Constraints on the ω- and σ-meson coupling constants with dibaryons

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The effect of narrow dibaryon resonances on basic nuclear matter properties and on the structure of neutron stars is investigated in mean-field theory and in relativistic Hartree approximation. The existence of massive neutron stars imposes constraints on the coupling constants of the ω and σ mesons with dibaryons. In the allowed region of the parameter space of the coupling constants, a Bose condensate of the light dibaryon candidates $d_1(1920)$ and $d'(2060)$ is stable against compression. This proves the stability of the ground state of heterophase nuclear matter with a Bose condensate of light dibaryons. [S0556-2813(97)00209-4]

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The prospect of observing the long-lived H particle predicted in 1977 by Jaffe [1] stimulated considerable activity in experimental dibaryon searches. It was proposed to investigate H-particle production in different reactions [2]. So far, experiments [3] have not provided an unambiguous signature for the H particle. The existence of the H particle remains an open question, which must eventually be settled by further experiments. Nonstrange dibaryons with exotic quantum numbers, which have a small width due to zero coupling to the NN channel, are promising candidates for experimental searches [4]. The data on pionic double charge exchange reactions on nuclei [5] exhibit a peculiar energy dependence at an incident total pion energy of 190 MeV, which can be interpreted [6] as evidence for the existence of a narrow d' dibaryon with quantum numbers $T=0$, $J^p=0^-$ and a total resonance energy of 2063 MeV. Recent experiments at TRIUMF (Vancouver) and CELSIUS (Uppsala) seem to support the existence of the d' dibaryon [7]. The properties of the d' dibaryon were recently analyzed in the constituent quark model [8]. A method for searching narrow, exotic dibaryon resonances in the $pp\rightarrow pp\gamma\gamma$ reaction is discussed in Ref. [9]. Recently, some indications for a $d_1(1920)$ dibaryon have been found in this reaction [10].

When the density of nuclear matter is increased beyond a critical value, production of dibaryons becomes energetically favorable. Dibaryons are Bose particles. Therefore, they condense in the ground state and form a Bose condensate [11,12]. An exactly solvable model for a one-dimensional Fermi system in which the fermions interact through a potential that leads to a resonance in the two-fermion channel is analyzed in Ref. [13]. When the density is increased beyond a critical value, the behavior of the system can be interpreted in terms of a Bose condensation of two-fermion resonances. In the limit of vanishing decay width, a dibaryon can be approximately described as an elementary field. The effect of narrow dibaryon resonances on nuclear matter has recently been analyzed in mean-field theory (MFT) [14,15].

Although a dibaryon Bose condensate does not exist in ordinary nuclei, dibaryons affect the properties of nuclei and nuclear matter through a Casimir effect: The presence of the background σ-meson mean field inside the nucleus modifies the nucleon and dibaryon masses. This in turn modifies the vacuum fluctuations of the nucleon and dibaryon fields and thus affects the energy density and pressure of the system. This effect, which is well known [16] for nucleons, can be evaluated within the relativistic Hartree approximation (RHA). We recall that in a loop expansion of quantum hadrodynamics (QHD), MFT corresponds to the lowest approximation (no loops), while the RHA corresponds to the one-loop approximation in the calculation of the equation of state for nuclear matter.

At zero temperature, a uniformly distributed system of bosons with an attractive potential is energetically unstable against compression and collapses [17]. In such a case, the long wave excitations (sound in the medium) have an imaginary dispersion law: The square of the sound velocity is negative $a_s^2<0$. The amplitude of these excitations increases with time, resulting in the instability of the system. It is necessary to analyze dispersion laws of other elementary excitations also. We shall see, however, that in MFT and the RHA only sound waves can generate an instability. The ground state of nuclear matter with a Bose condensate of dibaryons is either stable or unstable against small perturbations depending on whether the repulsive ω-meson exchange or the attractive σ-meson exchange interaction between dibaryons dominates.

In this paper, we investigate the possibility that dibaryon matter is unstable against compression. If dibaryon matter is unstable, the formation of dibaryons in nuclear matter would provide a possible mechanism for a phase transition to quark matter. If the central density of a massive neutron star exceeds a critical value for the formation of dibaryons, the neutron star could convert into a quark star, a strange star, or a black hole. Some of the observed pulsars are identified quite reliably with ordinary neutron stars [18]. From the requirement that dibaryon formation be not energetically favored at densities lower than the central density of neutron stars with a mass of $1.3M_{\odot}$, we derive constraints on the coupling constants of the ω and σ mesons with the possible $d_1(1920)$ and $d'(2060)$ dibaryons. We conclude that narrow dibaryons in this mass range can only form a Bose condensate that is stable against perturbations. The effect of the dibaryons on the stability and structure of neutron stars in

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different phenomenological models is analyzed in Refs. [12, 19]. Constraints on the binding energy of strange matter [20] derived from the existence of massive neutron stars are discussed in Ref. [21].

