*}\langle\sigma^*\rangle}{M_\Lambda} = -\left[1 + \frac{1}{2}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right)\right]\bar{\sigma}^*_\Lambda. \tag{41}$$

We have introduced

$$\bar{\sigma}_Y \equiv \frac{g_{YY\sigma}}{M_Y}\langle\sigma\rangle, \quad \bar{\sigma}^*_Y \equiv \frac{g_{YY\sigma^*}}{M_Y}\langle\sigma^*\rangle. \tag{42}$$

Solving Eqs. (40) and (41),

$$\frac{1}{2}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right) = -\frac{1-(1/2)\,\bar{\sigma}^*_\Lambda}{C_\Lambda}\,\bar{\sigma}_\Lambda, \tag{43}$$

$$\bar{S}_s^\Lambda = -\frac{2-(1/2)\,\bar{\sigma}_\Lambda}{C_\Lambda}\,\bar{\sigma}^*_\Lambda, \tag{44}$$

where

$$C_\Lambda = 2 + \frac{1}{2}\left(1-\bar{\sigma}^*_\Lambda\right)\bar{\sigma}_\Lambda. \tag{45}$$

Substituting Eqs. (43) and (44) into (39),

$$\bar{\sigma}^*_\Lambda = \frac{2\left(1-m^*_\Lambda\right) - \left[2 - \frac{1}{2}\left(1-m^*_\Lambda\right)\right]\bar{\sigma}_\Lambda}{2 - \frac{1}{2}\left(2+m^*_\Lambda\right)\bar{\sigma}_\Lambda}. \tag{46}$$

Because $\bar{\sigma}_\Lambda$ or the mean-field $\langle\sigma\rangle$ is a function of m^*_p and m^*_n, the $\bar{\sigma}^*_\Lambda$ or the mean-field $\langle\sigma^*\rangle$ is a function of m^*_p, m^*_n and m^*_Λ. The renormalized coupling constants of Λ are also expressed by m^*_p, m^*_n and m^*_Λ through

$$1 + \frac{1}{4}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right) + \frac{1}{2}\bar{S}_s^\Lambda = \frac{2-\bar{\sigma}^*_\Lambda}{C_\Lambda}, \tag{47}$$

$$1 + \frac{1}{2}\left(\bar{S}_u^\Lambda + \bar{S}_d^\Lambda\right) = \frac{2-(1/2)\,\bar{\sigma}_\Lambda}{C_\Lambda}. \tag{48}$$

Because the similar expressions to Eqs. (40) and (41) hold for Σ^0, its effective mass is expressed by m^*_p, m^*_n and m^*_Λ through

$$m^*_{\Sigma^0} = 1 + \bar{S}_{\Sigma^0} = 1 + \bar{S}_u^{\Sigma^0} + \bar{S}_d^{\Sigma^0} + \bar{S}_s^{\Sigma^0} = \frac{2-(3/2)\,\bar{\sigma}_\Sigma - 2\,\bar{\sigma}^*_\Sigma + \bar{\sigma}_\Sigma\bar{\sigma}^*_\Sigma}{C_{\Sigma^0}}, \tag{49}$$

where

$$C_{\Sigma^0} = 2 + \frac{1}{2}\left(1-\bar{\sigma}^*_\Sigma\right)\bar{\sigma}_\Sigma. \tag{50}$$

The renormalized coupling constants of Σ^0 are also expressed by m^*_p, m^*_n and m^*_Λ through

$$1 + \frac{1}{4}\left(\bar{S}_u^{\Sigma^0} + \bar{S}_d^{\Sigma^0}\right) + \frac{1}{2}\,\bar{S}_s^{\Sigma^0} = \frac{2 - \bar{\sigma}_\Sigma^*}{C_{\Sigma^0}}, \tag{51}$$

$$1 + \frac{1}{2}\left(\bar{S}_u^{\Sigma^0} + \bar{S}_d^{\Sigma^0}\right) = \frac{2 - (1/2)\,\bar{\sigma}_\Sigma}{C_{\Sigma^0}}. \tag{52}$$

The results for Σ^+ can be obtained by replacing $\bar{\sigma}_\Sigma$ in Eqs. (49)-(52) with $\bar{\sigma}_\Sigma + \bar{\delta}_\Sigma$, where

$$\bar{\delta}_Y \equiv \frac{g_{YY\delta}}{M_Y}\,\langle\delta_3\rangle\,. \tag{53}$$

Therefore the effective mass of Σ^+ becomes

$$m_{\Sigma^+}^* = 1 + \bar{S}_{\Sigma^+} = 1 + 2\,\bar{S}_u^{\Sigma^+} + \bar{S}_s^{\Sigma^+} = \frac{2 - (3/2)\left(\bar{\sigma}_\Sigma + \bar{\delta}_\Sigma\right) - 2\,\bar{\sigma}_\Sigma^* + \left(\bar{\sigma}_\Sigma + \bar{\delta}_\Sigma\right)\bar{\sigma}_\Sigma^*}{C_{\Sigma^+}}, \tag{54}$$

and the renormalized coupling constants are determined through

$$1 + \frac{1}{2}\left(\bar{S}_u^{\Sigma^+} + \bar{S}_s^{\Sigma^+}\right) = \frac{2 - \bar{\sigma}_\Sigma^*}{C_{\Sigma^+}}, \tag{55}$$

$$1 + \bar{S}_u^{\Sigma^+} = \frac{2 - (1/2)\left(\bar{\sigma}_\Sigma + \bar{\delta}_\Sigma\right)}{C_{\Sigma^+}}, \tag{56}$$

where

$$C_{\Sigma^+} = 2 + \frac{1}{2}\left(1 - \bar{\sigma}_\Sigma^*\right)\left(\bar{\sigma}_\Sigma + \bar{\delta}_\Sigma\right). \tag{57}$$

For Σ^- the contributions of $\bar{\delta}_\Sigma$ are negative. Thus the effective mass of Σ^- is

$$m_{\Sigma^-}^* = 1 + \bar{S}_{\Sigma^-} = 1 + 2\,\bar{S}_d^{\Sigma^-} + \bar{S}_s^{\Sigma^-} = \frac{2 - (3/2)\left(\bar{\sigma}_\Sigma - \bar{\delta}_\Sigma\right) - 2\,\bar{\sigma}_\Sigma^* + \left(\bar{\sigma}_\Sigma - \bar{\delta}_\Sigma\right)\bar{\sigma}_\Sigma^*}{C_{\Sigma^-}}, \tag{58}$$

and the renormalized coupling constants are determined through

$$1 + \frac{1}{2}\left(\bar{S}_d^{\Sigma^-} + \bar{S}_s^{\Sigma^-}\right) = \frac{2 - \bar{\sigma}_\Sigma^*}{C_{\Sigma^-}}, \tag{59}$$

$$1 + \bar{S}_d^{\Sigma^-} = \frac{2 - (1/2)\left(\bar{\sigma}_\Sigma - \bar{\delta}_\Sigma\right)}{C_{\Sigma^-}}, \tag{60}$$

where

$$C_{\Sigma^-} = 2 + \frac{1}{2}\left(1 - \bar{\sigma}_\Sigma^*\right)\left(\bar{\sigma}_\Sigma - \bar{\delta}_\Sigma\right). \tag{61}$$

Next, the effective mass of Ξ^0 is defined by

$$m_{\Xi^0}^* = 1 + \bar{S}_{\Xi^0}, \tag{62}$$

$$\bar{S}_{\Xi^0} = \bar{S}_u^{\Xi^0} + 2\,\bar{S}_s^{\Xi^0}, \tag{63}$$

where from Eqs.(21) and (22)

$$\bar{S}_u^{\Xi^0} = -\frac{g^*_{\Xi^0\Xi^0\sigma}\,\langle\sigma\rangle}{M_\Xi} - \frac{g^*_{\Xi^0\Xi^0\delta}\,\langle\delta_3\rangle}{M_\Xi} = -\left(1+\bar{S}_s^{\Xi^0}\right)\left(\bar{\sigma}_\Xi + \bar{\delta}_\Xi\right), \tag{64}$$

$$2\,\bar{S}_s^{\Xi^0} = -\frac{g^*_{\Xi^0\Xi^0\sigma^*}\,\langle\sigma^*\rangle}{M_\Xi} = -\left[1+\frac{1}{2}\left(\bar{S}_u^{\Xi^0}+\bar{S}_s^{\Xi^0}\right)\right]\bar{\sigma}^*_\Xi. \tag{65}$$

Solving Eqs. (64) and (65),

$$\bar{S}_u^{\Xi^0} = -\frac{2-(1/2)\,\bar{\sigma}^*_\Xi}{C_{\Xi^0}}\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right), \tag{66}$$

$$\bar{S}_s^{\Xi^0} = -\frac{1-\frac{1}{2}\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right)}{C_{\Xi^0}}\,\bar{\sigma}^*_\Xi, \tag{67}$$

where

$$C_{\Xi^0} = 2+\frac{1}{2}\left[1-\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right)\right]\bar{\sigma}^*_\Xi. \tag{68}$$

Substituting Eqs. (64) and (65) into (63), the effective mass of Ξ^0 becomes

$$m^*_{\Xi^0} = \frac{2-2\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right)-(3/2)\,\bar{\sigma}^*_\Xi+\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right)\bar{\sigma}^*_\Xi}{C_{\Xi^0}}. \tag{69}$$

The renormalized coupling constants are determined through

$$1+\bar{S}_s^{\Xi^0} = \frac{2-(1/2)\,\bar{\sigma}^*_\Xi}{C_{\Xi^0}}, \tag{70}$$

$$1+\frac{1}{2}\left(\bar{S}_u^{\Xi^0}+\bar{S}_s^{\Xi^0}\right) = \frac{2-\left(\bar{\sigma}_\Xi+\bar{\delta}_\Xi\right)}{C_{\Xi^0}}. \tag{71}$$

For Ξ^- the contributions of $\bar{\delta}_\Xi$ are negative. Thus the effective mass of Ξ^- is

$$m^*_{\Xi^-} = 1+\bar{S}_{\Xi^-} = 1+\bar{S}_d^{\Xi^-}+2\,\bar{S}_s^{\Xi^-} = \frac{2-2\left(\bar{\sigma}_\Xi-\bar{\delta}_\Xi\right)-(3/2)\,\bar{\sigma}^*_\Xi+\left(\bar{\sigma}_\Xi-\bar{\delta}_\Xi\right)\bar{\sigma}^*_\Xi}{C_{\Xi^-}}. \tag{72}$$

and the renormalized coupling constants are determined through

$$1+\bar{S}_s^{\Xi^-} = \frac{2-(1/2)\,\bar{\sigma}^*_\Xi}{C_{\Xi^-}}, \tag{73}$$

$$1+\frac{1}{2}\left(\bar{S}_d^{\Xi^-}+\bar{S}_s^{\Xi^-}\right) = \frac{2-\left(\bar{\sigma}_\Xi-\bar{\delta}_\Xi\right)}{C_{\Xi^-}}, \tag{74}$$

where

$$C_{\Xi^-} = 2+\frac{1}{2}\left[1-\left(\bar{\sigma}_\Xi-\bar{\delta}_\Xi\right)\right]\bar{\sigma}^*_\Xi. \tag{75}$$

