Physics Letters B 747 (2015) 43–47

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Hyperons in neutron stars

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ARTICLE INFO

Article history:
Received 23 January 2015
Received in revised form 13 March 2015
Accepted 18 March 2015
Available online 23 March 2015
Editor: W. Haxton

Keywords:
Neutron stars
Hyperons in matter
Equation of state
Dirac–Brueckner–Hartree–Fock approach

ABSTRACT

Using the Dirac–Brueckner–Hartree–Fock approach, the properties of neutron-star matter including hyperons are investigated. In the calculation, we consider both time and space components of the vector self-energies of baryons as well as the scalar ones. Furthermore, the effect of negative-energy states of baryons is partly taken into account. We obtain the maximum neutron-star mass of $2.08 M_\odot$, which is consistent with the recently observed, massive neutron stars. We discuss a universal, repulsive three-body force for hyperons in matter.

Neutron stars may be the most dense and exotic state of nuclear matter, and its core serves as a natural laboratory to investigate the nuclear matter whose density reaches several times higher than the normal nuclear-matter density, n_B^0 [1]. In fact, the recently observed, massive neutron stars, J1614-2230 (the mass of $1.97 \pm 0.04 M_\odot$, $M_\odot$: the solar mass) [2] and J0348+0432 ($2.01 \pm 0.04 M_\odot$) [3], have provided important information on the equation of state (EoS) for dense nuclear matter.

To understand these heavy objects, various nuclear models have been examined, in which relativistic mean-field theory (RMFT) is very popular and has been successfully applied to the dense nuclear matter [4]. However, in RMFT, nucleon (N)-nucleon short-range correlations in matter cannot be treated. In contrast, in the Dirac–Brueckner–Hartree–Fock (DBHF) approach, although the calculation is involved, one can consider the effects of the Pauli exclusion principle and short-range correlations.

Until now, several groups have performed the DBHF calculations not only in the region around n_B^0 but also in matter at higher densities (see Refs. [5–15]). However, so far there has not been any relativistic attempt to take account of the degrees of freedom of hyperons (Ys) as well as nucleons in dense matter. Because it is quite interesting to see how hyperons contribute to the EoS and to the maximum mass of neutron stars, it seems very urgent to perform the DBHF calculation for matter including hyperons.

In this Letter, we study such dense neutron-star matter using the DBHF approach. Here, we particularly pay attention to the following two points: (1) the space component of vector self-energy of baryon (B), Σ_B^V, is taken into account, because, although it is certainly small at low density, it is expected to be important in dense matter, (2) as in Refs. [13–15], we partly consider the effect of negative-energy states of baryons in the Bethe–Salpeter (BS) equation to remove the ambiguity in the relationship between the on-shell T-matrix for baryon–baryon scattering and the baryon self-energies [8–10]. Furthermore, when hyperons take place in matter, the effective masses of interacting two baryons become very different from each other, and thus we should treat the baryon-mass difference in the BS equation explicitly.

We now start with the self-energy of baryon in the rest frame of infinite, uniform nuclear matter. It is given by

$$\Sigma_B(k) = \Sigma_B^S(k) - \gamma_0 \Sigma_B^0(k) + \boldsymbol{\gamma} \cdot \boldsymbol{k} \Sigma_B^V(k), \quad (1)$$

where $\boldsymbol{k}$ (k) is the three (four) momentum of baryon. Here, $\Sigma_B^{S\,(0)\,[V]}$ is the scalar (zero-th component of vector) [space component of vector] part of baryon self-energy. Using these self-energies, the effective mass, M_B^*, the effective momentum, $\boldsymbol{k}_B^*$, and the effective energy, E_B^*, in matter are defined by

$$M_B^*(k) \equiv M_B + \Sigma_B^S(k), \quad \boldsymbol{k}_B^* \equiv \boldsymbol{k}[1 + \Sigma_B^V(k)],$$
$$E_B^*(k) \equiv \sqrt{\boldsymbol{k}_B^{*2} + M_B^{*2}(k)}, \quad (2)$$

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http://dx.doi.org/10.1016/j.physletb.2015.03.039

with M_B being the free baryon mass. Then, the baryon spinor states with positive or negative energy are respectively constructed as

$$\Phi_B(\boldsymbol{k}, s) = \sqrt{M_B^*(k) + E_B^*(k)} \begin{pmatrix} \chi_s, \\ \frac{\boldsymbol{k}_B^* \cdot \boldsymbol{\sigma}}{M_B^*(k)+E_B^*(k)} \chi_s \end{pmatrix}, \tag{3}$$

$$\Theta_B(\boldsymbol{k}, s) = \sqrt{M_B^*(k) + E_B^*(k)} \begin{pmatrix} \frac{\boldsymbol{k}_B^* \cdot \boldsymbol{\sigma}}{M_B^*(k)+E_B^*(k)} \chi_{-s} \\ \chi_{-s} \end{pmatrix}, \tag{4}$$

where $\boldsymbol{\sigma}$ is the Pauli matrix, and χ_s denotes a 2-component Pauli spinor.

