

# Effects of Dirac sea polarization on hadronic properties -A chiral SU(3) approach

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## Abstract

The effect of vacuum fluctuations on the in-medium hadronic properties is investigated using a chiral SU(3) model in the nonlinear realization. The effect of the baryon Dirac sea is seen to modify hadronic properties and in contrast to a calculation in mean field approximation it is seen to give rise to a significant drop of the vector meson masses in hot and dense matter. This effect is taken into account through the summation of baryonic tadpole diagrams in the relativistic Hartree approximation (RHA), where the baryon self energy is modified due to interactions with both the non-strange ( $\sigma$ ) and the strange ( $\zeta$ ) scalar fields.

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## 1. INTRODUCTION

The study of hot and dense matter is an important problem in strong interaction physics. In recent times there have been numerous experimental investigations, e.g. in the context of relativistic heavy ion collision experiments, to study how hadronic matter is modified under extreme conditions of high temperatures and/or densities [1, 2, 3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15]. The experimental observables from relativistic heavy ion collision are related to the medium modifications of the hadrons in the dynamically evolving strongly interacting matter (fireball) resulting from the nuclear collision. One of the explanations of the observed enhanced dilepton production in the low invariant mass regime is the medium modification of the vector mesons [1, 2, 3]. It was first conjectured to be a simple scaling law for the vector meson masses in the medium [6]. There have been also QCD sum rule calculations for studying the in-medium vector meson properties [8, 9, 10, 11]. In the Quantum Hadrodynamics framework [16], the vector meson masses were shown to have dominant contributions from the nucleon Dirac sea. The vector meson masses have very insignificant modifications in the hot and dense medium, when the contributions only from the Fermi sea are taken into account [17, 18, 19, 20, 21]. Recently, it was shown in a chiral SU(3) model [23, 24], that the Dirac sea polarization leads to a significant drop of vector meson masses in nuclear matter [25]. In the present investigation, we study the properties of hadrons in the hot hyperonic matter, taking into account the effects of the Dirac sea polarization [26, 27, 28].

We organize the paper as follows. In the section 2, we briefly recapitulate the SU(3) chiral model used in the present investigation. In section 3, we outline the mean field approximation for the study of hadronic properties. Section 4 discusses the inclusion of vacuum polarization effects using RHA. This is done by summing over the baryonic tadpole diagrams which includes couplings to both the nonstrange ( $\sigma$ ) and strange ( $\zeta$ ) scalar meson fields. In section 5 we discuss the medium modification of the vector meson masses due to their interaction with the baryons in the hot hadronic matter. The effect of vacuum polarization compared to the mean field approximation is studied. We discuss the results of the present investigation in section 6. Finally, in section 7, we summarize our findings and discuss possible improvements of the current approach.

## 2. THE HADRONIC CHIRAL $SU(3) \times SU(3)$ MODEL

In this section the various terms of the Lagrangian

$$\mathcal{L} = \mathcal{L}_{kin} + \sum_{W=X,Y,V,A,u} \mathcal{L}_{BW} + \mathcal{L}_{VP} + \mathcal{L}_{vec} + \mathcal{L}_0 + \mathcal{L}_{SB} \quad (1)$$

are discussed. The calculation is done within the framework of a relativistic quantum field theoretical model of baryons and mesons built on chiral symmetry and broken scale invariance [23] to describe strongly interacting nuclear matter. We adopt a nonlinear realization of the chiral symmetry which allows a simultaneous description of hyperon potentials and properties of finite nuclei [22, 23, 24]. This Lagrangian contains the baryon octet, the spin-0 and spin-1 meson multiplets as degrees of freedom.

$\mathcal{L}_{kin}$  is the kinetic energy term,  $\mathcal{L}_{BW}$  contains the baryon-meson interactions in which the baryon-spin-0 meson interaction terms generate the baryon masses.  $\mathcal{L}_{VP}$  describes the interactions of vector mesons with the pseudoscalar mesons (and with photons).  $\mathcal{L}_{vec}$  describes the dynamical mass generation of the vector mesons through coupling to the scalar fields and contains additionally quartic self-interactions of the vector fields.  $\mathcal{L}_0$  are the meson-meson interaction terms inducing the spontaneous breaking of chiral symmetry. It also includes a scale invariance breaking logarithmic potential.  $\mathcal{L}_{SB}$  describes the explicit symmetry breaking of  $U(1)_A$ ,  $SU(3)_V$  and the chiral symmetry.

### 2.1. The kinetic energy terms

An important property of the nonlinear realization of chiral symmetry is that all terms of the model-Lagrangian only have to be invariant under the  $SU(3)_V$  transformation in order to ensure chiral symmetry. This vector transformation depends in general on the pseudoscalar mesons and thus is local. Covariant derivatives have to be introduced for the kinetic terms in order to preserve chiral invariance [23]. The covariant derivative used in this case, reads:  $D_\mu = \partial_\mu + [\Gamma_\mu, \ ]$  with  $\Gamma_\mu = -\frac{i}{2}[u^\dagger \partial_\mu u + u \partial_\mu u^\dagger]$  where  $u = \exp\left[\frac{i}{\sigma_0} \pi^a \lambda^a \gamma_5\right]$  is the unitary transformation operator [23]. The pseudoscalar mesons are given as parameters of the symmetry transformation.

In summary, the kinetic energy terms read [23]

$$\begin{aligned} \mathcal{L}_{kin} = & i \text{Tr} \bar{B} \gamma_\mu D^\mu B + \frac{1}{2} \text{Tr} D_\mu X D^\mu X + \text{Tr}(u_\mu X u^\mu X + X u_\mu u^\mu X) + \frac{1}{2} \text{Tr} D_\mu Y D^\mu Y \\ & + \frac{1}{2} D_\mu \chi D^\mu \chi - \frac{1}{4} \text{Tr} \left( \tilde{V}_{\mu\nu} \tilde{V}^{\mu\nu} \right) - \frac{1}{4} \text{Tr} (F_{\mu\nu} F^{\mu\nu}) - \frac{1}{4} \text{Tr} (\mathcal{A}_{\mu\nu} \mathcal{A}^{\mu\nu}) . \end{aligned} \quad (2)$$

B denotes the baryon octet, X the scalar meson multiplet, Y the pseudoscalar chiral singlet,  $\tilde{V}^\mu$  ( $\mathcal{A}^\mu$ ) the renormalised vector (axial vector) meson multiplet with the field strength tensor

$\tilde{V}_{\mu\nu} = \partial_\mu \tilde{V}_\nu - \partial_\nu \tilde{V}_\mu$  ( $\mathcal{A}_{\mu\nu} = \partial_\mu \mathcal{A}_\nu - \partial_\nu \mathcal{A}_\mu$ ).  $F_{\mu\nu}$  is the electro-magnetic field tensor and  $\chi$  is the scalar, iso-scalar dilaton (glueball) -field.

## 2.2. Baryon-meson interaction

Except for the difference in Lorentz indices, the SU(3) structure of the spin-1/2 baryon-meson interaction terms are the same for all mesons. This interaction for a general meson field  $W$  has the form

$$\mathcal{L}_{BW} = -\sqrt{2}g_8^W (\alpha_W [\overline{B}OBW]_F + (1 - \alpha_W) [\overline{B}OBW]_D) - g_1^W \frac{1}{\sqrt{3}} \text{Tr}(\overline{B}OB) \text{Tr}W, \quad (3)$$

with  $[\overline{B}OBW]_F := \text{Tr}(\overline{B}OWB - \overline{B}OBW)$  and  $[\overline{B}OBW]_D := \text{Tr}(\overline{B}OWB + \overline{B}OBW) - \frac{2}{3} \text{Tr}(\overline{B}OB) \text{Tr}W$ . The different terms to be considered are those for the interaction of baryons with scalar mesons ( $W = X, \mathcal{O} = 1$ ), with vector mesons ( $W = \tilde{V}_\mu, \mathcal{O} = \gamma_\mu$  for the vector and  $W = \tilde{V}_{\mu\nu}, \mathcal{O} = \sigma^{\mu\nu}$  for the tensor interaction), with axial vector mesons ( $W = \mathcal{A}_\mu, \mathcal{O} = \gamma_\mu \gamma_5$ ) and with pseudoscalar mesons ( $W = u_\mu, \mathcal{O} = \gamma_\mu \gamma_5$ ), respectively. For the current investigation the following interactions are relevant.

### 2.2.1. Baryon-scalar meson interaction

This is the term generating the baryon masses through coupling of the baryons to the non-strange  $\sigma$  ( $\sim \langle \bar{u}u + \bar{d}d \rangle$ ) and the strange  $\zeta$  ( $\sim \langle \bar{s}s \rangle$ ) scalar quark condensate [23]. After insertion of the scalar meson matrix  $X$ , one obtains the baryon masses

$$\begin{aligned} m_N &= m_0 - \frac{1}{3}g_8^S(4\alpha_S - 1)(\sqrt{2}\zeta - \sigma), \\ m_\Lambda &= m_0 - \frac{2}{3}g_8^S(\alpha_S - 1)(\sqrt{2}\zeta - \sigma), \\ m_\Sigma &= m_0 + \frac{2}{3}g_8^S(\alpha_S - 1)(\sqrt{2}\zeta - \sigma), \\ m_\Xi &= m_0 + \frac{1}{3}g_8^S(2\alpha_S + 1)(\sqrt{2}\zeta - \sigma), \end{aligned} \quad (4)$$

with  $m_0 = g_1^S(\sqrt{2}\sigma + \zeta)/\sqrt{3}$ . The parameters  $g_1^S, g_8^S$  and  $\alpha_S$  can be used to fix the baryon masses to their experimentally measured vacuum values. It should be emphasized that the nucleon mass also depends on the *strange condensate*  $\zeta$ . This general case will be used in the present investigation, to study hot and strange hadronic matter. Recently the vector meson masses were investigated in nuclear matter[25], for the situation  $\alpha_S = 1$  and  $g_1^S = \sqrt{6}g_8^S$ , where the nucleon mass depends only on the non-strange quark condensate [23, 25].

