Possibility of s-wave pion condensates in neutron stars revisited

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We examine possibilities of pion condensation with zero momentum (s-wave condensation) in neutron stars by using the pion-nucleus optical potential $U$ and the relativistic mean field (RMF) models. We use low-density phenomenological optical potentials parameterized to fit deeply bound pionic atoms or pion-nucleus elastic scatterings. Proton fraction ($Y_p$) and electron chemical potential ($\mu_e$) in neutron star matter are evaluated in RMF models. We find that the s-wave pion condensation hardly takes place in neutron stars and especially has no chance if hyperons appear in neutron star matter and/or $b_1$ parameter in $U$ has density dependence.

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Pion condensation in neutron stars has a long history of study from the first suggestion in early 1970s by Migdal [1] and Sawyer [2]. The pion-nucleon interaction is attractive in p-wave, then the main interest along this line was to explore possible appearance of pionic excitations with zero energy and finite momentum, i.e. p-wave pion condensation [3] in nuclear matter. Possibilities of p-wave pion condensation in finite nuclei had been investigated extensively in 1970’s and 1980’s. Those possibilities were denied by the non-observation of anomalous angular momentum distribution in the inelastic excitation of the pionic quantum numbers [4]. Possibilities of p-wave pion condensation at high densities were also considered to be improbable based on the universal repulsion assumption, $g'_{N\Delta} \sim g'_{NN} \sim 0.6 - 0.8$ [5]. In 1990’s, new experiments on the Gamow-Teller giant resonances were performed, and the sum rule value including the $2p - 2h$ states was found to be around 90 % [6], suggesting that the transition to the $\Delta$ region is weak and $g'_{N\Delta} \sim 0.2$ would be smaller than $g'_{NN}$ [7]. In addition, microscopic variational calculation [8] suggests $\pi^0$ condensation in symmetric nuclear matter at high densities ($\rho_0 > 0.2$ fm$^{-3}$) generated from the $\Delta$ mediated three nucleon force. Thus at present, we cannot completely deny the possibility of pion condensation in dense matter, and it is necessary to examine all the ingredients of $\pi N$ and $\pi$-nucleus interactions with updated experimental and theoretical knowledge.

The study of the in-medium pion properties, especially the s-wave pion-nucleon interaction, has been recently developed in experiment. Precise observations of deeply bound atomic states of $\pi^-$ in Pb and Sn isotopes [9,10,11,12,13] and low energy pion-nucleus elastic scattering [13] provide us with detailed information of density dependent optical potentials at low densities. Theoretical calculations of in-medium pion self-energy have also experienced much progress based on chiral dynamics [14,15,16,17,18].

Together with these developments, it may be interesting to revisit pion condensation at high densities, such as in neutron stars. Motivated by the recent progress in the s-wave pion-nucleus interaction, we concentrate on the study of pion condensation with zero momentum (s-wave condensation), for simplicity. Even such a limited investigation would make progress of our understanding of dense matter. The s-wave pion condensation takes place in neutron stars with nuclear matter instability where the transition $n \rightarrow p\pi^-$ becomes energetically possible [19]. This happens with $\mu_n - \mu_p \geq E_\pi$, where $\mu_n$ and $\mu_p$ are the neutron and proton chemical potentials, respectively, and $E_\pi$ is the $\pi^-$ energy at rest. Due to $\beta$ equilibrium under the charge neutral and neutrino-less conditions, $\mu_n - \mu_p$ can be written as the electron chemical potential $\mu_e = \mu_n - \mu_p$. Since $nn$ interaction is more repulsive than $pn$, $\mu_n$ is pushed up then $\mu_e$ increases compared to Fermi gas value in neutron rich matter. For example, relativistic mean field (RMF) models suggest [19,20] that $\mu_e$ largely exceeds the in-vacuum $\pi^-$ mass at nuclear density $\rho_0 = (1 \sim 5)\rho_0$ in neutron stars, where $\rho_0$ is the saturation density.

In this paper, we examine whether the condition for the s-wave $\pi^-$ condensation, $E_\pi = \mu_e$, is satisfied in neutron stars for $E_\pi$ and $\mu_e$ obtained in our present knowledge of low density pion optical potentials and equation of state (EOS) of neutron star matter. For $E_\pi$, we use various pion optical potentials fitted so as to reproduce the pionic atom [12,21,22,23,24,25] and $\pi$-nucleus elastic scattering data [14]. Although the fitted optical potentials in normal nuclei may not be extrapolated to higher density and/or highly asymmetric nuclear matter, it is interesting to examine the present status and to think about next steps. For $\mu_e$, we adopt the results calculated in RMF with several parameter sets [26,27,28,29], which explain the bulk properties of nuclei such as the binding energy and the charge radius in a wide mass range. We also use proton fractions ($Y_p$) evaluated with RMF for calculating pion optical potentials$^1$. There are several theoretical works on the s-wave pion condensation at higher densities [30], while the connection to low density phenomena observed in experiments is not clear yet. Since we concentrate on the possibility of the s-wave $\pi^-$ condensate, we do not consider the double pole condition for $\pi^- \pi^+_s$ pair creation, which takes place with finite momentum.