In the conventional DBHF calculation, the baryon–baryon scattering is usually evaluated in the center of mass frame with respect to the interacting two baryons. In such cases, instead of Eqs. (3)–(4), the helicity spinors and the partial-wave decomposition are often used to solve the BS equation [5–11,13–15]. However, when Σ_B^V remains finite and $\boldsymbol{k} \neq \boldsymbol{k}_B^*$, although $\boldsymbol{k}$ and $\boldsymbol{k}_B^*$ are parallel with each other in the nuclear-matter rest frame, they are not in the center of mass frame. It is thus more convenient to perform the calculation with the standard spinors, Eqs. (3)–(4), in the nuclear-matter rest frame, rather than with the helicity spinors in the center of mass frame.

Furthermore, the inclusion of negative-energy states of baryon in the BS amplitude may be necessary to remove the ambiguity of the relationship between the reaction matrices for baryon–baryon scattering and the baryon self-energies [13–15]. Thus, we here define four reaction amplitudes

$$\begin{aligned} &T_{B'''B''B'B}(\boldsymbol{k}', \boldsymbol{k}, s''', s'', s', s; \boldsymbol{P}) \\ &\equiv \bar{\Phi}_{B'''}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}', s'''\right) \bar{\Phi}_{B''}\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}', s''\right) \\ &\quad \times \Gamma\, \Phi_{B'}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}, s'\right) \Phi_B\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}, s\right), \end{aligned} \tag{5}$$

$$\begin{aligned} &R_{B'''B''B'B}(\boldsymbol{k}', \boldsymbol{k}, s''', s'', s', s; \boldsymbol{P}) \\ &\equiv \bar{\Theta}_{B'''}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}', s'''\right) \bar{\Phi}_{B''}\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}', s''\right) \\ &\quad \times \Gamma\, \Phi_{B'}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}, s'\right) \Phi_B\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}, s\right), \end{aligned} \tag{6}$$

$$\begin{aligned} &O_{B'''B''B'B}(\boldsymbol{k}', \boldsymbol{k}, s''', s'', s', s; \boldsymbol{P}) \\ &\equiv \bar{\Phi}_{B'''}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}', s'''\right) \bar{\Phi}_{B''}\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}', s''\right) \\ &\quad \times \Gamma\, \Theta_{B'}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}, s'\right) \Phi_B\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}, s\right), \end{aligned} \tag{7}$$

$$\begin{aligned} &P_{B'''B''B'B}(\boldsymbol{k}', \boldsymbol{k}, s''', s'', s', s; \boldsymbol{P}) \\ &\equiv \bar{\Theta}_{B'''}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}', s'''\right) \bar{\Phi}_{B''}\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}', s''\right) \\ &\quad \times \Gamma\, \Theta_{B'}\left(\frac{1}{2}\boldsymbol{P} + \boldsymbol{k}, s'\right) \Phi_B\left(\frac{1}{2}\boldsymbol{P} - \boldsymbol{k}, s\right), \end{aligned} \tag{8}$$

where Γ represents the effective reaction operator, and these amplitudes satisfy the following, coupled BS equations

$$\begin{aligned} &T_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &= \bar{V}_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &\quad + \sum_{s''s'''B''B'''} \int \frac{d^3q}{(2\pi)^4} \bar{V}_{BB'B''B'''}(\boldsymbol{k}, \boldsymbol{q}, s, s', s'', s'''; \boldsymbol{P}) \\ &\quad \times Q_{B''B'''}(\boldsymbol{P}, \boldsymbol{q}) g_{Th\,B''B'''}(\boldsymbol{P}, \boldsymbol{q}) \\ &\quad \times T_{B''B'''BB'}(\boldsymbol{q}, \boldsymbol{k}, s''', s'', s, s'; \boldsymbol{P}), \end{aligned} \tag{9}$$

$$\begin{aligned} &R_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &= \bar{U}_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &\quad + \sum_{s''s'''B''B'''} \int \frac{d^3q}{(2\pi)^4} \bar{U}_{BB'B''B'''}(\boldsymbol{k}, \boldsymbol{q}, s, s', s'', s'''; \boldsymbol{P}) \\ &\quad \times Q_{B''B'''}(\boldsymbol{P}, \boldsymbol{q}) g_{Th\,B''B'''}(\boldsymbol{P}, \boldsymbol{q}) \\ &\quad \times T_{B''B'''BB'}(\boldsymbol{q}, \boldsymbol{k}, s''', s'', s, s'; \boldsymbol{P}), \end{aligned} \tag{10}$$

