Constraints to Coupling Constants of the $\omega$- and $\sigma$-Mesons with Dibaryons

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Abstract

The effect of narrow dibaryon resonances to nuclear matter and structure of neutron stars is investigated in the mean-field theory (MFT) and in the relativistic Hartree approximation (RHA). The existence of massive neutron stars imposes constraints to the coupling constants of the $\omega$- and $\sigma$-mesons with dibaryons. We conclude that the experimental candidates to dibaryons $d_1(1920)$ and $d'(2060)$ if exist form in nuclear matter a Bose condensate stable against compression. This proves stability of the ground state for nuclear matter with a Bose condensate of the light dibaryons.
The prospect to observe the long-lived H-particle predicted in 1977 by R. Jaffe [1] stimulated considerable activity in the experimental searches of dibaryons. It was proposed to examine the H-particle production in different reactions [2]. The experiments [3] did not give a positive sign for the H-particle, however, the existence of the H-particle remains an open question which must eventually be settled by experiment. The non-strange 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 pion double charge exchange (DCE) reactions on nuclei [5] exhibit a peculiar energy dependence, which can be interpreted [6] as evidence for the existence of a narrow d’ dibaryon with quantum numbers \( T = 0, J^P = 0^- \) and the total resonance energy of 2063 MeV. Recent experiments at TRIUMPF (Vancouver) and CELSIUS (Uppsala) seem to support the existence of the d’ dibaryon [7]. A method for searching narrow, exotic dibaryon resonances in the double proton-proton bremsstrahlung reaction is discussed in Ref. [8]. Recently, some indications for a d_1(1920) dibaryon in this reaction have been found [9].

When density of nuclear matter is increased beyond a critical value, production of dibaryons becomes energetically favorable. Dibaryons are Bose particles, so they condense in the ground state and form a Bose condensate [10, 11]. An exactly solvable model for a one-dimensional Fermi-system of fermions interacting through a potential leading to a resonance in the two-fermion channel is analyzed in Ref. [12]. The behavior of the system with increasing the density can be interpreted in terms of a Bose condensation of two-fermion resonances. The effect of narrow dibaryon resonances on nuclear matter in the mean field theory (MFT) is analyzed in Refs. [13, 14]. In the limit of vanishing decay width, a dibaryon can be approximately described as an elementary field.

Despite the dibaryon Bose condensate does not exist in ordinary nuclei, dibaryons affect properties of nuclear matter and the ordinary nuclei through a Casimir effect. Presence of the background \( \sigma \)-meson mean field inside of nuclei modifies the nucleon and dibaryon masses and in turn modifies the zero-point vacuum fluctuations of the nucleon and dibaryon fields. This effect contributes to the energy density and pressure. It can be evaluated within the relativistic Hartree approximation (RHA). For nucleons, this effect is well known [15]. In the loop expansion of quantum hadrodynamics (QHD), MFT corresponds to the lowest approximation (no loops), while RHA corresponds to the one-loop approximation in a calculation of the equation of state for nuclear matter.

At zero temperature, a uniformly distributed system of bosons with attractive potential is energetically unstable against compression and collapses [16]. In such a case, the long wave excitations (sound in the medium) have imaginary dispersion law: The square of the sound velocity is negative \( c^2_s < 0 \). The amplitude of these excitations increases with the time, providing instability of the system. It is necessary to analyze dispersion laws of other elementary excitations also. We shall see, however, that in MFT and RHA only sound waves can generate an instability. The ground state of nuclear matter with a Bose condensate of dibaryons is stable or unstable against small perturbations according as the repulsive \( \omega \)-meson exchange or the attractive \( \sigma \)-meson exchange is dominant between dibaryons.