The dibaryonic extension of the Walecka model [16] is obtained by including dibaryons in the Lagrangian density [14,15]

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

Here, Ψ is the nucleon field, ω_μ and σ are fields of the ω and σ mesons, $F_{\mu\nu}=\partial_\nu\omega_\mu-\partial_\mu\omega_\nu$ is the field-strength tensor, and φ is the dibaryon isoscalar-scalar (or isoscalar-pseudoscalar) field. Here, m_ω and m_σ are the ω- and σ-meson masses and $g_\omega,g_\sigma,h_\omega,h_\sigma$ are the coupling constants of the ω and σ mesons with nucleons (g) and dibaryons (h).

The σ-meson mean field σ_c determines the effective nucleon and dibaryon masses in the medium:

$$m_N^*=m_N+g_\sigma\sigma_c, \quad (2)$$

$$m_D^*=m_D+h_\sigma\sigma_c. \quad (3)$$

The nucleon scalar density in the RHA is defined by the expression [16]

$$\rho_{NS}=\langle\overline{\Psi}(0)\Psi(0)\rangle=\gamma\int\frac{d\mathbf{p}}{(2\pi)^3}\frac{m_N^*}{E^*(\mathbf{p})}\theta(p_F-|\mathbf{p}|) -4m_N^3\zeta(m_N^*/m_N), \quad (4)$$

where

$$4\pi^2\zeta(x)=x^3\ln x+1-x-\tfrac{5}{2}(1-x)^2+\tfrac{11}{2}(1-x)^3$$

and where the statistical factor γ is $\gamma=2$ for neutron matter and $\gamma=4$ for nuclear matter. The last term in Eq. (4) is due to the Dirac sea. This vacuum contribution renormalizes the scalar density.

Here, we investigate the properties of nuclear matter below the critical density for formation of dibaryons: i.e., the dibaryon condensate is zero, $\langle\varphi(0)\rangle=0$. The vacuum contribution to the scalar density of dibaryons can be found to be

$$2m_D^*\rho_{DS}=2m_D^*\langle\varphi(0)^*\varphi(0)\rangle=m_D^3\zeta(m_D^*/m_D). \quad (5)$$

It differs from the corresponding vacuum contribution to the nucleon scalar density in sign and statistical factor [one should replace $4(=2_S\times2_I)\rightarrow1$]. The self-consistency condition for the nucleon effective mass has the form

$$m_N^*=m_N-\frac{g_\sigma}{m_\sigma^2}(g_\sigma\rho_{NS}+h_s2m_D^*\rho_{DS}). \quad (6)$$

The renormalized vacuum contribution to the nucleon energy-momentum tensor has the form [16]

$$\langle T_{\mu\nu}^N(0)\rangle_{\rm vac}=-4g_{\mu\nu}m_N^4\eta(m_N^*/m_N), \quad (7)$$

where

$$16\pi^2\eta(x)=x^4\ln x+1-x-\tfrac{7}{2}(1-x)^2 +\tfrac{13}{3}(1-x)^3-\tfrac{25}{12}(1-x)^4.$$

For dibaryons, we get the expression

$$\langle T_{\mu\nu}^D(0)\rangle_{\rm vac}=g_{\mu\nu}m_D^4\eta(m_D^*/m_D). \quad (8)$$

The elementary excitations in nuclear matter with a Bose condensate of dibaryons correspond to nucleons and antinucleons, ω mesons, σ mesons, and dibaryons and antidibaryons. The dispersion laws for these quasiparticles can be found in Ref. [15]. The nucleon and antinucleon dispersion laws have the same form as in the vacuum with the replacement $m_N\rightarrow m_N^*$. Therefore, they cannot generate an instability of the system. The dispersion laws of ω mesons, σ mesons, and antidibaryons turn out to be real, too. The only possible source of instability is the dibaryon quasiparticle excitations. The latter are responsible for long wavelength perturbations of the system and are connected with the existence of sound in the medium.

The square of the sound velocity has the form [15]

$$a_s^2=\frac{\alpha}{1+\alpha}, \quad (9)$$

where

$$\alpha=2\rho_{DS}\frac{m_\sigma^2}{\widetilde{m}_\sigma^2}\left(\frac{h_\omega^2}{m_\omega^2}-\frac{h_\sigma^2}{m_\sigma^2}\right) \quad (10)$$

and $\widetilde{m}_\sigma^2=m_\sigma^2+2h_\sigma^2\rho_{DS}$. We see that $a_s^2>0$ for

$$\frac{h_\omega^2}{m_\omega^2}>\frac{h_\sigma^2}{m_\sigma^2}. \quad (11)$$

The validity of inequality (11) is a sufficient condition for the stability of the ground state of nuclear matter with a Bose condensate of dibaryons. Later, we show that a violation of the inequality (11) for light dibaryons is in contradiction with the existence of massive neutron stars.