The $\pi^-$ energy in uniform matter may be evaluated by $E_\pi = \sqrt{m_\pi^2 + 2m_\pi U}$ ($p = 0$) with the real part of the

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$^1$ A preliminary study has been done along the same line for limited combinations of pion potentials and an RMF parameter set [20].
TABLE I: Pion potential parameters. The upper four sets by the pionic atom are taken from Ref. [16], in which $b_0$ and $B_0^{Re}$ were readjusted to reproduce the recent data of the deeply bound $\pi^{-}$ states in Pb with fixing the other parameters as the original values given in Ref. [21] for T, [22] for BFG, [23] for SM and [24] for ET.

| system set | $b_0$       | $b_1$       | $B_0^{Re}$   | $\alpha$ |
|------------|-------------|-------------|-------------|---------|
| T          | -0.034      | -0.078      | 0           | 0       |
| pionic      | -0.025      | -0.085      | -0.021      | 0       |
| atom        | -0.027      | -0.12       | 0           | 0       |
| ET         | -0.020      | -0.0873     | -0.049      | 0       |
| NOG [25]   | -0.013      | -0.105      | 0           | 0       |
| KY [12]    | -0.0233     | -0.1473     | -0.019      | 0.367   |
| pion-nucleus | F-C[14] | -0.009      | -0.114      | -0.040  | 0.391   |
| scattering  | F-W[14]    | -0.009      | -0.081      | -0.040  | 0.391   |

energy-independent potential $U$ based on Ericson-Ericson parameterization [31]:

$$U = -\frac{2\pi}{m_\pi} \left[ (1 + \epsilon) (b_0 \rho_\pi + b_1 \delta \rho) + \left(1 + \frac{\epsilon}{2}\right) B_0^{Re} \rho^{(2)} \right]$$

with $\epsilon = m_\pi / M_N$, $\rho_\pi = \rho_\pi + \rho_\rho$, $\delta \rho = \rho_\pi - \rho_\rho = \rho_\rho (1 - 2Y_\rho)$ and a squared density $\rho^{(2)}$ defined below. This potential is related to the pion self-energy via $\Sigma_\pi = 2m_\pi (U + iW)$ with an imaginary potential $W$. The s-wave $\pi N$ potential parameters ($b_0$, $b_1$, $B_0^{Re}$) was determined with special care from precise measurements of pionic atom and pion-nucleus scattering data.

In Table I we summarize the parameter sets adopted here. The upper 6 sets were determined from the pionic atom data and the lower two from pion-nucleus scattering. For the former parameter sets except NOG, $b_0 = \bar{b}_0$ and $\rho^{(2)} = \bar{\rho}^2$ are used, whereas, for the latter, double scattering modifications were explicitly included by $b_0 = b_0 - 3(b_0^0 + 2b_1^0)k_F/(2\pi)$ with $k_F = (3\pi^2 \rho_\pi / 2)^{1/3}$, and $\rho^{(2)} = \bar{\rho}^2 - \delta \rho^2$ is used. For NOG, $b_0 = \bar{b}_0 + \delta b_0 - 3 (1 + \epsilon) \bar{b}_1^0 (2b_1^0 + 2\bar{b}_1^0) / (2\pi)$ with $\delta b_0 = -0.005\rho_\pi^{-3}$. In the parameter sets of KY and F-W, $b_1$ is assumed to have density dependence through $b_1 = b_1 (1 - \alpha \rho_\pi / \rho_0)$ [17, 23, 33] with finite $\alpha$ and $\rho_0 = 0.17$ fm$^{-3}$. This is a consequence of the pion wave function renormalization associated with energy dependence of the optical potential [17, 23] and the renormalization was performed at $E_\pi = m_\pi$. For more realistic calculations in dense nuclear matter, the wave function renormalization should be done at the in-medium pion mass and other parameters should be also renormalized.