Consequently, we can express all the relevant quantities by the three effective masses m_p^*, m_n^* and m_Λ^*. Their values are determined by extremizing Eq. (33) by them;

$$
\begin{aligned}
\frac{\partial}{\partial M_N^*}\left(\frac{\mathcal{E}}{\rho_T}\right) &= \sum_B \frac{\partial M_B^*}{\partial M_N^*}\frac{\rho_{BS}}{\rho_T} + \frac{m_\sigma^2\langle\sigma\rangle}{M_N\,\rho_T}\frac{\partial\langle\sigma\rangle}{\partial m_N^*} + \frac{m_\delta^2\langle\delta_3\rangle}{M_N\,\rho_T}\frac{\partial\langle\delta_3\rangle}{\partial m_N^*} + \frac{m_{\sigma^*}^2\langle\sigma^*\rangle}{M_N\,\rho_T}\frac{\partial\langle\sigma^*\rangle}{\partial m_N^*} \\
&+ \frac{\rho_T}{M_N\,m_\omega^2}\left(\sum_B f_B\,g_{BB\omega}^*\right)\left(\sum_B f_B\frac{\partial g_{BB\omega}^*}{\partial m_N^*}\right) \\
&+ \frac{\rho_T}{M_N\,m_\rho^2}\left(\sum_B I_{3B}f_B\,g_{BB\rho}^*\right)\left(\sum_B I_{3B}f_B\frac{\partial g_{BB\rho}^*}{\partial m_N^*}\right) \\
&+ \frac{\rho_T}{M_N\,m_\phi^2}\left(\sum_Y f_Y\,g_{YY\phi}^*\right)\left(\sum_Y f_Y\frac{\partial g_{YY\phi}^*}{\partial m_N^*}\right) = 0, \qquad (76)
\end{aligned}
$$

$$
\begin{aligned}
\frac{\partial}{\partial M_\Lambda^*}\left(\frac{\mathcal{E}}{\rho_T}\right) &= \sum_B \frac{\partial M_B^*}{\partial M_\Lambda^*}\frac{\rho_{BS}}{\rho_T} + \frac{m_{\sigma^*}^2\langle\sigma^*\rangle}{M_\Lambda\,\rho_T}\frac{\partial\langle\sigma^*\rangle}{\partial m_\Lambda^*} \\
&+ \frac{\rho_T}{M_\Lambda\,m_\omega^2}\left(\sum_B f_B\,g_{BB\omega}^*\right)\left(\sum_Y f_Y\frac{\partial g_{YY\omega}^*}{\partial m_\Lambda^*}\right) \\
&+ \frac{\rho_T}{M_\Lambda\,m_\rho^2}\left(\sum_B I_{3B}f_B\,g_{BB\rho}^*\right)\left(\sum_Y I_{3B}f_Y\frac{\partial g_{YY\rho}^*}{\partial m_\Lambda^*}\right) \\
&+ \frac{\rho_T}{M_\Lambda\,m_\phi^2}\left(\sum_Y f_Y\,g_{YY\phi}^*\right)\left(\sum_Y f_Y\frac{\partial g_{YY\phi}^*}{\partial m_\Lambda^*}\right) = 0, \qquad (77)
\end{aligned}
$$

where $m_N^* = m_p^*$ or m_n^* in Eq. (76). The ρ_T is the total baryon density, $f_B = \rho_B/\rho_T$ is the fraction of each baryon. The solutions of Eqs. (76) and (77) must satisfy the energy minimization conditions. It is tedious but straightforward task to calculate the derivatives in Eqs. (76) and (77) and so we will not present their expressions explicitly.

3 Numerical analyses

The present work treats the cold non-rotating β-equilibrated NS. The properties of NS matter in our EZM model are essentially determined by m_p^*, m_n^* and m_Λ^* that are the solutions of Eqs. (76) and (77). To solve these equations, the densities of every baryon are needed at fixed total baryon density ρ_T. They are determined from the chemical potential as

$$\mu_B = \left(k_{BF}^2 + M_B^{*\,2}\right)^{1/2} + V_B. \qquad (78)$$

The β-equilibrium condition requires

$$\mu_i = b_i \mu_n - q_i \mu_e, \tag{79}$$

where μ_i is the chemical potential of all the baryons and leptons (e^- and μ^-), and b_i and q_i are the corresponding baryon number and charge. There exist only two independent chemical potentials of neutron and electron. They are determined to satisfy

$$\rho_T = \sum_{\substack{B=p,n,\Lambda,\Sigma^+,\\ \Sigma^0,\Sigma^-,\Xi^0,\Xi^+}} \rho_B, \tag{80}$$

and the charge neutral condition,

$$\sum_{i=B,l} q_i \rho_i = 0. \tag{81}$$

Since the right hand side of Eq. (78) includes M_B^*, Eqs. (76)-(81) have to be solved self-consistently.

For numerical calculations of the NS matter, we determine the free meson-baryon coupling constants. The $NN\sigma$ and $NN\omega$ coupling constants are fixed to reproduce the nuclear matter saturation. We assume the saturation energy of $-15.75\,\mathrm{MeV}$ at the saturation density $0.16\,\mathrm{fm}^{-3}$. The values $(g_{NN\sigma}/m_\sigma)^2 = 16.9\,\mathrm{fm}^2$ and $(g_{NN\omega}/m_\omega)^2 = 12.5\,\mathrm{fm}^2$ are obtained. The effective nucleon mass and the incompressibility of saturated nuclear matter are $m_N^* = 0.605$ and $K = 302\,\mathrm{MeV}$. The $NN\delta$ and $NN\rho$ coupling constants are fixed to the values by the Bonn A potential in Ref. [13]. The detailed explanations of these meson-nucleon coupling constants have been given in Refs. [22] and [23].

Next, the meson-hyperon coupling constants have to be determined. Unfortunately, at present there is little reliable information on the nucleon-hyperon (NY) and hyperon-hyperon (YY) interactions. In this work we investigate four sets of the coupling constants. First, all the meson-hyperon coupling constants are related to the meson-nucleon coupling constants by the SU(6) symmetry. Hereafter this choice will be referred to as the coupling set 1. Because the QMC model in Ref. [21] has employed this coupling set, the comparison between our EZM and the QMC model is interesting.

The other three coupling sets are determined by using both empirical method and the SU(6) symmetry. Namely, the $YY\omega$, $YY\delta$, $YY\rho$ and $YY\phi$ coupling constants are determined by the SU(6) symmetry. The $YY\sigma$ coupling constants are determined to give the hyperon potentials in saturated nuclear matter $U_Y^{(N)}(\rho_{nm})$ [29,30],

$$U_\Lambda^{(N)}(\rho_{nm}) = -28\,\mathrm{MeV}, \quad U_\Sigma^{(N)}(\rho_{nm}) = 30\,\mathrm{MeV} \quad \text{and} \quad U_\Xi^{(N)}(\rho_{nm}) = -18\,\mathrm{MeV}. \tag{82}$$

In our model they are given by

$$U_Y^{(N)}\left(\rho_{nm}\right) = -\, g^*_{YY\sigma}\left\langle\sigma\right\rangle_{NM} + \, g^*_{YY\omega}\left\langle\omega_0\right\rangle_{NM}\,, \tag{83}$$

where $\langle\sigma\rangle_{NM}$ and $\langle\omega_0\rangle_{NM}$ are the mean-fields in saturated nuclear matter. These choices of the coupling constants are the most plausible at present and so are common to the coupling sets 2, 3 and 4.

One of the differences between the three sets is in the choices of $YY\sigma^*$ coupling constants. In the set 2 they are adjusted [16] so that the potential of a single hyperon, embedded in a bath of Ξ matter at ρ_{nm}, becomes

$$U_\Xi^{(\Xi)}\left(\rho_{nm}\right) = U_\Lambda^{(\Xi)}\left(\rho_{nm}\right) = -40\,\mathrm{MeV}, \tag{84}$$

and $g_{\Sigma\Sigma\sigma^*} = g_{\Lambda\Lambda\sigma^*}$ is assumed. The resulting coupling constants predict the following potential of a single hyperon embedded in a bath of Λ matter:

$$U_\Xi^{(\Lambda)}\left(\rho_{nm}/2\right) \approx U_\Lambda^{(\Lambda)}\left(\rho_{nm}/2\right) \approx -20\,\mathrm{MeV}. \tag{85}$$

It is noted that this rather strong attractive $\Lambda\Lambda$ interaction is based on the old data of double Λ hypernuclei.

However the recent discovery of ${}^{\;6}_{\Lambda\Lambda}$He in the KEK-E373 experiment [12] has presented a question on such strong $\Lambda\Lambda$ attraction. This new result was reproduced by the calculation [7] based on the three-body Faddeev equation and the most recent Nijmegen soft-core potential NSC97 [31,32]. In fact the NRBHF calculation [33] using NSC97 model predicted no binding of pure Λ matter. Thus, in the coupling set 3, we implement the NSC97f potential in our EZM model by adjusting the $YY\sigma^*$ coupling constants to reproduce the hyperon binding energy curves in Fig. 2 of Ref. [33]. This implementation is according to Ref. [34] but we have reduced $U_\Sigma^{(N)}\left(\rho_{nm}\right)$ in Eq. (82) to 20 MeV, so that the binding energy of pure Σ matter agrees with the result of Ref. [33] more precisely than Ref. [34]. Although the calculation of Ref. [7] used the NSC97e potential, we have adopted the model f because the deep binding of pure Ξ matter by the model e predicted in Ref. [33] cannot be reproduced by the RMF model using the reasonable coupling constant.

The NRBHF calculation using NSC97 potential in Ref. [33] predicted relatively deep binding of pure Σ matter. Because the $U_\Sigma^{(N)}\left(\rho_{nm}\right)$ is taken to be repulsive, this means that the above implementation of the NSC97 potential leads to much stronger $\Sigma\Sigma\sigma^*$ coupling constant than the value derived from SU(6) symmetry. It is not clear whether such a strong coupling is physically rational or not. Therefore, in the coupling set 4, we take the same coupling constant for $\Sigma\Sigma\sigma^*$ as $\Lambda\Lambda\sigma^*$ according to the SU(6) symmetry. The values of coupling constants of each set are summarized in Table 1.

Figure 1 calculates the EOSs of β-equilibrated NS matter. The dashed, dashed-dotted, solid and dotted curves are the results using the meson-hyperon coupling sets 1, 2, 3 and 4 respectively. Taking these EOSs as inputs, we integrate the so-called Tolman-Oppenheimer-Volkov (TOV) equation [35]. For the outer region of NS, we use the EOSs by Feynman-Metropolis-Teller, Baym-Pethick-Sutherland and Negele-Vautherin from Ref. [36]. Figures 2 and 3 show the gravitational masses of neutron stars in unit of the solar mass as functions of the central energy density $\mathcal{E}_C$ and the radius R. The curves are the same as those in Fig. 1, but the coupling set 2 is excluded because it is unphysical as discussed below. Figures 4 to 7 show the fraction of baryons and leptons in β-equilibrated NS matter using the coupling sets 1 to 4 respectively.

Now we are going to discuss the results. It is first noted that there are the critical densities above which we cannot find β-equilibrium state. Their values are $\rho_T = 0.887\,\mathrm{fm}^{-3}$, $0.867\,\mathrm{fm}^{-3}$, $0.691\,\mathrm{fm}^{-3}$ and $0.998\,\mathrm{fm}^{-3}$ for the coupling sets 1 to 4. This is not surprising because in our EZM model the vector potential V_B in the chemical potential (78) depends directly on the effective mass M_B^* while in the other RMF models it depends only on the each baryon density ρ_B. (See Eq. (29).) The constraint of the β-equilibrium condition on the self-consistency equations of the effective masses or the meson mean-fields are much severer in the EZM model than the other models.

The critical density will be no trouble if it is higher than the central baryon density ρ_T in NS with maximum mass. The meson-hyperon coupling set 1 however cannot satisfy this condition. The maximum mass appears at the end point of the dashed curve in Fig. 2 or 3, that is, at the critical density $\rho_T = 0.887\,\mathrm{fm}^{-3}$. Furthermore its value $M_{ns} = 1.402\,M_\odot$ is lower than the canonical value $1.44\,M_\odot$. This is because the coupling set 1 produces the softer EOS than the other sets in the region of $250\,\mathrm{MeV{\cdot}fm}^{-3} < \mathcal{E} < 500\,\mathrm{MeV{\cdot}fm}^{-3}$. The softness is due to the abundance of Σ^- above $\rho_T = 0.26\,\mathrm{fm}^{-3}$ as seen in Fig. 4. In the other coupling sets the hyperons appear at higher densities $\rho_T \approx 0.4\,\mathrm{fm}^{-3}$ as seen in Figs. 5, 6 and 7 and there are no Σ's in the sets 2 and 4. The early abundance of Σ^- and the appearance of Σ^0 and Σ^+ in the set 1 are due to its stronger $\Sigma\Sigma\sigma$ coupling constant, which produces relatively deep attractive potential $U_\Sigma^{(N)}(\rho_{nm})$ while in the other coupling sets the repulsive potentials are assumed. It also leads to the delayed appearance and scarcity of Ξ hyperons and so the abundance of Λ hyperons. At present the fact that there are no observations of Σ hypernuclei strongly supports the positive value of $U_\Sigma^{(N)}(\rho_{nm})$. In addition the detailed theoretical analysis of Σ^- atomic data [37] and (π^-, K^+) inclusive spectra [38] predict a repulsive Σ-nucleus optical potential in the nuclear interior.

Consequently, the possibility of the coupling set 1 with strong attractive $N\Sigma$ interaction has been denied in our EZM model. On the contrary, the nonlinear Walecka and QMC models [21] predicted the consistent results of NSs with observations even if all the meson-hyperon coupling constants are derived from the SU(6) symmetry. Furthermore both the models work well at much higher density, where the RMF model of baryon matter cannot be expected to be valid, than our critical density. In this respect we can expect that the EZM model is much sensitive to the choice of the meson-hyperon coupling constants and so selects reasonable values of them. This expectation will be further examined in the other coupling sets.