The physical meaning of the inequality (11) can be clarified by considering the interaction energy of uniformly distributed dibaryon matter $\rho_{DV}(\mathbf{x})=\rho_{DV}=$const:

$$W=\frac{1}{2}\int d\mathbf{x}_1d\mathbf{x}_2\rho_{DV}(\mathbf{x}_1)\rho_{DV}(\mathbf{x}_2)V(|\mathbf{x}_1-\mathbf{x}_2|). \quad (12)$$

The Yukawa potential $V(r)$ for two dibaryons has the form

$$V(r)=\frac{h_\omega^2}{4\pi}\frac{e^{-m_\omega r}}{r}-\frac{h_\sigma^2}{4\pi}\frac{e^{-m_\sigma r}}{r}. \quad (13)$$

Integration gives

$$W=\frac{1}{2}N_D\rho_{DV}\left(\frac{h_\omega^2}{m_\omega^2}-\frac{h_\sigma^2}{m_\sigma^2}\right), \quad (14)$$

where N_D is the number of dibaryons. When the dibaryon density ρ_{DV} increases, the energy (for $a_s^2>0$) increases also, the pressure is positive, and therefore the system is stable.

At present, the coupling constants of the mesons with dibaryons are not known with sufficient precision to draw a definite conclusion concerning the stability of dibaryon mat-

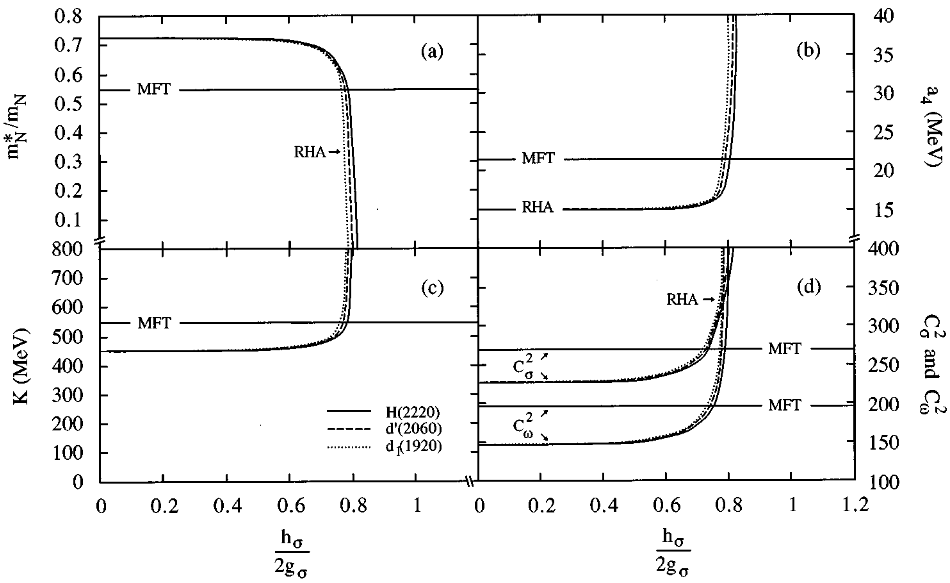

FIG. 1. The effective nucleon mass (a), asymmetry coefficient (b), incompressibility of nuclear matter at saturation density (c), and the coupling constants $C_s^2 = g_s^2(m_N/m_s)^2$ and $C_\omega^2 = g_\omega^2(m_N/m_\omega)^2$ in the RHA, versus the ratio $h_\sigma/(2g_\sigma)$ of σ-dibaryon (h_σ) and σ-nucleon (g_σ) coupling constants. The MFT results are shown for comparison. The coupling constants are fixed by fitting the minimum position and depth of the energy per baryon number at the saturation density of nuclear matter. The solid, dashed, and dotted curves correspond to the dibaryons H(2220), d'(2060), and d_1(1920), respectively. A satisfactory description of the properties of nuclear matter at the saturation density is possible for $h_\sigma/(2g_\sigma) < 0.8$.

ter. Given that the ratio $h_\sigma/(2g_\sigma)$ between the σ-meson couplings with dibaryons and nucleons is fixed, one can extract the ω- and σ-meson couplings with nucleons by fitting the nuclear matter binding energy $E/A - m_N = -15.75$ MeV at the empirical equilibrium density $\rho_0 = 0.148$ fm^{-3}. The empirical equilibrium density ρ_0, which is determined from the density in the interior of ^{208}Pb [16], corresponds to the equilibrium Fermi wave number $k_F = 1.30$ fm^{-1}.