The phenomenological potentials are determined with a fixed pion energy. The energy dependence of the optical potential can be estimated by theoretical calculations. The s-wave in-medium pion self-energy $\Sigma_\pi(E_\pi)$, equivalent to the optical potential, was derived based on the chiral perturbation theory in Ref. [15] within a linear density approximation. The in-medium pion mass is obtained by solving $m^2_\pi - m^2_\pi - \Sigma(m^2_\pi) = 0$, which automatically takes account of the wave function renormalization. Here we use the following self-energy, abbreviated as MOW [16].

$$\Sigma(m^2_\pi) = c_1 \frac{4\rho_\rho}{f^2} m^2_\pi - \frac{2\rho_\rho}{f^2} \frac{m^2_\pi}{c_2 + c_3} \frac{\rho^2}{8m_N} + \frac{m^2_\pi \rho^2}{2f^2}$$

with $c_1 = -0.81$, $c_2 = 3.20$, $c_3 = -4.66$ in units of GeV$^{-1}$ and $f = 88$ MeV. We also consider the $\pi^{-}$ self-energy calculated by an in-medium chiral perturbation theory in $O(p^3)$ discussed in Ref. [17] (KKW), which reproduces well the energies and widths of deeply bound $\pi^{-}$ atomic states in Pb.

Let us see the model dependence of the pion energy $E_\pi$. In Fig. 1 we show $E_\pi$ in the cases of symmetric ($Y_\rho = 0.5$) and asymmetric ($Y_\rho = 0.2$) nuclear matter. The proton fraction $Y_\rho = 0.2$ is a typical value in neutron star matter obtained in RMF, as shown later in Fig. 2. The negative signs of the coefficients ($b_0$, $b_1$, $B_0^{Re}$) imply that $\pi^{-}$ feels repulsive potential in nuclear matter. The phenomenological pion potentials in symmetric nuclear matter agree well with each other at low densities below $\rho_0$, whereas, in asymmetric nuclear matter with $Y_\rho = 0.2$, we have 50-100 MeV ambiguities at $\rho_\rho \sim \rho_0$ in order to fix the large ambiguity of the potentials in asymmetric nuclear matter, it is very interesting to obtain pionic atom and scattering data in neutron rich nuclei [34].

RMF models have been developed to describe bulk properties of nuclei and nuclear matter with the mean field via the meson fields. We here adopt the RMF models, NL1 [26], NL3 [27], TM1 [28], SCL [29], having the Lagrangian in the following form:

$$L = L_{\text{tree}} + \bar{\psi} \left[ g_\sigma \sigma - g_\omega \omega - g_\rho \rho \right] \psi + \frac{c_\omega}{4} \omega^4 - V_\sigma(\sigma),$$

$$V_\sigma = \left\{ \begin{array}{ll}
\frac{1}{3} g_3 \sigma^3 + \frac{1}{4} g_4 \sigma^4 & \text{NL1, NL3, TM1} \\
-\alpha_0 f_{\text{SCL}}(\sigma / f_\pi) & \text{SCL} \end{array} \right. ,$$

where $\psi$, $\sigma$, $\omega$, $\rho$ represent nucleon and $\sigma$-, $\omega$- and $\rho$-meson fields, respectively, and $f_{\text{SCL}}(x) = \log(1-x) + x + x^2 / 2$. The model parameters are summarized in Table II. These RMF models describe the binding energies of heavy semi-double magic nuclei well, and are expected to give reasonable EOS of nuclear matter. We have solved the $\beta$ equilibrium condition
TABLE II: RMF parameters. In SCL, \(g_4\) and \(g_4\) are from the expansion of \(f_{\text{SCL}}\).

| Model   | \(g_0N\) | \(g_0N\) | \(g_\pi N\) | \(g_\pi N\) | \(g_\pi N\) | \(g_\pi N\) | \(c_\omega\) | \(m_\sigma\) (MeV) | \(m_\omega\) (MeV) | \(m_\rho\) (MeV) |
|---------|----------|----------|-------------|-------------|-------------|-------------|-------------|-----------------|-----------------|-----------------|
| NL1 [26] | 10.138   | 13.285   | 4.976       | 2401.9      | -36.265     | 0           | 492.25      | 795.359         | 763             |
| NL3 [27] | 10.217   | 12.868   | 4.474       | 2058.35     | -28.885     | 0           | 508.194     | 782.501         | 763             |
| TM1 [28] | 10.0289  | 12.6139  | 4.6322      | 1426.466    | 0.6183      | 71.3075     | 511.198     | 783             | 770             |
| SCL [29] | 10.08    | 13.02    | 4.40        | 1255.88     | 13.504      | 200         | 502.63      | 783             | 770             |

FIG. 2: (Color online) RMF results of proton fraction \(Y_p = \rho_p/\rho_B\), upper panel), electron chemical potential \(\mu_e\) and energy per baryon \(E/B\), lower right panel).

in cold neutron star matter,

\[
\mu_e = \mu_n - \mu_p, \quad \rho_e = \rho_p, \quad (1)
\]

\[
\rho_{n,p} = \frac{M_N^2 + k_p^2 + g_{\omega B}^2 + g_{\rho B}^2}{3}, \quad (2)
\]

where \(M_N^2 = M_N - g_\sigma^n\) represents the effective mass of nucleon. As shown in Fig. 2, calculated values of \(E/B\), \(Y_p\) and \(\mu_e\) in neutron star matter are consistent at low densities (\(\rho_B < \rho_0\)), since meson-baryon coupling constants are well determined by the binding energies of heavy-nuclei. Significant differences are found in \(E/B\) at higher densities, where the mesons have large expectation values and the self-interaction terms \((\sigma^3, \sigma^4, \omega^4)\) contribute to \(E/B\) considerably. While we have small differences in \(Y_p\) and \(\mu_e\), the model dependence is smaller compared with those in \(E/B\) and \(\mu_e\).