The outstanding feature of the meson-hyperon coupling set 2 is much stronger $\Lambda\Lambda\sigma^*$ coupling constant than the other sets as seen in Table I. It leads to strong $\Lambda\Lambda$ attraction, which is revealed as much abundance of Λ in Fig. 5 compared with Figs. 6 and 7. The EOSs by the coupling sets 2, 3 and 4 in Fig. 1 are almost the same below $\mathcal{E} = 400\,\mathrm{MeV{\cdot}fm^{-3}}$. However the pressure of the set 2 depletes above $\rho_T = 0.424\,\mathrm{fm^{-3}}$. This is accompanied by the abundance of Ξ^- hyperons above $\rho_T = 0.4\,\mathrm{fm^{-3}}$. They appear immediately after the abundance of Λ hyperons because of the stronger $\Xi\Xi\sigma^*$ coupling than the other hyperons. We further find a kink around $\mathcal{E} = 700\,\mathrm{MeV{\cdot}fm^{-3}}$. It is accompanied by the abundance of Ξ^0 hyperons above $\rho_T = 0.65\,\mathrm{fm^{-3}}$. These features of the EOS by the set 2 are essentially due to the strong $\Lambda\Lambda$ attraction. For the coupling sets 3 and 4 with weak $\Lambda\Lambda$ interaction, we can see in Figs. 6 and 7 that the abundance of Λ is restrained after appearances of the other hyperons and so the Ξ^- is the most rich hyperon at high densities. On the contrary, the strong $\Lambda\Lambda$ attraction allows the Λ to be the most rich hyperon even after the appearances of Ξ's. Therefore the abundances of Ξ's in Fig. 5 lead to the rapid increase of σ^* mean-field owing to the strong $\Lambda\Xi$ interactions. Consequently, the pressure depletes or the EOS is softened according to Eq. (34).

Although the pressure P of the set 2 depletes above $\rho_T = 0.424\,\mathrm{fm^{-3}}$, the Gibbs free energy per baryon $\bar{G}$, which is given by only the baryochemical potential as $\bar{G} = (\mathcal{E}+P)/\rho_T = \mu_n$ owing to the Gibbs-Duhem relation and the charge neutrality, remains increasing. As the result we cannot find an evidence of the first-order phase transition or a point of intersection in the correlation between P and $\bar{G}$. It is therefore concluded that the EOS by the meson-hyperon coupling set 2 is unphysical and so the strong $\Lambda\Lambda$ attraction is ruled out in our EZM model. (For this reason, the calculations using the set 2 are not shown in Figs. 2 and 3.) If the strong $\Lambda\Lambda$ attraction suggested by the old data of double Λ hypernuclei is true, our result requires the transition to a new phase as the de-confined quark matter below $\rho_T = 0.424\,\mathrm{fm^{-3}}$. On the contrary, the nonlinear Walecka model [39] and the straightforward extension of the original ZM model [20] work well until much higher densities than our critical density even if the strong $\Lambda\Lambda$ attraction is employed. It has been again found that the EZM model selects the meson-hyperon coupling constants.

Because the EZM model has denied the possibility of strong $\Lambda\Lambda$ attraction, its consistency with new ${}^{6}_{\Lambda\Lambda}\mathrm{He}$ data by the KEK-E373 experiment [12] has to be examined. The meson-hyperon coupling sets 3 and 4 are appropriate for this purpose. Their $\Lambda\Lambda\sigma^*$ coupling constant is much weaker than the set 2 and is consistent to the NSC97 potential, which was found [7] to be able to reproduce new ${}^{6}_{\Lambda\Lambda}\mathrm{He}$ data. The difference between the set 3 and 4 is in their $\Sigma\Sigma\sigma^*$ coupling constants. (Although their $\Sigma\Sigma\sigma$ coupling constants in Table I are different from each other, the results of the set 4 are not altered if its $\Sigma\Sigma\sigma$ coupling constant is the same as the set 3.) The set 3 has much stronger $\Sigma\Sigma\sigma^*$ coupling than the other sets. Within the mean-field theory the $\Sigma\Sigma$ interaction has no effect on the $\Lambda\Lambda$ channel. Thus both the set 3 and 4 are consistent to new ${}_{\Lambda\Lambda}^{6}$He data.

The EOSs by the sets 3 and 4, the solid and dotted curves in Fig. 1, are almost the same below $\mathcal{E} = 700\,\mathrm{MeV{\cdot}fm^{-3}}$. They have two kinks around $\mathcal{E} = 400\,\mathrm{MeV{\cdot}fm^{-3}}$ and $\mathcal{E} = 700\,\mathrm{MeV{\cdot}fm^{-3}}$. The first is due to the abundance of Ξ^- hyperons above $\rho_T = 0.41\,\mathrm{fm^{-3}}$ while the second is due to the appearance of Ξ^0 hyperons above $\rho_T = 0.661\,\mathrm{fm^{-3}}$ in the set 3 and $\rho_T = 0.678\,\mathrm{fm^{-3}}$ in the set 4. Because the Ξ hyperons have the strangeness $S = -2$, their abundances mean the increase of σ^* mean-field and so soften the EOS. In Fig. 6 the Σ^- hyperons appear immediately after Ξ^-. As the result of Σ^- abundance, the EOS by the set 3 is slightly softer than the set 4 above $\mathcal{E} = 400\,\mathrm{MeV{\cdot}fm^{-3}}$. We have also seen the abundance of Σ^- hyperons in the coupling set 1 (Fig. 4). It was due to the strong attractive $N\Sigma$ interactions. Thus the Σ^- appears first rather than Λ. On the other hand, the abundance of Σ^- in the coupling set 3 is due to the strong attractive $\Lambda\Sigma$ and $\Sigma\Xi$ interactions. It therefore occurs after the abundances of Λ and Ξ hyperons.

It is seen that the results by the sets 3 and 4 in Figs. 2 and 3 are almost the same. It therefore seems that the maximum NS mass cannot resolve the uncertainty in the Σ^- coupling. The same fact has been also found [40] in the nonlinear Walecka model. We can however see that the result by the set 3 reaches the critical density just before maximum NS mass. In this respect the strong $\Sigma\Sigma\sigma^*$ coupling constant in the set 3 cannot be regarded to be reasonable and so the existence of Σ^- hyperons in NSs is ruled out. Consequently, the EZM model allows only the meson-hyperon coupling set 4 with weak $\Lambda\Lambda$ and $\Sigma\Sigma$ interactions. Its maximum mass is $M_{ns} = 1.615\,M_\odot$ with the radius $R = 12.98\,\mathrm{km}$ at the central density just above $\rho_T = 0.691\,\mathrm{fm^{-3}}$. There exist only Λ and Ξ^- as hyperons in NSs.

At present all the confirmed masses of NSs [5,41-44] lie in a narrow range $M_{ns} = 1.35 \pm 0.1\,M_\odot$. The X-ray pulsar Vela X-1 is often referred to [45] as one exception with $M_{ns} \approx 1.8\,M_\odot$. However its value has been claimed to be close to $1.4\,M_\odot$ in Ref. [46]. Even if the larger mass of Vela will be confirmed, the effect of NS rotation raises [47] the masses of Fig. 2 considerably. The maximum mass by the set 4 is therefore reasonable. Although the corresponding radius is much larger than the values from the other EOSs [2], it satisfies both the upper and lower limits derived by the analyses of quasi-periodic oscillations [2] and glitches [48] from X-ray pulsars. The central baryon density of maximum NS is merely $\rho_T \simeq 0.7\,\mathrm{fm^{-3}}$ being much smaller than the nonlinear Walecka and QMC models [21], and the neutrons are dominant to the other baryons. Nevertheless the proton fraction can exceed the critical value $f_p \simeq 0.148$ for the rapid cooling by direct URCA process in the core of neutron stars $\rho_T \geq 0.465\,\mathrm{fm^{-3}}$. The result is common to all the RMF models. In summary we can see that the NSs predicted by the EZM model using the coupling set 4 have rather unique properties as compared with the other models.

4 Conclusions

We have extended the modified ZM model of dense baryon matter, which was previously developed in terms of a constituent quark picture of baryons, to take into account the contributions of isovector and strange mesons for studying NS equations-of-state. The resulting EZM model contains the renormalized meson-baryon coupling constants in the mean-fields. They are effectively density dependent through the effective masses of each baryon.

For calculating the EOS of NS matter, we consider four sets of meson-hyperons coupling constants. The first set is based on pure SU(6) symmetry of baryon octet. The second is derived in terms of empirical hyperon potential in nuclear matter and Ξ bath. It is based on old data of double hypernuclei. The third set mimics the NSC97f YY potential. The forth set is obtained by reducing the $\Sigma\Sigma\sigma^*$ coupling constant in the third set to the same value of the $\Lambda\Lambda\sigma^*$ coupling constant according to the SU(6) symmetry. The set 1 is characterized by strong attractive $N\Sigma$ interaction while the others have repulsive $N\Sigma$ interactions. The set 2 is characterized by strong attractive $\Lambda\Lambda$ interaction. The set 3 has weak $\Lambda\Lambda$ but strong attractive $\Sigma\Sigma$ interactions while the set 4 has much weaker $\Sigma\Sigma$ interaction than the set 3.

There have been found the critical densities above which the β-equilibrium cannot be arrived. Using the coupling set 1, we have a maximum NS mass at the critical density. Furthermore the mass is lower than the canonical value $1.44\,M_{\odot}$. This is because that the EOS is softened by the abundance of Σ^- hyperons owing to the strong $N\Sigma$ attraction. The EOS by the set 2 has been found to be unphysical because of outstanding depletion of pressure accompanied with no first-order phase transitions. It has been attributed to the abundance of Λ hyperons arising from the strong $\Lambda\Lambda$ attraction. The results of the NS mass sequences by the sets 3 and 4 are almost the same. However the EOS by the set 3 also has a maximum NS mass at the critical density as in the set 1. Because only the difference between the sets 3 and 4 is in $\Sigma\Sigma$ interaction, the result means that the strong $\Sigma\Sigma$ attraction is not reasonable.

In conclusion the strong attractive $N\Sigma$, $\Lambda\Lambda$ and $\Sigma\Sigma$ interactions have been ruled out. The result is consistent to the most recent information on hyperon interactions from the experimental and theoretical investigations of hypernuclei. It is found that the EZM model is sensitive to the choice of meson-hyperon coupling constants. In this respect the EZM model is contrasted to the other RMF models. We believe that the physically reasonable model of dense baryon matter favors only the physically reasonable values of meson-baryon coupling constants.

Appendix: derivation of the renormalized coupling constants from relativistic SU(6) model of baryons

In section 2 we derived the renormalized meson-baryon coupling constants in the EZM model using intuitive schematic methods. Here they are derived again more rigorously using the relativistic SU(6) model of baryons.

In the relativistic SU(6) model [49-52], the wave function of baryon octet $\Psi_{(\alpha p)(\beta q)(\gamma r)}$ has three Dirac indices (α, β, γ) and three SU(3)-spin (or flavor) indices (p, q, r), and is fully symmetric in both these indices. (Anti-symmetric tensor representing color singlet state is omitted because the present model does not take into account the color degrees of freedom explicitly.) It is presented by

$$\Psi_{(\alpha p)(\beta q)(\gamma r)}(k) = \frac{1}{2\sqrt{6}} \left\{ \varepsilon_{pqs} \mathcal{B}_{[\alpha\beta]\gamma,}{}^{s}_{r}(k) + \varepsilon_{qrs} \mathcal{B}_{[\beta\gamma]\alpha,}{}^{s}_{p}(k) + \varepsilon_{rps} \mathcal{B}_{[\gamma\alpha]\beta,}{}^{s}_{q}(k) \right\}, \tag{A1}$$

where k is the four-momentum of the baryon and ε_{pqr} is a totally asymmetric tensor of rank 3. $\tilde{B}^{q}_{p}$ is the 3×3 matrix representing the octet:

$$\mathcal{B}^{q}_{p} = \begin{pmatrix} \Sigma^0/\sqrt{2} + \Lambda/\sqrt{6} & \Sigma^+ & P \\ \Sigma^- & -\Sigma^0/\sqrt{2} + \Lambda/\sqrt{6} & N \\ \Xi^- & \Xi^0 & -2\Lambda/\sqrt{6} \end{pmatrix}. \tag{A2}$$

Each wave function of these octet members is expressed by the third-rank Bargmann-Wigner wave function $B_{[\alpha\beta]\gamma}$ [53]:

$$\mathcal{B}_{[\alpha\beta]\gamma,}{}^{q}_{p} = \left(B_{[\alpha\beta]\gamma}\right)^{q}_{p}, \tag{A3}$$

$$B_{[\alpha\beta]\gamma}(k) = (1/M)\left[(\not{k} + M)\gamma_5 C\right]_{\alpha\beta} u_\gamma(k), \tag{A4}$$

where C is the charge conjugate operator and u_γ is the Dirac spinor. We have assumed that all the octets have the same mass M. At the final stage of calculation, it is replaced by each mass of baryons. This prescription phenomenologically breaks the SU(3) invariance.