In Fig. 1, we show how relevant nuclear matter observables are influenced by the presence of dibaryons in the Walecka model. The dependence of the effective nucleon mass on the ratio $h_\sigma/(2g_\sigma)$ at the equilibrium density is shown in Fig. 1(a). In Figs. 1(b) and 1(c) we show the dependence of the incompressibility $K = 9\rho_0(\partial^2\varepsilon/\partial\rho^2)|_{\rho=\rho_0}$ and the asymmetry coefficient a_4 on the ratio $h_\sigma/(2g_\sigma)$. In Fig. 1(d), the values $C_s^2 = g_\sigma^2(m_N/m_\sigma)^2$ and $C_\omega^2 = g_\omega^2(m_N/m_\omega)^2$ are plotted. Note that $h_\sigma/(2g_\sigma) = 0$ is equivalent to the RHA without dibaryons. For comparison, we give the MFT results. In MFT, dibaryons do not influence the properties of nuclear matter below the critical density for the formation of a dibaryon Bose condensate. The RHA results do not sensitively depend on the mass of the dibaryon.

When the ratio $h_\sigma/(2g_\sigma)$ approaches the value 0.8, the system of equations does not yield physical solutions and the empirical equilibrium properties of nuclear matter can no longer be reproduced. For $x \to 1$, $\zeta(x) = O((1-x)^4)$, and the zero-point contributions to the scalar density of nucleons and dibaryons, which have opposite signs, are comparable for $4g_\sigma^4/m_N \approx h_\sigma^4/m_D$. Dibaryon effects become large for $h_\sigma/(2g_\sigma) \approx 0.5(4m_D/m_N)^{1/4} \approx 0.84$. The greater the dibaryon mass, the greater the upper limit of the critical ratio $h_\sigma/(2g_\sigma)$. This effect is seen in Fig. 1.

In this work, we study the physical implications of the effective Lagrangian (1) describing nucleon and dibaryon degrees of freedom. Other baryons can be included in the QHD framework in a similar way. Their effect is quite small. We have checked that the Casimir effect caused by the inclusion of other octet baryons ($2\times 8 = 16$ degrees of freedom) shifts the critical value h_σ only by about 25% if one assumes that the σ meson is an SU(3)$_f$ singlet. The inclusion of octet and decuplet baryons ($2\times 8 + 4\times 10 = 56$ degrees of freedom) with a universal sigma-meson coupling constant increases the critical value of $h_\sigma/(2g_\sigma)$ to 1.2.

The saturation curve for nuclear matter is shown in Fig. 2. The equation of state (EOS) in the RHA is softer than in MFT. The contributions of the vacuum zero-point fluctuations of nucleons and dibaryons partially cancel each other, and therefore the inclusion of dibaryons makes the EOS stiffer. This is shown in Fig. 2, where the RHA with dibaryons (dashed curve) lies above the RHA calculation without

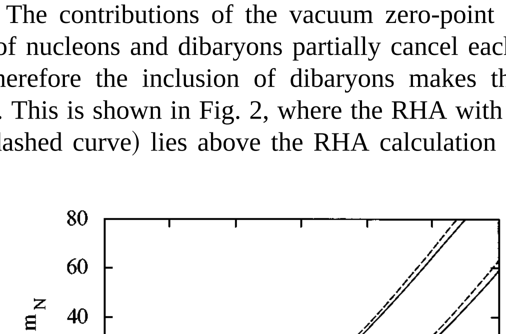

FIG. 2. Saturation curve for nuclear matter in the RHA: without dibaryons (solid line) and with the inclusion of H dibaryons (dashed line) for $h_\sigma/(2g_\sigma) = 0.6$.

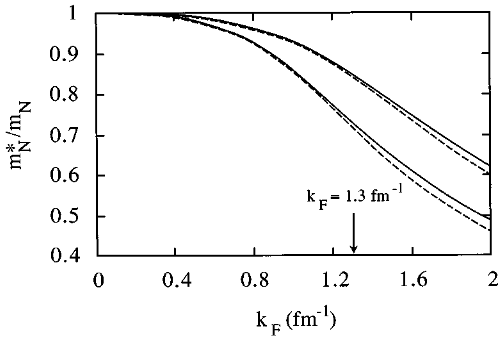

FIG. 3. The effective nucleon mass versus the Fermi momentum of nucleons in nuclear matter (upper curves) and in neutron matter (lower curves). The solid lines correspond to the RHA without dibaryons, and the dashed lines correspond to the RHA with inclusion of H dibaryons.

dibaryons (solid line). The same effect is seen in Fig. 3, the dependence of the effective nucleon mass on the Fermi wave number k_F being shown for the case of the H particle with $h_\sigma/(2g_\sigma)=0.6$. In MFT, the effective nucleon mass decreases faster than in the RHA. As a result of the partial cancellation of the nucleon and dibaryon contributions to the vacuum scalar density, the dashed lines lie below the solid lines.

