Hot Nuclear Matter in the Quark Meson Coupling Model with Dilatons

I. Zakout and H. R. Jaqaman

Department of Physics, Bethlehem University, P.O. Box 9, Bethlehem, Palestine

Abstract

We study hot nuclear matter in an explicit quark model based on a mean field description of nonoverlapping nucleon bags bound by the self-consistent exchange of scalar and vector mesons as well as the glueball field. The glueball exchange as well as a realization of the broken scale invariance of quantum chromodynamics is achieved through the introduction of a dilaton field. The calculations also take into account the medium-dependence of the bag constant. The effective potential with dilatons is applied to nuclear matter. The nucleon properties at finite temperature as calculated here are found to be appreciably different from cold nuclear matter. The introduction of the dilaton potential improves the shape of the saturation curve at $T=0$ and is found to affect hot nuclear matter significantly.

PACS: 21.65.+f, 24.85.+p, 12.39Ba
I. INTRODUCTION

The description of nuclear phenomena in relativistic mean-field theory has been successfully formulated using the hadronic degrees of freedom \cite{1,2}. However, due to the observations which revealed the medium modification of the internal structure of the nucleon \cite{3}, it has become essential to explicitly incorporate the quark-gluon degrees of freedom while respecting the established model based on the hadronic degrees of freedom in nuclei. It was in this spirit that Guichon proposed the quark-meson coupling model which describes nuclear matter as a collection of non-overlapping MIT bags interacting through the self-consistent exchange of mesons in the mean field approximation with the meson fields directly coupled to the quarks \cite{4–7}. It has been further suggested that including a medium-dependent bag constant may be essential for the success of relativistic nuclear phenomenology \cite{6}. This modification has been recently applied in a study of the properties of nuclear matter at finite temperature \cite{8} where it was found that the bag constant decreases appreciably above a critical temperature $T_c \approx 200\,\text{MeV}$ indicating the onset of quark deconfinement.

On the other hand, it is important to incorporate the breaking of the classical scale invariance of Quantum Chromodynamics (QCD) resulting from the QCD trace anomaly \cite{9}. It has been determined that it is possible to incorporate the broken global chiral and scale symmetries of QCD into an effective lagrangian through the introduction of a scalar, chiral isoscalar glueball-dilaton field $\phi$ \cite{9–15}. There have also been attempts to introduce the dilaton field in an appropriate way in order to obtain reasonable values for the compression modulus of nuclear matter \cite{12,13,16}. It has been noted that the scale invariant term which leads to an omega meson, after symmetry breaking, is strongly favored to be of the form $\omega^\mu \omega^\nu \frac{\phi^2}{\phi_0^2}$ by the bulk properties of nuclei \cite{13}. The effective lagrangian

$$\mathcal{L} = \mathcal{L}_0 - V_G,$$

consists of a chiral and scale invariant part $\mathcal{L}_0$ and an explicitly scale-breaking potential $V_G$. The divergence of the scale current in QCD is given by the trace of the improved energy
momentum tensor and so the scalar potential \( V_G \) was chosen to reproduce, via Noether’s theorem, the effective trace anomaly \[ \theta^\mu_\mu = 4V_G(\Phi) - \sum_i \Phi_i \frac{\partial V_G}{\partial \Phi_i} = 4\epsilon_{\text{vac}} \left( \frac{\phi}{\phi_0} \right)^4, \] (1.2)
where \( \Phi_i \) runs over the scalar fields \( \{\sigma, \vec{\pi}, \phi\} \) and \( \epsilon_{\text{vac}} \) is the vacuum energy \[ 12,13 \]. The proportionality \( \theta^\mu_\mu \propto \left( \frac{\phi}{\phi_0} \right)^4 \) is suggested by the form of the QCD trace anomaly. In the mean field approximation, we have \( <\vec{\pi}> = 0 \) for symmetric nuclear matter so that the isovector pionic field has no effect on the properties of nuclear matter in this approximation.

The outline of the present paper is as follows. In Sect. II, we formalize the quark-meson coupling model with dilatons at finite temperature. This is essentially an extension of our earlier work \[ 8 \] where the dilaton field was not included. Finally in Sect. III, we present our results and conclusions.