The free meson-baryon vertex function Γ_{BBM} in the SU(6) model is given by [52]

$$(-i)\,\Gamma_{BBM} = \sum_{i=0}^{8} G^{i}_{R}\, \bar{\Psi}^{ABC} \left(\hat{O}^{R}\lambda^{i}\right)^{A'}_{A} \Psi_{A'BC}\, \Phi^{i}_{R}, \tag{A5}$$

where $A = (\alpha p)$ *etc.*, λ^i is the Gell-Mann matrix and $\hat{O}^R$ is $\mathbf{1}$ or γ^μ for the scalar $(R = S)$ or vector $(R = V)$ meson. Under the SU(3) invariance the coupling constant G^i_R is given by $G^0_R = G_{1R}$ for meson singlet Φ^0_R and $G^i_R = G_{8R}$ $(i = 1-8)$ for meson octet $\Sigma^{8}_{i=1}\Phi^{i}_{R}\lambda^{i}$.

The medium correction to Eq. (A5) in the EZM model is defined by

$$(-i)\,\Delta\Gamma_{BBM} = \sum_{i=0}^{8} G^{i}_{R}\, \bar{\Psi}^{ABC} \left(\hat{O}^{R}\lambda^{i}\right)^{A'}_{A} (S)^{B'}_{B}\, \Psi_{A'B'C}\, \Phi^{i}_{R}, \tag{A6}$$

where

$$S = \frac{1}{M}\begin{pmatrix} S_u & 0 & 0 \\ 0 & S_d & 0 \\ 0 & 0 & S_s \end{pmatrix} \equiv \begin{pmatrix} \bar{S}_u & 0 & 0 \\ 0 & \bar{S}_d & 0 \\ 0 & 0 & \bar{S}_s \end{pmatrix}. \tag{A7}$$

$S_{u,d,s}$ are the scalar potentials of u, d and s-quarks. As for the mass M we have assumed the same values for all the octet members and will replace them by each value for baryons at the final stage of calculation. Substituting Eq. (A1) into (A6) and using the Bianchi identity

$$\varepsilon^{pqs}\mathcal{B}^r_s + \varepsilon^{qrs}\mathcal{B}^p_s + \varepsilon^{rps}\mathcal{B}^q_s = 0, \tag{A8}$$

after some manipulations we obtain

$$\begin{aligned}
(-i)\,\Delta\Gamma_{BBM} \;=\;& (1/24)\Bigg\{ \bar{S}_q \,\mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\beta\gamma]\alpha} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} [G_R\Phi_R]\left(\mathcal{B}_{[\beta\gamma]\alpha'} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)\right] \\
&- \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\beta\gamma]\alpha} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} [G_R\Phi_R]\left(\mathcal{B}_{[\beta\gamma]\alpha'} - \mathcal{B}_{[\alpha'\beta]\gamma}\right) S\right] \\
&- \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\beta\gamma]\alpha} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} [G_R\Phi_R]\, S\left(\mathcal{B}_{[\gamma\alpha']\beta} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)\right] \\
&+ \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\beta\gamma]\alpha} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)[G_R\Phi_R]\right]\left(\hat{O}^R\right)_\alpha^{\alpha'} \mathrm{Tr}\left[\left(\mathcal{B}_{[\gamma\alpha']\beta} - \mathcal{B}_{[\alpha'\beta]\gamma}\right) S\right] \\
&- \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\gamma\alpha]\beta} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} S\,[G_R\Phi_R]\left(\mathcal{B}_{[\beta\gamma]\alpha} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)\right] \\
&+ \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\gamma\alpha]\beta} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right) S\right]\left(\hat{O}^R\right)_\alpha^{\alpha'} \mathrm{Tr}\left[\left(\mathcal{B}_{[\beta\gamma]\alpha'} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)[G_R\Phi_R]\right] \\
&+ \sqrt{6}\, G_R^0\, \Phi_R^0\, \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\gamma\alpha]\beta} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} S\left(\mathcal{B}_{[\gamma\alpha']\beta} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)\right] \\
&- \mathrm{Tr}\left[\left(\bar{\mathcal{B}}^{[\gamma\alpha]\beta} - \bar{\mathcal{B}}^{[\alpha\beta]\gamma}\right)\left(\hat{O}^R\right)_\alpha^{\alpha'} S\left(\mathcal{B}_{[\gamma\alpha']\beta} - \mathcal{B}_{[\alpha'\beta]\gamma}\right)[G_R\Phi_R]\right]\Bigg\},
\end{aligned} \tag{A9}$$

where $\bar{S}_q \equiv \bar{S}_u + \bar{S}_d + \bar{S}_s$ and $[G_R\Phi_R] = \Sigma_{i=0}^{8} G_R^i \Phi_R^i \lambda^i$.

Then we calculate the traces of the SU(3)-spin matrices for the vector mesons by assuming the $\phi - \omega$ ideal mixing,

$$[G_V\Phi_V] = G_{8V}\begin{pmatrix} \omega + \rho_3 & 0 & 0 \\ 0 & \omega - \rho_3 & 0 \\ 0 & 0 & -\sqrt{2}\phi \end{pmatrix} \equiv G_{8V}\,[\Phi_V]\,. \tag{A10}$$

For brevity we omit $\hat{O}^V = \gamma^\mu$ and the relevant indices of the Dirac spinors. The results are

$$\begin{aligned}\mathrm{Tr}\left(\bar{\mathcal{B}}[\Phi_V]\mathcal{B}\right) &= \bar{P}\,\omega P + \bar{N}\omega N + \frac{1}{3}\bar{\Lambda}\omega\,\Lambda + \bar{\Sigma}^+\omega\,\Sigma^+ + \bar{\Sigma}^0\omega\,\Sigma^0 + \bar{\Sigma}^-\omega\,\Sigma^- \\ &\quad - \frac{2\sqrt{2}}{3}\,\bar{\Lambda}\,\phi\,\Lambda - \sqrt{2}\;\bar{\Xi}^0\phi\,\Xi^0 - \sqrt{2}\;\bar{\Xi}^-\phi\,\Xi^-, \end{aligned} \tag{A11}$$

$$\begin{aligned}\mathrm{Tr}\left(\bar{\mathcal{B}}[\Phi_V]\mathcal{B}S\right) &= \bar{S}_s\left(\bar{P}\omega P + \bar{N}\omega N\right) + \frac{1}{6}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Lambda}\omega\,\Lambda \\ &\quad + \bar{S}_d\bar{\Sigma}^+\omega\,\Sigma^+ + \frac{1}{2}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Sigma}^0\omega\,\Sigma^0 + \bar{S}_u\bar{\Sigma}^-\omega\,\Sigma^- \\ &\quad - \frac{2\sqrt{2}}{3}\,\bar{S}_s\,\bar{\Lambda}\,\phi\,\Lambda - \sqrt{2}\left(\bar{S}_d\,\bar{\Xi}^0\phi\,\Xi^0 + \bar{S}_u\,\bar{\Xi}^-\phi\,\Xi^-\right) \\ &\quad + \bar{S}_s\left(\bar{P}\rho_3 P - \bar{N}\rho_3 N\right) + \bar{S}_d\,\bar{\Sigma}^+\rho_3\,\Sigma^+ - \bar{S}_u\,\bar{\Sigma}^-\rho_3\,\Sigma^-, \end{aligned} \tag{A12}$$

$$\begin{aligned}\mathrm{Tr}\left(\bar{\mathcal{B}}[\Phi_V]\,S\mathcal{B}\right) &= \mathrm{Tr}\left(\bar{\mathcal{B}}S\,[\Phi_V]\mathcal{B}\right) = \bar{S}_u\bar{P}\,\omega\,P + \bar{S}_d\bar{N}\,\omega\,N + \frac{1}{6}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Lambda}\omega\,\Lambda \\ &\quad + \bar{S}_u\,\bar{\Sigma}^+\omega\,\Sigma^+ + \frac{1}{2}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Sigma}^0\omega\,\Sigma^0 + \bar{S}_d\,\bar{\Sigma}^-\omega\,\Sigma^- \\ &\quad - \frac{2\sqrt{2}}{3}\,\bar{S}_s\bar{\Lambda}\,\phi\,\Lambda - \sqrt{2}\,\bar{S}_s\left(\bar{\Xi}^0\phi\,\Xi^0 + \bar{\Xi}^-\phi\,\Xi^-\right) \\ &\quad + \bar{S}_u\bar{P}\rho_3 P - \bar{S}_d\,\bar{N}\rho_3 N + \bar{S}_u\,\bar{\Sigma}^+\rho_3\,\Sigma^+ - \bar{S}_d\,\bar{\Sigma}^-\rho_3\,\Sigma^-, \end{aligned} \tag{A13}$$

$$\mathrm{Tr}\left(\bar{\mathcal{B}}[\Phi_V]\right) \times \mathrm{Tr}\left(\mathcal{B}S\right) = \mathrm{Tr}\left(\bar{\mathcal{B}}S\right) \times \mathrm{Tr}\left(\mathcal{B}[\Phi_V]\right) = \frac{1}{3}\left(\bar{S}_u + \bar{S}_d - 2\bar{S}_s\right)\left(\bar{\Lambda}\,\omega\,\Lambda + \sqrt{2}\,\bar{\Lambda}\,\phi\,\Lambda\right), \tag{A14}$$

$$\begin{aligned}\mathrm{Tr}\left(\bar{\mathcal{B}}S\mathcal{B}[\Phi_V]\right) &= \frac{1}{6}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Lambda}\,\omega\,\Lambda + \bar{S}_u\bar{\Sigma}^+\omega\,\Sigma^+ + \frac{1}{2}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Sigma}^0\omega\,\Sigma^0 + \bar{S}_d\,\bar{\Sigma}^-\omega\,\Sigma^- \\ &+ \bar{S}_s\left(\bar{\Xi}^0\omega\,\Xi^0 + \bar{\Xi}^-\omega\,\Xi^-\right) - \sqrt{2}\,\bar{S}_u\bar{P}\,\phi\,P - \sqrt{2}\,\bar{S}_d\bar{N}\,\phi\,N - \frac{2\sqrt{2}}{3}\,\bar{S}_s\bar{\Lambda}\,\phi\Lambda \\ &- \bar{S}_u\,\bar{\Sigma}^+\rho_3\,\Sigma^+ + \bar{S}_d\,\bar{\Sigma}^-\rho_3\,\Sigma^- - \bar{S}_s\left(\bar{\Xi}^0\rho_3\,\Xi^0 - \bar{\Xi}^-\rho_3\,\Xi^-\right), \end{aligned} \tag{A15}$$

$$\begin{aligned}\mathrm{Tr}\left(\bar{\mathcal{B}}S\mathcal{B}\right) &= \bar{S}_u\bar{P}P + \bar{S}_d\bar{N}N + \frac{1}{6}\left(\bar{S}_u + \bar{S}_d + 4\bar{S}_s\right)\bar{\Lambda}\Lambda + \bar{S}_u\,\bar{\Sigma}^+\Sigma^+ \\ &\quad + \frac{1}{2}\left(\bar{S}_u + \bar{S}_d\right)\bar{\Sigma}^0\Sigma^0 + \bar{S}_d\,\bar{\Sigma}^+\Sigma^+ + \bar{S}_s\left(\bar{\Xi}^0\Xi^0 + \bar{\Xi}^-\Xi^-\right), \end{aligned} \tag{A16}$$

where the $\Lambda\Lambda\rho_3$ and $\Sigma^0\Sigma^0\rho_3$ couplings and the $\Lambda\Sigma^0\omega\,(\phi,\rho_3)$ couplings have been suppressed.

