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Chiral symmetry and strangeness content in nuclear physics parametrized by a medium dependent bag constant

R. Aguirre

Departamento de Fisica, Facultad de Ciencias Exactas, Univ. Nacional de La Plata, Argentina

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Abstract

Non-perturbative QCD vacuum effects at finite density are parametrized by means of a bag constant B. It is extracted from a Nambu–Jona-Lasinio model with two or three flavors. The parameter B is used in an effective bag-like model of baryons to study the nuclear phenomenology. We examine the nucleon structure and the thermodynamical properties of symmetric nuclear matter, particular attention is paid to the symmetry energy and to the eventual phase transition to deconfined quark matter. An alternative sketch of the binding mechanism of symmetric nuclear matter emerges within this approach. It is also found that the inclusion of strangeness content in B is crucial for an appropriate description.
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The low energy regime of QCD is very difficult to manage due to its non-perturbative nature. This feature gives rise to different methods, like sum rules and effective models, to avoid dealing with the full theory in the description of the hadronic phenomenology. These approaches intend to include the symmetries of QCD implicitly into the final results.

A simple shortcut procedure to take care of non-perturbative vacuum contributions is the introduction of the so-called bag constant B, in both confined and deconfined quark phases. The former bag models used it to account for the energy required to create a unit volume of the Weyl–Wigner vacuum into the complex QCD ground state. In Ref. [1] two different scales were associated to B, the higher one is related to the condensation of gluons and consequently with the recovery of the scale symmetry of QCD. The lower one is indicative of the restoration of chiral symmetry. Approximate calculations for these scales there given are $B^{1/4} = 245$ MeV and $B^{1/4} = 140$ MeV, respectively.

Another significant aspect in the definition of the bag parameter is its hypothetical variation with the medium properties of the hadronic matter. Speculation about the medium dependence of B has been motivated from different areas and ad-hoc phenomenological parametrizations have been proposed [2–6]. It is expected that the gradual recovery of fundamental symmetries with increasing density and/or tempera-

E-mail address: aguirre@venus.fisica.unlp.edu.ar (R. Aguirre).

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ture modifies the vacuum structure and the value of B accordingly.

In this work we attempt to extract a bag parameter that varies with the baryonic density but carrying the signature of the chiral symmetry and, ultimately to examine its influence over the nuclear dynamics at finite density. For this purpose we use two rather different theoretical descriptions of the high density regime of the strong interacting matter and we combine them to describe in a simplified manner the influence of the QCD vacuum changes over the gross properties of the nuclear matter. We have selected the Nambu–Jona-Lasinio (NJL) model [7] to study the modification of the QCD vacuum with increasing baryonic density and to extract a medium-dependent bag constant. This parameter is then inserted into the Quark Meson Coupling (QMC) model [8,9] to describe the nucleon substructure. It must be mentioned that each one of these issues have been independently investigated in the past, but we intend to put them together in a novel combination. For instance, a density dependent bag constant has been proposed within the QMC [2], but it was initially regarded as a way to restore certain aspects of the nuclear phenomenology instead of a manifestation of the underlying QCD vacuum. This fact explains why the parametrization of B has been made in terms of either the scalar σ meson mean field or the effective nucleon mass.

Adopting the viewpoint of the early bag models, we expect that B would carry significative information about the fundamental symmetries and would transfer it to the hadronic sector. Thus, as previously mentioned, the recovery of chiral and scale symmetries with increasing density and/or temperature would have remarkable effects over B.

The formulation of chiral invariant theories was certainly one of the major successes in the development of the bag models [10]. However, QMC abandoned this requirement since it was thought as a way to match the bag picture of the hadron structure with the Quantum Hadrodynamics [11] treatment of the nuclear many-body problem. The coupling to a scalar σ meson plays a key role in this description, therefore the explicit breaking of chiral symmetry seems to be an inherent feature of the QMC formulation. In the following we shall see how a bag constant respecting the chiral invariance in its definition, could modify this situation.