The effect of including RHA is similar to the results obtained in the Walecka model [17, 26].

In the present investigation, however, in the summing over baryon tadpoles the effects of coupling of baryons to both scalar fields ( $\sigma$  and  $\zeta$ ) have to be taken into account for the baryonic Dirac sea in RHA.

### 2.2.2. Baryon-vector meson interaction

In analogy to the baryon-scalar meson coupling there exist two independent baryon-vector meson interaction terms corresponding to the F-type (antisymmetric) and D-type (symmetric) couplings. Here we will use the symmetric coupling because from the universality principle [29] and the vector meson dominance model one can conclude that the antisymmetric coupling should be small. We realize this assumption by setting  $\alpha_V = 1$  for all fits. Additionally we decouple the strange vector field  $\phi_\mu \sim \bar{s}\gamma_\mu s$  from the nucleon by setting  $g_1^V = \sqrt{6}g_8^V$  and the remaining baryon-vector meson interaction reads

$$\mathcal{L}_{BV} = -\sqrt{2}g_8^V \left\{ [\bar{B}\gamma_\mu B V^\mu]_F + \text{Tr}(\bar{B}\gamma_\mu B) \text{Tr} V^\mu \right\}. \quad (5)$$

Note that in this limit all coupling constants are fixed once  $g_8^V$  is specified [23]. This is done by fitting the nucleon- $\omega$  coupling to the energy density at nuclear matter saturation ( $E/A = -16$  MeV). With the above choice, the vector meson- baryon couplings reduce to those from the additive quark model given as

$$\begin{aligned} g_{\Lambda\omega} &= \frac{2}{3}g_{N\omega} = g_{\Sigma\omega} = 2g_{\Xi\omega} \\ g_{\Lambda\phi} &= -\frac{\sqrt{2}}{3}g_{N\omega} = g_{\Sigma\phi} = 2g_{\Xi\phi} \end{aligned} \quad (6)$$

## 2.3. Meson-meson interaction

### 2.3.1. Spin-0 potential

The Lagrangian describing the interaction for the scalar mesons,  $X$ , and pseudoscalar singlet,  $Y$ , is given as [23]

$$\mathcal{L}_0 = -\frac{1}{2}k_0\chi^2 I_2 + k_1(I_2)^2 + k_2 I_4 + 2k_3\chi I_3, \quad (7)$$

with  $I_2 = \text{Tr}(X + iY)^2$ ,  $I_3 = \det(X + iY)$  and  $I_4 = \text{Tr}(X + iY)^4$ . The scalar glueball field  $\chi$  is introduced to satisfy the QCD trace anomaly i.e. nonvanishing energy-momentum tensor  $\Theta_\mu^\mu = (\beta_{QCD}/2g)\langle G_{\mu\nu}^a G^{a,\mu\nu} \rangle$ , where  $G_{\mu\nu}^a$  is the gluon field tensor.

A scale breaking potential [33]

$$\mathcal{L}_{\text{scalebreak}} = -\frac{1}{4}\chi^4 \ln \frac{\chi^4}{\chi_0^4} + \frac{\delta}{3}\chi^4 \ln \frac{I_3}{\det\langle X \rangle_0} \quad (8)$$

is introduced. This yields

$$\theta_\mu^\mu = 4\mathcal{L} - \chi \frac{\partial \mathcal{L}}{\partial \chi} - 2\partial_\mu \chi \frac{\partial \mathcal{L}}{\partial(\partial_\mu \chi)} = \chi^4 \quad (9)$$

and allows for the identification of the  $\chi$  field with the gluon condensate  $\Theta_\mu^\mu = (1 - \delta)\theta_\mu^\mu = (1 - \delta)\chi^4$ . Finally the term  $\mathcal{L}_\chi = -k_4\chi^4$  generates a phenomenologically consistent finite vacuum expectation value. We shall use the frozen glueball approximation i.e. assume  $\chi = \langle 0|\chi|0 \rangle \equiv \chi_0$ , since the variation of  $\chi$  in the medium is rather small [23].

### 2.3.2. Vector mesons masses

The Lagrangian for the vector meson interaction is written as [30, 31, 32]

$$\begin{aligned} \mathcal{L}_{vec} = & \frac{1}{2}m_V^2 \frac{\chi^2}{\chi_0^2} \text{Tr}(\tilde{V}_\mu \tilde{V}^\mu) + \frac{1}{4}\mu \text{Tr}(\tilde{V}_{\mu\nu} \tilde{V}^{\mu\nu} X^2) \\ & + \frac{1}{12}\lambda_V \left( \text{Tr}(\tilde{V}^{\mu\nu}) \right)^2 + 2(\tilde{g}_4)^4 \text{Tr}(\tilde{V}_\mu \tilde{V}^\mu)^2. \end{aligned} \quad (10)$$

The vector meson fields,  $\tilde{V}_\mu$  are related to the renormalized fields by  $V_\mu = Z_V^{1/2} \tilde{V}_\mu$  with  $V = \omega, \rho, \phi$  [25]. The masses of  $\omega, \rho$  and  $\phi$  are fitted by tuning  $m_V, \mu$  and  $\lambda_V$ . The vector meson masses have contributions from the quartic self-interaction, and we get in the frozen glueball approximation

$$\begin{aligned} m_\omega^{*2} &= m_\omega^2 + 12g_4^4 \omega^2, \\ m_\rho^{*2} &= m_\rho^2 + 12g_4^4 \frac{Z_\rho}{Z_\omega} \omega^2, \\ m_\phi^{*2} &= m_\phi^2 + 24g_4^4 \frac{Z_\phi^2}{Z_\omega^2} \phi^2, \end{aligned} \quad (11)$$

with  $g_4 = \sqrt{Z_\omega} \tilde{g}_4$  as the renormalized coupling. Since the quartic vector-interaction contributes only in the medium, the coupling  $g_4$  cannot be unambiguously fixed. It is fitted, so that the compressibility lies in the desired region between 200 – 300 MeV in the mean field approximation. Note that the  $N - \omega$  as well as the  $N - \rho$  - couplings are also affected by the redefinition of the fields with the corresponding renormalised coupling constants as  $g_{N\omega} \equiv 3g_V^8 \sqrt{Z_\omega}$  and  $g_{N\rho} \equiv g_V^8 \sqrt{Z_\rho}$ .

## 2.4. Explicit chiral symmetry breaking

The explicit symmetry breaking term is given as [23]

$$\mathcal{L}_{SB} = \text{Tr} A_p (u(X + iY)u + u^\dagger(X - iY)u^\dagger) \quad (12)$$

with  $A_p = 1/\sqrt{2}\text{diag}(m_\pi^2 f_\pi, m_\pi^2 f_\pi, 2m_K^2 f_K - m_\pi^2 f_\pi)$  and  $m_\pi = 139$  MeV,  $m_K = 498$  MeV. This choice for  $A_p$  together with the constraints  $\sigma_0 = -f_\pi$ ,  $\zeta_0 = -\frac{1}{\sqrt{2}}(2f_K - f_\pi)$ , for the VEV on the scalar condensates assure that the PCAC-relations of the pion and kaon are fulfilled. With  $f_\pi = 93.3$  MeV and  $f_K = 122$  MeV we obtain  $|\sigma_0| = 93.3$  MeV and  $|\zeta_0| = 106.56$  MeV.

### 3. THE MEAN FIELD APPROXIMATION

The Lagrangian density in the mean field approximation [23] consists of the following terms

$$\mathcal{L}_{BX} + \mathcal{L}_{BV} = - \sum_i \bar{\psi}_i [g_{i\omega}\gamma_0\omega + g_{i\phi}\gamma_0\phi + m_i^*] \psi_i \quad (13)$$

$$\mathcal{L}_{vec} = \frac{1}{2}m_\omega^2 \frac{\chi^2}{\chi_0^2} \omega^2 + g_4^4 \omega^4 + \frac{1}{2}m_\phi^2 \frac{\chi^2}{\chi_0^2} \phi^2 + g_4^4 \left(\frac{Z_\phi}{Z_\omega}\right)^2 \phi^4 \quad (14)$$

$$\begin{aligned} \mathcal{V}_0 &= \frac{1}{2}k_0\chi^2(\sigma^2 + \zeta^2) - k_1(\sigma^2 + \zeta^2)^2 - k_2\left(\frac{\sigma^4}{2} + \zeta^4\right) - k_3\chi\sigma^2\zeta \\ &+ k_4\chi^4 + \frac{1}{4}\chi^4 \ln \frac{\chi^4}{\chi_0^4} - \frac{\delta}{3}\chi^4 \ln \frac{\sigma^2\zeta}{\sigma_0^2\zeta_0} \end{aligned} \quad (15)$$

$$\mathcal{V}_{SB} = \left(\frac{\chi}{\chi_0}\right)^2 \left[ m_\pi^2 f_\pi \sigma + (\sqrt{2}m_K^2 f_K - \frac{1}{\sqrt{2}}m_\pi^2 f_\pi)\zeta \right], \quad (16)$$

where  $m_i^* = -g_{\sigma i}\sigma - g_{\zeta i}\zeta$  is the effective mass of the baryon of type  $i$ , with  $i = N, \Sigma, \Lambda, \Xi$ . The thermodynamical potential of the grand canonical ensemble,  $\Omega$ , per unit volume  $V$  at given chemical potential  $\mu$  and temperature  $T$  can be written as

$$\begin{aligned} \frac{\Omega}{V} &= -\mathcal{L}_{vec} - \mathcal{L}_0 - \mathcal{L}_{SB} - \mathcal{V}_{vac} + \sum_i \frac{\gamma_i}{(2\pi)^3} \int d^3k E_i^*(k) \left( f_i(k) + \bar{f}_i(k) \right) \\ &- \sum_i \frac{\gamma_i}{(2\pi)^3} \mu_i^* \int d^3k \left( f_i(k) - \bar{f}_i(k) \right). \end{aligned} \quad (17)$$

