Neutron stars with isovector scalar correlations

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Neutron stars with the isovector scalar $\delta$-field are studied in the framework of the relativistic mean field (RMF) approach in a pure nucleon plus lepton scheme. The $\delta$-field leads to a larger repulsion in dense neutron-rich matter and to a definite splitting of proton and neutron effective masses. Both features are influencing the stability conditions of the neutron stars. Two parametrizations for the effective nonlinear Lagrangian density are used to calculate the nuclear equation of state (EOS) and the neutron star properties, and compared to correlated Dirac-Brueckner results. We conclude that in order to reproduce reasonable nuclear structure and neutron star properties within a RMF approach a density dependence of the coupling constants is required.

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A Relativistic Mean Field (RMF) approach to nuclear matter with the coupling to an isovector scalar field, a virtual $a_0(980)$ $\delta$-meson, has been studied for the asymmetric nuclear matter at low densities, including its linear response \[\textsuperscript{1,2,3}\], and for heavy ion collisions at intermediate energies where larger density and momentum regions can be probed \[\textsuperscript{4,5,6}\]. In this work we extend the analysis of the contribution of the $\delta$-field in dense asymmetric matter to the influence on neutron star properties. We also want to test which effective interaction is more appropriate for the description of dense matter, including the symmetric part. In order to make a comparison, we use the two parametrizations for the effective nonlinear Lagrangian density to study the strongly isospin-asymmetric matter at high density regions.

A Lagrangian density of the interacting many-particle system consisting of nucleons, isoscalar (scalar $\sigma$, vector $\omega$), and isovector (scalar $\delta$, vector $\rho$) mesons is the starting point of the RMF theory

$$
\mathcal{L} = \bar{\psi} \left( i \gamma_\mu \partial^\mu - (M - g_\sigma \phi - g_\delta \bar{\delta} \cdot \vec{\delta}) - g_\omega \gamma_\mu \omega^\mu - g_\rho \gamma_\mu \vec{b}_\mu \cdot \vec{b} \right) \psi \\
+ \frac{1}{2} \left( \partial_\mu \phi \partial^\mu \phi - m_\phi^2 \phi^2 \right) - U(\phi) + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \frac{1}{2} m_\rho^2 \vec{b}_\mu \cdot \vec{b} \\
+ \frac{1}{2} \left( \partial_\mu \vec{\delta} \cdot \partial^\mu \vec{\delta} - m_\delta^2 \vec{\delta}^2 \right) - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{4} \vec{G}_{\mu\nu} \vec{G}^{\mu\nu},
$$

where $\phi$ is the $\phi$-meson field, $\omega_\mu$ is the $\omega$-meson field, $\vec{b}_\mu$ is $\rho$ meson field, $\vec{\delta}$ is the isovector scalar field of the $\delta$-meson. $F_{\mu\nu} \equiv \partial_\mu \omega_\nu - \partial_\nu \omega_\mu$, $\vec{G}_{\mu\nu} \equiv \partial_\mu \vec{b}_\nu - \partial_\nu \vec{b}_\mu$, and the $U(\phi)$ is a nonlinear potential of $\sigma$ meson: $U(\phi) = \frac{1}{4} a \phi^4 + \frac{1}{4} b \phi^4$.

The field equations in a mean field approximation (MFA) are
energy-momentum tensor. The energy density has the form
\[ \epsilon = \gamma \rho + \omega \gamma^0 \omega_0 - \rho \gamma^0 \tau_3 b_0 \psi = 0, \]
\[ m^2 \phi^2 + b \phi^3 = g_s \gamma^0 \psi = g_s \rho_s, \]
\[ m^2_\omega \omega_0 = g_\omega \gamma^0 \psi = g_\omega \rho, \]
\[ m^2_\omega b_0 = g_\omega \gamma^0 \tau_3 \psi = g_\omega \rho_3, \]
\[ m^2_\delta \delta_3 = g_5 \gamma^0 \tau_3 \psi = g_5 \rho_{s3}, \]
where \( \rho_3 = \rho_p - \rho_n \) and \( \rho_{s3} = \rho_{sp} - \rho_{sn} \). \( \rho \) and \( \rho_s \) are the baryon and the scalar densities, respectively. Neglecting the derivatives of mesons fields, the energy-momentum tensor in \( MFA \) is given by
\[ T_{\mu \nu} = i \bar{\psi} \gamma^\mu \partial_\mu \psi + \frac{1}{2} m^2_s \phi^2 + U(\phi) + \frac{1}{2} m^2_\omega \omega^2 - \frac{1}{2} m^2_\omega \omega_0 \lambda - \frac{1}{2} m^2_\delta \delta_3 | g_{\mu \nu}. \]