Our starting point is the original meaning of B, i.e., the energy difference between the fundamental state with explicit realization of the symmetries and that with symmetries spontaneously broken [12]. For this purpose we consider the NJL for both $SU(2)$ and $SU(3)$ versions of the chiral symmetries [13–15]. In the last case the flavor degeneracy is removed introducing a heavy strange quark.

The Lagrangian density of the model is [15]

$$\mathcal{L} = \bar{q}(i\not{\partial} - m)q + \mathcal{L}_4 + \mathcal{L}_6 \tag{1}$$

with $q^t = (u, d, s)$ and $m = \mathrm{diag}(m_u, m_d, m_s)$ is the current mass matrix that explicitly breaks chiral invariance. The four quarks interaction term $\mathcal{L}_4$ is invariant under $U_L(3) \times U_R(3)$ transformations

$$\mathcal{L}_4 = G\big[(\bar{q}\lambda_k q)^2 + (\bar{q}i\gamma_5\lambda_k q)^2\big], \tag{2}$$

where λ_k, $k = 0, \ldots, 8$ are the Gell-Mann matrices with $\lambda_0 = \sqrt{2/3}\,\mathrm{diag}(1, 1, 1)$.

The six quarks term removes the redundant $U_A(1)$ symmetry

$$\mathcal{L}_6 = -K\big\{\det\big[\bar{q}_i(1+\gamma_5)q_j\big] + \det\big[\bar{q}_i(1-\gamma_5)q_j\big]\big\}. \tag{3}$$

Due to the development of the quark–antiquark condensates the current quark masses acquires its large constituent values M_k, given by the self-consistent equations [15]

$$M_k = m_k - 4G\langle\bar{q}_k q_k\rangle + 2K\langle\bar{q}_i q_i\rangle\langle\bar{q}_j q_j\rangle, \tag{4}$$

with $i \neq k \neq j$, and the condensates are given by

$$\begin{aligned}\langle\bar{q}_k q_k\rangle &= C_k \\ &= -i \lim_{x_0' \to x_0+} \int \frac{d^4p}{(2\pi)^4} e^{-ip.(x'-x)}\,\mathrm{Tr}\,G_k(p) \\ &= -\frac{N_c}{2\pi^2} M_k\bigg[\Lambda E_\Lambda - p_k E_{p_k} \\ &\qquad - M_k^2 \ln\bigg(\frac{\Lambda + E_\Lambda}{p_k + E_{p_k}}\bigg)\bigg],\end{aligned} \tag{5}$$

where we have used the quark propagator

$$\begin{aligned}G_k(p) &= (\not{p} + M_k) \\ &\quad\times\bigg[\frac{1}{p^2 - M_k^2 + i\epsilon} \\ &\qquad + 2\pi i\,\delta\big(p^2 - M_k^2\big)\,\theta(p_k - p)\bigg]\end{aligned} \tag{6}$$

and $E_x = \sqrt{x^2 + M_k^2}$. We have used a non-covariant 3-momentum cutoff Λ, and explicitly included finite density effects using the Fermi momentum p_k. Therefore, the ground state is obtained filling the quark Fermi shell plus the Dirac sea divergent contribution. The contribution of k-flavor n_k to the total baryon density is related to its Fermi momentum by

$$n_k = \frac{N_c}{3}\frac{p_k^3}{3\pi^2}. \tag{7}$$

Once Eqs. (4) and (5) have been solved the energy density for the fundamental state is evaluated by taking the statistical average of the Hamiltonian density

$$\varepsilon_{QM} = \frac{N_c}{8\pi^2}\sum_k \bigl(F(M_k, p_k) - F(M_k, \Lambda)\bigr) + 2G\bigl(C_u^2 + C_d^2 + C_s^2\bigr) - 4KC_uC_dC_s, \tag{8}$$

with $F(M, x) = xE_x(E_x^2 + x^2) - M^4 \ln((x + E_x)/M)$.