**II. THE QUARK MESON COUPLING MODEL WITH DILATONS AT FINITE TEMPERATURE**

The quark field \( \psi_q(\vec{r}, t) \) inside the bag satisfies

\[
\left[ i\gamma^\mu \partial_\mu - \left( m_q^0 \frac{\phi}{\phi_0} - g_\sigma^0 \sigma \right) - g_\omega^0 \omega \gamma^0 \right] \psi_q(\vec{r}, t) = 0, \]

(2.1)

where \( m_q^0 \) is the current quark mass while \( g_\sigma^0 \) and \( g_\omega^0 \) are the quark couplings with the scalar \( \sigma \) and vector \( \omega \) mesons, respectively, and \( \frac{\phi}{\phi_0} \) is the scale invariance breaking.

The single-particle quark and antiquark energies in units of \( R^{-1} \), where \( R \) is the bag radius, are given by \[ 7 \]

\[ \epsilon_{n\kappa} = \Omega^{n\kappa} \pm g_\omega^0 \omega R \] (2.2)

where \( \Omega^{n\kappa} = \sqrt{x_{n\kappa}^2 + R^2 m_q^2} \) and \( m_q^* = \frac{\phi}{\phi_0} m_q^0 - g_\sigma^0 \sigma \) is the effective quark mass. The quark momentum \( x_{n\kappa} \) in the state characterized by specific values of \( n \) and \( \kappa \) is determined by the boundary condition at the bag surface.
\[ i \gamma \cdot n \psi^{\text{nk}}_q |_R = \psi^{\text{nk}}_q |_R, \]  
\[ (2.3) \]

which reduces for the ground state \[3\] to

\[ j_0(x_{nk}) = \beta^{\text{nk}}_q j_1(x_{nk}) \]  
\[ (2.4) \]

where

\[ \beta^{\text{nk}}_q = \sqrt{\Omega^{\text{nk}} - R m_q^* \over \Omega^{\text{nk}} + R m_q^*} \]  
\[ (2.5) \]

and the coefficient \( \kappa = -1 \) for the \( s \)-state.

The total energy from the quarks and antiquarks reads

\[ E_{\text{tot}} = 3 \sum_{n_{k}} \Omega^{n_{k}}_R \left[ {1 \over e^{(\epsilon^{n_{k}}_e/R-\mu_q)/T} + 1} + {1 \over e^{(\epsilon^{n_{k}}_e/R+\mu_q)/T} + 1} \right]. \]  
\[ (2.6) \]

The bag energy now becomes

\[ E_{\text{bag}} = E_{\text{tot}} - \left( Z \Phi {\Phi_0 \over R} + {4\pi \over 3} R^3 \Phi_0 \Phi B(\sigma) \right) \]  
\[ (2.7) \]

where the \( \Phi \) dependence is suggested by the asymptotic behavior of the bag energy in the purely hadronic models \[13\].

\[ E_{\text{bag}} \propto {\phi \over \phi_0} M_N. \]  
\[ (2.8) \]

The medium effects are taken into account by adopting the bag parameter \( B(\sigma) = B_0 \exp \left( -{4g^{\text{B}}_\sigma \over \phi_0 M_N} \right) \) where \( g^{\text{B}}_\sigma \) is an additional fitting parameter \[3\]. The spurious center-of-mass motion in the bag is subtracted to obtain the effective nucleon mass \[17\]

\[ M_N^* = \sqrt{E_{\text{bag}}^2 - <p_{\text{cm}}^2>>}. \]  
\[ (2.9) \]

The spurious center-of-mass average momentum squared is given by

\[ <p_{\text{cm}}^2> = {x^2 \over R^2}, \]  
\[ (2.10) \]

where
\[<x^2> = 3 \sum_{n\kappa} x_{n\kappa}^2 \left[ \frac{1}{e^{(\epsilon_{+n}/R-\mu_q)/T} + 1} + \frac{1}{e^{(\epsilon_{-n}/R+\mu_q)/T} + 1} \right] \]  

(2.11)
is written in terms of the sum of the quark and antiquark Fermi distribution functions since the center of mass motion does not distinguish between a quark and an antiquark. The quark chemical potential \(\mu_q\), assuming that there are three quarks in the nucleon bag, is determined through the constraint

\[n_q = 3 = 3 \sum_{n\kappa} \left[ \frac{1}{e^{(\epsilon_{+n}/R-\mu_q)/T} + 1} - \frac{1}{e^{(\epsilon_{-n}/R+\mu_q)/T} + 1} \right]. \]  

(2.12)
The temperature-dependent bag radius \(R\) is obtained through the condition

\[\frac{\partial M_N^*}{\partial R} = 0. \]  