Noting $G_{0V} = -G_{8V}$ and $\Phi_V^0 = \sqrt{1/3}\,\phi - \sqrt{2/3}\,\omega$, the medium correction to $pp\omega$ coupling is

$$
\begin{aligned}
(-\,i)\,\Delta\Gamma_{pp\omega} \;=\;& \frac{1}{24}\left(\bar{S}_u+\bar{S}_d\right) G_{8V}\left(\bar{P}^{[\beta\gamma]\alpha}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}\left(P_{[\beta\gamma]\alpha'}-P_{[\alpha'\beta]\gamma}\right)\omega_{\mu}\\
&-\frac{1}{24}\,\bar{S}_u\,G_{8V}\left(\bar{P}^{[\beta\gamma]\alpha}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}\left(P_{[\gamma\alpha']\beta}-P_{[\alpha'\beta]\gamma}\right)\omega_{\mu}\\
&-\frac{1}{24}\,\bar{S}_u\,G_{8V}\left(\bar{P}^{[\gamma\alpha]\beta}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}\left(P_{[\beta\gamma]\alpha'}-P_{[\alpha'\beta]\gamma}\right)\omega_{\mu}\\
&+\frac{1}{12}\,\bar{S}_u\,G_{8V}\left(\bar{P}^{[\gamma\alpha]\beta}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}\left(P_{[\gamma\alpha']\beta}-P_{[\alpha'\beta]\gamma}\right)\omega_{\mu}. \qquad \text{(A17)}
\end{aligned}
$$

Using the Bianchi identity

$$
B_{[\alpha\beta]\gamma}+B_{[\beta\gamma]\alpha}+B_{[\gamma\alpha]\beta}=0, \qquad \text{(A18)}
$$

and after some manipulations, Eq. (A17) becomes

$$
\begin{aligned}
(-i)\,\Delta\Gamma_{pp\omega} \;=\;& \frac{1}{24}\,\bar{S}_u\,G_{8V}\left[-7\bar{P}^{[\beta\gamma]\alpha}\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}P_{[\alpha'\beta]\gamma}+2\,\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}P_{[\alpha'\gamma]\beta}\right]\omega_{\mu}\\
&+\frac{1}{24}\,\bar{S}_d\,G_{8V}\left[-5\,\bar{P}^{[\beta\gamma]\alpha}\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}P_{[\alpha'\beta]\gamma}+4\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}P_{[\alpha'\gamma]\beta}\right]\omega_{\mu}. \qquad \text{(A19)}
\end{aligned}
$$

Here we note

$$
\bar{B}^{[\beta\alpha]\gamma}\left(p'\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}B_{[\alpha'\beta]\gamma}\left(p\right)=8\,\bar{u}\left(p'\right)\left[\gamma^{\mu}-i\,\sigma^{\mu\nu}q_{\nu}\left(2M\right)\right]u\left(p\right), \qquad \text{(A20)}
$$

$$
\bar{B}^{[\beta\alpha]\gamma}\left(p'\right)\left(\gamma^{\mu}\right)_{\alpha}^{\alpha'}B_{[\alpha'\gamma]\beta}\left(p\right)=4\,\bar{u}\left(p'\right)\left[\left(1+q^2/\left(4M^2\right)\right)\gamma^{\mu}-\,i\,\sigma^{\mu\nu}q_{\nu}/M\right]u\left(p\right), \qquad \text{(A21)}
$$

where $u\left(p\right)$ is the Dirac spinor for each baryon and $q=p'-p$. Therefore, at the limit of $q\rightarrow 0$ being equivalent to the mean-field approximation, Eq. (A19) becomes

$$
(-i)\,\Delta\Gamma_{pp\omega}\;=-\left(2\bar{S}_u+\bar{S}_d\right)G_{8V}\,\bar{\psi}_p\,\gamma^{\mu}\,\psi_p\,\omega_{\mu}. \qquad \text{(A22)}
$$

Utilizing the above-mentioned prescription $\bar{S}_{u(d)}\rightarrow\bar{S}^N_{u(d)}$ and noting $g_{NN\omega}=3G_{8V}$, we have

$$
(-i)\,\Delta\Gamma_{pp\omega}=-\left(1/3\right)\bar{S}_p\,g_{NN\omega}\,\bar{\psi}_p\,\gamma^{\mu}\,\psi_p\,\omega_{\mu}, \qquad \text{(A23)}
$$

where S_p is the scalar potential of proton:

$$
2\,\bar{S}^N_u+\bar{S}^N_d=\bar{S}_p=S_p/M. \qquad \text{(A24)}
$$

Therefore the effective $pp\omega$ coupling becomes

$$
\Gamma^*_{pp\omega}=\Gamma_{pp\omega}+\Delta\Gamma_{pp\omega}=(-\,i)\,g^*_{NN\omega}\,\bar{\psi}_p\,\gamma^{\mu}\,\psi_p\,\omega_{\mu}, \qquad \text{(A25)}
$$

$$g^*_{NN\omega} = \left(1 + \frac{1}{3}\bar{S}_p\right) g_{NN\omega} = \left[1 + \frac{1}{3}\left(m^*_p - 1\right)\right] g_{NN\omega}. \tag{A26}$$

This is exactly Eq. (1). Similarly, the effective $nn\omega$ coupling (2) is derived. As the result of the $\phi - \omega$ ideal mixing, we naturally have $\Delta\Gamma_{NN\phi} = 0$.

Next, the medium correction to $\Lambda\Lambda\omega$ coupling is

$$\begin{aligned}
(-i)\,\Delta\Gamma_{\Lambda\Lambda\omega} &= \frac{1}{144}\left(\bar{S}_u + \bar{S}_d + 2\bar{S}_s\right) G_{8V}\left(\bar{\Lambda}^{[\beta\gamma]\alpha} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\beta\gamma]\alpha'} - \Lambda_{[\alpha'\beta]\gamma}\right)\omega_\mu \\
&- \frac{1}{144}\left(\bar{S}_u + \bar{S}_d - 4\bar{S}_s\right) G_{8V}\left(\bar{\Lambda}^{[\beta\gamma]\alpha} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\gamma\,\alpha']\beta} - \Lambda_{[\alpha'\beta]\gamma}\right)\omega_\mu \\
&- \frac{1}{144}\left(\bar{S}_u + \bar{S}_d - 4\bar{S}_s\right) G_{8V}\left(\bar{\Lambda}^{[\gamma\,\alpha]\beta} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\beta\gamma]\alpha'} - \Lambda_{[\alpha'\beta]\gamma}\right)\omega_\mu \\
&+ \frac{1}{144}\left(\bar{S}_u + \bar{S}_d + 8\bar{S}_s\right) G_{8V}\left(\bar{\Lambda}^{[\gamma\alpha]\beta} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\gamma\,\alpha']\beta} - \Lambda_{[\alpha'\beta]\gamma}\right)\omega_\mu, \qquad \text{(A27)} \\
&= -\frac{1}{16}\left(\bar{S}_u + \bar{S}_d + 2\bar{S}_s\right) G_{8V}\,\bar{\Lambda}^{[\beta\alpha]\gamma}\,(\gamma^\mu)^{\alpha'}_{\alpha}\,\Lambda_{[\alpha'\beta]\gamma}\,\omega_\mu, \qquad \text{(A28)} \\
&= -\left[\frac{1}{2}\left(\bar{S}_u + \bar{S}_d\right) + \bar{S}_s\right] G_{8V}\,\bar{\psi}_\Lambda\,\gamma^\mu\,\psi_\Lambda\,\omega_\mu. \qquad \text{(A29)}
\end{aligned}$$

Using the prescription $\bar{S}_{u,d,s} \to \bar{S}^\Lambda_{u,d,s}$ and noting $g_{\Lambda\Lambda\omega} = 2G_{8V}$, the effective $\Lambda\Lambda\omega$ coupling becomes

$$\Gamma^*_{\Lambda\Lambda\omega} = (-i)\,g^*_{\Lambda\Lambda\omega}\,\bar{\psi}_\Lambda\,\gamma^\mu\,\psi_\Lambda\,\omega_\mu, \tag{A30}$$

$$g^*_{\Lambda\Lambda\omega} = \left[1 + \frac{1}{4}\left(\bar{S}^\Lambda_u + \bar{S}^\Lambda_d\right) + \frac{1}{2}\,\bar{S}^\Lambda_s\right] g_{\Lambda\Lambda\omega}. \tag{A31}$$

This is nothing but Eq. (9).

Then the medium correction to $\Lambda\Lambda\phi$ coupling is

$$\begin{aligned}
(-i)\,\Delta\Gamma_{\Lambda\Lambda\phi} &= -\frac{\sqrt{2}}{36}\left(\bar{S}_u + \bar{S}_d\right) G_{8V}\left(\bar{\Lambda}^{[\beta\gamma]\alpha} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\beta\gamma]\alpha'} - \Lambda_{[\alpha'\beta]\gamma}\right)\phi_\mu \\
&- \frac{\sqrt{2}}{72}\left(\bar{S}_u + \bar{S}_d\right) G_{8V}\left(\bar{\Lambda}^{[\beta\gamma]\alpha} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\gamma\alpha']\beta} - \Lambda_{[\alpha'\beta]\gamma}\right)\phi_\mu \\
&- \frac{\sqrt{2}}{72}\left(\bar{S}_u + \bar{S}_d\right) G_{8V}\left(\bar{\Lambda}^{[\gamma\alpha]\beta} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\beta\gamma]\alpha'} - \Lambda_{[\alpha'\beta]\gamma}\right)\phi_\mu \\
&- \frac{\sqrt{2}}{144}\left(\bar{S}_u + \bar{S}_d\right) G_{8V}\left(\bar{\Lambda}^{[\gamma\alpha]\beta} - \bar{\Lambda}^{[\alpha\beta]\gamma}\right)(\gamma^\mu)^{\alpha'}_{\alpha}\left(\Lambda_{[\gamma\alpha']\beta} - \Lambda_{[\alpha'\beta]\gamma}\right)\phi_\mu, \qquad \text{(A32)}
\end{aligned}$$

$$= -\frac{\sqrt{2}}{8}\left(\bar{S}_u+\bar{S}_d\right) G_{8V}\left[-\bar{\Lambda}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} \Lambda_{[\alpha'\beta]\gamma}+\bar{\Lambda}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} \Lambda_{[\alpha'\gamma]\beta}\right] \phi_\mu, \quad \text{(A33)}$$

$$= \frac{\sqrt{2}}{2}\left(\bar{S}_u+\bar{S}_d\right) G_{8V}\, \bar{\psi}_\Lambda\, \gamma^\mu\, \psi_\Lambda\, \phi_\mu. \quad \text{(A34)}$$

Using the prescription $\bar{S}_{u(d)} \rightarrow \bar{S}^\Lambda_{u(d)}$ and noting $g_{\Lambda\Lambda\phi} = -\sqrt{2}\,G_{8V}$, the effective $\Lambda\Lambda\phi$ coupling becomes

$$\Gamma^*_{\Lambda\Lambda\phi} = (-i)\, g^*_{\Lambda\Lambda\phi}\, \bar{\psi}_\Lambda\, \gamma^\mu\, \psi_\Lambda\, \phi_\mu, \quad \text{(A35)}$$

$$g^*_{\Lambda\Lambda\phi} = \left[1+\frac{1}{2}\left(\bar{S}^\Lambda_u+\bar{S}^\Lambda_d\right)\right] g_{\Lambda\Lambda\phi}. \quad \text{(A36)}$$

This is none other than Eq. (10).

The medium correction to $pp\rho$ coupling is

$$\begin{aligned}(-i)\,\Delta\Gamma_{pp\rho} = {} & \frac{1}{24}\left(\bar{S}_u+\bar{S}_d\right) G_{8V}\left(\bar{P}^{[\beta\gamma]\alpha}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(P_{[\beta\gamma]\alpha'}-P_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3} \\ & -\frac{1}{24}\bar{S}_u\, G_{8V}\left(\bar{P}^{[\beta\gamma]\alpha}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(P_{[\gamma\alpha']\beta}-P_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3} \\ & -\frac{1}{24}\bar{S}_u\, G_{8V}\left(\bar{P}^{[\gamma\alpha]\beta}-\bar{P}^{[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(P_{[\beta\gamma]\alpha'}-P_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3}, \end{aligned} \quad \text{(A37)}$$

$$\begin{aligned} = {} & \frac{1}{24}\bar{S}_u\, G_{8V}\left[-3\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} P_{[\alpha'\beta]\gamma}+6\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} P_{[\alpha'\gamma]\beta}\right]\rho_{\mu 3} \\ & +\frac{1}{24}\bar{S}_d\, G_{8V}\left[-5\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} P_{[\alpha'\beta]\gamma}+4\bar{P}^{[\beta\alpha]\gamma}\left(\gamma^\mu\right)_\alpha^{\alpha'} P_{[\alpha'\gamma]\beta}\right]\rho_{\mu 3}, \end{aligned} \quad \text{(A38)}$$

$$= -\bar{S}_d\, G_{8V}\, \bar{\psi}_p\, \gamma^\mu\, \psi_p\, \rho_{\mu 3}. \quad \text{(A39)}$$

Using the prescription $\bar{S}_d \rightarrow \bar{S}^N_d$ and noting $g_{NN\rho} = G_{8V}$, the effective $pp\rho$ coupling becomes

$$\Gamma^*_{pp\rho} = (-i)\, g^*_{pp\rho}\, \bar{\psi}_p\, \gamma^\mu\, \psi_p\, \rho_{\mu 3}, \quad \text{(A40)}$$

$$g^*_{pp\rho} = \left(1+\bar{S}^N_d\right) g_{NN\rho}. \quad \text{(A41)}$$

This is just equivalent to Eq. (15) in Ref. [25] or Eq. (3).