In Fig. 4 we show the critical densities for the formation of dibaryons $H(2220)$, $d'(2060)$, and $d_1(1920)$ in nuclear and neutron matter versus the ratio $h_\sigma/(2g_\sigma)$, both in MFT and the RHA. For each dibaryon, three different values for the ω-dibaryon coupling are assumed, namely, $h_\omega/h_\omega^{\max}=1$, 0.8, and 0.6. Here, $h_\omega^{\max}=h_\sigma m_\omega/m_\sigma$ is the maximum value for the ω-dibaryon coupling constant for which which the inequality (11) is violated. Note that $h_\omega/h_\omega^{\max}=1$ corresponds to $a_s^2=0$ and $h_\omega/h_\omega^{\max}=0.8$ and 0.6 correspond to $a_s^2<0$. In the RHA, dibaryons occur at higher densities. The coupling constant h_ω determines the energy of dibaryons in the positive ω-meson mean field. The greater h_ω, the greater the density that is required to make the production of dibaryons energetically favorable. This effect is seen in Fig. 4: The solid lines $h_\omega/h_\omega^{\max}=1$ lie above the long-dashed and dashed lines $h_\omega/h_\omega^{\max}=0.8$ and 0.6, respectively.

If dibaryon matter is unstable against compression, production of dibaryons with increasing density results in instability of neutron stars with a subsequent phase transition to quark matter and the conversion of neutron stars to quark stars, strange stars, or black holes. In such a case, the maximum masses of neutron stars are determined by the mass and the coupling constants of the mesons with the lightest dibaryon. In Fig. 5 we show the minimal neutron star masses for which dibaryons can occur.

The MFT and RHA EOS for neutron matter at supranuclear densities are matched smoothly with the Baum-Bethe-Pethick EOS [22] at densities $\rho_{\rm drip}<\rho<0.8\rho_0$ where $\rho_{\rm drip}=4.3\times10^{11}$ g/cm^3 and then with the Baum-Pethick-Sutherland EOS [23] at densities $\rho<\rho_{\rm drip}$. The maximum neutron star masses are sensitive to the value of the equilibrium Fermi wave number. If we chose $k_F=1.42$ fm^{-1} instead of $k_F=1.30$ fm^{-1}, the maximum masses in MFT (without dibaryons) are reduced from $3M_\odot$ down to $2.6M_\odot$ [16]. The choice $k_F=1.30$ fm^{-1} provides less stringent and therefore more conservative constraints for the meson-dibaryon coupling constants. We do not show the results for the $d_1(1920)$ dibaryon, because its condensation starts al-

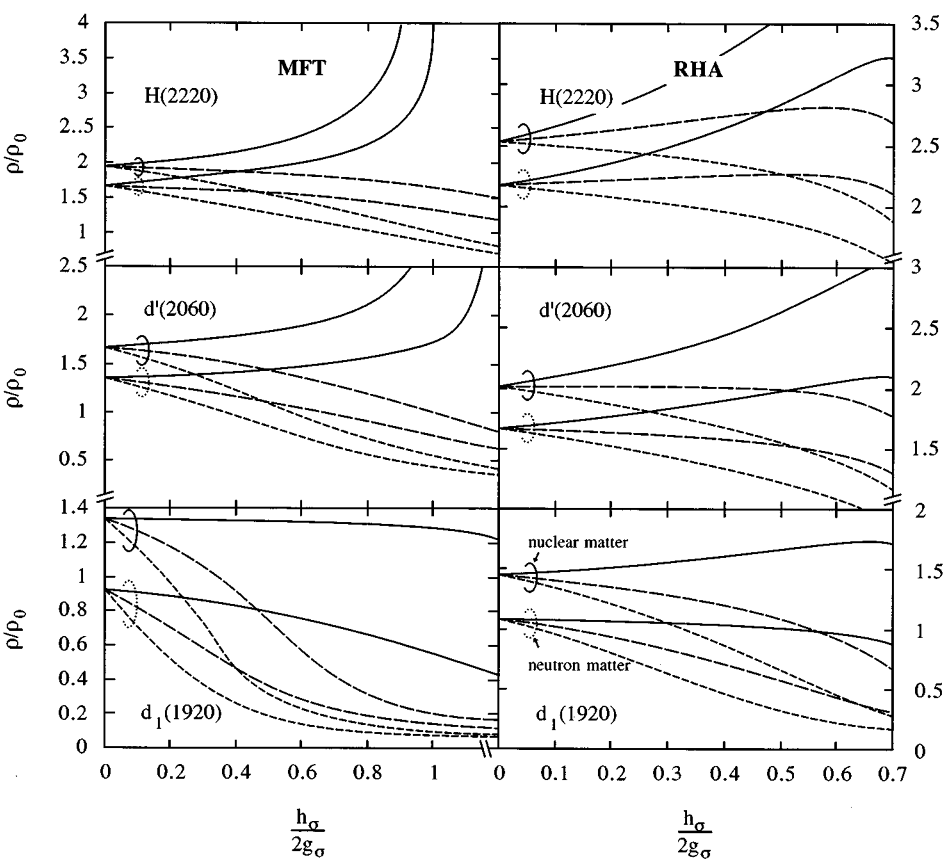

FIG. 4. Critical densities for the formation of $H(2220)$, $d'(2060)$, and $d_1(1920)$ dibaryons in nuclear and neutron matter in both MFT and the RHA, versus the σ-dibaryon coupling constant for $h_\omega/h_\omega^{\max}=1$, 0.8, and 0.6 (the solid, longed-dashed, and dashed lines, respectively). Here, $h_\omega^{\max}=h_\sigma m_\omega/m_\sigma$ is the maximum value of the ω-dibaryon coupling constant for which dibaryon matter is unstable against compression (see text).