A definite constant ε_0 must be added to Eq. (8) in order to get zero energy when no valence quarks are present. The pressure is evaluated using the thermodynamical relation

$$P_{QM} = \sum_k \mu_k n_k - \varepsilon_{QM}, \tag{9}$$

where $\mu_k = \partial\varepsilon_{QM}/\partial n_k$ is the chemical potential for quarks of flavor k.

As above mentioned we take B as the difference between Eq. (8) and the energy of the state with zero quark condensates, but at the same baryonic density. This definition is in agreement with that used in [16]. In this way a density dependent bag parameter has been obtained, able to be used in both hadron and quark matter since the parametrization is made in terms of only a common conserved charge. It is expected that the flavor composition of matter affects the density dependence of B, therefore, we examine three different situations (a) symmetric $u - d$ quark matter in the $SU(2)$ version of the model, (b) symmetric $u - d$ quark matter in the $SU(3)$ version, that is the Dirac sea contribution of strange quarks is included, and (c) electrically neutral quark matter coexisting with leptons in equilibrium by electroweak decay, a situation of interest in astrophysical applications. Thus we have the constraints for the quark Fermi momenta $p_u = p_d = p$, $p_s = 0$, $n = 2N_c p^3/(9\pi^2)$ for the items (a), (b) and $\mu_s = \mu_d$, $\mu_u + \mu_e = \mu_d$, $0 = N_c(2p_u^3 - p_d^3 - p_s^3)/3 - p_e^3$, $n = N_c(p_u^3 + p_d^3 + p_s^3)/(9\pi^2)$ for the case (c). Here n stands for the baryon number density and μ_e for the electron chemical potential.

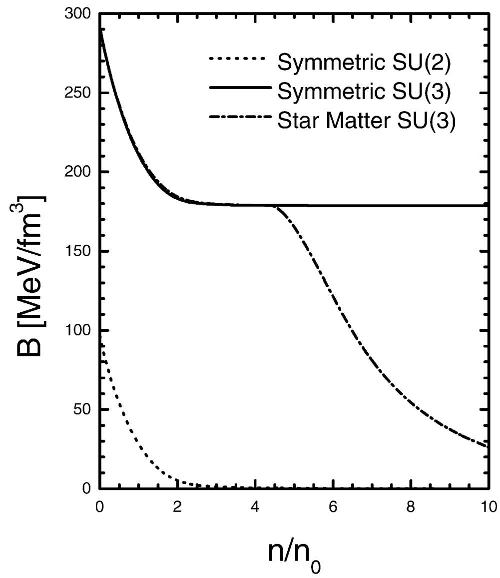

Fig. 1. The bag parameter in terms of the normalized baryon number density.

For practical calculations we have used $m_u = m_d = 5.5$ MeV, $m_s = 140.7$ MeV and the model parameters $\Lambda = 602.3$ MeV, $G\Lambda^2 = 1.835$, $K\Lambda^5 = 12.36$ MeV which were obtained by reproducing the π-, K- and η'-meson masses together with the constant f_π [14], for cases (b) and (c). For the item (a) instead, we have used the parameters given in [15] for the Hartree results in the 3-momentum cutoff approach, i.e., $\Lambda = 653$ MeV, $G\Lambda^2 = 2.14$.

The density dependence obtained in our calculations is shown in Fig. 1. The $SU(2)$ and $SU(3)$ schemes differ essentially in all the range, while within the three flavor treatment the cases (b) and (c) are similar for low and medium densities. The items (a) and (c) exhibit a steep descent at sufficiently high densities, in contrast with case (b) that keeps a large constant value. The cause of this behavior is that the strange condensate does not melt in (b) as the non-strange ones do. In turn this is due to the lack of the valence strange quark contribution, as confirmed by the sudden splitting of the curves (b) and (c) at $n \simeq 4.5\,n_0$ shown in Fig. 1. The fast decrease of the graph corresponding to (c) coincides with the turning up of the strange quark into the Fermi sea. Therefore, the large asymptotic value shown by the $SU(3)$ approach for the symmetric $u-d$ matter can be assigned to the constraint of null strange Fermi momentum.