Here the vacuum energy (the potential at  $\rho = 0$ ) has been subtracted in order to obtain a vanishing vacuum energy.  $\gamma_i$  is the spin-isospin degeneracy factor for baryon,  $i$  and  $E_i^*(k) = \sqrt{k_i^2 + m_i^{*2}}$  and  $\mu_i^* = \mu_i - g_{i\omega}\omega$  are the effective single particle energy and effective chemical potential respectively.  $f_i$  and  $\bar{f}_i$  are the thermal distribution functions for the baryons and antibaryons given as

$$f_i(k) = \frac{1}{e^{\beta(E_i^*(k) - \mu_i^*)} + 1}, \quad \bar{f}_i(k) = \frac{1}{e^{\beta(E_i^*(k) + \mu_i^*)} + 1}. \quad (18)$$

The mesonic field equations are determined by minimizing the thermodynamical potential

$$\left. \frac{\partial(\Omega/V)}{\partial\sigma} \right|_{MFT} = k_0\chi^2\sigma - 4k_1(\sigma^2 + \zeta^2)\sigma - 2k_2\sigma^3 - 2k_3\chi\sigma\zeta - 2\frac{\delta\chi^4}{3\sigma} +$$

$$+ m_\pi^2 f_\pi + \sum_i \frac{\partial m_i^*}{\partial \sigma} \rho_i^s = 0, \quad (19)$$

$$\begin{aligned} \left. \frac{\partial(\Omega/V)}{\partial \zeta} \right|_{MFT} &= k_0 \chi^2 \zeta - 4k_1(\sigma^2 + \zeta^2)\zeta - 4k_2 \zeta^3 - k_3 \chi \sigma^2 - \frac{\delta \chi^4}{3\zeta} + \\ &+ \left[ \sqrt{2} m_K^2 f_K - \frac{1}{\sqrt{2}} m_\pi^2 f_\pi \right] + \sum_i \frac{\partial m_i^*}{\partial \zeta} \rho_i^s = 0, \end{aligned} \quad (20)$$

$$\left. \frac{\partial(\Omega/V)}{\partial \omega} \right|_{MFT} = -m_\omega^2 \omega - 4g_4^4 \omega^3 + \sum_i g_{i\omega} \rho_i = 0, \quad (21)$$

$$\left. \frac{\partial(\Omega/V)}{\partial \phi} \right|_{MFT} = -m_\phi^2 \phi - 8g_4^4 \left( \frac{Z_\phi}{Z_\omega} \right)^2 \phi^3 + \sum_i g_{i\phi} \rho_i = 0. \quad (22)$$

In the above,  $\rho_i^s$  and  $\rho_i$  are the scalar and vector densities for the baryons at finite temperature given as,

$$\rho_i^s = \gamma_i \int \frac{d^3 k}{(2\pi)^3} \frac{m_i^*}{E_i^*} (f_i(k) + \bar{f}_i(k)), \quad \rho_i = \gamma_i \int \frac{d^3 k}{(2\pi)^3} (f_i(k) - \bar{f}_i(k)). \quad (23)$$

The energy density and the pressure follow from the Gibbs-Duhem relation,  $\epsilon = \Omega/V + \mu_i \rho_i + TS$  and  $p = -\Omega/V$ .

#### 4. THE RELATIVISTIC HARTREE APPROXIMATION

The relativistic Hartree approximation takes into account the effects from the Dirac sea through evaluating the baryonic tadpole diagrams. The interacting propagator for baryon of type  $i$  is given by the Schwinger-Dyson equation

$$G_i^H(p) = G_i^0(p) + G_i^0(p) \Sigma_i G_i^H(p), \quad (24)$$

with  $G_i^0(p)$  as the free propagator and  $\Sigma_i(p)$  as the self-energy consisting of the scalar and vector parts as

$$\Sigma_i = \Sigma_i^S - \gamma^\mu \Sigma_{i\mu}^V. \quad (25)$$

The formal solution of the Schwinger-Dyson equation is

$$[G_i^H(p)]^{-1} = \gamma \cdot \bar{p} - m_i^* \quad (26)$$

or equivalently,

$$\begin{aligned} G_i^H(p) &= (\gamma^\mu \bar{p}_\mu + m_i^*) \left\{ \frac{1}{\bar{p}^2 - m_i^{*2} + i\epsilon} \right. \\ &\quad \left. + \frac{\pi i}{E_i^*(p)} \left[ \frac{\delta(\bar{p}^0 - E_i^*(p))}{e^{\beta(E_i^*(p) - \mu_i^*)} + 1} + \frac{\delta(\bar{p}^0 + E_i^*(p))}{e^{\beta(E_i^*(p) + \mu_i^*)} + 1} \right] \right\} \\ &\equiv (G_i^H)^F(p) + (G_i^H)^D(p) \end{aligned} \quad (27)$$



where  $E_i^*(p) = \sqrt{\mathbf{p}^2 + m_i^{*2}}$ ,  $\bar{p} = p + \Sigma_i^V$  and  $m_i^* = m_i + \Sigma_i^S$ .

In the present investigation for the study of hot hyperonic matter, the baryons couple to both the non-strange ( $\sigma$ ) and strange ( $\zeta$ ) scalar fields, so that we have the scalar self energy as

$$\Sigma_i^S = -(g_{\sigma i}\tilde{\sigma} + g_{\zeta i}\tilde{\zeta}), \quad (28)$$

where  $\tilde{\sigma} = \sigma - \sigma_0$ ,  $\tilde{\zeta} = \zeta - \zeta_0$ . The scalar self-energy  $\Sigma_i^S$  can also be written as

$$\begin{aligned} \Sigma_i^S &= i \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) \int \frac{d^4\mathbf{p}}{(2\pi)^4} \text{Tr}[G_i^F(p) + G_i^D(p)] e^{ip^0\eta} \\ &\equiv (\Sigma_i^S)^F + (\Sigma_i^S)^D. \end{aligned} \quad (29)$$

$(\Sigma_i^S)^D$  is the density dependent part and is identical to the mean field contribution

$$\begin{aligned} (\Sigma_i^S)^D &= - \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) \times \gamma_i \int \frac{d^3\mathbf{p}}{(2\pi)^3} \frac{m_i^*}{E_i^*(p)} (f_i(p) + \bar{f}_i(p)) \\ &= - \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) \rho_i^S. \end{aligned} \quad (30)$$

The Feynman part  $(\Sigma_i^S)^F$  of the scalar part of the self-energy is divergent. We adopt a dimensional regularization scheme to extract the convergent part by performing the integration in  $n$  dimensions. Finally one takes the limit  $n \rightarrow 4$  to extract the divergence which is rendered finite by adding the appropriate counter terms. After regularization, we finally get

$$(\Sigma_i^S)^F = \frac{\gamma_i}{8\pi^2} \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) [m_i^{*3}\Gamma(1 - n/2) + 2m_i^{*3} \ln m_i^*], \quad (31)$$

with  $m_i^* = m_i - g_{\sigma i}\tilde{\sigma} - g_{\zeta i}\tilde{\zeta}$ . We renormalize the Feynman part of the self energy, as given by the first term on the rhs. of (29) by adding the counter term

$$(\Sigma_i^S)_{CTC} = - \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) \sum_{n=0}^3 \frac{1}{n!} (g_{\sigma i}\tilde{\sigma} + g_{\zeta i}\tilde{\zeta})^n \beta_{n+1}^i. \quad (32)$$

The coefficients  $\beta_n^i$ 's are fixed from a set of renormalization conditions. The first term in (32) ensures that the tadpole contribution vanishes in vacuum. The second term fixes the masses for the scalar fields  $\sigma$  and  $\zeta$  fields to their vacuum values, in addition to ensuring that there are no mixed terms (of the type  $\sigma\zeta$ ) introduced due to such RHA contributions. The last two terms in (32) are chosen so that there are no cubic or quartic interaction contributions in the scalar meson fields arising from contributions in RHA in vacuum. Explicitly, the coefficients  $\beta_n^i$ 's are given as

$$\beta_1^i = \frac{\gamma_i}{8\pi^2} [m_i^3\Gamma(1 - n/2) + 2m_i^3 \ln m_i^*],$$

$$\begin{aligned}
\beta_2^i &= -\frac{\gamma_i}{8\pi^2}[3m_i^2\Gamma(1-n/2) + 2m_i^2(1 + 3\ln m_i)], \\
\beta_3^i &= \frac{\gamma_i}{8\pi^2}[6m_i\Gamma(1-n/2) + 10m_i + 12m_i\ln m_i], \\
\beta_4^i &= -\frac{\gamma_i}{8\pi^2}[6\Gamma(1-n/2) + 22 + 12\ln m_i].
\end{aligned} \tag{33}$$

