If sufficiently light dibaryon resonances exist, a Bose condensate of dibaryons can occur in nuclear matter before the quark-hadron phase transition. Within a relativistic mean-field model we show that heterophase nuclear-dibaryon matter is for a wide set of parameters energetically more favorable than normal nuclear matter. Production of dibaryons is, however, relatively suppressed as compared to estimates based on the model of non-interacting nucleons and dibaryons.

The possibility of existence of dibaryon resonances was investigated in the last two decades experimentally and theoretically. The most promising candidates for experimental searches are those dibaryons which have a small width. In 1977 R.Jaffe predicted the existence of a loosely bound dihyperon, $H$, with a mass just below the threshold for two-lambda decay. Calculations of the $H$-particle mass within QCD-inspired models [2-4] showed that the existence of a dihyperon near the $\Lambda\Lambda$-threshold is plausible. Dibaryons with exotic quantum numbers, which have a small width due to zero coupling to the $NN$-channel, are of special interest [5-7]. A method for searching narrow, exotic dibaryon resonances in the double proton-proton bremsstrahlung reaction is discussed in Ref. [8]. Data from pion double charge exchange (DCE) reaction on nuclei [9, 10] exhibit a peculiar energy dependence at the total pion energy of 190 $MeV$ that can be interpreted as evidence for the existence of a narrow $d'$ dibaryon with quantum numbers $T = 0$, 1.
$J^P = 0^-$. Recent experiments at TRIUMPF (Vancouver) and CELSIUS (Uppsala) seem to support the existence of the $d'$ dibaryon [12].

The properties of nuclear matter with admixture of multiquark clusters are discussed in Ref. [13]. A dibaryon Bose condensate in interiors of neutron stars decreases the maximum masses of neutron stars [14]. In a recent paper [15] an exactly solvable model for a one-dimensional system of fermions interacting through a potential, which leads to a resonance in the two-fermion channel, is constructed. The behavior of this system can be interpreted in terms of a Bose condensation of the two-fermion resonances.

There is no dibaryon condensate in ordinary nuclei. From this one can conclude that the masses of dibaryons coupled to the $NN$-channel should be greater than

$$m_D > 2\mu_N = 2(m_N + \varepsilon_F) = 1.96 \text{ GeV}$$

where $\mu_N$ is the chemical potential of nucleons and $\varepsilon_F = 40 \text{ MeV}$ is the Fermi energy of nucleons in nuclei. Here, the ideal gas approximation for nucleons and dibaryons and the assumption that the shell model potential for dibaryons is twice as deep as the one for nucleons have been used.

The $d'$ dibaryon is coupled to the $NN\pi$ channel only. In the nuclear medium, the reaction $nd' \leftrightarrow nnp$ is possible. In nuclei the equilibrium condition for the chemical potentials has the form $\mu_n + m_D = 2\mu_n + \mu_p$. Since $\mu_p \approx \mu_n$, we arrive at the same inequality (1).

A Bose condensate of dibaryons can presumably be formed at high densities when relativistic effects for nucleons become important. In order to describe such a system, we should go beyond non-relativistic many-body theory. The relativistic field-theoretical Walecka model [16] is known to be very successful in describing properties of infinite nuclear matter and of ordinary nuclei throughout the periodic table. In this paper we study the influence of narrow dibaryon resonances on nuclear matter in the framework of the Walecka model in the mean-field approximation.

The Lagrangian of the model contains nucleons interacting through $\omega$- and $\sigma$-meson exchanges. We add to the Lagrangian dibaryons interacting with nucleons and each other through $\omega$- and $\sigma$-meson exchanges also. Inclusion of dibaryons entails uncertainties connected to the lack of reliable information on dibaryon masses and coupling constants. However, many conclusions can be drawn on quite general grounds without knowing precise values for the newly added parameters. The Lagrangian density is given by

$$\mathcal{L} = \bar{\Psi}(i\partial_\mu \gamma_\mu - m_N - g_\sigma \sigma - g_\omega \omega_\mu \gamma_\mu)\Psi + \frac{1}{2}(\partial_\mu \sigma)^2 - \frac{1}{2}m_\sigma^2 \sigma^2 - \frac{1}{4}F_{\mu\nu}^2 + \frac{1}{2}m_\omega^2 \omega^2 + (\partial_\mu - ih_\omega \omega_\mu)\varphi^* (\partial_\mu + ih_\omega \omega_\mu)\varphi - (m_D + h_\sigma \sigma)^2 \varphi^* \varphi + \mathcal{L}_c.$$  