The resulting B parameter can be used to evaluate nuclear matter properties. For this purpose we have selected a bag-like model that provides a satisfactory description of a wide range of nuclear phenomena, the QMC model [8,9]. Within the QMC baryons are regarded as non-overlapping spherical bags, where three valence quarks are confined. These quarks interact with its surrounding media by the interchange of light σ and ω-mesons, as suggested by the successful Quantum Hadrodynamics models [11]. Inside the bag the quark fields obey the mean-field equation

$$\left(i\gamma^{\mu}\partial_{\mu} - g_{\omega}^{k}\gamma^{0}\omega_{0} - m_{k}^{*}\right)q_{k}(x) = 0, \tag{10}$$

where $m_k^* = m_k - g_\sigma^k \sigma$, since within the bag there is only a small breakdown of chiral symmetry we use the current value of the quark masses m_k. In Eq. (10) σ and ω_0 stand for the mean field values of the meson fields evaluated in the dense hadronic medium and $g_{\sigma,\omega}^k$ are the quark–meson coupling constants. The solutions of this equation is [9]

$$q_k(r,t) = \mathcal{N}^{-1/2}\frac{e^{-i\varepsilon_k t}}{\sqrt{4\pi}}\begin{pmatrix} j_0(x_k r/R) \\ i\beta_k\vec{\sigma}.\hat{r}\, j_1(x_k r/R)\end{pmatrix}\chi^k, \tag{11}$$

where χ^k is the quark spinor and

$$\varepsilon_k = \frac{\Omega_k}{R} + g_\omega^k \omega_0,$$
$$\mathcal{N} = R^3\left[2\Omega_k(\Omega_k - 1) + R m_k^*\right]\frac{j_0^2(x_k)}{x_k^2},$$
$$\beta_k = \left[\frac{\Omega_k - R m_k^*}{\Omega_k + R m_k^*}\right]^{1/2} \tag{12}$$

with $\Omega_k = [x_k^2 + (R m_k^*)^2]^{1/2}$. The eigenvalue x_k is solution of the equation

$$j_0(x_k) = \beta_k j_1(x_k) \tag{13}$$

which arises from the condition of zero quark current through the bag surface.
In this model the ground state bag energy is identified with the baryon mass M^*,

$$M^* = \frac{\sum_k N_k \Omega_k - z_0}{R} + \frac{4}{3}\pi B R^3, \tag{14}$$

where N_k is the number of quarks of flavor k inside the bag. As usual in the bag model, two phenomenological corrective terms are added to the quark energy: the one containing z_0 amends the single particle spectrum and the other, proportional to the bag volume, represents the energy required to create a Weyl–Wigner bubble in the complex QCD vacuum. The last one has acquired a density dependence in our approach.

The bag radius adjusts to maintain the confinement volume in equilibrium with it surrounding media. Thus once B is given and z_0 fixed to reproduce the baryon mass at zero baryon density, the bag radius R is determined by the equilibrium condition $dM^*/dR = 0$.

Density dependent bag parameters have been proposed previously within the QMC description [2–4]. Formerly it was regarded as a way to restore the phenomenology of $\sigma - \omega$ fields at the saturation nuclear density $n_0 = 0.15\ \mathrm{fm}^{-3}$ [2]. However, the density dependence of B was guessed by means of the effective nucleon mass or by the σ-mean field, thus its relation with the underlying theory was not clear. Instead a phenomenological relation between B and the equation of state of nuclear matter was the proposal in [3], and restriction to experimental electron scattering analysis was imposed in [4]. The decrease of B at $0.7n_0$ is of 55% and 22% for curves (a) and (b), (c), respectively. In the last case it is only a few percent above the upper bound of 10–17% established in [4]. Based on different grounds, the necessity of a variable bag parameter was claimed in [6] in order to describe appropriately the structure of neutron stars.