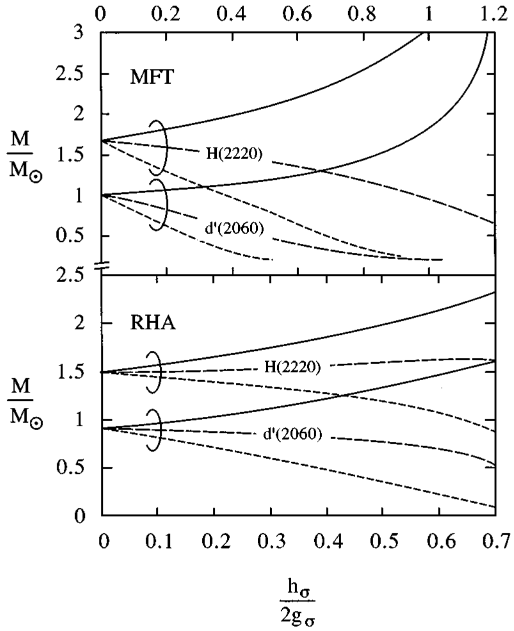

FIG. 5. Lowest neutron star masses, for which dibaryon formation becomes energetically favorable, versus the σ-dibaryon coupling constant for $h_\omega/h_\omega^{\max}=1$, 0.8, and 0.6 (the solid, longed-dashed, and dashed lines, respectively). Here, $h_\omega^{\max}$ is the maximum value of the ω-dibaryon coupling constant for which dibaryon matter is unstable against compression ($a_s^2<0$). The results are shown for both MFT and the RHA and for $H(2220)$ and $d'(2060)$ dibaryons. In the case of the $d_1(1920)$ dibaryon, the resulting neutron star masses are very small ($<0.2M_\odot$).

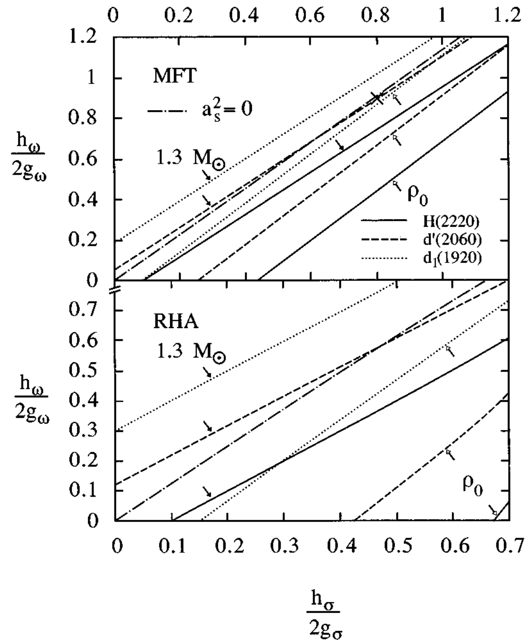

FIG. 6. Parameter space for the σ- and ω-dibaryon coupling constants in MFT and the RHA. The dash-dotted line $a_s^2=0$ divides the parameter space into two parts. The upper left part corresponds to dibaryon matter stable against compression (square of the sound velocity is positive $a_s^2>0$), and the lower right part corresponds to dibaryon matter unstable against compression ($a_s^2<0$). The solid, dashed, and dotted curves restrict the parameter space for the meson coupling constants with the H particle and the $d'(2060)$ and $d_1(1920)$ dibaryons. In the left upper region, dibaryon formation is energetically unfavorable in ordinary nuclei (curves marked by arrows with white arrowheads) and in neutron stars of a mass $1.3M_\odot$ (curves marked by arrows with black arrowheads). In the lower right parts of the parameter space, dibaryons could appear, respectively, in ordinary nuclei and in massive neutron stars. The cross refers to the H-particle coupling constants with the mesons, determined from the adiabatic H-H potential [26].

ready at a density $\rho\approx\rho_0$ and results in the conversion of neutron stars with very low masses $M<0.2M_\odot$ into quark stars.

In Fig. 6 we show the parameter space for the coupling constants of dibaryons with σ and ω mesons. As mentioned in the beginning, our discussion is restricted to the region $a_s^2<0$ for which dibaryon matter is unstable against compression. The requirement of stability of normal nuclear matter at the saturation density allows us to get constraints on the coupling constants. The corresponding curves (straight lines in the MFT case) marked by arrows with white arrowheads restrict the parameter space of the coupling constants from below. The dotted line in the MFT case, which refers to the $d_1(1920)$ dibaryon, is very close to the dash-dotted line $a_s^2=0$. For low values h_σ, the dotted line lies above the line $a_s^2=0$. In this case, dibaryon matter unstable against compression cannot exist. The window in parameter space for unstable $d_1(1920)$-dibaryon matter is, however, much larger in the case of the RHA.