$$\begin{aligned} &O_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &= \bar{W}_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &\quad + \sum_{s''s'''B''B'''} \int \frac{d^3q}{(2\pi)^4} \bar{V}_{BB'B''B'''}(\boldsymbol{k}, \boldsymbol{q}, s, s', s'', s'''; \boldsymbol{P}) \\ &\quad \times Q_{B''B'''}(\boldsymbol{P}, \boldsymbol{q}) g_{Th\,B''B'''}(\boldsymbol{P}, \boldsymbol{q}) \\ &\quad \times O_{B''B'''BB'}(\boldsymbol{q}, \boldsymbol{k}, s''', s'', s, s'; \boldsymbol{P}), \end{aligned} \tag{11}$$

$$\begin{aligned} &P_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &= \bar{Z}_{BB'BB'}(\boldsymbol{k}, \boldsymbol{k}, s, s', s, s'; \boldsymbol{P}) \\ &\quad + \sum_{s''s'''B''B'''} \int \frac{d^3q}{(2\pi)^4} \bar{U}_{BB'B''B'''}(\boldsymbol{k}, \boldsymbol{q}, s, s', s'', s'''; \boldsymbol{P}) \\ &\quad \times Q_{B''B'''}(\boldsymbol{P}, \boldsymbol{q}) g_{Th\,B''B'''}(\boldsymbol{P}, \boldsymbol{q}) \\ &\quad \times O_{B''B'''BB'}(\boldsymbol{q}, \boldsymbol{k}, s''', s'', s, s'; \boldsymbol{P}), \end{aligned} \tag{12}$$

with $\bar{V}, \bar{U}, \bar{W}$ and $\bar{Z}$ being the anti-symmetrized matrices of one-boson-exchange (OBE) interaction [16] with respect to the positive- and negative-energy states (as seen in Eqs. (5)–(8)).

In Eqs. (9)–(12), $Q_{BB'}$ is the Pauli exclusion operator for baryons B and B', and $g_{ThBB'}$ denotes the Thompson's two-particle propagator [17]. The seven arguments in the four reaction amplitudes, T, R, O, P, are as follows: from left to right, the first variable represents the final (or intermediate) relative three-momentum; the second, the initial (or intermediate) relative three-momentum; the third and fourth are for the spins of the final (or intermediate) two baryons, each of which is up (+) or down (−); the fifth and sixth, the spins of the initial (or intermediate) two baryons; and the last one is the total three-momentum of interacting two baryons. We note that the negative-energy states appear only in the initial and/or final states of the BS amplitudes, and they are not included in the intermediate states, because, in the realistic baryon–baryon potentials such as the Bonn potentials, the negative-energy states are usually not considered [13,14].

The ladder-approximated, coupled BS equations can be numerically solved in the nuclear-matter rest frame. To reduce the number of variables and make the present calculation feasible, we here average the azimuthal angle in the spinors, Eqs. (3)–(4), namely we replace $E_B^*(1/2\boldsymbol{P} \pm \boldsymbol{k})$ by the averaged one, $\frac{1}{2\pi}\int d\phi E_B^*(1/2\boldsymbol{P} \pm \boldsymbol{k})$. We have checked that this change does not lead any large numerical error in our final results.

Given the reaction amplitudes, we can calculate the following components [14]

$$\begin{aligned} \Sigma_{\Phi\Phi}^B(k) &\equiv \bar{\Phi}_B(k,+)\, \Sigma_B(k)\, \Phi_B(k,+) \\ &= 2M_B^*(k)\Sigma_B^S(k) - 2E_B^*(k)\Sigma_B^0(k) + 2\boldsymbol{k}\cdot\boldsymbol{k}_B^*\Sigma_B^V(k), \end{aligned} \tag{13}$$

$$\begin{aligned} \Sigma_{\Theta\Phi}^B(k) &\equiv \bar{\Theta}_B(k,+)\, \Sigma_B(k)\, \Phi_B(k,-) \\ &= 2|\boldsymbol{k}_B^*|(k)\Sigma_B^0(k) - 2|\boldsymbol{k}|E_B^*(k)\Sigma_B^V(k), \end{aligned} \tag{14}$$

$$\begin{aligned} \Sigma_{\Theta\Theta}^B(k) &\equiv \bar{\Theta}_B(k,+)\, \Sigma_B(k)\, \Theta_B(k,+) \\ &= -2M_B^*(k)\Sigma_B^S(k) - 2E_B^*(k)\Sigma_B^0(k) + 2\boldsymbol{k}\cdot\boldsymbol{k}_B^*\Sigma_B^V(k), \end{aligned} \tag{15}$$