The equation of state (EOS) for nuclear matter at \( T=0 \) is straightforwardly obtained from the energy-momentum tensor. The energy density has the form
\[ \epsilon = \frac{1}{2} \sum_{i=n,p} \int \frac{d^3k}{(2\pi)^3} E_i^*(k) + \frac{1}{2} m^2_s \phi^2 + U(\phi) + \frac{1}{2} m^2_\omega \omega^2 + \frac{1}{2} m^2_\omega \omega_0 + \frac{1}{2} m^2_\delta \delta_3 + \frac{1}{2} m^2_\delta \delta_3, \]
and the pressure is
\[ p = \frac{1}{3} \sum_{i=1}^3 \frac{1}{3} \int \frac{d^3k}{(2\pi)^3} E_i^*(k) - \frac{1}{2} m^2_s \phi^2 - U(\phi) + \frac{1}{2} m^2_\omega \omega^2 + \frac{1}{2} m^2_\omega \omega_0 - \frac{1}{2} m^2_\delta \delta_3, \]
where \( E_i^* = \sqrt{k^2 + M_i^2}, i=p,n \). The nucleon effective masses are, respectively
\[ M_p^* = M - g_\sigma \phi - g_\delta \delta_3, \]
and
\[ M_n^* = M - g_\sigma \phi + g_\delta \delta_3. \]

The nucleon chemical potentials \( \mu_i \) are given in terms of the vector meson mean fields
\[ \mu_i = \sqrt{k_{F_i}^2 + M_i^2}, + g_\omega \omega_0 + g_\rho b_0 \] (proton, neutron),
where the Fermi momentum \( k_{F_i} \) of the nucleon is related to its density, \( k_{F_i} = (3\pi^2 \rho_i)^{1/3} \).

Since we are interested in the effects of the Nuclear Equation of State we will consider only pure nucleonic (+lepton) neutron star structures, i.e. without strangeness bearing baryons and even deconfined quarks, see the recent nice review [3] and the ref. [4]. In particular we will use two models for the neutron star composition, pure neutron and \( \beta \)-stable matter. In the latter case we limit the constituents to be neutrons, protons and electrons. Then the composition is determined by the request
of charge neutrality and $\beta$-equilibrium. The $(n, p, e^−)$ matter is indeed the most important $\beta$-stable nucleon + lepton matter at low temperature.

The chemical potential equilibrium condition for the $(n, p, e^−)$ system can be written as

$$\mu_e = \mu_n - \mu_p . \quad (9)$$

The charge neutrality condition is

$$\rho_e = \rho_p = X_p \rho , \quad (10)$$

where $X_p = Z/A = \rho_p/\rho$ is the proton fraction (asymmetry parameter $\alpha = 1 - 2X_p$), and $\rho$ is the total baryon density. The electron density $\rho_e$ in the ultrarelativistic limit for non-interacting electrons can be denoted as a function of its chemical potential

$$\rho_e = \frac{1}{3\pi^2} \mu_e^3 , \quad (11)$$

where $\mu_e = \sqrt{k_{F_e}^2 + m_e^2}$. The $X_p$ can be obtained by using Eqs.(8), (9), (10) and (11). The $X_p$ is related to the nuclear symmetry energy $E_{sym}$

$$3\pi^2 \rho X_p - [4E_{sym}(\rho)(1 - 2X_p)]^3 = 0. \quad (12)$$

In presence of a coupling to an isovector-scalar $\delta$-meson field, the expression for the symmetry energy has a simple transparent form, see [2, 3]:

$$E_{sym}(\rho) = \frac{1}{6} \frac{k_{F_e}^2}{E_{F_e}} + \frac{1}{2} [f_\delta - f_{\delta^*}(\frac{M^*}{E_{F_e}^*})^2] \rho , \quad (13)$$

where $M^* = M - g_s \phi$ and $E_{F_e}^* = \sqrt{k_{F_e}^2 + M^*}$. The $E_{sym}$ and the EOS for the $\beta$-stable $(npe^−)$ matter at $T=0$ can be estimated by using the obtained values of $X_p$. Equilibrium properties of the neutron stars can be finally studied by solving Tolmann-Oppenheimer-Volkov (TOV) equations [9, 10] inserting the derived nuclear EOS as an input.