Next we shall examine the behavior of symmetric nuclear matter at medium and high densities under the varying B-parameter that includes the signatures of the chiral symmetry of the strong interaction, considering two or three flavors contributions.

The thermodynamical properties in the QMC model are evaluated regarding nucleons as quasi-particles with single particle spectrum $\epsilon_p = \sqrt{p^2 + M^{*2}} + g_\omega \omega_0$, where the effective nucleon mass must be evaluated from Eq. (14). The energy density and the pressure for symmetric nuclear matter are, therefore [9]

$$\varepsilon_{NM} = \frac{1}{4\pi^2}F\left(M^*, p_F\right) + \frac{1}{2}\left(m_\sigma^2\sigma^2 + m\omega^2\omega^2\right), \tag{15}$$
$$P_{NM} = n\epsilon_{p_F} - \varepsilon_{NM}. \tag{16}$$

In Eq. (15) the energy of mesons has been included, p_F is the nucleon Fermi momentum related with the baryon density $n = 2p_F^3/(3\pi^2)$, and conditions of homogeneity and static equilibrium has been assumed. The symmetry energy of nuclear matter is a very significative quantity to understand a variety of nuclear processes at medium and high densities [17]. For example, it has a crucial role in the dynamics and structure of neutron stars, as well as in the composition of matter in heavy-ion collisions. However, its behavior far from the normal nuclear density is poorly known experimentally, and theoretical predictions range over a wide and sometimes contradictory qualitative results. For a linear nucleon-ρ meson coupling the energy symmetry for symmetric nuclear matter is

$$\epsilon_{\text{sym}} = \frac{1}{8}\left(\frac{g_\rho}{m_\rho}\right)^2 n + \frac{p_F^2}{6\sqrt{p_F^2 + M^{*2}}}. \tag{17}$$

The mesons mean field values have not been determined yet, they can be obtained minimizing the energy density with respect to σ and ω. This leads to

$$\sigma = -\frac{1}{2\pi^2 m_\sigma^2}\frac{\partial M^*}{\partial \sigma} \times \left[p_F\sqrt{p_F^2 + M^{*2}} - M^{*2}\sinh^{-1}\left(p_F/M^*\right)\right], \tag{18}$$

$$\omega = \frac{g_\omega}{m_\omega^2} n. \tag{19}$$

The nucleon–meson couplings are expressed in terms of the quark–meson couplings as $g_\sigma = 3g_\sigma^k$, $g_\omega = 3g_\omega^k$. Numerical values can be assigned reproducing the conditions of nuclear matter at the saturation density n_0. Meaningfully it is found in the $SU(3)$ treatment that, for a certain value of the binding energy, one can do calculations without the σ-meson coupling, i.e., $g_\sigma^k = 0$. Thus we choose the values $P_{NM}(n_0) = 0$, $\varepsilon_{NM}(n_0) - M_0 = -16.807$ MeV for the pressure and the binding energy at the normal density. If the vector–isovector meson ρ is also considered, then the corresponding coupling constant $g_\rho = g_\rho^k$ can be deduced equating the symmetry energy $\epsilon_{\text{sym}} = (\partial\epsilon_N/\partial n)$ at n_0 to its phenomenological value $\epsilon_{\text{sym}} = 34$ MeV. The coupling constants obtained under these assumptions are shown in Table 1.