(2.13)
The baryon-antibaryon thermal distribution functions are given by

\[f_B(\mu_B) = \frac{1}{e^{(\epsilon^*(\vec{k})-\mu_B^*)/T} + 1} \]  

(2.14)
and

\[\overline{T}_B(\mu_B) = \frac{1}{e^{(\epsilon^*(\vec{k})+\mu_B^*)/T} + 1} \]  

(2.15)
with \(\epsilon^*(\vec{k}) = \sqrt{\vec{k}^2 + M_N^*}^2\) and the effective baryon chemical potential \(\mu_B^* = \mu_B - g_\omega \omega\). The baryon chemical potential, \(\mu_B\), is obtained from the nontrivial solution of

\[\rho_B = \frac{\gamma}{(2\pi)^3} \int d^3k \left( f_B(\mu_B) - \overline{T}_B(\mu_B) \right). \]  

(2.16)
However, for a given chemical potential \(\mu_B\), the baryon density \(\rho_B\) can be calculated easily. Subsequently, the vector mean field reads

\[\omega = \frac{1}{m_\omega (\phi/\phi_0)^2 g_\omega \rho_B}. \]  

(2.17)
The total energy density at finite temperature and baryon density \(\rho_B\) reads

\[
\epsilon = \frac{\gamma}{(2\pi)^3} \int d^3k \sqrt{k^2 + M_N^*} \left[ f_B(\mu_B) + \overline{T}_B(\mu_B) \right]
+ \frac{1}{2} m_\omega \left( \frac{\phi}{\phi_0} \right)^2 \omega^2 + \frac{1}{2} m_\sigma \left( \frac{\phi}{\phi_0} \right)^2 \sigma^2 + U[\phi],
\]  

(2.18)
with the dilaton potential $U[\phi]$ given by $U[\phi] = V_G[\phi] - V_G[\phi_0]$ where

$$V_G[\phi] = \epsilon_{\text{vac}} \left( \frac{\phi}{\phi_0} \right)^4 \left[ \log \left( \frac{\phi^4}{\phi_0^4} \right) - 1 \right].$$

Here $\epsilon_{\text{vac}}$ is the dilaton potential constant and corresponds to the vacuum energy. The pressure of nuclear matter reads

$$P = \frac{1}{3} \frac{\gamma}{(2\pi)^3} \int d^3 k \, \frac{\vec{k}^2}{\epsilon^*(k)} \left[ f_B(\mu_B) + \overline{f}_B(\mu_B) \right] + \frac{1}{2} m^2 \omega \left( \frac{\phi}{\phi_0} \right)^2 - \frac{1}{2} m^2 \sigma \left( \frac{\phi}{\phi_0} \right)^2 - U[\phi].$$

The scalar and dilaton fields are determined by the extremization of the pressure through the conditions

$$\frac{\partial P}{\partial \sigma} \bigg|_{\sigma} = \left( \frac{\partial P}{\partial M^*_N} \right)_{\mu_B,T} \left( \frac{\partial M^*_N}{\partial \sigma} \right) + \left( \frac{\partial P}{\partial \sigma} \right)_{M^*_N} = 0,$$

and

$$\frac{\partial P}{\partial \chi} \bigg|_{\phi/\phi_0} = \left( \frac{\partial P}{\partial M^*_N} \right)_{\mu_B,T} \left( \frac{\partial M^*_N}{\partial \chi} \right) + \left( \frac{\partial P}{\partial \chi} \right)_{M^*_N} = 0,$$

for the $\sigma$ and $\chi = \frac{\phi}{\phi_0}$ fields, respectively. This extremization is carried out while taking into consideration the full coupling of the scalar mean fields to the internal quark structure of the bag by means of the solution of the point-like Dirac equation with the required boundary condition of confinement at the surface of the bag as suggested by Refs. [5, 8]. The variation of the effective chemical potential $\mu^*_B$ is taken as

$$\frac{\partial \mu^*_B}{\partial M^*_N} = -g_\omega \frac{\partial \omega}{\partial M^*_N}$$

where

$$\frac{\partial \omega}{\partial M^*_N} = -\frac{g_\omega}{(\phi/\phi_0)^2 m^2 \omega \gamma (2\pi)^3} \int d^3 k \, \frac{M^*_N}{\epsilon} \left( f_B(1 - f_B) - \overline{f}_B(1 - \overline{f}_B) \right) + \frac{1}{1 + \frac{g_\omega}{(\phi/\phi_0)^2 m^2 \omega \gamma (2\pi)^3} \int d^3 k \, f_B(1 - f_B) + \overline{f}_B(1 - \overline{f}_B)}.$$