The medium correction to $\Sigma^+\Sigma^+\rho$ coupling is

$$\begin{aligned}(-i)\,\Delta\Gamma_{\Sigma^+\Sigma^+\rho} = {} & \frac{1}{24}\left(\bar{S}_u+\bar{S}_s\right) G_{8V}\left(\bar{\Sigma}^{+[\beta\gamma]\alpha}-\bar{\Sigma}^{+[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(\Sigma^+_{[\beta\gamma]\alpha'}-\Sigma^+_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3} \\ & -\frac{1}{24}\bar{S}_u\, G_{8V}\left(\bar{\Sigma}^{+[\beta\gamma]\alpha}-\bar{\Sigma}^{+[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(\Sigma^+_{[\gamma\alpha']\beta}-\Sigma^+_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3} \\ & -\frac{1}{24}\bar{S}_u\, G_{8V}\left(\bar{\Sigma}^{+[\gamma\alpha]\beta}-\bar{\Sigma}^{+[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(\Sigma^+_{[\beta\gamma]\alpha'}-\Sigma^+_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3} \\ & +\frac{1}{24}\bar{S}_u\, G_{8V}\left(\bar{\Sigma}^{+[\gamma\alpha]\beta}-\bar{\Sigma}^{+[\alpha\beta]\gamma}\right)\left(\gamma^\mu\right)_\alpha^{\alpha'}\left(\Sigma^+_{[\gamma\alpha']\beta}-\Sigma^+_{[\alpha'\beta]\gamma}\right)\rho_{\mu 3}, \end{aligned} \quad \text{(A42)}$$

$$= \frac{1}{24}\left(\bar{S}_u + \bar{S}_s\right) G_{8V} \left[-5\bar{\Sigma}^{+[\beta\alpha]\gamma} \left(\gamma^\mu\right)_\alpha^{\alpha'} \Sigma^+_{[\alpha'\beta]\gamma} + 4\,\bar{\Sigma}^{+[\beta\alpha]\gamma} \left(\gamma^\mu\right)_\alpha^{\alpha'} \Sigma^+_{[\alpha'\gamma]\beta} \right] \rho_{\mu 3}, \tag{A43}$$

$$= -\left(\bar{S}_u + \bar{S}_s\right) G_{8V}\, \bar{\psi}_{\Sigma^+} \gamma^\mu \psi_{\Sigma^+} \rho_{\mu 3}. \tag{A44}$$

Using the prescription $\bar{S}_{u(s)} \to \bar{S}^{\Sigma^+}_{u(s)}$ and noting $g_{\Sigma\Sigma\rho} = 2G_{8V}$, the effective $\Sigma^+\Sigma^+\rho$ coupling becomes

$$\Gamma^*_{\Sigma^+\Sigma^+\rho} = (-i)\, g^*_{\Sigma^+\Sigma^+\rho}\, \bar{\psi}_{\Sigma^+} \gamma^\mu \psi_{\Sigma^+} \rho_{\mu 3}, \tag{A45}$$

$$g^*_{\Sigma^+\Sigma^+\rho} = \left[1 + \frac{1}{2}\left(\bar{S}^{\Sigma^+}_u + \bar{S}^{\Sigma^+}_s \right) \right] g_{\Sigma\Sigma\rho}. \tag{A46}$$

This is no less than Eq. (19).

The other couplings can be derived similarly. The scalar meson couplings are derived by analogies with the corresponding vector mesons.

References

[1] S. Balberg, I. Lichtenstadt and G.B. Cook, Astrophys. J. Suppl. **121** (1999) 515.

[2] J.M. Lattimer and M. Prakash, Astrophys. J. **550** (2001) 426.

[3] M. Baldo, G.F. Burgio and H.-J. Schulze, Phys. Rev. **C61** (2000) 055801, [arXiv:nucl-th/9912066].

[4] I. Vidaña, A. Polls, A. Ramos, L. Engvik and M. Hjorth-Jensen, Phys. Rev. **C62** (2000) 035801, [arXiv:nucl-th/0004031].

[5] J.M. Weisberg and J.H. Taylor, Proceedings of August 2002 meeting "Radio Pulsars" in Chania, Crete, ASP Conference Series Vol. CS-302 (2003), M. Bailes, D.J. Nice and S.E. Thorsett, eds. [arXiv:astro-ph/0211217].

[6] C. Albertus, J.E. Amaro and J. Nieves, Phys. Rev. Lett. **89** (2002) 032501, [arXiv:nucl-th/0110046].

[7] I.N. Filikhin and A. Gal, Phys. Rev. **C65** (2002) 041001(R); Nucl. Phys. **A707** (2002) 491, [arXiv:nucl-th/0203036]; I.N. Filikhin, A. Gal and V.M. Suslov, Phys. Rev. **C68** (2003) 024002, [arXiv:nucl-th/0303028]; Nucl. Phys. **A743** (2004) 194, [arXiv:nucl-th/0406049];

[8] K.S. Myint, S. Shinmura and Y. Akaishi, Eur. Phys. J. **A16** (2003) 21, [arXiv:nucl-th/0209090].

[9] T. Yamada, Phys. Rev. **C69** (2004) 044301, [arXiv:nucl-th/0403016].

[10] I. Vidaña, A. Ramos and A. Polls, Phys. Rev. **C70** (2004) 024306, [arXiv:nucl-th/0307096].

[11] H. Nemura, S. Shinmura, Y. Akaishi and K.H. Myint, arXiv:nucl-th/0407033.

[12] H. Takahashi *et al.*, Phys. Rev. Lett. **87** (2001) 212502.

[13] R. Brockmann and R. Machleidt, Phys. Rev. **C42** (1990) 1965.

[14] B.D. Serot and J.D. Walecka, *Advances in Nuclear Physics*, Vol. **16** (Plenum, New York, 1986).

[15] S. Gmuca, Nucl. Phys. **A547** (1992) 447.

[16] J. Schaffner, C.B. Dover, A. Gal, C. Greiner and H. Stőker, Phys. Rev. Lett. **71** (1993) 1328.

[17] F. Hofmann, C.M. Keil and H. Lenske, Phys. Rev. **C64** (2001) 025804, [arXiv:nucl-th/0008038].

[18] J. Zimanyi and S.A. Moszkowski, Phys. Rev. **C42** (1990) 1416.

[19] P.A.M. Guichon, Phys. Lett. **B200** (1988) 235; K. Saito and A.W. Thomas, Phys. Lett. **B327** (1994) 9; Phys. Rev. **C52** (1995) 2789.

[20] I. Bednarek, M. Keska and R. Manka, Phys. Rev. **C68** (2003) 035805, [arXiv:nucl-th/0212065].

[21] S. Pal, M. Hanauske, I. Zakout, H. Stőker and W. Greiner, Phys. Rev. **C60** (1999) 015802, [arXiv:nucl-th/9905010].

[22] K. Miyazaki, CERN Document Server (CDS) ext-2003-062 revised by Mathematical Physics Preprint Archive (mp_arc) 05-178.

[23] K. Miyazaki, CERN Document Server (CDS) ext-2003-077 revised by Mathematical Physics Preprint Archive (mp_arc) 05-190.

[24] M. Barranco, R.J. Lombard, S. Marcos and S.A. Moszkowski, Phys. Rev. **C44** (1991) 178.

[25] N.K. Glendenning, F. Weber and S.A. Moszkowski, Phys. Rev. **C45** (1992) 844.

[26] K. Miyazaki, CERN Document Server (CDS) ext-2004-010 revised by Mathematical Physics Preprint Archive (mp_arc) 05-199.

[27] K. Miyazaki, CERN Document Server (CDS) ext-2004-036 revised by Mathematical Physics Preprint Archive (mp_arc) 05-206.

[28] K. Miyazaki, CERN Document Server (CDS) ext-2004-071 revised by Mathematical Physics Preprint Archive (mp_arc) 05-216.

[29] N.K. Glendenning and S.A. Moszkowski, Phys. Rev. Lett. **67** (1991) 2414.

[30] J. Schaffner-Bielich, M. Hanauske, H. Stőker and W. Greiner, Phys. Rev. Lett. **89** (2002) 171101, [arXiv:astro-ph/0005490].

[31] V.G.J. Stoks and Th.A. Rijken, Phys. Rev. **C59** (1999) 3009, [arXiv:nucl-th/9901028].

[32] Th.A. Rijken, V.G.J. Stoks and Y. Yamamoto, Phys. Rev. **C59** (1999) 21, [arXiv:nucl-th/9807082].

[33] V.G.J. Stoks and T.-S.H. Lee, Phys. Rev. **C60** (1999) 024006, [arXiv:nucl-th/9901030].

[34] J. Schaffner-Bielich and A. Gal, Phys. Rev. **C62** (2000) 034311, [arXiv:nucl-th/0005060].

[35] W.D. Arnett and R.L. Bowers, Astrophys. J. Suppl. **33** (1977) 415.

[36] V. Canuto, Ann. Rev. Astr. Ap. **12** (1974) 167; **13** (1975) 335.

[37] J. Mareš, E. Friedman, A. Gal and B.K. Jennings, Nucl. Phys. **A594** (1995) 311, [arXiv:nucl-th/9505003].

[38] M. Kohno, Y. Fujiwara, Y. Watanabe, K. Ogata and M. Kawai, arXiv:nucl-th/0410073.

[39] J. Schaffner and I.N. Mishustin, Phys. Rev. **C53** (1996) 1416, [arXiv:nucl-th/9506011].

[40] N.K. Glendenning, Phys. Rev. **C64** (2001) 025801, [arXiv:nucl-th/0009082].

[41] S.E. Thorsett and D. Chakrabarty, Astrophys. J. **512** (1999) 288.

[42] I.H. Stairs, S.E. Thorsett, J.H.Taylor and A. Wolszczan, Astrophys. J. **581** (2002) 501, [arXiv:astro-ph/0208357].

[43] M. Bailes, S.M. Ord, H.S. Knight and A.W. Hotan, Astrophys. J. **595** (2003) L49. [arXiv:astro-ph/0307468].

[44] A.G. Lyne *et al.*, Science **303** (2004) 1153, [arXiv:astro-ph/0401086].

[45] M.H. van Kerkwijk, arXiv:astro-ph/0403489.

[46] D. Stickland, C. Lloyd and A. Radziun-Woodham, MNRAS **286** (1997) L21.

[47] M. Salgado, S. Bonazzola, E. Gourgoulhon and P. Haensel, A&A **291** (1994) 155.

[48] B. Link, R.I. Epstein and J.M. Lattimer, Phys. Rev. Lett. **83** (1999) 3362.

[49] A. Salam, R. Delbourgo and J. Strathdee, Proc. R. Soc. London **A284** (1965) 146; A. Salam, R. Delbourgo, M.A. Rashid and J. Strathdee, Proc. R. Soc. London **A285** (1965) 312.

[50] B. Sakita and K.C. Wali, Phys. Rev. **89** (1965) B1355.

[51] F. Hussain, J.G. Kőrner and G. Thompson, Ann. of Phys. **206** (1991) 334.

[52] K. Miyazaki, Prog. Theor. Phys. **89** (1993) 639.

[53] Y. Takahashi, *An Introduction to Field Quantization* (Pergamon Press, Oxford, 1969); W. Greiner, *Relativistic Quantum Mechanics* (Springer-Verlag, Berlin, 1990).

Table 1: The four sets of meson-hyperon coupling constants used in the calculations of NS matter. The derivations of each value are discussed in section 3.