Table 1
Coupling constants for the QMC model for the cases considered: (a) and (b) correspond to medium dependent B obtained from a Nambu–Jona-Lasinio model with $SU(2)$ and $SU(3)$ chiral symmetries, respectively, (d1) and (d2) correspond to the constant B fixed at the zero density value of (a) and (b), respectively. In the last two columns the zero density value of the bag parameter and the bag radius are shown

Case	g_σ	g_ω	g_ρ	$B_0^{1/4}$ (MeV)	R_0 (fm)
(a)	11.97	16.46	8.43	165.06	0.83
(b)	0.00	6.00	9.53	217.59	0.57
(d1)	17.39	8.26	9.17	165.06	0.83
(d2)	18.30	9.23	9.04	217.59	0.57

Numerical calculations have been developed taking $m_\sigma = 2.787$ fm^{-1}, $m_\omega = 3.968$ fm^{-1}, and $m_\rho = 3.902$ fm^{-1} for the meson masses, $M_p = 4.755$ fm^{-1}, $M_n = 4.761$ fm^{-1} for the nucleon masses at zero baryon density. For the sake of comparison we have also carried out calculations with the standard fixed B treatment (in the following labelled case (d)).

Some results for the in-medium nucleon properties are displayed in Fig. 2. The neutron mass is uniformly decreasing, the falling at n_0 is approximately 10% in the $SU(3)$ treatment, 20–25% in the constant B option and of an excessively high 40% in the $SU(2)$ case. The behavior of the bag radius in terms of the baryon density is shown in Fig. 2, the cases (b) and (d) provide a radius with a bounded variation. In particular, for the approach (b) an increasing radius is obtained whose growth saturates 15% over its zero density value. Instead for case (a) a monotonously increasing radius is obtained, exceeding the limit of validity of the model. The large increment of the bag radius yields a violation of the non-overlapping bag hypothesis at $n \simeq n_0$, thus we may conclude that the two flavor description of the bag parameter is not able of true physical consideration. In the three flavor case this assumption is also broken down but at the higher density $n \simeq 6n_0$.

The results for the pressure and the symmetry energy are shown in Fig. 3, it can be seen that the softest equation of state at high densities corresponds to (d). The nuclear compressibility $\kappa = 9\,\partial P_{NM}/\partial n$ takes on the values 310 and 335 MeV/fm^3 at $n = n_0$ for the (d) and (b) approaches, respectively. For the symmetry energy all the results are very similar for

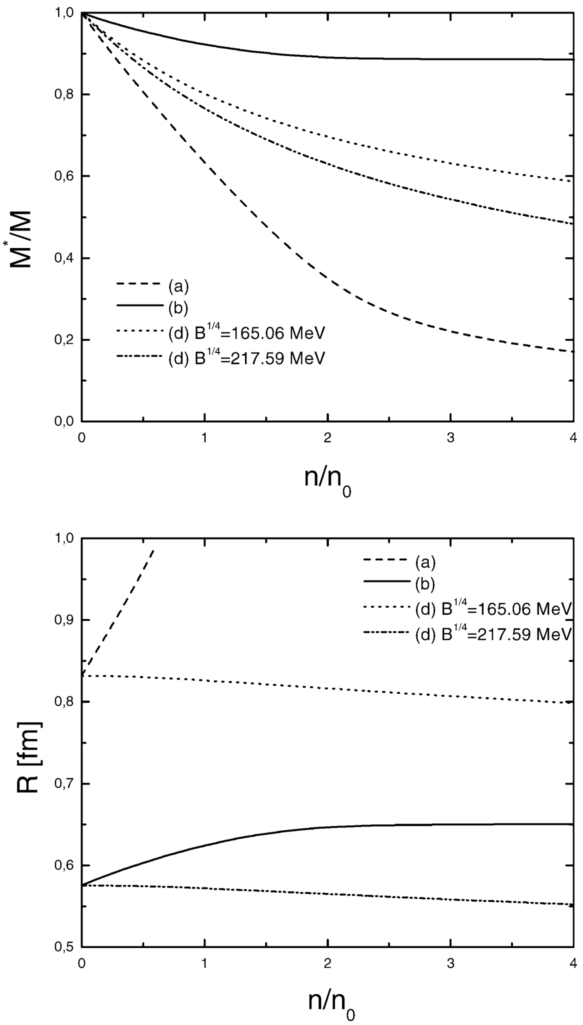

Fig. 2. In-medium nucleon properties as functions of the normalized baryon number density. In the upper figure the effective neutron mass relative to its value at zero baryon number density, in the lower figure the equilibrium bag radius. The line convention is explained in each figure, the labels (a), (b) and (d) have the same meaning as in the text. In the last case we have taken two different constant values for B obtained at zero baryon density in the treatments (a) and (b).