**III. RESULTS AND CONCLUSIONS**

We have used the quark meson coupling to study nuclear matter at zero and finite temperatures. Glueball exchange and the breaking of scale invariance are taken into account...
by introducing the dilaton field. We have adopted the fitting parameters that are used in the earlier calculations, in the absence of the dilaton potential, which fit the saturation properties of nuclear matter [1,7]. The bag parameters $B_0^{1/4} = 188.1$ MeV and $Z_0 = 2.03$ are chosen to reproduce the free nucleon mass $M_N$ at its experimental value 939 MeV and bag radius $R_0 = 0.60$ fm. The current quark mass $m_q$ is taken equal to zero. For $g^2_\sigma = 1$, the values of the vector meson coupling and the parameter $g^B_\sigma$ as fitted from the saturation properties of nuclear matter, are given as $g^2_\omega/4\pi = 5.24$ and $g^{B2}_\sigma/4\pi = 3.69$.

For specific values of the temperature and $\mu_B$, the thermodynamic potential is given in terms of the effective nucleon mass which depends on the bag radius $R$, the quark chemical potential $\mu_q$ and the mean fields $\sigma$ and $\omega$ as well as the dilaton field $\phi/\phi_0$. For given values of the mean fields $\sigma$ and $\omega$ as well as the dilaton field $\phi/\phi_0$, the quark chemical potential $\mu_q$ and the bag radius $R$ are determined using the conditions Eqns.(2.12) and (2.13), respectively. We then determine the values of $\sigma$ and $\phi/\phi_0$ by minimizing the thermodynamic potential $\Omega$ through the conditions of Eqns.(2.21) and (2.22), respectively, together with the self-consistency condition [7] for the $\omega$ mean field as given in Eq.(2.17). For given values of the temperature and baryon chemical potential $\mu_B$, we have calculated the different thermodynamic quantities.

The breaking of scale invariance is tested by studying the behavior of the dilaton scale $\chi = \left(\frac{\phi}{\phi_0}\right)$. When the glueball field is frozen, the dilaton scale takes the value $\chi = 1$ [15,16]. We display the dilaton scale $\chi$ as a function of the baryon density $\rho_B$ for different values of temperature in Fig. 1. The dilaton potential constant is taken as $\epsilon_{vac} = (250$ MeV$)^4$. The dilaton field for cold nuclear matter (T=0) increases weakly at small baryon density (for values $\rho_B \leq 0.12$ fm$^{-3}$). When $\rho_B$ reaches 0.12 fm$^{-3}$, it starts to decrease. Furthermore, it takes values $\chi < 1$ for baryon densities $\rho_B > 0.22$ fm$^{-3}$. The behavior is quite different for hot nuclear matter where the dilaton scale $\chi$ is monotonically increasing with $\rho_B$ at all temperatures. It is interesting to note that in contrast to the results of calculations involving only baryonic degrees of freedom [15] where the dilaton scale at finite temperature takes values $\chi < 1$ it here takes values $\chi > 1$. Moreover, the dilaton scale is seen to attain

7
values $\chi > 1$ at zero baryon density for temperatures $T \geq T_c \approx 200$ MeV. This is related to the phase transition seen in the earlier calculations without the dilaton field when the system becomes a dilute gas of baryons in a sea of baryon-antibaryon pairs and quark deconfinement sets in. Therefore at the phase transition, the scale invariance is broken even at zero baryon density.

The effective nucleon mass $M_N^*$ decreases monotonically with the baryon density $\rho_B$ as can be seen in Fig. 2. It also increases with temperature for all values of $\rho_B$ up to the critical temperature and then suddenly starts to decrease rapidly with temperature for $T \geq 200$ MeV at low baryon density $\rho_B$. This rapid fall of $M_N^*$ with increasing temperature is also related to the above-mentioned phase transition and was observed in the earlier calculation. The general behavior of the pressure with temperature and baryon density is also found to be similar to that obtained in the earlier calculation in the absence of the dilaton field. For example, it takes a nonzero value at $\rho_B = 0$ for temperatures $T \geq 200$ MeV.

The bag radius $R$ as a function of the baryon density $\rho_B$ is displayed in Fig. 3 for several values of temperature. The bag radius $R$ increases monotonically with the increase in the baryon density and decreases when the temperature is increased. However, when the temperature reaches 200 MeV, the bag radius $R$ suddenly starts to increase with temperature for low baryon densities. For sufficiently high temperatures, for instance $T = 240$ MeV, the bag radius $R$ tends to be approximately constant with very little variation with the baryon density $\rho_B$.