In this paper, we investigate the hypothesis that the dibaryon matter is unstable against compression. In such a case, formation of dibaryons in nuclear matter can be treated as a possible mechanism for a phase transition into the quark matter. If central density of a massive neutron stars exceeds a critical value for formation of dibaryons, the
neutron star should 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 [17]. From the requirement that the dibaryon formation is not energetically favorable at densities lower than the central density of neutron stars with a mass $1.3M_\odot$, we derive constraints to the coupling constants of the mesons and dibaryons $d_1(1920)$ and $d'(2060)$ and conclude that narrow dibaryons in this mass range can form a Bose condensate stable against perturbations only. The effect of the dibaryons to stability and structure of neutron stars in different phenomenological models is analyzed in Refs. [11, 18]. Constraints to the binding energy of strange matter [19] from the existence of massive neutron stars are discussed in Ref. [20].

The dibaryonic extension of the Walecka model [13] is obtained by including dibaryons to the Lagrangian density [13, 14]

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

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 $\sigma$-meson mean field $\sigma_c$ 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 expression [13]

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

where

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

The last term in Eq.(4) occurs after the renormalization of the scalar density. Here, $\gamma = 2$ for neutron matter and $\gamma = 4$ for nuclear matter.

We investigate here properties of the nuclear matter below the critical density for formation of dibaryons, so the dibaryon condensate is zero $\langle |\varphi(0)| \rangle = 0$. The vacuum contribution to the scalar density of dibaryons can be found to be

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

It differs from the nucleon term in the sign and in the statistical factor (one should replace 4 ($= 2_s \times 2_t$) → 1). 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 [13]

$$< T_{\mu\nu}^N(0) >_{\text{vac}} = -4g_{\mu\nu}m_N^4\eta(m_N^*/m_N) \quad (7)$$
where
\[ 16\pi^2 \eta(x) = x^4 \ln x + 1 - x - \frac{7}{2} (1 - x)^2 + \frac{13}{3} (1 - x)^3 - \frac{25}{12} (1 - x)^4. \]

For dibaryons, we get expression
\[ < T_{\mu\nu}^D(0) >_{vac} = g_{\mu\nu} m_D^4 \eta(m_D^*/m_D). \] (8)

The elementary excitations in nuclear matter with a Bose condensate of dibaryons correspond to nucleons and antinucleons, \( \omega \)-mesons, \( \sigma \)-mesons, and dibaryons and antidibaryons. The dispersion laws for these quasiparticles are found in Ref. [14]. The nucleon and antinucleon dispersion laws have the same form as in the vacuum with a replacement \( m_N \rightarrow m_N^* \) and therefore cannot generate an instability of the system. The dispersion laws of \( \omega \)-mesons, \( \sigma \)-mesons, and antidibaryons turn out to be real also. The possible source of the instability are the dibaryon quasiparticle excitations only, which are responsible for a long wave perturbations of the system and connected with existence of the sound in the medium.

The square of the sound velocity has the form [14]
\[ a_s^2 = \frac{\alpha}{1 + \alpha} \] (9)

where
\[ \alpha = 2 \rho DS \frac{m_\omega^2 h_\omega^2}{m_\sigma^2} - \frac{h_\sigma^2}{m_\sigma^2} \] (10)

and \( \tilde{m}_\sigma^2 = m_\sigma^2 + 2 h_\sigma^2 \rho DS \). We see that \( a_s^2 > 0 \) for
\[ \frac{h_\omega^2}{m_\omega^2} > \frac{h_\sigma^2}{m_\sigma^2}. \] (11)

The validity of inequality (11) is sufficient condition for stability of the ground state of nuclear matter with a Bose condensate of dibaryons.