This yields additional contributions from the Dirac sea to the baryon self energy, as

$$\begin{aligned}
(\Sigma_i^S)^{F finite} &= (\Sigma_i^S)^F + (\Sigma_i^S)_{CTC} = \frac{\gamma_i}{4\pi^2} \left( \frac{g_{\sigma i}^2}{m_\sigma^2} + \frac{g_{\zeta i}^2}{m_\zeta^2} \right) \\
&\times \left[ m_i^{*3} \ln \left( \frac{m_i^*}{m_i} \right) + m_i^2(m_i - m_i^*) - \frac{5}{2}m_i(m_i - m_i^*)^2 + \frac{11}{6}(m_i - m_i^*)^3 \right].
\end{aligned} \tag{34}$$

The above corresponds to the following counter terms in the Lagrangian as

$$\mathcal{L}_{CTC} = \sum_i \sum_{n=1}^4 \frac{\beta_n^i}{n!} (g_{\sigma i} \tilde{\sigma} + g_{\zeta i} \tilde{\zeta})^n, \tag{35}$$

where the sum extends over all the baryon species  $i$ . The energy density can then be evaluated to be

$$\epsilon_{RHA} = \epsilon_{MFT} + \Delta\epsilon, \tag{36}$$

with the contribution to the energy density from the Dirac sea as

$$\begin{aligned}
\Delta\epsilon &= -\sum_i \frac{\gamma_i}{16\pi^2} \left[ m_i^{*4} \ln \left( \frac{m_i^*}{m_i} \right) + m_i^3(m_i - m_i^*) - \frac{7}{2}m_i^2(m_i - m_i^*)^2 \right. \\
&\quad \left. + \frac{13}{3}m_i(m_i - m_i^*)^3 - \frac{25}{12}(m_i - m_i^*)^4 \right].
\end{aligned} \tag{37}$$

The field equations for the scalar meson fields are modified to

$$\left. \frac{\partial(\Omega/V)}{\partial\Phi} \right|_{RHA} = \left. \frac{\partial(\Omega/V)}{\partial\Phi} \right|_{MFT} + \sum_i \frac{\partial m_i^*}{\partial\Phi} \Delta\rho_i^s = 0 \quad \text{with } \Phi = \sigma, \zeta, \tag{38}$$

where the additional contribution to the nucleon scalar density is given as

$$\Delta\rho_i^s = -\frac{\gamma_i}{4\pi^2} \left[ m_i^{*3} \ln \left( \frac{m_i^*}{m_i} \right) + m_i^2(m_i - m_i^*) - \frac{5}{2}m_i(m_i - m_i^*)^2 + \frac{11}{6}(m_i - m_i^*)^3 \right]. \tag{39}$$

These make a refitting of some of the parameters necessary. First we have to account for the change in the energy and the pressure, i.e.  $g_{N\omega}$  and  $\chi_0$  have to be refitted. Due to a change in  $\chi_0$  the parameters  $k_0, k_2$  and  $k_4$  must be adapted to ensure that the vacuum equations for  $\sigma, \zeta$  and  $\chi$  have minima at the vacuum expectation values of the fields. Table I shows the parameters corresponding to the Mean-field and the Hartree approximations.

Parameter	Mean Field		Hartree	
$g_4$	2.7	0	2.7	0
$k_1$	1.4	1.4	1.4	1.4
$g_{N\omega}$	13.24	10.74	10.61	9.45
$g_{N\rho}$	5.04	4.1	4.04	3.6
$\chi_0$	405.7	435.8	437.4	454.2
$k_3$	-2.57	-1.95	-1.91	-1.56
$k_0$	2.33	2.02	2	1.86
$k_2$	-5.55	-5.55	-5.55	-5.55
$k_4$	-0.23	-0.23	-0.23	-0.24
$m_N^*/m_N(\rho_0)$	0.62	0.7	0.73	0.76
$m_\sigma$	485	578	583	637

TABLE I: Parameters for the Mean-Field and the Hartree Fit

## 5. VECTOR MESON MASSES IN THE MEDIUM

### 5.1. Mass modifications of $\omega$ and $\rho$ -mesons

We now examine how the Dirac sea effects discussed in section 4 modify the masses of the vector mesons ( $\omega$  and  $\rho$ ) due to their interaction with the in-medium nucleons. From (5), the interaction can be written as

$$\mathcal{L}_{NV} = g_{N\omega}\omega_\mu\bar{\psi}_N\gamma^\mu\psi_N + g_{N\rho}\vec{\rho}_\mu\bar{\psi}_N\gamma^\mu\vec{\tau}\psi_N, \quad (40)$$

in terms of the renormalized couplings  $g_{N\omega}$  and  $g_{N\rho}$ . Furthermore a tensor coupling is introduced:

$$\mathcal{L}_{\text{tensor}} = -\frac{g_{NV}\kappa_V}{2m_N} [\bar{\psi}_N\sigma_{\mu\nu}\tau^a\psi_N\partial^\nu V_a^\mu], \quad (41)$$

where  $(g_{NV}, \kappa_V) = (g_{N\omega}, \kappa_\omega)$  or  $(g_{N\rho}, \kappa_\rho)$  for  $V_a^\mu = \omega^\mu$  or  $\rho_a^\mu$ ,  $\tau_a = 1$  or  $\vec{\tau}$ ,  $\vec{\tau}$  being the Pauli matrices. The vector meson self energy is expressed in terms of the full nucleon propagator (27) and is given by

$$\Pi_V^{\mu\nu}(k) = -\gamma_I g_{NV}^2 \frac{i}{(2\pi)^4} \int d^4p \text{Tr} \left[ \Gamma_V^\mu(k) G^H(p) \Gamma_V^\nu(-k) G^H(p+k) \right], \quad (42)$$

where  $\gamma_I = 2$  is the isospin degeneracy factor for nuclear matter, and,  $\Gamma_V^\mu(k) = \gamma^\mu\tau_a - (\kappa_V/2m_N)\sigma^{\mu\nu}\tau_a$  represents the meson-nucleon vertex function obtained from (40) and (41).

For the  $\omega$  meson, the tensor coupling is generally small in comparison to the vector coupling to the nucleons [18] and is neglected in the present calculations.

We then have the meson-nucleon vertex functions as

$$\begin{aligned}\Gamma^\mu(k) &= \gamma_\mu, & \text{for } \omega, \\ \Gamma^\mu(k) &= \gamma_\mu + \frac{i\kappa_\rho}{2m_N}\sigma^{\mu\alpha}k_\alpha & \text{for } \rho.\end{aligned}\quad (43)$$

After insertion of  $G^H(p)$ , the vector meson self energy separates into

$$\Pi^{\mu\nu}(k) = \Pi_F^{\mu\nu}(k) + \Pi_D^{\mu\nu}(k), \quad (44)$$

where  $\Pi_F^{\mu\nu}(k)$  is the vacuum polarization and describes the correction to the meson propagator due to coupling to baryon-antibaryon excitations of the Dirac sea and  $\Pi_D^{\mu\nu}(k)$  describes coupling to the particle-hole excitations of the Fermi sea. The Feynman part of the self energy,  $\Pi_F^{\mu\nu}(k)$ , is divergent and has to be renormalized. After dimensional regularization to separate the divergent part and using the subtraction procedure described in [17, 18, 34], we obtain the following  $\Pi_F^{\mu\nu}(k)$  for the  $\omega$  and  $\rho$  mesons

$$\Pi_F^\omega(k^2) \equiv \frac{1}{3}\text{Re}(\Pi_F^{\text{ren}})^\mu{}_\mu = -\frac{g_{N\omega}^2}{\pi^2}k^2 I_1 \quad (45)$$

$$\Pi_F^\rho(k^2) = -\frac{g_{N\rho}^2}{\pi^2}k^2 \left[ I_1 + \frac{m_N^* \kappa_\rho}{2m_N} I_2 + \frac{1}{2} \left( \frac{\kappa_\rho}{2m_N} \right)^2 (k^2 I_1 + m_N^* I_2) \right], \quad (46)$$

where

$$I_1 = \int_0^1 dz z(1-z) \ln \left[ \frac{m_N^{*2} - k^2 z(1-z)}{m_N^2 - k^2 z(1-z)} \right], \quad (47)$$

$$I_2 = \int_0^1 dz \ln \left[ \frac{m_N^{*2} - k^2 z(1-z)}{m_N^2 - k^2 z(1-z)} \right]. \quad (48)$$

The renormalization condition which has been used to obtain the  $\omega$ -self energy is  $\Pi_F^\omega(k^2)(m_N^* \rightarrow m_N) = 0$  ensuring the vanishing of the vector self energy in vacuum. Due to the tensor interaction in (43), the vacuum self energy for the  $\rho$ -meson is not renormalizable. We have employed a phenomenological subtraction procedure [18] to extract the finite part using the condition  $\partial^n \Pi_F^\rho(k^2)/\partial(k^2)^n|_{m_N^* \rightarrow m_N} = 0$ , with  $n = 0, 1, 2, \dots, \infty$ .

The contribution from the Fermi sea is given as

$$\Pi^D(k_0, \mathbf{k} \rightarrow 0) = -\frac{4g_{NV}^2}{\pi^2} \int p^2 dp F(|\mathbf{p}|, m_N^*) [f_{FD}(\mu^*, T) + \bar{f}_{FD}(\mu^*, T)], \quad (49)$$

with

$$F(|\mathbf{p}|, m_N^*) = \frac{1}{\epsilon^*(p)(4\epsilon^*(p)^2 - k_0^2)} \left[ \frac{2}{3}(2|\mathbf{p}|^2 + 3m_N^{*2}) + k_0^2 \left\{ 2m_N^* \left( \frac{\kappa_V}{2m_N} \right) + \frac{2}{3} \left( \frac{\kappa_V}{2m_N} \right)^2 (|\mathbf{p}|^2 + 3m_N^{*2}) \right\} \right], \quad (50)$$

where  $\epsilon^*(p) = (\mathbf{p}^2 + m_N^{*2})^{1/2}$  is the effective energy of the nucleon. The effective mass of the vector meson at rest is then obtained by solving the equation, with  $\Pi = \Pi_F + \Pi_D$ ,

$$k_0^2 - m_V^2 + \text{Re}\Pi(k_0, \mathbf{k} = 0) = 0. \quad (51)$$

## 5.2. Mass modification of $\phi$ - meson

The  $\phi$ -meson, which does not couple to the nucleons, is modified due to the hyperon-antihyperon excitations in the relativistic Hartree approximation. The mass of the  $\phi$ -meson in the medium is obtained as a solution of the dispersion relation given as [26],

$$k_0^2 - m_V^2 + \sum_i \text{Re}\Pi_i(k_0, \vec{k} = 0) = 0. \quad (52)$$

Using the hyperon- $\phi$  couplings as given in (6) the mass of the  $\phi$ -meson is calculated.