With the conservative assumption that pulsars with a mass $1.3M_\odot$ are ordinary neutron stars, the constraints on the meson-dibaryon coupling constants can be further improved. The corresponding curves (straight lines in the MFT case) are shown in Fig. 6. We see that the dotted and dashed lines lie above or very close to the line $a_s^2=0$. This means that for $d_1(1920)$ and $d'(2060)$ dibaryons, dibaryon matter unstable against compression cannot exist [owing to the very small window for $d'(2060)$ at higher values of h_σ]. For the H particle, there is a window in parameter space between the line $a_s^2=0$ and the solid curves marked by arrows with black arrowheads that corresponds to dibaryon matter unstable against compression.

The H-particle interaction was studied in the nonrelativistic quark cluster model [24,25], which simultaneously describes the NN-phase shifts. The coupling constants of the mesons with the H particle can be fixed by fitting the depth and the position of the minimum of the HH-adiabatic potential [26] to give $h_\omega/(2g_\omega)=0.89$ and $h_\sigma/(2g_\sigma)=0.80$. These values are marked on Fig. 6 with a cross. These estimates are used in the MFT calculations [14,15]. They correspond to unstable dibaryon matter and are in the allowed region of the parameter space for the H particle. The energetically favorable compression of H matter can lead to the formation of absolutely stable strange matter [20] and the conversion of neutron stars to strange stars [27].

The MFT and RHA EOS of the extended Walecka model are both very stiff. In soft NN interaction models such as the Reid potential (for a review of nuclear matter models see [18]), the central density of neutron stars is much larger than in stiff models. Therefore, the conditions for the occurrence of new forms of nuclear matter are more favorable. From Figs. 5 and 6, we see that the softer RHA EOS produces lower upper limits on the neutron star masses and, respectively, more stringent constraints on the meson-dibaryon coupling constants as compared to the stiffer MFT EOS, despite the fact that in the RHA dibaryons occur at higher densities (see Fig. 4). One can assume that this effect is of general validity and that softer EOS give more stringent constraints on the meson-dibaryon coupling constants. Therefore, we consider the constraints given in Fig. 6 as rather conservative.

In conclusion, we have shown that the assumption of instability of $d_1(1920)$- and $d'(2060)$-dibaryon matter against compression is in contradiction with the hypothesis that pulsars of a mass of $1.3M_{\odot}$ are ordinary neutron stars. This conclusion is valid for all narrow dibaryons with the same quantum numbers in the same mass range. On the other hand, the H particle is sufficiently heavy, its condensation starts at higher densities, and constraints on meson-dibaryon coupling constants are not as stringent. One cannot exclude the possibility that H particles form in nuclear matter a condensate that is unstable against compression. Finally, the ω- and σ-meson coupling constants with $d_1(1920)$ and $d'(2060)$ dibaryons should obey the inequality (11). The corresponding coupling constants with the H particle should lie above the solid curves in Fig. 6.

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[1] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977).

[2] M. Sano, M. Wakai, and H. Bando, Phys. Lett. B **224**, 359 (1989); F. S. Rotondo, Phys. Rev. D **47**, 3871 (1993); A. T. M. Aerts and C. B. Dover, Phys. Rev. Lett. **24**, 1752 (1982), Phys. Rev. D **28**, 1752 (1982); **29**, 433 (1984).

[3] S. Aoki *et al.*, Phys. Rev. Lett. **65**, 1729 (1990); A. Rusek *et al.*, *ibid.* **52**, 1580 (1995).

[4] P. J. Mulders, A. T. Aerts, and J. J. de Swart, Phys. Rev. D **21**, 2653 (1980); L. A. Kondratyuk, B. V. Martemyanov, and M. G. Schepkin, Yad. Fiz. **45**, 1252 (1987) [Sov. J. Nucl. Phys. **45**, 776 (1987)].

[5] R. Bilger *et al.*, Phys. Lett. B **269**, 247 (1991); R. Bilger, H. A. Clement, and M. G. Schepkin, Phys. Rev. Lett. **71**, 42 (1993); H. Clement, M. Schepkin, G. J. Wagner, and O. Zaboronsky, Phys. Lett. B **337**, 43 (1994).

[6] B. V. Martemyanov and M. G. Schepkin, Pis'ma Zh. Eksp. Teor. Fiz. **53**, 132 (1991) [JETP Lett. **53**, 139 (1991)].

[7] R. Meyer *et al.*, contribution to PANIC 1996 (unpublished); W. Brodowski *et al.*, Z. Phys. C **355**, 5 (1996).