Here, $\Psi$ is the nucleon field, $\omega_\mu$ and $\sigma$ are fields of the $\omega$- and $\sigma$-mesons, $F_{\mu\nu} = \partial_\mu \omega_\nu - \partial_\nu \omega_\mu$, $\varphi$ is the dibaryon isoscalar-scalar (or isoscalar-pseudoscalar) field. The values $m_\omega$ and
\( m_\sigma \) are the \( \omega \)- and \( \sigma \)-meson masses and the values \( g_\omega, g_\sigma, h_\omega, h_\sigma \) are coupling constants of the \( \omega \)- and \( \sigma \)-mesons with nucleons (\( g \)) and dibaryons (\( h \)).

The term \( \mathcal{L}_c \) describes conversion of dibaryons into nucleons. The \( H \)-particle is coupled to the \( NN \)-channel through a double weak decay and \( \mathcal{L}_c = O(G_F^2) \). For the nonstrange \( d_1 \) dibaryon \(^{[17]}\) and the \( d' \) dibaryon, we neglect possible virtual transitions \( \text{e.g.} \) to the \( NN\sigma \) channel. The on-shell couplings for these dibaryons are small. The exotic \( d_1 \) dibaryon decays to the \( NN\gamma \)-channel only, and so \( \mathcal{L}_c = O(\alpha) \). The \( d' \) dibaryon decays to the \( NN\pi \) channel. Due to Adler’s consistency condition \(^{[18]}\) \( \mathcal{L}_c \propto \partial_\mu \pi \). In the mean-field approximation \( \partial_\mu \pi = 0 \), and the term \( \mathcal{L}_c \) does not modify the mean-field equations. In what follows we set \( \mathcal{L}_c = 0 \). The effect of a small term \( \mathcal{L}_c \) reduces to providing a chemical equilibrium with respect to transitions between dibaryons and nucleons \(^{[19]}\).

The field operators can be expanded in \( c \)-numbers and operator parts: \( \omega_\mu = g_{\mu 0} \omega_c + \hat{\omega}_\mu \), \( \sigma = \sigma_c + \hat{\sigma}, \varphi = \varphi_c + \hat{\varphi} \), and \( \varphi^* = \varphi_c^* + \hat{\varphi}^* \). The \( \sigma \)-meson mean field determines the effective nucleon and dibaryon masses in the medium: \( m_N^2 = m_N + g_\sigma \sigma_c \) and \( m_D^2 = m_D + h_\sigma \sigma_c \). The baryon number current has the form \( j_\mu^B = j_\mu^N + 2j_\mu^D \) where \( j_\mu^N = \bar{\Psi} \gamma_\mu \Psi \) and \( j_\mu^D = \varphi^* i \partial_\mu \varphi - 2h_\omega \omega_\mu \varphi^* \varphi \). The \( \omega \)-mesons are coupled to the current \( j_\mu^\omega = g_\omega j_\mu^N + h_\omega j_\mu^D \).

The nucleon vector and scalar densities are defined by the expectation values \( \rho_{\text{NV}} = < \bar{\Psi}(0) \gamma_0 \Psi(0) > \) and \( \rho_{\text{NS}} = < \bar{\Psi}(0) \Psi(0) > \). The scalar density of the dibaryon condensate is given by \( \rho_{\text{DS}}^* = |< \varphi(0) >|^2 \). The time evolution of the dibaryon condensate \( \varphi \)-field is determined by the dibaryon chemical potential \( \mu_D \)

\[
\varphi_c(t) = e^{-i \mu_D t} \sqrt{\rho_{\text{DS}}^*}.
\]

(3)

The vector density of the dibaryon condensate is given by \( \rho_{\text{DV}}^* = 2\mu_D^* \rho_{\text{DS}}^* \) where \( \mu_D = \mu_D^* + h_\omega \omega_c \).