## 6. RESULTS AND DISCUSSIONS

In this section, we discuss the findings of the present investigation for the properties of the hadrons properties in a chiral SU(3) model. We investigate how the hadron properties in mean field approximation are modified due to the effect of the vacuum polarizations in relativistic Hartree approximation. Figure 1 shows the temperature and density dependence of the pressure. The effect from the Dirac sea of the baryons through the relativistic Hartree approximation is seen to lead to a softening of the equation of state for the hot hyperonic matter. In figures 2 and 3 the baryon masses are shown as functions of temperature for zero baryon density. Accounting for the baryonic Dirac sea effects is seen to give rise to higher values for the baryon masses in the medium. This is related to the fact that these contributions lead to a softening of the equation of state as illustrated in the figure 1. However, note that the masses stay almost constant up to a temperature of around 160 MeV above which there is a drop of the masses in the hot matter. The modification of the hyperon masses is smaller as compared to the nucleon mass because of their stronger coupling to the strange condensate  $\zeta$ , which shows much weaker temperature dependence than the non-strange condensate  $\sigma$  as illustrated in 4. Especially, the  $\Xi$  mass is seen to stay

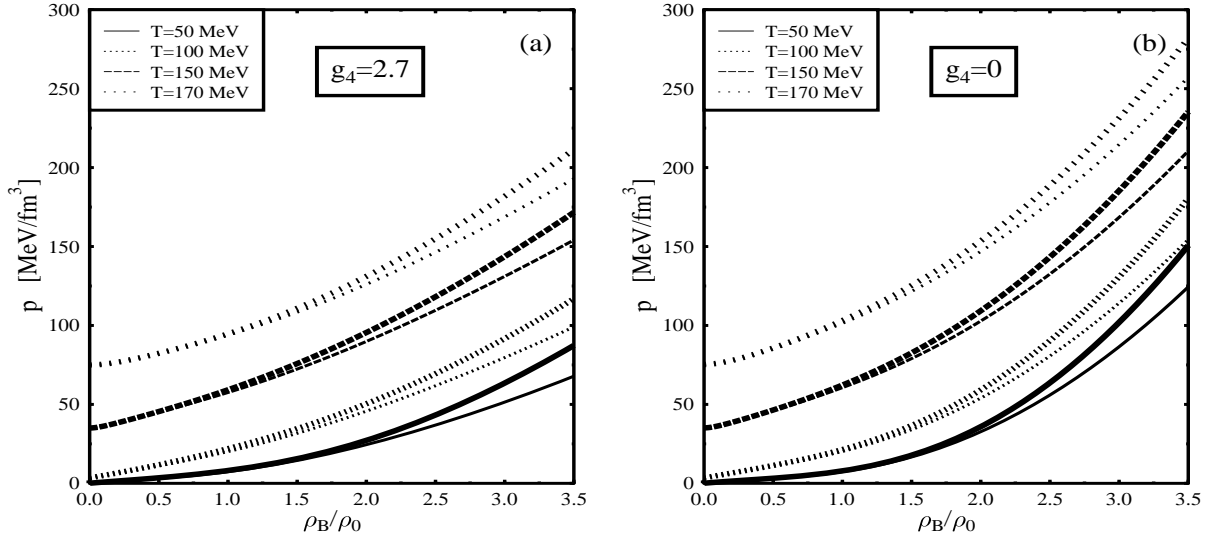


FIG. 1: The equation of state in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  for the vector self-interaction coupling (a)  $g_4=2.7$  and (b)  $g_4=0$ .

almost constant even up to a temperature of about 180 MeV. The presence of quartic self interaction for the vector field enhances the mass modification as seen in figures 2 and 3. This is because the vector field strength is attenuated leading to a larger effective chemical potential and hence a larger thermal distribution function for the baryon. As a consequence the contribution to baryon scalar self energy from the medium dependent part as given by (30) becomes larger. This explains the smaller baryon mass for  $g_4 \neq 0$  in the mean field approximation. The qualitative features remain the same with additional contributions from  $N\bar{N}$  fluctuations in RHA. Figures 5 and 6 illustrate the temperature and density dependence of the non-strange ( $\sigma$ ) and strange ( $\zeta$ ) scalar fields. The effective baryon masses as functions of density, for different temperatures and zero net strangeness are shown in figures 7 -10. Again, higher baryon masses are predicted if we take the quantum corrections of the Dirac sea into account. This behaviour mirrors itself in the density dependence of the scalar fields (figures 5-6).

Note that the masses of the baryons at finite densities first increases with temperature up to around 170 MeV and then decreases. Such a behaviour of the nucleon mass increasing with temperature was also observed earlier within the framework of the Walecka model by Ko and Li [35] in a mean field calculation. This behaviour of the baryon self energy, given by (30) in the mean field approximation can be understood in the following manner. The

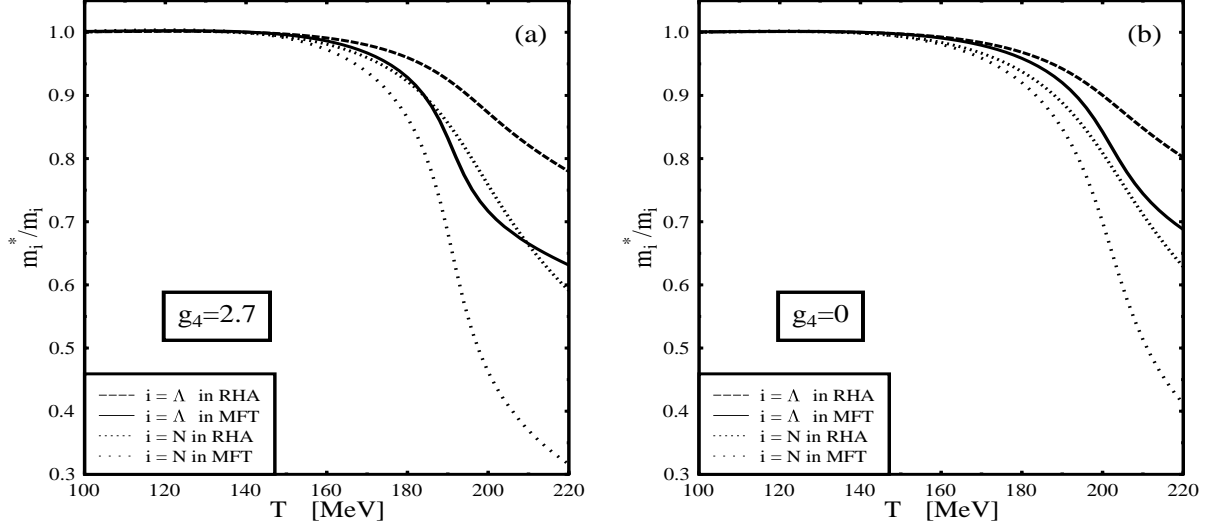


FIG. 2: Effective nucleon and  $\Lambda$  masses as functions of temperature in the mean field (MFT) and in RHA for  $f_s = 0$  and  $\rho_B = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

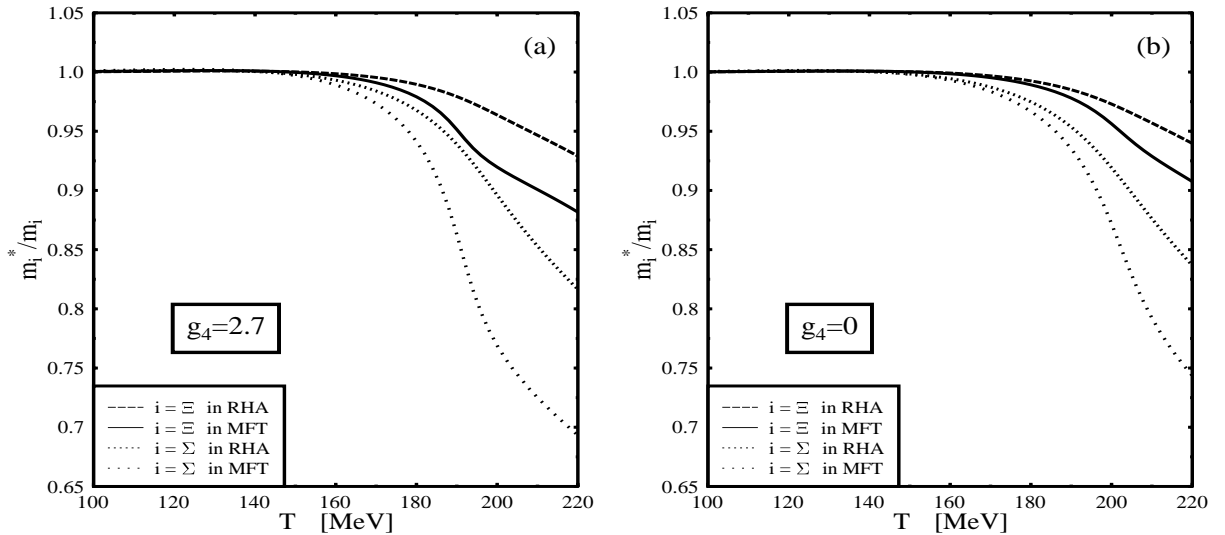


FIG. 3: Effective  $\Sigma$  and  $\Xi$  masses as functions of temperature in the mean field (MFT) and in RHA for  $f_s = 0$  and  $\rho_B = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