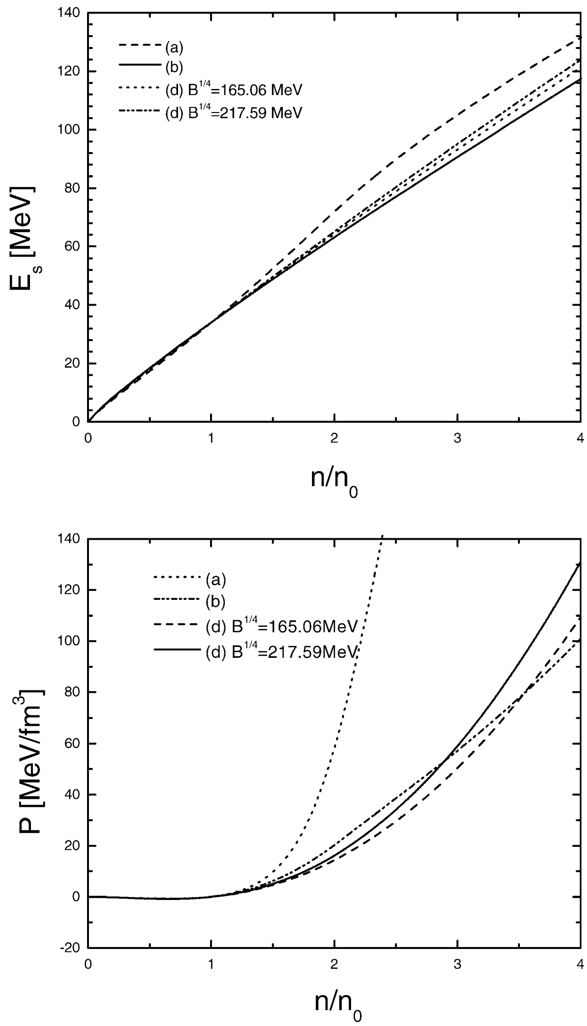

Fig. 3. Thermodynamical properties of symmetric nuclear matter as functions of the normalized baryon number density. In the upper figure the symmetry energy and the pressure in the lower panel. For the line convention see the explanation in Fig. 2.

low densities and differences becomes sensible only beyond $n/n_0 \sim 2.5$.

Thermodynamical quantities are similar in the approaches (b) and (d), corresponding greater pressures and lower symmetry energy to the instance (b). Instead nucleon mass and radius have distinguishing behaviors in all the range of densities.

The medium dependent B description has an upper bound of applicability given by the point where bags start to overlap, unless new physical aspects arise before this situation is reached. This could be the case if a deconfinement phase transition take place. To examine this possibility we have compared the equation of state obtained previously with that corresponding to quark matter. The Gibbs criterium for phase transitions $P_{NM} = P_{QM}$ and $\mu_n = 2\mu_d + \mu_u$ is adopted, and variable or fixed B is used. Thus a first order phase transition is found only for the case (a), but at so high density that QMC model assumptions have been violated. Not transition at all is found for the cases (b).

In this Letter we have studied the effects of including the chiral symmetry of QCD and the strange degree of freedom in the definition of the bag parameter B. This has been obtained from a Nambu–Jona-Lasinio model with two or three flavors, at finite baryon density. Since baryon number is a conserved charge of both quark and hadron matter, the quantity B obtained can be regarded as medium-dependent and applied to examine nucleon properties and the nuclear matter equation of state. Thus our results for B are based on a dynamical effective model of QCD, in contrast with previous calculations that used *ad hoc* parametrizations in terms of the σ-meson mean field. The description in terms of only non-strange quarks must be discarded, i.e., the bag parameter must include the strange vacuum contribution even for the description of ordinary nuclear matter. Within the $SU(3)$ scheme we have found that the behavior of B strongly depends on the composition of matter at medium and high densities.