The existence of a dibaryon condensate depends on the values of the coupling constants of dibaryons with the \( \omega \)- and \( \sigma \)-mesons. The \( \omega \)- and \( \sigma \)-meson coupling constants \( h_\omega \) and \( h_\sigma \) enter the dibaryon-dibaryon Yukawa potential

\[
V(r) = \frac{h_\omega^2}{4\pi} e^{-m_\omega r} - \frac{h_\sigma^2}{4\pi} e^{-m_\sigma r}.
\]

(4)

The interaction energy for dibaryons in the condensate is for a constant density distribution \( \rho_D(x) = \rho_D \) equal to

\[
W = \frac{1}{2} \int d\mathbf{x}_1 d\mathbf{x}_2 \rho_D(\mathbf{x}_1) \rho_D(\mathbf{x}_2) V(|\mathbf{x}_1 - \mathbf{x}_2|) = 2\pi N_D \rho_D \left( \frac{h_\omega^2}{4\pi m_\omega^2} - \frac{h_\sigma^2}{4\pi m_\sigma^2} \right)
\]

(5)

where \( N_D \) is the total number of dibaryons. A negative \( W \) would imply instability of the system against compression. The value \( W \) is positive and the system is stable for

\[
\frac{h_\omega^2}{4\pi m_\omega^2} > \frac{h_\sigma^2}{4\pi m_\sigma^2}.
\]

(6)
In a nonrelativistic theory for systems of interacting bosons \[20\] and in the model considered, the requirement of stability is equivalent to the requirement of a positive value for the square of the sound velocity \((a_s^2 > 0)\).

The \(H\)-particle interactions are studied on the basis of the non-relativistic quark cluster model \[3, 22-24\] which is successful in describing the \(NN\)-phase shifts. The calculation of the interaction integral (3) with the adiabatic \(HH\)-potential \[24\] gives a negative energy, so the \(H\)-dibaryon condensate is probably unstable against compression. The coupling constants of the mesons with the \(H\)-particle can be fixed by fitting the depth and the position for minimum of the \(HH\)-potential to give \(h_\omega^2 = 603.7\) and \(h_\sigma^2 = 279.2\).

The coupling constants of the mesons with the \(d_1\) and \(d'\) dibaryons are unknown. The \(\omega\)- and \(\sigma\)-mesons interact with nucleons and pions inside the dibaryon. For dibaryons decaying into the \(NN\)-channel, the \(\sigma - D\) and \(\omega - D\) couplings are in nonrelativistic approximation two times greater than for nucleons: \(h_\omega = 2g_\omega\) and \(h_\sigma = 2g_\sigma\). The scalar charge is, however, suppressed by the Lorentz factor. For standard parameters of the Walecka model \[16\], \(m_\sigma = 520\) MeV, \(g_\omega^2 = 190.4\), and \(g_\sigma^2 = 109.6\), the inequality (6) becomes \(98.85(\frac{g_\omega}{2g_\omega})^2 GeV^{-2} > 129.0(\frac{g_\sigma}{2g_\sigma})^2 GeV^{-2}\). With these additive estimates, the inequality (6) is not fulfilled. The exchange current contributions to the meson couplings with dibaryons, which violate additivity, are analysed in Ref. \[21\]. At present no definite conclusions concerning the stability of the \(d_1\) and \(d'\) dibaryon matter can be drawn.

The mean-field solutions of the equations of motion corresponding to the Lagrangian density (2) are obtained by neglecting the operator parts of the meson fields. For the \(\omega\)- and \(\sigma\)-meson mean fields, we get the following expressions

\[
\omega_c = \frac{g_\omega \rho_{NV} + h_\omega 2\mu_D^* \rho_{DS}^*}{m_\omega^2},
\]

\[
\sigma_c = -\frac{g_\sigma \rho_{NS} + h_\sigma 2m_D^* \rho_{DS}^*}{m_\sigma^2}.
\]

Substituting expression (3) into the equation of motion for the dibaryon field \(\varphi\), we get \(\mu_D^* = m_D^*\). The nucleon and dibaryon chemical potentials have the form \(\mu_N = E_F^* + g_\omega \omega_c\) and \(\mu_D = m_D^* + h_\omega \omega_c\), where \(E_F^* = \sqrt{m_N^* + k_F^2}\) is the Fermi energy of nucleons with the effective mass \(m_N^*\).