The physical meaning of the inequality (11) can be clarified by considering the interaction energy of a uniformly distributed dibaryon matter \( \rho_{DV}(x_1) = \rho_{DV} = \text{constant} \):
\[ W = \frac{1}{2} \int d\mathbf{x}_1 d\mathbf{x}_2 \rho_{DV}(\mathbf{x}_1) \rho_{DV}(\mathbf{x}_2) V(|\mathbf{x}_1 - \mathbf{x}_2|). \] (12)

The Yukawa potential \( V(r) \) for two dibaryons has the form
\[ V(r) = \frac{h_\omega^2 e^{-m_\omega r}}{4\pi r} - \frac{h_\sigma^2 e^{-m_\sigma r}}{4\pi r}. \] (13)

The integration gives
\[ W = \frac{1}{2} N_D \rho_{DV}(\frac{h_\omega^2}{m_\omega^2} - \frac{h_\sigma^2}{m_\sigma^2}) \] (14)

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

At present, the coupling constants of the mesons with dibaryons are not known with precision good enough to draw a definite conclusion concerning the stability of the dibaryon matter.
Here, we show that violation of the inequality (11) for light dibaryons is in contradiction with the existence of massive neutron stars.

Given that the ratio $h_\sigma/(2g_\sigma)$ between the $\sigma$-meson couplings with dibaryons and nucleons is fixed, one can find the $\omega$- and $\sigma$-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$^{-3}$ determined from the density in the interior of $^{208}$Pb $^{15}$. It corresponds to the equilibrium Fermi wavenumber $k_F = 1.3$ fm$^{-1}$. The dependence of the effective nucleon mass at the equilibrium density on the ratio $h_\sigma/(2g_\sigma)$ is shown on Fig.1 (a). In Figs.1 (b) and (c) we give 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 $C^2_\omega = g'^2_\omega(m_N/m_\omega)^2$ and $C^D_\omega = g'^2_\omega(m_N/m_\omega)^2$ are plotted. Notice that $h_\sigma/(2g_\sigma) = 0$ is equivalent to RHA with no dibaryons. For comparison, we give the results of MFT where the effect of dibaryons to the nuclear matter below the critical density for occurrence of a Bose condensate of dibaryons is absent. The results RHA for different dibaryons are not much different. 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) shift the critical value $h_\sigma$ only by about 25% if one assumes that the $\sigma$-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.

When the ratio $h_\sigma/(2g_\sigma)$ approaches a value 0.8, the system of equations blows up and the empirical equilibrium properties of the nuclear matter can no longer be reproduced. When $x \rightarrow 0 \ z(x) = O(x^4)$, so the zero-point contributions to the scalar density of nucleons and dibaryons, which have the opposite signs, are comparable for $4g'^2_\omega/m_N \approx h_\sigma^4/m_D$. The dibaryon effects become large for $h_\sigma/(2g_\sigma) \approx 0.5(4m_D/m_N)^{1/4} \approx 0.84$. The greater dibaryon mass, the greater the upper limit to the ratio $h_\sigma/(2g_\sigma)$. This effect is seen in Fig.1.

The saturation curve is shown in Fig.2 and the effective nucleon mass dependence on the Fermi wavenumber $k_F$ is shown in Fig.3 for $h_\sigma/(2g_\sigma) = 0.6$ in case of the $H$-particle. EOS in RHA is softer than in MFT. The contributions of the vacuum zero-point fluctuations of nucleons and dibaryons partially cancel each other, so the inclusion of the dibaryons makes the EOS stiffer. In Fig.2, we see that the dashed curves corresponding to RHA with dibaryons lie above the solid lined corresponding to the RHA with no dibaryons. The same effect is seen on Fig.3. In MFT, the nucleon effective mass decreases with the density faster then in RHA. Due to the partial compensation 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 occurrence of dibaryons $H(2220)$, $d'(2060)$, and $d_1(1920)$ in MFT and RHA in nuclear and neutron matter versus the ratio $h_\omega/(2g_\sigma)$ for $h_\omega/h_\omega^{max} = 1, 0.8$, and 0.6 where $h_\omega^{max} = h_\sigma m_\omega/m_\sigma$ is the maximum value for the $\omega$-meson coupling constant with dibaryons at which the inequality (11) is violated ($h_\omega/h_\omega^{max} = 1$ corresponds to $a_2^\sigma = 0$ and $h_\omega/h_\omega^{max} = 0.8$ and 0.6 correspond to $a_2^\sigma < 0$). In RHA, dibaryons occur at higher densities. The coupling constant $h_\omega$ determines the energy of dibaryons in the positive $\omega$-meson mean field. The greater the $h_\omega$, the greater the density is required to make 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.
\[ h_\omega / h_\omega^{\text{max}} = 0.8 \text{ and } 0.6, \text{ respectively.} \]

When the dibaryon matter is unstable against compression, production of dibaryons with increasing the density results to instability of neutron stars with subsequent phase transition into the quark matter and conversion of neutron stars into 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 in which dibaryons can occur.