For the hadronic phase we used the QMC model of confined quarks coupled to mesons, that allows us to extract the confinement volume medium-dependence. From the comparison of the items (b) and (d) we conclude that the case (b) does not need of the σ meson to appropriately describe the properties of nuclear matter. Therefore, the saturation mechanism differs from those of Quantum Hadrodynamics, the attractive contribution is provided by the interplay between many-body effects and the modification of the hadronic phase vacuum. Avoiding a direct coupling between quarks and the σ-meson opens new possibilities such as a chiral extension of the QMC, or could eliminate the unpleasant features reported in [18].

We have not found any first order phase transition from symmetric nuclear matter to quark matter, but this could be a failure of the approach that could be mended including quark correlations between neighboring bags in the high density regime of the hadronic phase [19].

Acknowledgements

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References

[1] C. Adami, G.E. Brown, Phys. Rep. 234 (1993) 1.
[2] X. Jin, B.K. Jennings, Phys. Lett. B 374 (1996) 13;
X. Jin, B.K. Jennings, Phys. Rev. C 54 (1996) 1427;
H. Muller, B.K. Jennings, Nucl. Phys. A 626 (1997) 966.
[3] R. Aguirre, M. Schvellinger, Phys. Lett. B 400 (1997) 245.
[4] D.H. Lu, K. Tsushima, A.W. Thomas, A.G. Williams, Nucl. Phys. A 634 (1998) 443.
[5] D. Blaschke, H. Grigorian, G. Poghosyan, C.D. Roberts, S. Schmidt, Phys. Lett. B 450 (1999) 207.
[6] G.F. Burgio, M. Baldo, P.K. Sahu, A.B. Santra, H.-J. Schulze, Phys. Lett. B 526 (2002) 19.
[7] Y. Nambu, G. Jona-Lasinio, Phys. Rev. 122 (1961) 345;
Y. Nambu, G. Jona-Lasinio, Phys. Rev. 124 (1961) 246.
[8] P.A.M. Guichon, Phys. Lett. B 200 (1988) 235.
[9] K. Saito, A.W. Thomas, Phys. Lett. B 327 (1994) 9;
K. Saito, A.W. Thomas, Phys. Rev. C 51 (1995) 2757.
[10] A.W. Thomas, Adv. Nucl. Phys 13 (1984) 1.
[11] B.D. Serot, J.D. Walecka, Int. J. Mod. Phys. E 6 (1997) 515.
[12] S. Li, R.S. Bhalerao, R.K. Bhaduri, Int. J. Mod. Phys. 6 (1991) 501.
[13] T. Kunihiro, T. Hatsuda, Phys. Lett. B 206 (1988) 385;
V. Bernard, R.L. Jaffe, U.-G. Meissner, Nucl. Phys. B 308 (1988) 753.
[14] P. Rehberg, S.P. Klevansky, J. Hufner, Phys. Rev. C 53 (1996) 410.
[15] S.P. Klevansky, Rev. Mod. Phys. 64 (1992) 649.
[16] M. Buballa, M. Oertel, Phys. Lett. B 457 (1999) 261.
[17] B.-A. Li, Phys. Rev. Lett. 88 (2002) 192701;
B.-A. Li, Nucl. Phys. A 708 (2002) 365.
[18] S.W. Hong, B.K. Jennings, Phys. Rev. C 64 (2001) 038203.
[19] K. Saito, K. Tsushima, A.W. Thomas, LANL report, nucl-th/9901084;
R. Aguirre, A.L. De Paoli, LANL report, nucl-th/0211038.