The isovector coupling constants, $\rho$-field and $\rho + \delta$ cases, are fixed from the symmetry energy at saturation and from Dirac-Brueckner estimations, see the detailed discussions in refs.[2, 3].

In order to make a comparison, two parameter sets for the isoscalar part are used. The first, Set $A$, is more suitable at high densities where it appears closer to various Dirac-Brueckner predictions. In fact recently this interaction has been used with success to describe reaction observables in RMF-transport simulations of relativistic heavy ion collisions, where high densities and momenta are reached [4, 5, 6]. The second, Set $B$, is taken from the NL3 parametrization [11], obtained by fitting properties of symmetric nuclear matter at saturation density and of finite nuclei. In ref.[5] it has been shown that the good description of finite nuclei, even exotic, is kept also when the isovector-scalar channel is included, normally not present in the NL3 Lagrangian.
The coupling constants, \( f_i \equiv g_i^2/m_i^2 \), \( i = \sigma, \omega, \rho, \delta \), and the two parameters of the \( \sigma \) self-interacting terms: \( A \equiv a/g_\sigma^3 \) and \( B \equiv b/g_\sigma^4 \), are reported in Table 1. The corresponding properties of nuclear matter are listed in Table 2.

### Table 1. Parameter sets.

| Parameter \( f_i\) \((fm^2)\) | \( \text{Set A} \) | \( \text{Set B} \) |
|-------------------------------|--------------------|--------------------|
| \( f_\sigma \) |
| \( f_\omega \) |
| \( f_\rho \) |
| \( f_\delta \) |
| \( A \) \((fm^{-1})\) |
| \( B \) |

### Table 2. Saturation properties of nuclear matter.

| Parameter \( \rho_0 \) \((fm^{-3})\) | \( E/A \) (MeV) | \( K \) (MeV) | \( E_{sym} \) (MeV) | \( M^*/M \) |
|---------------------------------|----------------|------------|----------------|------------|
| \( \rho_0 \)         | 0.16     | 0.148     | 240.0         | 31.3       |
| \( E/A \)             | -16.0    | -16.299   | 271.7         | 33.7       |
| \( K \)               | 31.3     | 33.7      | 31.3          | 33.7       |
| \( M^*/M \)           | 0.75     | 0.60      | 0.75          | 0.60       |

We first use the two parametrizations to calculate the scalar and the vector potentials, and the binding energy \( E/A \) for symmetric nuclear matter \((\alpha=0.0)\) as a function of baryon density. The results are presented in Fig.1.

The relativistic Dirac-Brueckner-Hartree-Fock (DBHF) approach is a microscopic model describing many-body system with correlations, that has been extensively used to study the nuclear matter properties \[12, 13, 14, 15, 16, 17\]. In order to make a comparison, in Fig.1 we also report the different results within the DBHF approaches, the relativistic Dirac-Brueckner calculations by Brockmann and Machleidt \[12\], denoted as DBBM, the Dirac-Brueckner T-matrix calculations \[13\], denoted as DBT, and the Dirac-Brueckner results by ter Haar and Malfliet \[14\], denoted as DBHM.

From Fig.1(a) we see that the scalar and vector potentials for the isoscalar channels, given by the \( \text{Set B} \) (i.e. the NL3 \( \sigma \omega \) couplings) in low density regions are consistent with the correlated relativistic predictions, while the results given by the \( \text{Set A} \) are in better agreement in high density regions.

The dotted line in Fig.1(b) denotes the EOS of symmetric nuclear matter given by the \( \text{Set B} \), in full overlap with the solid line given by the NL3 interaction. It shows a nice agreement with correlated relativistic predictions at low densities but is clearly too repulsive with increasing density. At variance Fig.1(b) also shows that, as expected, the EOS of symmetric nuclear matter given by the \( \text{Set A} \) is more consistent with that given by Dirac-Brueckner calculations, in particular for the

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FIG. 1: (a) Scalar and vector potentials vs the baryon density; (b) binding energy as a function of the baryon density for symmetric nuclear matter. See text.

$DBT$ and $DBBM$ estimations, in high density regions up to about three-four times normal density. Such very different behaviors at high baryon density, in connection to the large difference in nucleon effective masses, will strongly influence the neutron star structure.

Relativistic heavy-ion collisions can provide crucial information about the EOS of nuclear matter. The investigation of nuclear EOS at high densities is one of driving force for study of heavy-ion reactions. The authors use different $DBHF$ approaches to study the collective flow of heavy-ion collisions. It is shown that the softer $DBT$ choice is in better agreement with experimental data of relativistic collisions, at least up to a few $AGeV$ beam energies, where densities up to $2.5\rho_0$ can be reached in the interacting zone. In general very accurate analyses of relativistic collisions data favor the predictions of a softer EOS at high densities. We remark from Fig.1(b) that for symmetric matter the $DBT$ is quite close to our $Set A$ parametrization, at least up to about $3\rho_0$.