The MFT and RHA EOS for neutron matter at supranuclear densities are matched smoothly with the BBP EOS [21] at densities \( \rho_{\text{drip}} < \rho < 0.8 \rho_0 \) where \( \rho_{\text{drip}} = 4.3 \times 10^{11} \text{g/cm}^3 \) and then with the BPS EOS [22] at densities \( \rho < \rho_{\text{drip}} \). The maximum neutron star masses are sensitive to the value of the equilibrium Fermi wavenumber. If we chose \( k_F = 1.42 \text{ fm}^{-1} \) instead of \( k_F = 1.3 \text{ fm}^{-1} \), the maximum masses in MFT (with no dibaryons) are reduced from \( 3M_\odot \) down to \( 2.6M_\odot \) [15]. The choice \( k_F = 1.3 \text{ fm}^{-1} \) provides less stringent (therefore more conservative) constraints to the meson-dibaryon coupling constants. We do not show the results for the \( d_1(1920) \) dibaryon, since its condensation starts at a density \( \rho \approx \rho_0 \) providing conversion to quark stars of neutron stars with very low masses \( M < 0.2M_\odot \).

In Fig. 6 we show the parameter space for the coupling constants of dibaryons with the mesons. Our discussion and the validity of our arguments are restricted to the region \( a_s^2 < 0 \) in which the dibaryon matter is unstable against compression. The requirement of stability of the normal nuclear matter at the saturation density allows to get constraints to the coupling constants. The corresponding curves (straight lines in the MFT case) marked by arrows with the white ends restrict from below the parameter space of the coupling constants. The dotted line in the MFT case, which refers to the \( d_1(1920) \) dibaryon, is very close to the dashed-dotted line \( a_s^2 = 0 \). For low values \( h_\sigma \) the dotted line lies above the line \( a_s^2 = 0 \). In such a case, the dibaryon matter unstable against compression cannot exist. The window in the parameter space for the unstable dibaryon matter for \( d_1(1920) \) is, however, much greater for RHA.

If a conservative assumption is used, namely, that pulsars with a mass \( 1.3M_\odot \) are ordinary neutron stars, the constraints to the meson coupling constants with dibaryons can further be improved. The corresponding curves (straight lines in the MFT case) are shown on Fig.6. We see that the dotted and dashed lines lie above or very close to the line \( a_s^2 = 0 \). It means that for \( d_1(1920) \) and \( d'(2060) \) dibaryons, the dibaryon matter unstable against compression cannot exist (owing to a very small window for \( d'(2060) \) at higher values of the \( h_\sigma \)). For the H-particle, there is a window in the parameter space between the line \( a_s^2 = 0 \) and the solid curves marked by arrows with the black ends, which corresponds to the dibaryon matter unstable against compression. The H-particle interaction were studied in the non-relativistic quark cluster model [23, 24] successful in describing 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 [25] 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 [13, 14]. They correspond to the unstable dibaryon matter and are in the allowed region of the parameter space for the H-particle. The energetically favorable compression of the H-matter can lead to the formation of the absolutely stable strange matter [19], producing conversion of neutron stars to strange stars [26].

MFT and RHA EOS both are very stiff. In the soft models like the Reid one (for
a review of the nuclear matter models see [17], the central density of neutron stars is much greater than in stiff models. Respectively, conditions for 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 to the neutron star masses and, respectively, more stringent constraints to the meson-dibaryon coupling constants as compared to the stiffer MFT EOS, despite in RHA dibaryons occur at higher densities (see Fig.4). One can assume that this effect is of the general validity and that softer EOS like the Reid one give more stringent constraints to the meson-dibaryon coupling constants. We consider therefore the constraints given in Fig.6 as the conservative ones.