As we can see from Eq.(2), \(\mu_n - \mu_p\) is modified from the Fermi gas value with \(M^*\) by the \(\rho\) meson, whose coupling with nucleons is well constrained by nuclear binding energies, and higher order terms of the \(\rho\) meson are not included in the RMF models under consideration. As a result, model dependence of the isospin dependent potential, \(g_{\rho B}\), is around 10 MeV at \(\rho_B = 0.8\) fm\(^{-3}\). In Fig. 2, we also show the results of some RMF models including hyperons (TM1-SM [35] and IOTSY [20]). With hyperons, the proton fraction and electron chemical potential significantly decrease.

Now let us compare the electron chemical potential \(\mu_e\) and the pion energy \(E_\pi\) as functions of \(\rho_B\) (Fig. 3). The RMF results of TM1 are adopted for the proton fraction \(Y_p\) to evaluate \(E_\pi\). Results with hyperons denoted by IOTSY are also shown. The left panel of Fig. 3 shows the comparison of \(\mu_e\) and \(E_\pi\) obtained from the pionic atom data, namely, Tauscher (T), Batty-Friedman-Gal (BFG), Seki-Masutani (SM), Ericson-Tauscher (ET) and Kienle-Yamazaki (KY). In these potentials, density dependence of the potential parameters is not taken into account except for KY. The right panel of Fig. 3 compares \(\mu_e\) with \(E_\pi\) obtained from the pion-nuclear scattering data and the theoretical calculations.

We find that \(E_\pi\) obtained from the potentials with density-independent \(b_1\) are very close to \(\mu_e\) at high densities \(\rho_B > 0.3 \) fm\(^{-3}\), and in further dense nuclear matter, \(\mu_e\) exceeds \(E_\pi\) obtained with some of the parameter sets. We could have possibility for the \(s\)-wave pion condensation to take place in dense neutron star matter. However, since inclusion of hyperons makes \(\mu_e\) suppressed, \(E_\pi\) is found to be larger than \(\mu_e\) (IOTSY) in most cases. Thus the \(s\)-wave pion condensation would not take place, if hyperons could participate in neutron star matter.

We note that even more attractive hyperon potentials make \(\mu_e\) smaller (TM1-SM). Also in non-relativistic variational treatments [8, 36], symmetry energy and \(\mu_e\) are generally smaller at \(\rho_B > \rho_0\) than in relativistic models. In the case of the density-dependent \(b_1\), which is a consequence of the renormalization of the pion wave function and is required to explain the pionic atom data of Sn isotopes [12], the pion self-energies are more repulsive. Nevertheless, it is important to note that, as already mentioned, the renormalization of the wave function has been done only for the \(b_1\) parameter in linear \(\delta\rho\). Thus, for more quantitative discussion, it is necessary to improve the phenomenological and theoretical pion optical potentials in a consistent way, for instance as done in...
Ref. [18]. It is also desired to include the effects of short-range and tensor correlations on $\mu_p$ under $\beta$-equilibrium [8] in relativistic frameworks [37].

In summary, we have discussed the in-medium pion energy in the context of possibility for the $s$-wave pion condensation to take place in neutron stars. We have compared the in-medium pion energies determined from pionic atom or pion-nucleus scattering data with the electron chemical potential evaluated in relativistic mean field (RMF) models, using the RMF result of the proton fraction. With our present limited knowledge of the in-medium pion properties obtained in experiments, we could conclude that the $s$-wave pion condensation would not take place in neutron stars with hyperons. It is certainly necessary to investigate in-medium pion self-energy theoretically in more elaborated prescription to go beyond nuclear density. Especially energy dependence of the pion self-energy should be treated in more proper ways for higher densities. At the same time, experimental observations of pionic atoms and scattering in neutron rich nuclei are essential to fix ambiguities in the pion optical potentials in asymmetric nuclear matter. Precise knowledge of pion self-energies at high density is also important to study finite temperature process such as black hole formations, where the temperature can be as large as $T = 70$ MeV [38]. At such high temperatures, pion contribution could be significant depending on the in-medium pion mass.

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