For the isovector channels one can see from Eqs.(6) and (7) that the presence of the $\delta$-field leads to
proton and neutron effective mass splitting. In Fig. 2 we present the baryon density dependence of n, p effective masses for different proton fractions by the two parameter sets. The solid lines in Fig. 2 are the nucleon effective mass for symmetric nuclear matter \((X_p = 0.5)\).

Fig. 2(a) shows that the proton and the neutron effective masses given by the Set A decrease slowly with increasing baryon density, at variance with the Set B case that presents a much faster decrease, Fig. 2(b). This main difference between the two, A and B, parametrizations, is actually coming from the isoscalar part. When coupled to the splitting due to the isovector \(\delta\)-field it will have large effects on the n-star equilibrium features.

The density dependence of symmetry energy for the two parameter sets is reported in Fig. 3. For both cases we see a similar behavior of \(E_{sym}\) at sub-saturation densities for \(NL\rho\) and \(NL\rho\delta\) models. With increasing baryon density \(\rho\), however, the differences arising from the presence of the \(\delta\)-meson in the isovector channel become more pronounced for both A and B, parametrizations. This is due to
FIG. 3: Symmetry energy vs the baryon density at T=0 MeV by Set A and Set B, respectively. In the insert the corresponding proton fraction, see text.

the quenching factor \((M^*/E_F^*)^2\) for the attractive \(\delta\) contribution in Eq. (13) at high density regions, see refs. 2, 3.

We note that in spite of the same isovector coupling constants, at high density the Set B gives a larger symmetry energy. This is related to the larger contribution of the kinetic term in the r.h.s. of Eq.(13) due to the faster decrease of the effective nucleon mass. This represents a nice example of how the isovector part of the nuclear EOS can be influenced by the isoscalar channels due to the Fermi correlations.

In the inserts of Fig.3 we show the corresponding proton fractions \(X_p\) at \(\beta\)-equilibrium, Eq. (12). Due to the stiffer nature of the symmetry energy in the \(NL\rho\delta\) cases in both parametrizations we see a decrease of the \(\rho_{Urca}\), i.e. of the baryon density corresponding to the value \(X_p = 1/9\) that makes possible a direct \(Urca\) process, see 7.

At this point we can derive the predictions for equilibrium properties of neutron stars just solving
FIG. 4: Mass of the neutron star as a function of the central density of the neutron star by Set A and Set B, respectively.

The main results are presented in the Figs. 4 and 5. Fig 4 displays the neutron-star mass as a function of the central density of the star given by the two parameter sets, for the two compositions, pure neutron and the $\beta$-equilibrium ($npe^-$) matter.

Fig 4(a) and (b) both show that the maximum masses of the $\beta$-equilibrium star are smaller than in the pure neutron case due to the presence of a proton fraction and so of a smaller symmetry repulsion. Consistently the corresponding central densities are larger. A related effect is that the mass of the ($npe^-$) star given by the $NL\rho\delta$ model decreases more quickly with increasing density than that given by the $NL\rho$ choice, i.e., we have a lower density instability onset. This is just because the introduction of the $\delta$ coupling increases the equilibrium proton fraction at high baryon densities, see the inserts in Fig. 3.

The larger stiffness of the symmetry energy in the $NL\rho\delta$ cases, in both parametrizations, can be directly seen in the fact that the corresponding curves are always above the ones without the $\delta$
coupling. This implies larger maximum masses and smaller central densities. As a consequence we systematically see that the results for the \((npe^-)\) composition in the \(NL\rho\delta\) models are approximately equivalent to the ones for the pure neutron matter in the \(NL\rho\) choices.

Fig.4(b) presents a qualitatively new feature of the \(Set\ B\) results: the lack of solution (maximum mass) in the \(\beta\)-equilibrium case for the \(NL\rho\delta\) model. The fast decrease of the neutron effective mass, see Fig.2(b), prevents the chemical potential equilibrium condition for the \((npe^-)\) matter to be satisfied at densities around \(3\rho_0\).

Fig.5 reports the correlation between neutron star mass and radius given by the two parameter sets, respectively, for the two cases, pure neutron and \((npe^-)\) matter. Fig.5 shows that the contribution of the \(\delta\)-field to the neutron stars at high density regions is quite remarkable. In particular we note that we systematically have larger masses and radii and lower central densities, as expected from the larger symmetry pressure.