The self-consistency condition for the effective nucleon mass can be transformed to a form equivalent to that in the standard Walecka model:

\[
m_N^* = \tilde{m}_N - \frac{g_\omega^2}{m_\omega^2} \rho_{NS},
\]

where \(\tilde{m}_N = m_N (\rho_{\rho_{DV}}^{c,\text{max}} / \rho_{\rho_{DV}}^c)^{\text{max}}\) and \(\rho_{\rho_{DV}}^{c,\text{max}} = m_N m_\sigma^2 / (g_\sigma h_\sigma) = 0.1507(\frac{2g_\sigma}{h_\sigma}) f m^{-3}\). If the densities \(\rho_{TV}\) and \(\rho_{\rho_{DV}}^c\) are fixed, equation (9) allows to find the effective nucleon...
mass $m^*_N$. Solutions to Eq. (1) exist for arbitrary total density $\rho_{TV}$ and dibaryon density $\rho_{DV}^* < \rho_{DV}^{c,max}$ when the value $\tilde{m}_N$ is positive.

It is clear that the fraction of dibaryons should increase when the difference $2\mu_N - \mu_D$ is positive and $\rho_{DV}^* = 0$. If the difference $2\mu_N - \mu_D$ is negative and the system consists of dibaryons only, production of nucleons is energetically favorable. If the difference $2\mu_N - \mu_D = 0$ and increases with the dibaryon fraction, small fluctuations take the system away from equilibrium. These types of states are unstable.

There are three stable cases: (i) The homophase nuclear matter: $2\mu_N - \mu_D < 0$ and $\rho_{DV}^* = 0$; (ii) the homophase dibaryon matter: $2\mu_N - \mu_D > 0$ and $2\rho_{DV}^* = \rho_{TV}$; (iii) the heterophase nuclear-dibaryon matter: $2\mu_N - \mu_D = 0$ and $\frac{d(2\mu_N - \mu_D)}{d\rho_{DV}^*}|_{\rho_{TV}} < 0$. Small fluctuations around these states lead the system back to the equilibrium points.

In Fig.1 we show the critical density for occurrence of a Bose condensate of non-strange dibaryons for $h\omega = 2g_\omega$ as a function of the $\sigma$-meson coupling constant with dibaryons. The critical density is determined from the equation $2\mu_N - \mu_D = 0$.

In Fig. 2 (a) we show the nucleon effective mass $m^*_N$ versus the dibaryon fraction $2\rho_{DV}^*/\rho_{TV}$ for $h\omega = 2g_\omega$ and $h_\sigma = 1.6g_\sigma$. The behavior of $m^*_N$ as a function of the dibaryon fraction does not depend on the $m_D$, since the dibaryon mass does not enter the self-consistency condition (9) directly. In Fig.2 (b) we show the difference for the chemical potentials versus the dibaryon fraction.

Mean-field solutions exist at all densities $\rho_{TV}$ for sufficiently small densities of dibaryons, $\rho_{DV}^* < \rho_{DV}^{c,max}$. This means that we can always investigate the stability of normal nuclear matter with respect to dibaryon condensate formation. Small total baryon number densities correspond to a stable equilibrium of type (i), at higher densities a stable equilibrium (iii) occurs. When the density $\rho_{TV}$ is high, dibaryon production is energetically favorable. The mean-field solutions disappear, however, before the system reaches an equilibrium.

In Fig. 3 we show the energy per nucleon and the pressure versus the total baryon number density for some possible dibaryons. The effect of zero compressibility for heterophase nuclear-dibaryon matter present in the ideal gas approximation reveals itself through the softening of the equation of state (EOS). Notice that the pressure of the heterophase system obeys the basic inequality $\partial p/\partial \rho_{TV} \geq 0$ of statistical mechanics. One can verify that in the model considered, the hydrostatic pressure coincides with the thermodynamic pressure. The $H$-particles are formed at a lower density, since the $\omega - H$ coupling constant $h_\omega/(2g_\omega) = 0.89$ is relatively small. The energy of the $H$-particles in the positive $\omega$-meson mean field is lower, so the production of the $H$-particles is energetically more favourable.

The qualitative estimates based on a model for non-interacting nucleons and dibaryons
show that in normal nuclear matter a dibaryon Bose condensate does not exist provided
the inequality (1) is satisfied. A more accurate estimate can be made on the basis of the
relativistic mean-field model (2). From the requirement of absence of a dibaryon Bose
condensate for \( \rho_{TV} \leq \rho_0 = 0.15 \, fm^{-3} \), we get for \( h_\omega = 2g_\omega \) a constraint
\[
  m_D > 1.89 \, GeV. \tag{10}
\]
This constraint is valid provided that the dibaryon matter is stable against compression.
In such a case, the \( d_1 \) resonance \( [17] \) with a mass \( m_D = 1.92 \, GeV \) does not affect properties
of ordinary nuclei.