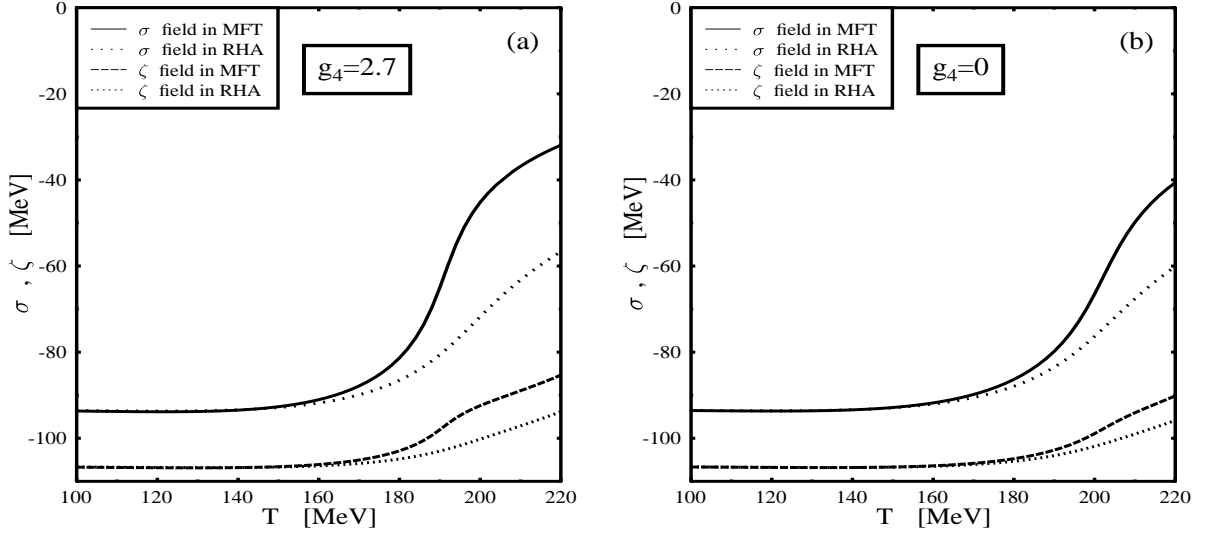


FIG. 4: The scalar fields as functions of temperature in the mean field (MFT) and in RHA for  $f_s = 0$  and  $\rho_B = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

thermal distribution functions have the effect of increasing the self energy as given by (30) (hence decreasing the masses). However, at finite densities, for increasing temperatures, there are contributions also from higher momenta thereby increasing the denominator of the integrand in the rhs of (30) (and hence increasing the value for the effective baryon mass). The competing effects arising from the thermal distribution functions and the contributions from higher-momenta states give rise to the observed behaviour of the effective baryon masses with temperature at finite densities. For temperatures of about 170 MeV one can see that the scalar fields have nonzero fluctuations from the vacuum values even at zero density, which indicates the existence of baryon-antibaryon pairs in the thermal bath. This behaviour for nuclear matter at finite temperatures has also been known in the literature [36, 37]. This leads to the baryon masses above this temperature to be different from the vacuum values at zero baryon density.

The vector meson masses ( $\omega$  and  $\rho$ ) as modified by the interaction to the nucleons in the thermal medium are plotted in figures 11 and 12. The effective vector meson masses arising due to nucleon-antinucleon loop in the relativistic Hartree approximation are compared to the mean field case. For the nucleon-rho couplings, the vector and tensor couplings as obtained from the N-N forward dispersion relation [18, 20, 38] are used. The medium modified vector meson masses are plotted in figures 11 and 12 with and without the quartic



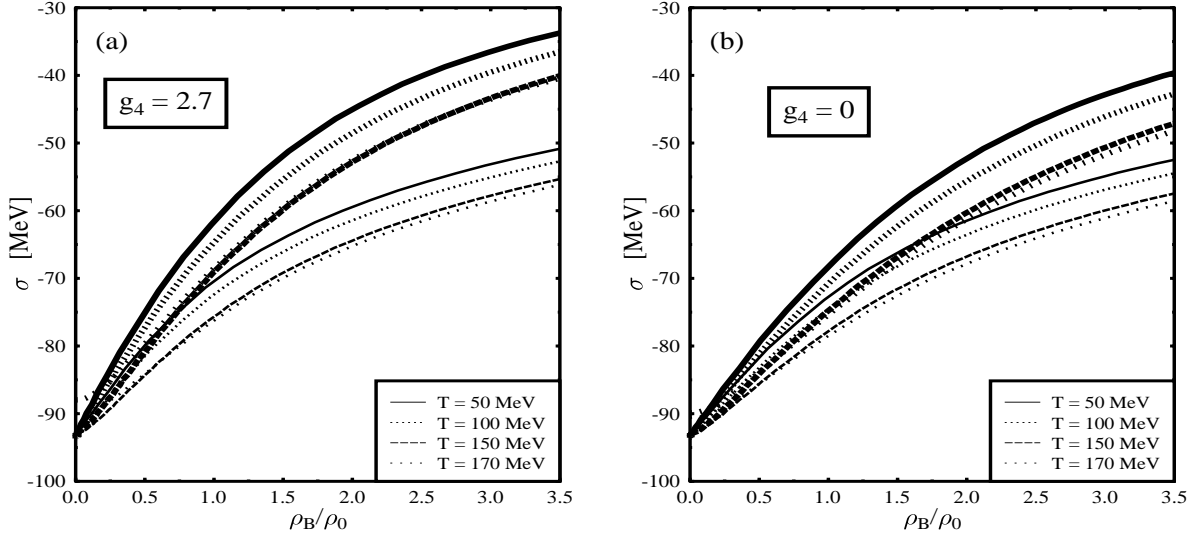


FIG. 5: The non-strange condensate as a function of density in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

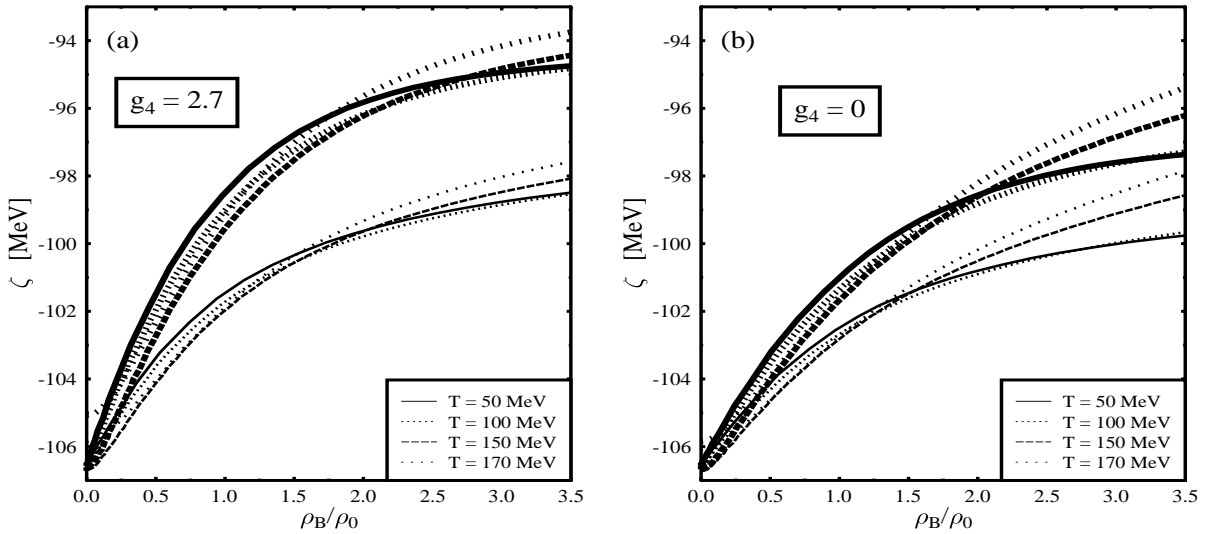


FIG. 6: The strange condensate as a function of density in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

interaction for the vector fields. The increase of the nucleon masses with temperature at finite densities is reflected as an increase in the vector meson masses as was also seen in Ref. [35]. If we switch off the quartic self interaction, the vector meson masses have no density and temperature dependence in MFT as seen from (11). In RHA, a significant reduction of the  $\omega$  and  $\rho$  masses due to the Dirac sea polarization is found up to around nuclear saturation density. At higher densities, the density dependent part of the vector meson self energy, describing the interaction with the Fermi sea fluctuations starts to be more dominating, leading to increasing masses instead. But in the case of the  $\omega$  mass, above 100 MeV the effect of the Fermi sea polarization seems not to be sufficient in order to overturn the original decreasing tendency. For  $g_4 \neq 0$  the vector meson masses increase monotonically with density in the mean field case. The mean field value of the  $\omega$  field increases with density and consequently due to (11) the vector meson masses increase.

In RHA, the masses drop up to  $\rho_0$  but a finite value for  $g_4$  leads to a modified high density behaviour. Above,  $\rho_0$ , the masses increase with density at all temperatures. It is seen that the density dependence dominates over the temperature dependence.