Coupling constants	Set 1	Set 2	Set 3	Set 4
$g_{\Lambda\Lambda\omega}$	$(2/3)\,g_{NN\omega}$			
$g_{\Sigma\Sigma\omega}$	$(2/3)\,g_{NN\omega}$			
$g_{\Xi\Xi\omega}$	$(1/3)\,g_{NN\omega}$			
$g_{\Lambda\Lambda\sigma}$	$(2/3)\,g_{NN\sigma}$	$0.604\,g_{NN\sigma}$		
$g_{\Sigma\Sigma\sigma}$	$(2/3)\,g_{NN\sigma}$	$0.461\,g_{NN\sigma}$	$0.485\,g_{NN\sigma}$	$0.461\,g_{NN\sigma}$
$g_{\Xi\Xi\sigma}$	$(1/3)\,g_{NN\sigma}$	$0.309\,g_{NN\sigma}$		
$g_{\Lambda\Lambda\phi}$	$-(\sqrt{2}/3)\,g_{NN\omega}$			
$g_{\Sigma\Sigma\phi}$	$-(\sqrt{2}/3)\,g_{NN\omega}$			
$g_{\Xi\Xi\phi}$	$-(2\sqrt{2}/3)\,g_{NN\omega}$			
$g_{\Lambda\Lambda\sigma^*}$	$(\sqrt{2}/3)\,g_{NN\sigma}$	$0.690\,g_{NN\sigma}$	$0.52\,g_{NN\sigma}$	
$g_{\Sigma\Sigma\sigma^*}$	$(\sqrt{2}/3)\,g_{NN\sigma}$	$0.690\,g_{NN\sigma}$	$0.96\,g_{NN\sigma}$	$0.52\,g_{NN\sigma}$
$g_{\Xi\Xi\sigma^*}$	$(2\sqrt{2}/3)\,g_{NN\sigma}$	$1.221\,g_{NN\sigma}$	$1.28\,g_{NN\sigma}$	
$g_{\Sigma\Sigma\rho}$	$2\,g_{NN\rho}$			
$g_{\Xi\Xi\rho}$	$g_{NN\rho}$			
$g_{\Sigma\Sigma\delta}$	$2\,g_{NN\delta}$			
$g_{\Xi\Xi\delta}$	$g_{NN\delta}$			

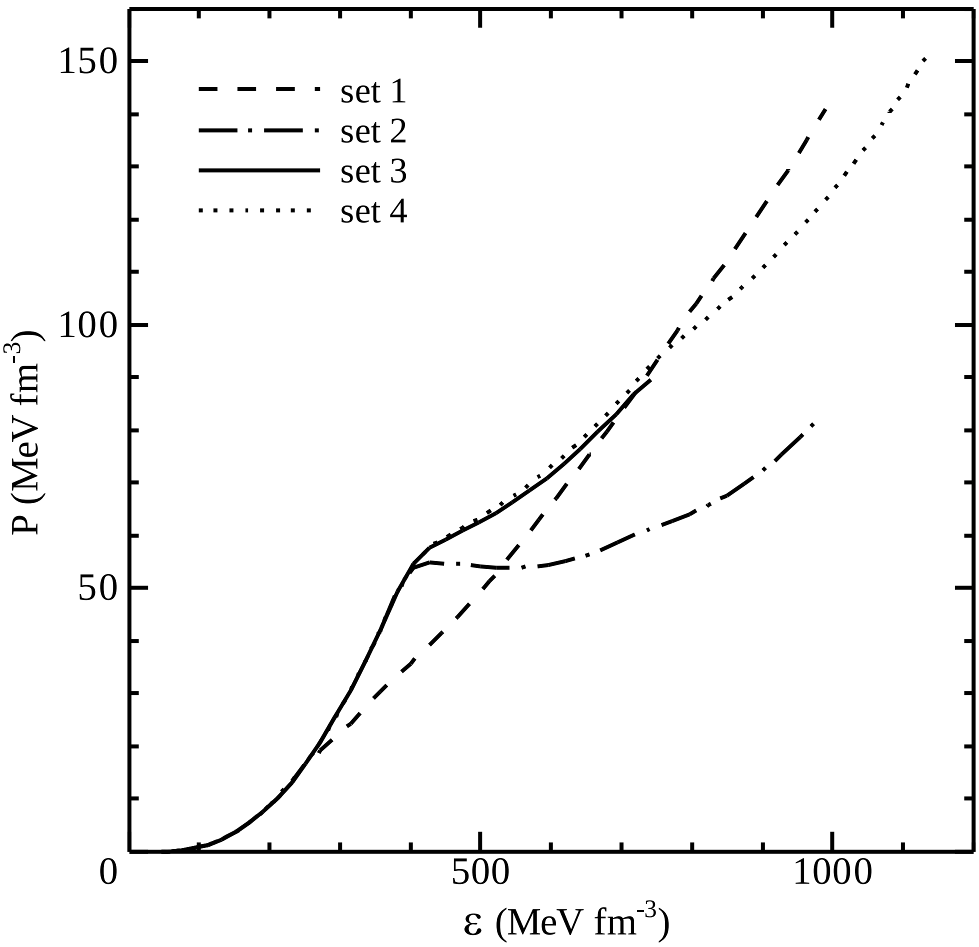


Figure 1: The EOSs of β-equilibrated NS matter. The dashed, dashed-dotted, solid and dotted curves are the results using the coupling sets 1, 2, 3 and 4 respectively.

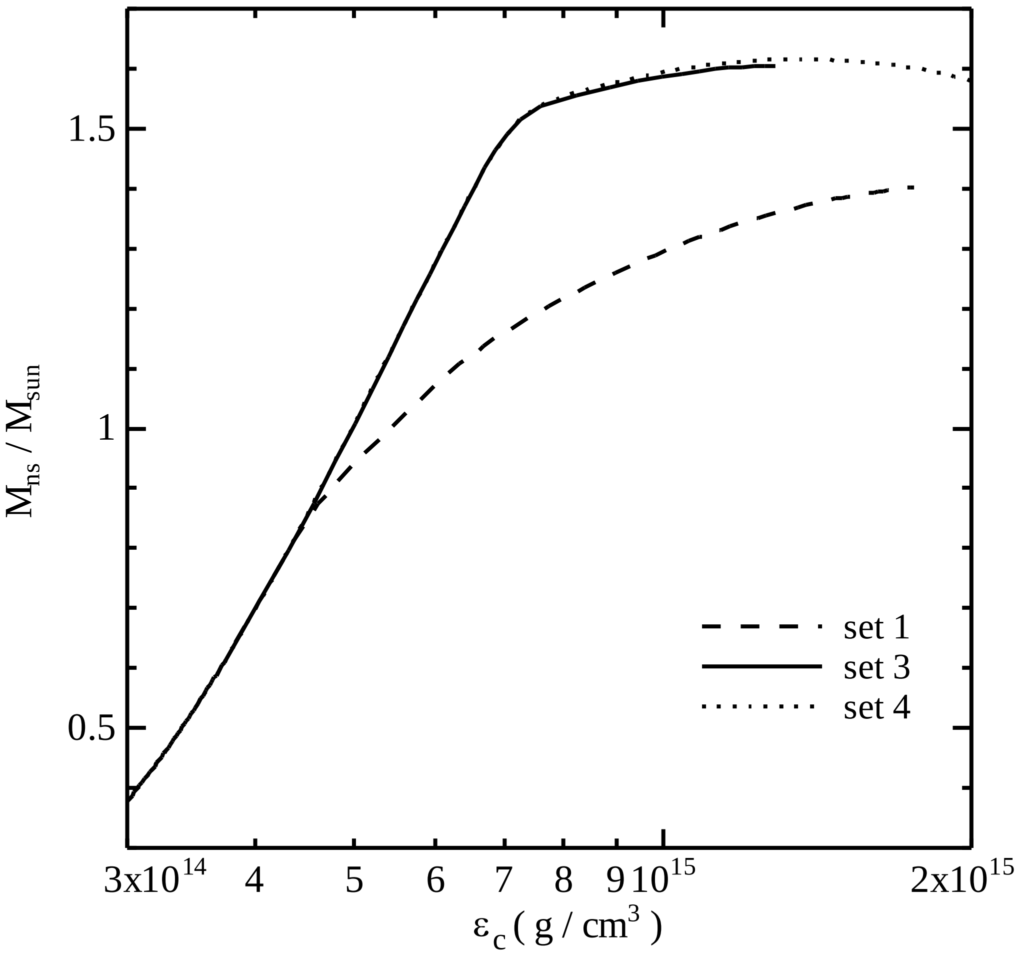


Figure 2: The mass sequences of non-rotating cold β-equilibrated NS using the meson-hyperon coupling sets 1, 3 and 4.

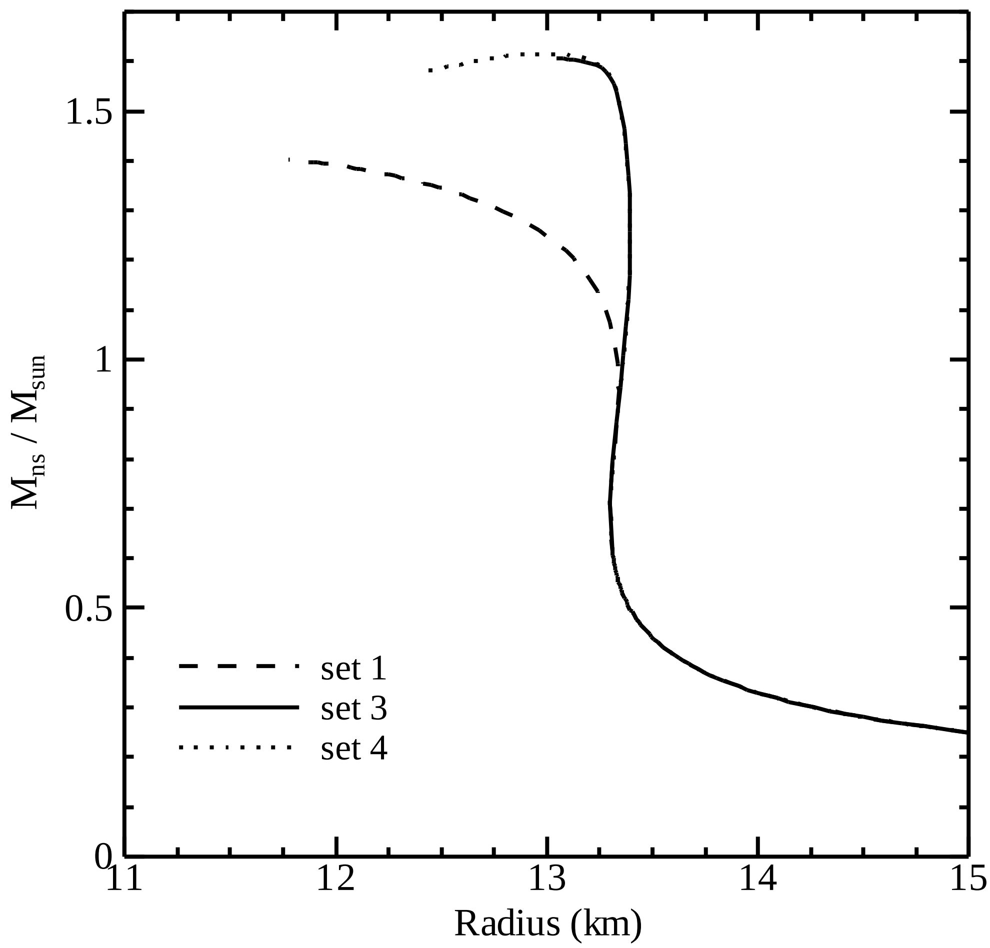


Figure 3: The mass-radius relations of non-rotating cold β-equilibrated NS using the meson-hyperon coupling sets 1, 3 and 4.