Phase transitions of nuclear matter to strange quark matter \( [26, 27] \) have been widely
discussed in the literature (for a recent review see \( [28] \)). Dense nuclear matter with
a dibaryon Bose condensate can exist as an intermediate state below the quark-gluon
phase transition. This is the case when dibaryon matter is stable against compression.
If dibaryon matter is unstable, the creation of dibaryons can be a possible mechanism
for the phase transition to quark matter. The energetically favourable compression of
\( H \)-matter can lead to the formation of strange matter. It would be interesting to check
astrophysical data for the presence of a dibaryon condensate in the interiors of massive
neutron stars as well as possible signatures of their instability caused by dibaryons.

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FIGURE CAPTIONS

Fig.1. The critical density for occurrence of a Bose condensate of dibaryons versus the $\sigma$-meson coupling constant $h_\sigma$ for $m_D = 1.88$ GeV ($= 2m_N$; the long dashed curve No. 1), $1.96$ GeV (the solid curve No. 2), etc. with a step $80$ MeV. The results for dibaryons $d_1(1920)$ and $d'(2060)$ are also shown (the dashed curves). The dibaryon matter is stable against compression when the square of the sound velocity $a_s$ is positive. This is the case for $h_\sigma/(2g_\sigma) < 0.8754$. The value $\rho_0 = 0.15$ fm$^{-3}$ is the saturation density for nuclear matter. For $2m_N \leq m_D \leq 1.89$ GeV, we start at zero density from a heterophase nuclear-dibaryon matter. With increasing the density, the matter can be transformed to a homophase nuclear matter and then again to a heterophase nuclear-dibaryon matter. For $m_D > 1.89$ GeV, we start at zero density from a homophase nuclear matter which converts with increasing the density (at $\rho_{TV} > \rho_0$ for $h_\sigma/(2g_\sigma) < 0.8754$) to a heterophase nuclear-dibaryon matter. The occurrence of the $H$-particles is denoted by the cross.

Fig.2. (a) The effective nucleon mass $m^*_N$ in GeV versus the dibaryon fraction $2\rho^c_{DV}/\rho_{TV}$ in the heterophase nuclear-dibaryon matter. The results do not depend on the dibaryon mass. (b) The difference $2\mu_N - \mu_D$ between the two nucleon chemical potentials and the dibaryon chemical potential versus the dibaryon fraction $2\rho^c_{DV}/\rho_{TV}$. The results are given for the total baryon densities 1, 2, 3, 4, 5, and 6 times greater than the saturation density. The normal nuclear matter is stable when $2\mu_N - \mu_D < 0$ and $\rho^c_{DV} = 0$. An intersection of a curve with a negative slope with the horizontal line $2\mu_N - \mu_D = 0$ indicates occurrence of a stable equilibrium in the heterophase nuclear-dibaryon matter. The results are given for $m_D = 1.96$ GeV. The dibaryon mass does not enter the self-consistency condition (9) and enters linearly in the difference $2\mu_N - \mu_D$, so the curves for other dibaryon masses can be obtained by vertical parallel displacements. The results for $m_D = 2.06$ GeV ($d'$-dibaryon) can be obtained e.g. by a 100 MeV negative shift, etc.

Fig.3. The pressure (left scale) and the energy per baryon (right scale) versus the total baryon number density $\rho_{TV} = \rho_{NV} + 2\rho^c_{DV}$ for normal nuclear matter (solid lines) and for the heterophase nuclear-dibaryon matter (dashed lines) for the $d_1(1920)$ and $d'(2060)$ dibaryons at $h_\omega = 2g_\omega$ and $h_\sigma/(2g_\sigma) = 0.8$. The dibaryon condensates occur at $\rho_{TV}/\rho_0 = 2.05$ and $3.15$ for the $d_1$ and $d'$. The dibaryon Bose condensation softens the EOS for nuclear matter. The $H$-dibaryons are formed at $\rho_{TV}/\rho_0 = 2.74$ when the meson coupling constants are fixed by fitting the adiabatic $HH$-potential [24]. In this case the $H$-matter is unstable against compression, providing a transition to strange matter.