The medium modification for the vector meson  $\phi$  in the hyperonic matter is plotted in figure 13. One observes that the strange meson  $\phi$  has smaller mass modifications compared to the  $\omega$  and  $\rho$  mesons. This is due to the fact that  $\phi$  meson does not couple to the nucleons and also, the hyperon masses are rather insensitive to the changes in the baryon densities as may be seen in figures 8 - 10. One might note here that unlike the  $\omega$  and  $\rho$  vector mesons, where the presence of the quartic self interaction for the vector fields attenuates the drop of the meson masses in relativistic Hartree approximation, the  $\phi$ -mass has a larger drop when  $g_4 \neq 0$ . The reason for this is that the nucleon- $\omega$  coupling is larger in the presence of this interaction as can be seen from table I. Hence the hyperon- $\phi$  couplings, which are related to  $g_{N\omega}$  through the relation (6) are higher in value. This gives rise to a larger drop of the  $\phi$  mass due to the hyperon-antihyperon excitations in relativistic Hartree approximation. The nonzero  $g_4$  gives rise to an increasing contribution to the  $\phi$  mass as proportional to the quadratic strength of the  $\phi$ -field unlike the case of  $\omega$  and  $\rho$  mesons where this is proportional to the strength of the  $\omega$ -field. Since the strength of the  $\phi$  field is much smaller than that of the  $\omega$ , this increase in the  $\phi$ -mass remains small compared to the drop due to the hyperon sea. The strange meson ( $\phi$ ) mass modification observed as small compared to the  $\omega$  and  $\rho$  meson masses is in line with the earlier observations [8, 18, 39].

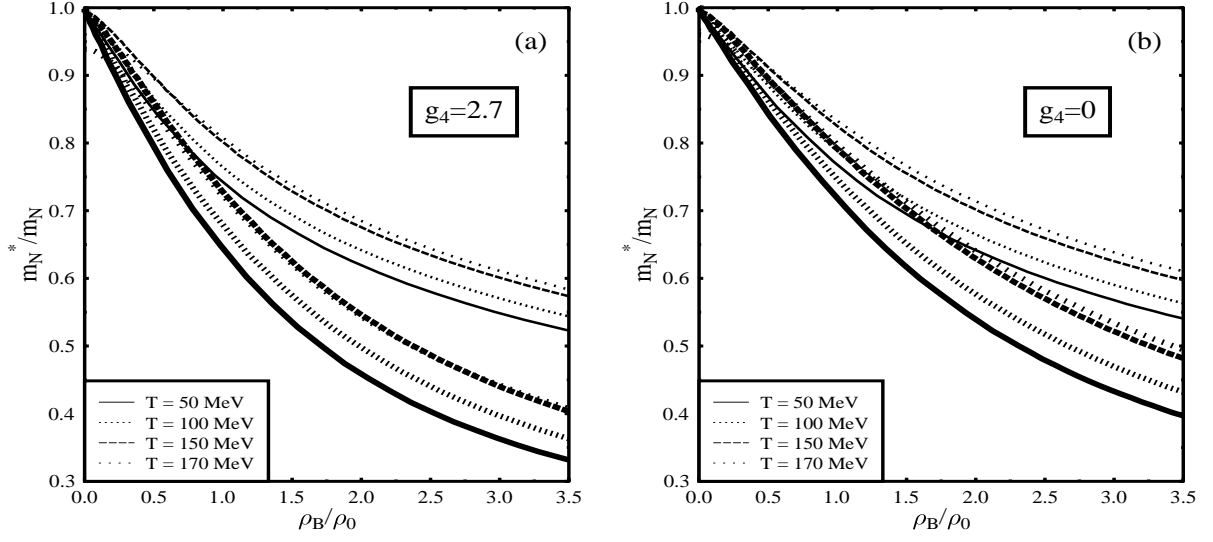


FIG. 7: Effective nucleon mass as a function of density in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

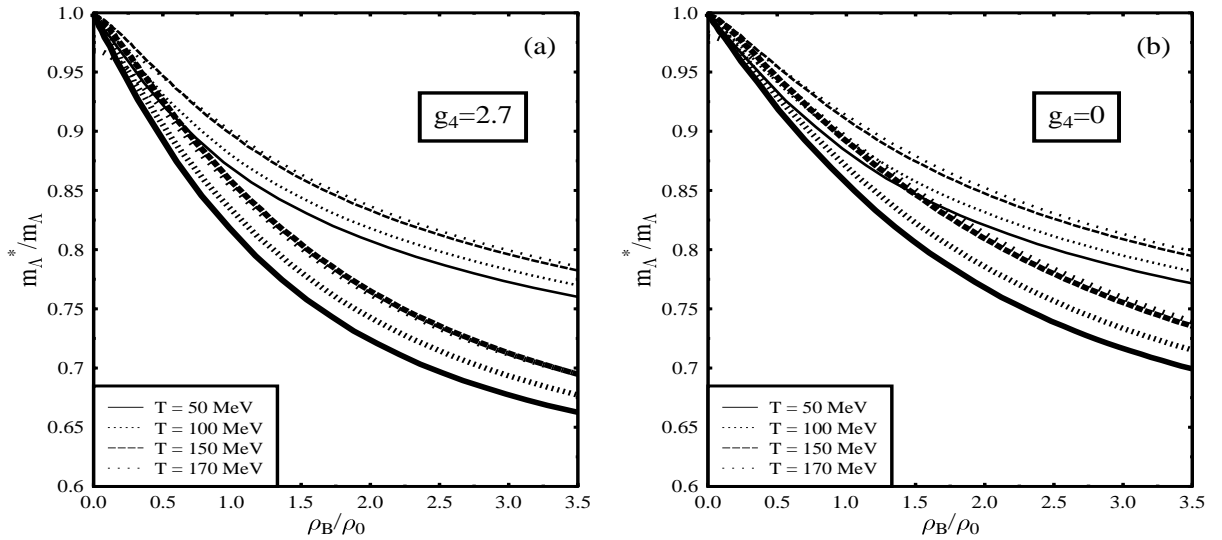


FIG. 8: Effective  $\Lambda$  mass as a function of density in the mean field (MFT) and in RHA for different temperatures and  $f_s = 0$  for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

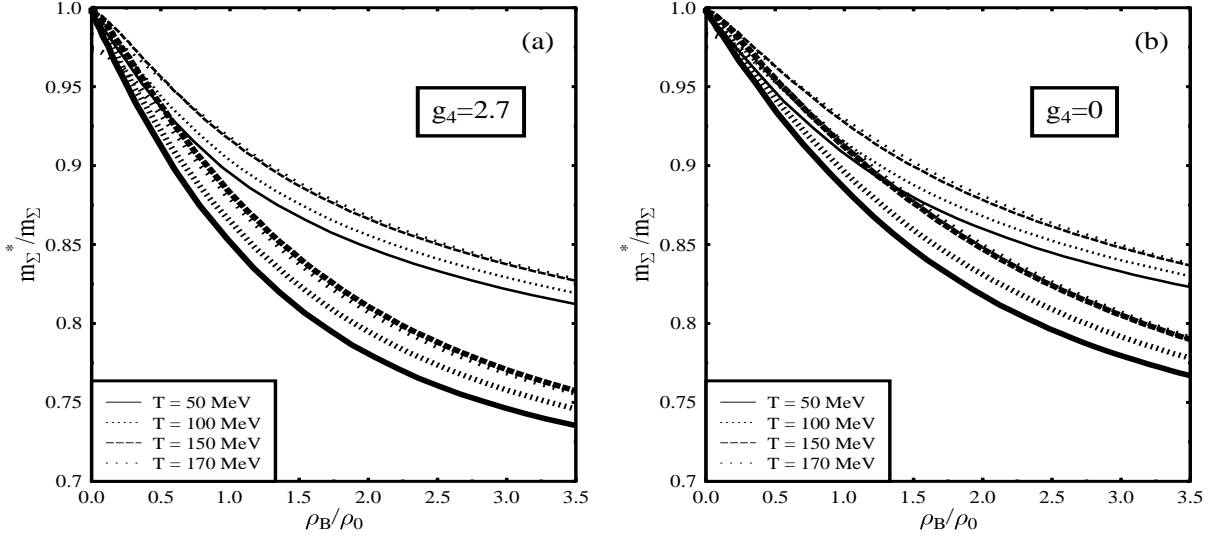


FIG. 9: Effective  $\Sigma$  mass as a function of density in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

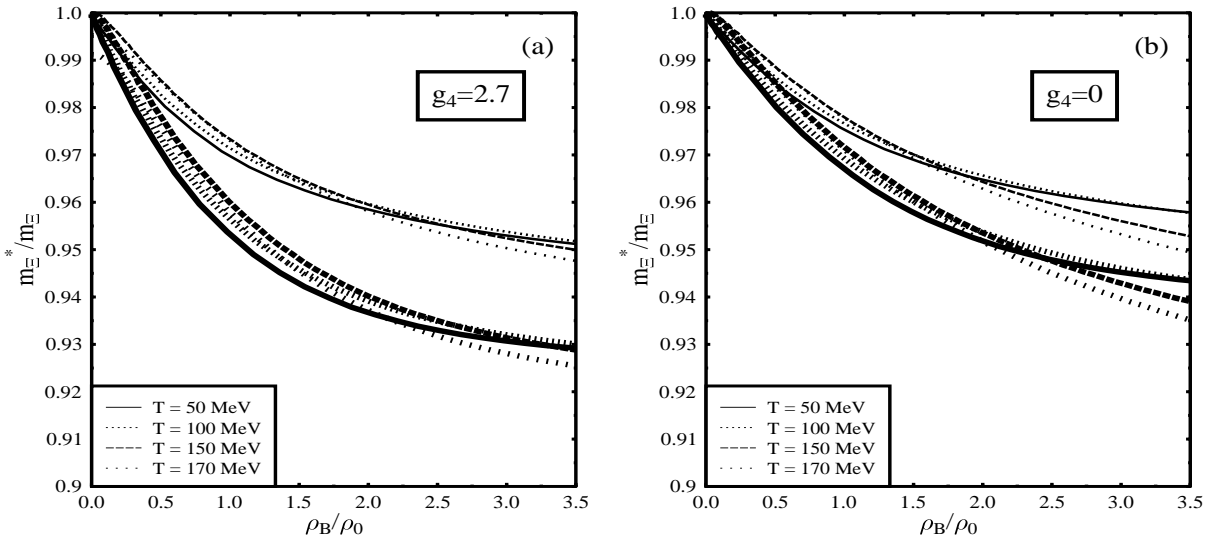


FIG. 10: Effective  $\Xi$  masses as a function of density in the mean field (MFT) (thick lines) and in RHA (thin lines) for different temperatures and  $f_s = 0$  for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