In the conclusion, we showed that the hypothesis of instability of the dibaryon matter against compression is in contradiction with the hypothesis that pulsars of a mass $1.3M_\odot$ are ordinary neutron stars for dibaryons $d_1(1920)$ and $d'(2060)$. This conclusion is valid for all narrow dibaryons with the same quantum numbers in the same mass range. The $H$-particle is sufficiently heavy, its condensation starts at higher densities, respectively, constraints to the meson-dibaryon coupling constants are not much stringent, allowing a possibility for the $H$-particle to form a condensate unstable against compression with subsequent formation of the strange matter. The meson coupling constants with dibaryons $d_1(1920)$ and $d'(2060)$ should obey the inequality (11). The meson coupling constants with the $H$-particle should lie above the solid curves on Fig.6.

The authors are grateful to B. V. Martemyanov for discussions of the results. One of the authors (M. I. K.) acknowledges hospitality of Institute for Theoretical Physics of University of Tuebingen, Alexander von Humboldt Stiftung for support with a Forschungsstipendium and DFG-RFBR for Grant No. Fa-67/20-1.
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The effective nucleon mass (a), the asymmetry coefficient (b), the incompressibility of nuclear matter (c) at the saturation density, and the coupling constants $C^2_s = g^2_s(m_N/m_s)^2$ and $C^2_\omega = g^2_\omega(m_N/m_\omega)^2$ in the RHA versus the ratio $h_\sigma/(2g_\sigma)$ between the $\sigma$-meson coupling constants with dibaryons and nucleons. The MFT results are shown for comparison. The coupling constants are fixed by fitting the minimum 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 $d_1(1920)$. The reasonable description of the properties of the nuclear matter at the saturation density is possible for $h_\sigma/(2g_\sigma) < 0.8$.

Fig. 2 Saturation curve for nuclear matter in RHA with no dibaryons (solid line) and in RHA with inclusion of the H-dibaryon (dashed line) for $h_\sigma/(2g_\sigma) = 0.6$.

Fig. 3 The effective nucleon mass versus the fermi momentum of nucleons in the nuclear matter (upper curves) and in the neutron matter (lower curves). The solid lines correspond to RHA with no dibaryons, the dashed lines correspond to RHA with inclusion of the H-dibaryon.

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

Fig. 5 The minimum neutron star masses, in which the dibaryon formation becomes energetically favorable, versus the $\sigma$-meson coupling constant with dibaryons for $h_\omega/h_\omega^{max} = 1$, 0.8, and 0.6 (the solid, longed-dashed, and dashed lines, respectively) where $h_\omega^{max}$ is the maximum value for the $\omega$-meson coupling constant at which the dibaryon matter is unstable against compression ($a^2_\sigma < 0$). The results are given in MFT and RHA and for two dibaryons $H(2220)$ and $d'(2060)$. The minimum neutron star masses for the $d_1(1920)$ are very small ($< 0.2M_\odot$).

Fig. 6 Parameter space for the coupling constants of the $\sigma$- and $\omega$-mesons with dibaryons in MFT and RHA. The dashed-dotted line $a^2_\sigma = 0$ divides the parameter space into two parts. The upper left part of the parameter space corresponds to the dibaryon matter stable against compression (square of the sound velocity is positive $a^2_\sigma > 0$), the lower right part corresponds to the nuclear matter unstable against compression ($a^2_\sigma > 0$). The solid, dashed, and dotted curves constrain from below the regions in which the dibaryon formation is energetically not favorable in ordinary nuclei (curves marked by arrows with white ends) and at the density equal to the central density of a $1.3M_\odot$ mass neutron star (curves marked by arrows with black ends). The cross refers to the H-particle coupling constants with the mesons, determined from the adiabatic potential [25].