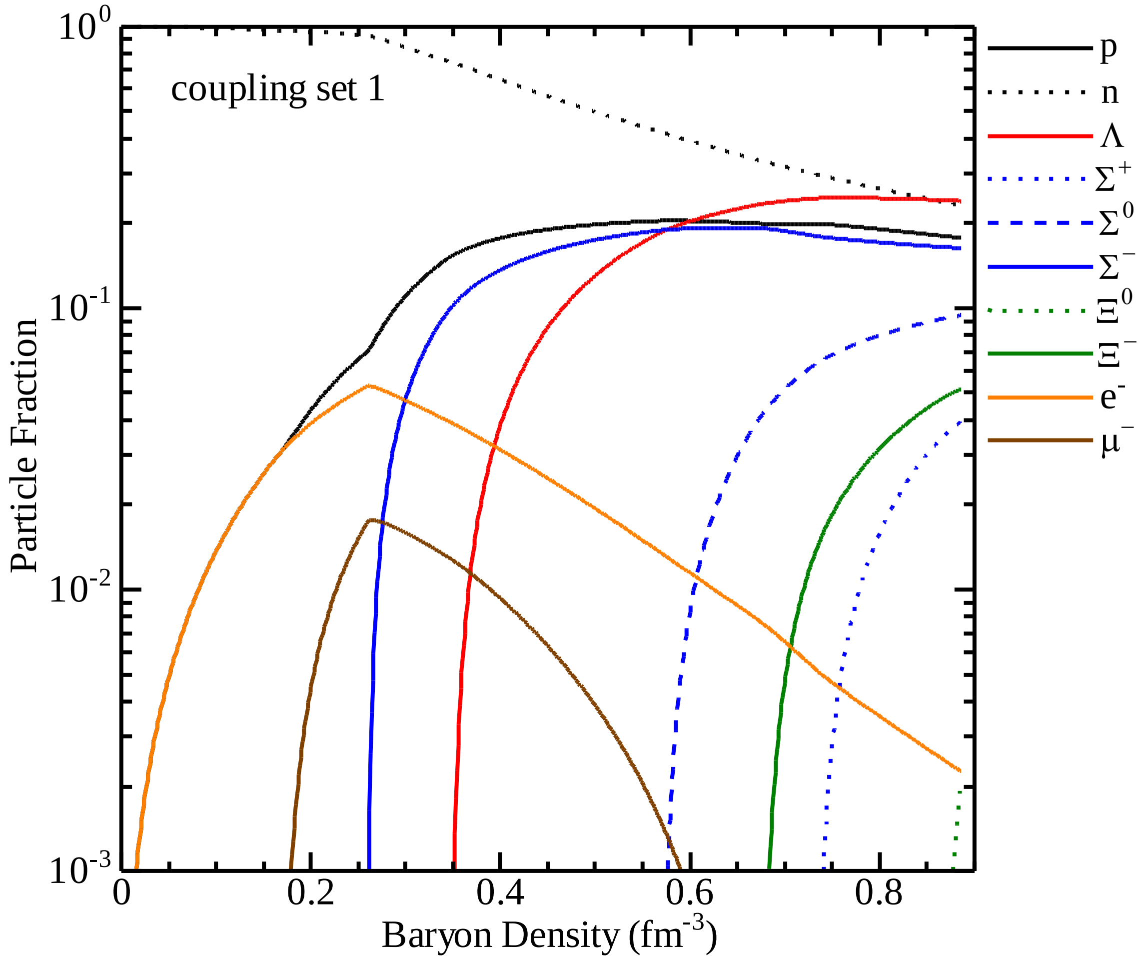


Figure 4: The fractions of baryons and leptons in the β-equilibrated NS matter using the meson-hyperon coupling set 1.

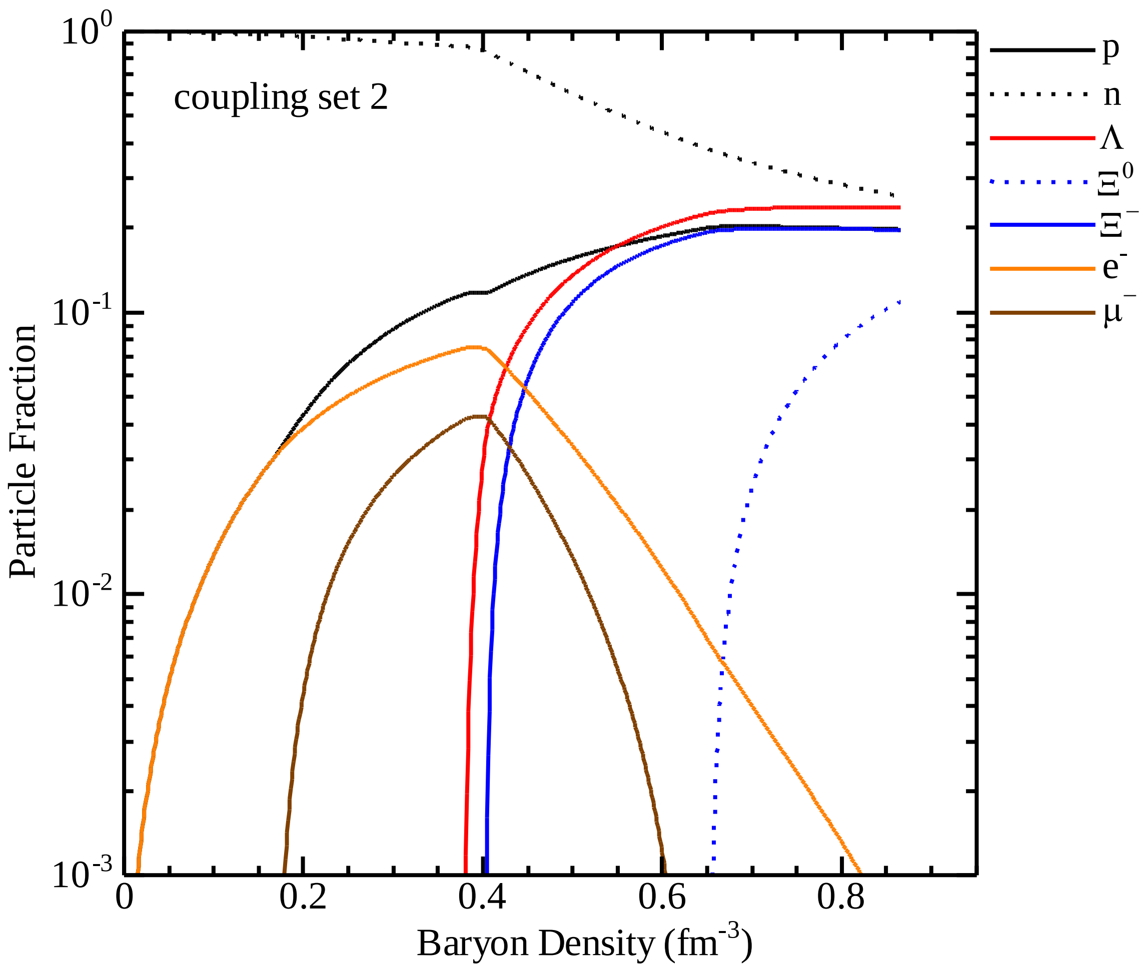


Figure 5: The same as Fig. 4 but using the meson-hyperon coupling set 2.

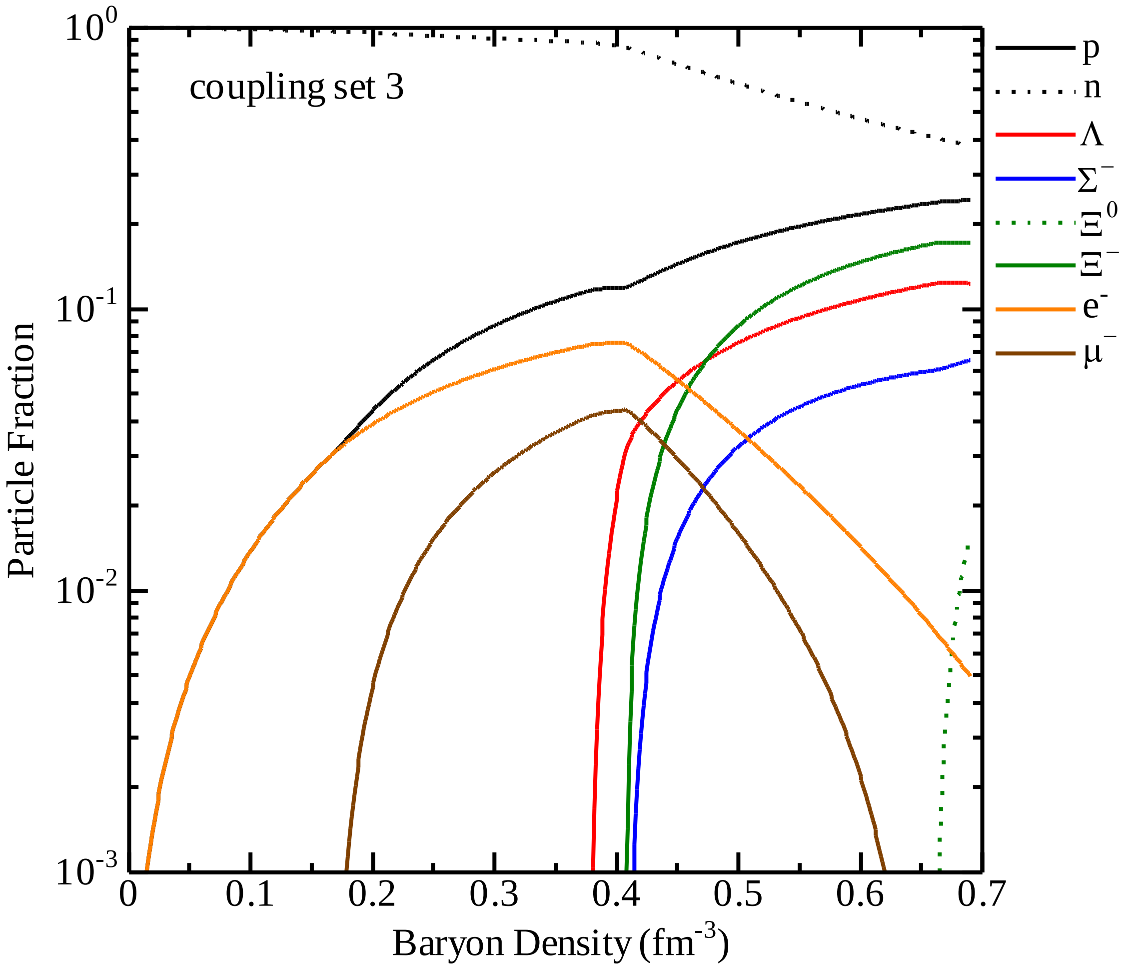


Figure 6: The same as Fig. 4 but using the meson-hyperon coupling set 3.

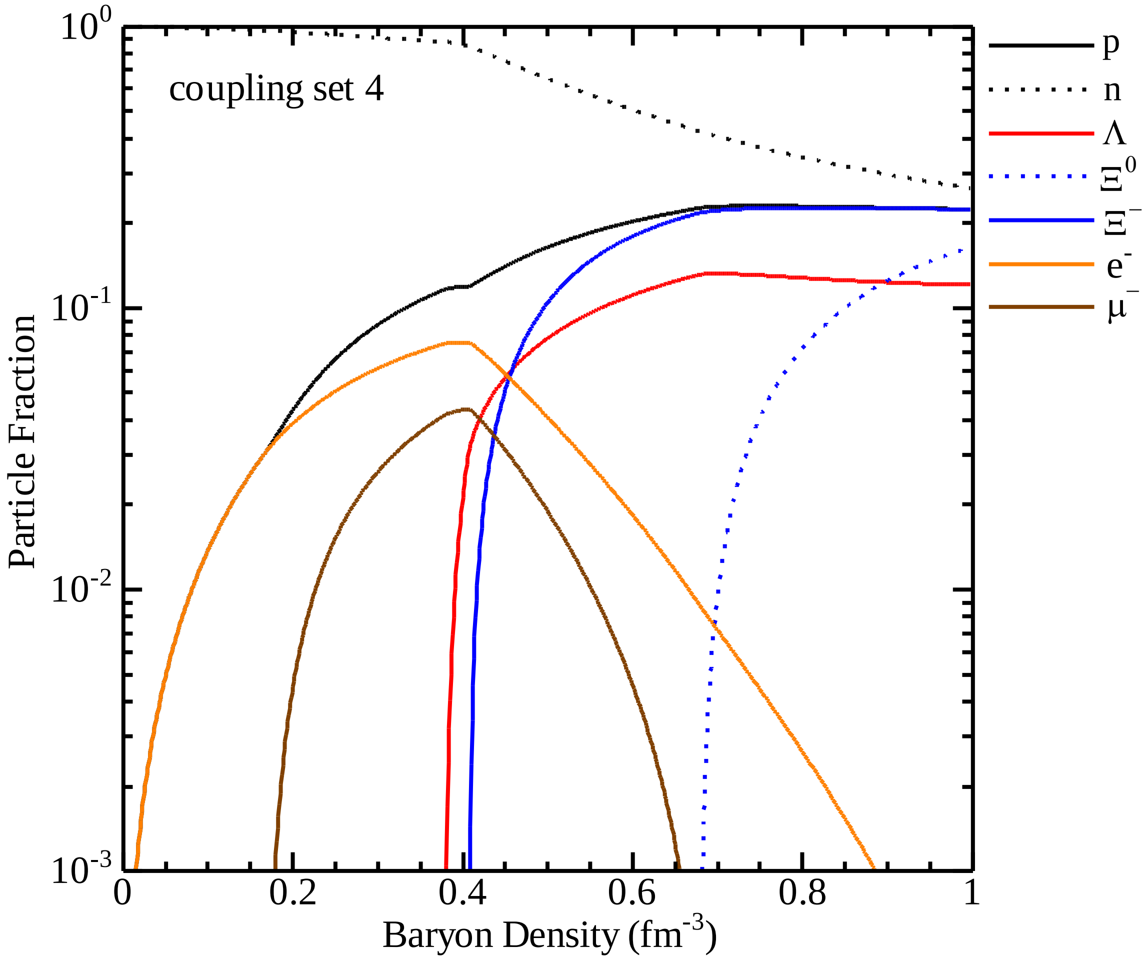


Figure 7: The same as Fig. 4 but using the meson-hyperon coupling set 4.