All the estimated maximum masses and the corresponding central densities and radii of the neutron
Table 3. Maximum mass, corresponding radius and central density of the star given by the two parameter sets.

| Neutron star Properties | Set A | Set B |
|-------------------------|-------|-------|
|                         | NLρ  | NLρδ | NLρ  | NLρδ |
| pure neutron            |       |       |       |
| \((M_s/M_\odot)_{max}\) | 2.24  | 2.56  | 2.88  | 3.12  |
| \(\rho_c/\rho_0\)      | 5.75  | 4.65  | 4.12  | 3.55  |
| \(R(km)\)              | 11.44 | 12.56 | 13.70 | 14.69 |
| (npe) matter            |       |       |       |
| \((M_s/M_\odot)_{max}\) | 2.14  | 2.30  | 2.83  |       |
| \(\rho_c/\rho_0\)      | 6.77  | 6.49  | 4.54  |       |
| \(R(km)\)              | 10.80 | 11.33 | 13.25 |       |

stars are reported in Table 3.

We note that the difference between the two, \(A\) and \(B\), parametrizations, in the neutron star predictions, is largely due to the isoscalar structure of the interactions, see Fig.1. The \(B\) parametrization is much stiffer at high density regions and this leads to differences on the neutron star masses, radii and central densities. The comparison between the results given by the two sets shows that the set \(B\) (i.e. the \(NL3\) forces) can be good at low densities below saturation density and it has serious problems at high densities, while the set \(A\) is a good choice for \(EOS\) of nuclear matter at larger density regions, which is consistent with that pointed out by refs. \[19, 20\]. In a sense all that just shows that we need density-dependent \(RMF\) parametrizations, the \(B\)-type at low densities and \(A\)-type at high densities. This has been already emphasized in the work \[21\]. While results of Density Dependent (\(DD\)) parametrizations and of the \(NL3\) forces agree very well below the saturation density, the \(EOS\) of \(DD\)-interactions at supra-normal densities shows a much softer behavior, similar to \(DBHF\) calculations and in better agreement with Heavy Ion Collision (\(HIC\)) data, see also ref. \[5\].

In any case our calculations show that the \(\delta\)-field provides significant contributions to the neutron star structure for the stiffness of the symmetry energy and the neutron/proton mass splitting in high density regions. The proton fraction in the \(\beta\)-equilibrium matter is much larger than that in the no \(\delta\)-field case. Moreover we can see from Fig.4(b) and Fig.5(b) that we do not have solutions in the \((npe^-)\) case for the set \(B\) with \(NL\rho\delta\) isovector interaction. This represents a quite dramatic effect of the splitting of neutron/proton effective masses at high densities, on top of the larger nucleon mass decrease of the \(B\) parametrizations, see Fig.2. The neutron chemical potential is then not able to satisfy the \(\beta\)-equilibrium conditions of the \((npe^-)\) matter. The contribution of the \(\delta\)-field for strongly isospin-asymmetric dense matter is important and it cannot be neglected.

With reference to neutron star properties we study the \(EOS\) for dense asymmetric matter in the \(RMF\) frame with two different parameter sets for the Lagrangian density. The set \(B\) is close to the \(NL3\) parametrization which has been proposed to describe finite nuclei properties. Since the proton and the neutron effective masses, especially the neutron effective mass, decrease quickly with increasing baryon density, the set \(B\) cannot provide the \(EOS\) needed in high density regions for the \((npe^-)\) star. This means that in general the \(B\) parametrization seems to have serious problems at high densities, as already remarked from \(HIC\) studies \[19\]. Though the \(B\) parametrization can be good
for the EOS at low densities, below and around saturation, it is too stiff in high density regions, in particular there is no solution for the \((npe^-)\) case with the \(\delta\)-field. So the set \(B\) is not suitable for the case of dense matter. We require a softer EOS at high densities, and indeed the softer Dirac-Brueckner predictions, in particular the DBT one, are in better agreement with relativistic collisions data. Our \(A\) parametrization, including the \(\delta\)-isovector-scalar field, is quite close to the DBT and it has been shown to lead to good predictions in transport simulations for heavy ion collisions at intermediate energies \([4, 5, 6]\). It appears then quite appropriate for the nucleonic part of the approach to neutron star properties. In this respect we note that quite extended neutron star structure calculations have been recently performed just using our Set A Lagrangian \([22]\).

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