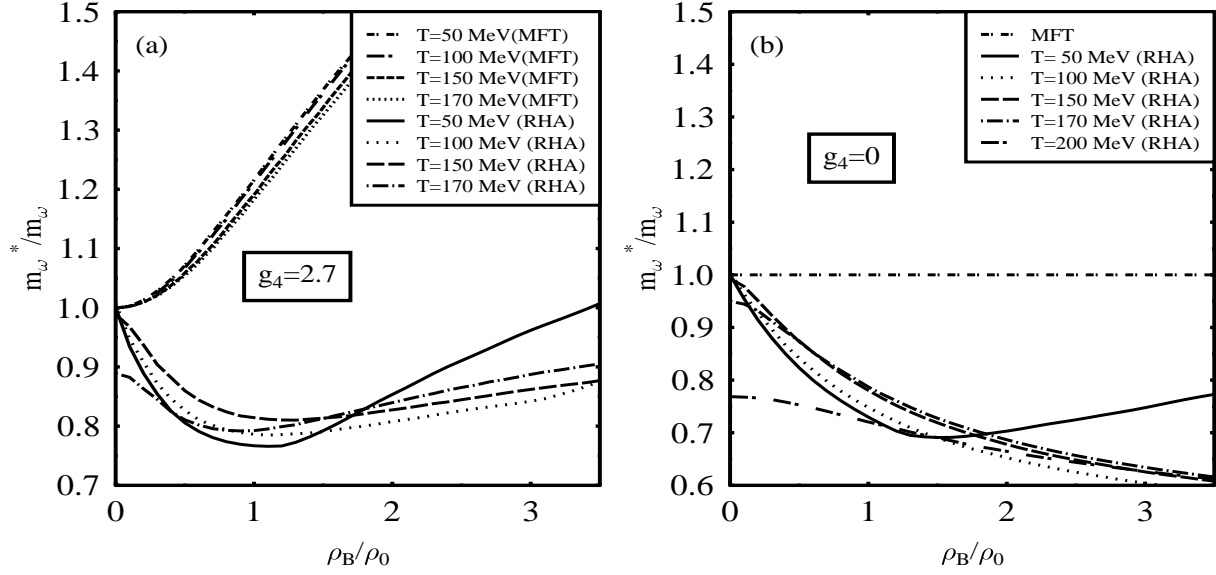


FIG. 11: Effective  $\omega$  mass as a function of density in the mean field (MFT) and in RHA for different temperatures and  $f_s = 0$  and for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

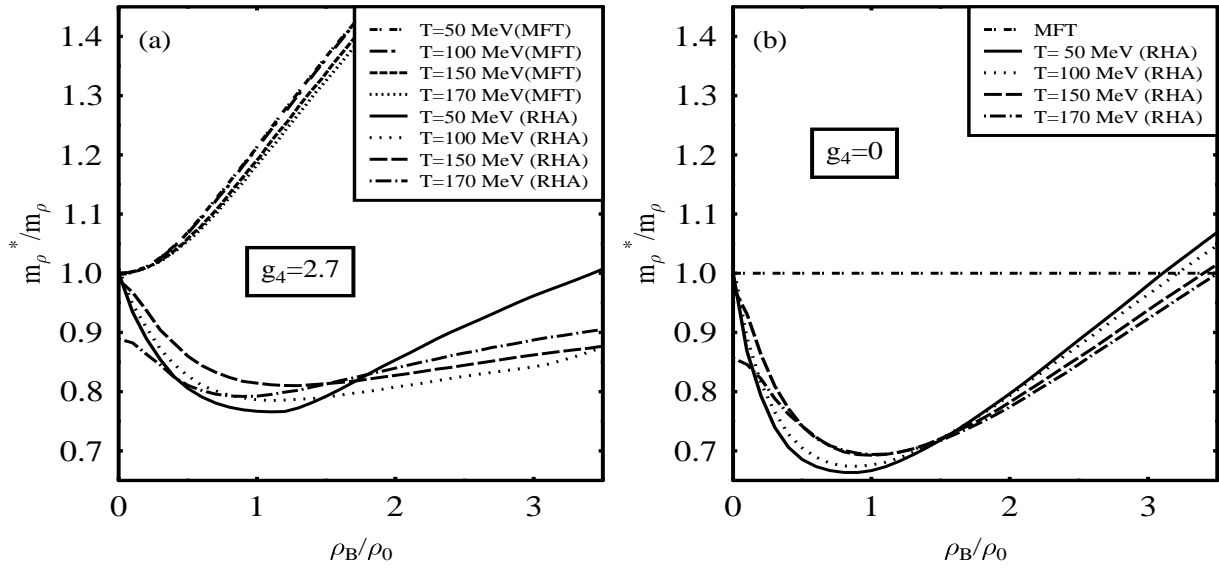


FIG. 12: Effective  $\rho$  mass as a function of density in the mean field (MFT) and in RHA for different temperatures and  $f_s = 0$  for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

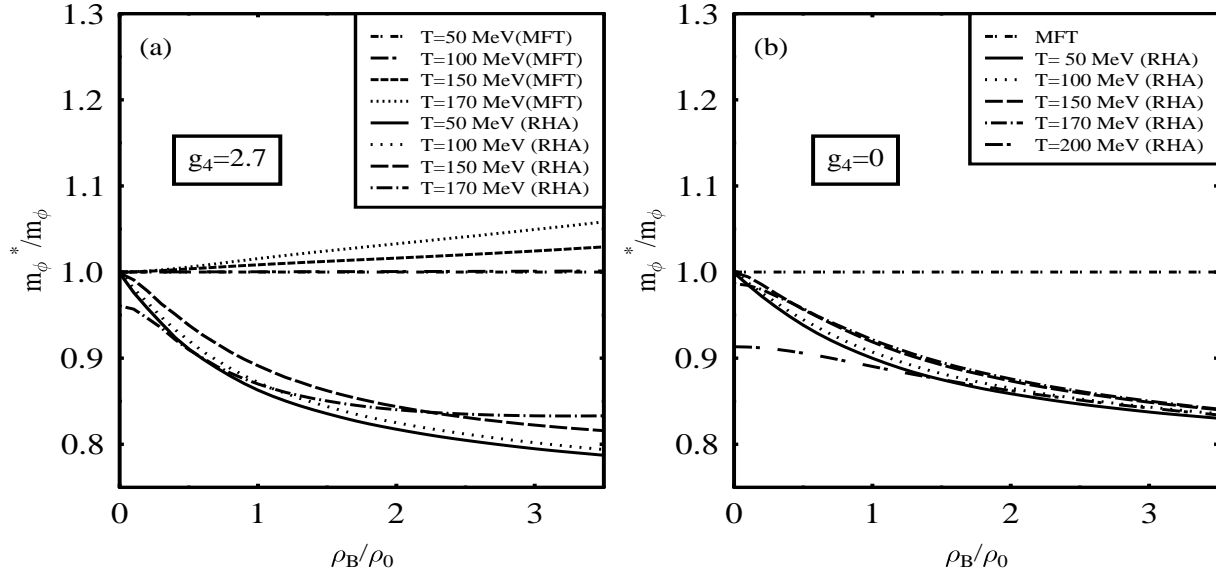


FIG. 13: Effective  $\phi$  meson mass in the mean field approximation and including the Hartree contributions for (a)  $g_4=2.7$  and (b)  $g_4=0$ .

## 7. SUMMARY

To summarize, in the present work, we have considered a chiral SU(3) model for the description of the hot and strange hadronic matter. The effect of the baryonic vacuum polarizations has been taken into account in the relativistic Hartree approximation for study of the hadronic properties. The coupling of the baryons to both nonstrange and strange scalar fields modifies the scalar self energy and is taken into account while summing over the baryonic tadpole diagrams in the relativistic Hartree approximation. The vector meson ( $\omega$  and  $\rho$ ) masses are calculated in the thermal medium arising from the nucleon-antinucleon loop and are seen to have large drops due to the Dirac sea contribution. However, the strange vector meson,  $\phi$  mass is seen to be much less modified due to the baryon Dirac sea, as compared to the nonstrange vector mesons. The vector meson properties are directly linked to the dilepton spectra in relativistic heavy ion collision experiments. It will thus be worth investigating how the dilepton spectra are modified due to the medium modification of the vector mesons. The hadron properties in the medium also modify the number densities and hence would modify the particle ratios observed in relativistic heavy ion collision experiments. These have been studied in the mean field approximation in the present chiral SU(3) model [40] and it will be interesting to examine the effects of vacuum polarisations

on the particle ratios. These and related problems are under investigation.

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