[8] L. Ya. Glozman, A. J. Buchmann, and Amand Faessler, J. Phys. G **20**, L49 (1994); Georg Wagner, L. Ya. Glozman, A. J. Buchmann, and Amand Faessler, Nucl. Phys. **A594**, 263 (1995); A. J. Buchmann, Georg Wagner, K. Tsushima, L. Ya. Glozman, and Amand Faessler, Prog. Part. Nucl. Phys. **36**, 383 (1995); K. Itonaga, A. J. Buchmann, Georg Wagner, and Amand Faessler, Nucl. Phys. **A609**, 422 (1996); A. J. Buchmann, Georg Wagner, and Amand Faessler, Phys. Rev. C (submitted).

[9] S. B. Gerasimov and A. S. Khrykin, Mod. Phys. Lett. A **8**, 2457 (1993).

[10] A. S. Khrykin, πN Newslett. **10**, 67 (1995); contribution to PANIC 96 (unpublished); V. M. Abazov *et al.*, Report No. JINR-E1-96-104, 1996.

[11] A. M. Baldin *et al.*, Dokl. Acad. Sci. USSR, **279**, 602 (1984); A. V. Chizov *et al.*, Nucl. Phys. **A449**, 660 (1986).

[12] M. I. Krivoruchenko, JETP Lett. **46**, 3 (1987).

[13] A. J. Buchmann, Amand Faessler, and M. I. Krivoruchenko, Ann. Phys. (N.Y.) **254**, 109 (1997).

[14] Amand Faessler, A. J. Buchmann, M. I. Krivoruchenko, and B. V. Martemyanov, Phys. Lett. B **391**, 255 (1997).

[15] Amand Faessler, A. J. Buchmann, M. I. Krivoruchenko, and B. V. Martemyanov (unpublished).

[16] J. D. Walecka, Ann. Phys. (N.Y.) **83**, 491 (1974); S. A. Chin and J. D. Walecka, Phys. Lett. **52B**, 24 (1974); S. A. Chin, Ann. Phys. (N.Y.) **108**, 301 (1977); B. D. Serot, Rep. Prog. Phys. **55**, 1855 (1992).

[17] A. A. Abrikosov, L. P. Gorkov, and I. E. Dzyaloshinski, *Methods of Quantum Field Theory in Statistical Physics* (Prentice-Hall, Englewood Cliffs, NJ, 1963).

[18] S. L. Shapiro and S. A. Teukolsky, *Black Holes, White Dwarfs, and Neutron Stars: The Physics of Compact Objects* (Wiley, New York, 1983).

[19] R. Tamagaki, Prog. Theor. Phys. **85**, 321 (1991); A. V. Olinto, P. Haensel, and J. Frieman, Report No. FERMILAB-PUB-91-176-A, 1991.

[20] E. Witten, Phys. Rev. D **30**, 272 (1984); E. Farhi and R. L. Jaffe, *ibid.* **30**, 3279 (1984); A. V. Olinto, in Relativistic Aspects of Nuclear Physics—1991, Proceedings of the International Conference, Rio de Janeiro, Brazil, 1991 (unpublished), p. 187; C. Greiner and J. Schaffner, Int. J. Mod. Phys. E **5**, 239 (1996).

[21] M. I. Krivoruchenko and B. V. Martemyanov, Astrophys. J. **378**, 628 (1991); Nucl. Phys. **B24**, 134c (1991).

[22] G. Baum, H. A. Bethe, and C. J. Pethick, Nucl. Phys. **A175**, 225 (1971).

[23] G. Baum, C. J. Pethick, and P. Sutherland, Astrophys. J. **170**, 299 (1971).

[24] U. Straub, Zong-Ye Zhang, K. Bräuer, Amand Faessler, and S. B. Khadkikar, Phys. Lett. B **200**, 241 (1988); Amand Faessler and U. Straub, Prog. Part. Nucl. Phys. **24**, 323 (1985); U. Straub, Z. Y. Zhang, K. Bräuer, Amand Faessler, S. B. Khadkikar, and G. Lübeck, Nucl. Phys. **A483**, 686 (1988).

[25] M. Oka, K. Shimizu, and K. Yazaki, Nucl. Phys. **A464**, 700 (1987); T. Sakai, A. Buchmann, and K. Yazaki, *ibid.* **A543**, 661 (1992); T. Sakai, Ph.D. thesis, University of Tokyo, 1994; T. Sakai, K. Shimizu, and K. Yazaki, Nucl. Phys. **A594**, 247 (1995).

[26] T. Sakai, J. Mori, A. J. Buchmann, K. Shimizu, and K. Yazaki, Research Center for Nuclear Physics (RCNP), Osaka, Annual Report 1995, p. 118; see also, Carl Dover, memorial volume Nucl. Phys. A (to be published).

[27] A. V. Olinto, Phys. Lett. B **192**, 71 (1987); Nucl. Phys. **B24**, 103c (1991); B. V. Martemyanov, Pis'ma Astron. Zh. **20**, 588 (1994) [Sov. Astron. Lett. **20**, 499 (1994)].