Fig. 4 shows the scalar field $\sigma$ as a function of the baryon density $\rho_B$. The scalar field $\sigma$ increases with $\rho_B$ and, at first, tends to decrease with temperature until the temperature reaches 200 MeV when it suddenly starts to increase for low baryon densities. Furthermore, the scalar field $\sigma$ takes nonzero values at and above the critical temperature $T_c = 200$ MeV. This behavior was also observed in the earlier calculations with a frozen glueball and is related to the phase transition and the onset of quark deconfinement.

To study further the effect of the dilaton field on the quark meson coupling model, we investigate the effect of the value of the dilaton potential constant, $\epsilon_{\text{vac}}$, on the properties
of cold and hot nuclear matter. In Fig. 5 we plot the dilaton scale $\chi$ as a function of $\rho_B$ for several values of $\epsilon_{\text{vac}}$. The cold nuclear matter case is displayed in Fig. 5(a) where it is seen that the dilaton field increases with $\rho_B$ for low baryon densities until it reaches its maximum value at $\rho_B = 0.12$ fm$^{-3}$ and then starts to decrease. The curves of the dilaton scale $\chi$ for several values of $\epsilon_{\text{vac}}$ all intersect with the frozen glueball scale ($\chi = 1$) at $\rho_B = 0.24$ fm$^{-3}$ and then continue to decrease below $\chi = 1$, for baryon densities $\rho_B > 0.24$ fm$^{-3}$. It is seen that for the lowest value $\epsilon_{\text{vac}} = (200\text{MeV})^4$ the variation in $\chi$ with $\rho_B$ is very dramatic and tend to very small values for $\rho_B > 0.35$ fm$^{-3}$ where the solution becomes unstable indicating that such low values of $\epsilon_{\text{vac}}$ are physically unacceptable. For larger values of $\epsilon_{\text{vac}}$ the variation becomes less and less dramatic. For sufficiently large values, for instance, $\epsilon_{\text{vac}} = (800\text{MeV})^4$, it tends to act almost as a frozen glueball field and the dilaton scale becomes equal to one. The variation of $\chi$ with $\epsilon_{\text{vac}}$ for hot nuclear matter at the temperature $T = 200$ MeV is displayed in Fig. 5(b). The dilaton scale $\chi$ increases monotonically with $\rho_B$ with $\chi$ always $> 1$ and is found to decrease with $\epsilon_{\text{vac}}$. Its variation with density becomes very steep for $\epsilon_{\text{vac}} = (200\text{MeV})^4$ while it tends to be constant for $\epsilon_{\text{vac}} = (800\text{MeV})^4$. It is thus seen that the behavior of the dilaton scale $\chi$ is completely different for cold and hot nuclear matter for low values of the dilaton potential constant $\epsilon_{\text{vac}}$. However, for sufficiently large values of $\epsilon_{\text{vac}}$, the dilaton field changes very slightly and cannot be distinguished from the case without dilatons for both cold and hot nuclear matter.

The dependence of the effective nucleon mass $M_N^*$ on the value of $\epsilon_{\text{vac}}$ is displayed in Figs. 6(a) and 6(b) for cold and hot nuclear matter, respectively. In the case of cold nuclear matter, the nucleon mass $M_N^*$ does not seem to be affected by the value of $\epsilon_{\text{vac}}$. However, for the hot nuclear matter case at a temperature $T = 200$ MeV, $M_N^*$ increases slightly with $\epsilon_{\text{vac}}$. In both cases, $M_N^*$ decreases with $\rho_B$ as already observed in Fig. 2 as well as in the calculations without the dilaton [8,18].

The variation of the bag radius $R$ with the value of $\epsilon_{\text{vac}}$ is displayed in Fig. 7(a) for cold nuclear matter. It is seen that $R$ increases monotonically with the baryon density $\rho_B$
for all $\epsilon_{\text{vac}}$ values and it increases weakly with $\epsilon_{\text{vac}}$ for low baryon densities $\rho_B < 0.24 \text{fm}^{-3}$. However, when $\rho_B = 0.24 \text{ fm}^{-3}$ the bag radius takes the same value for all values of $\epsilon_{\text{vac}}$. When $\rho_B$ exceeds 0.24 $\text{ fm}^{-3}$, the bag radius $R$ starts to decrease with the increase in the value of $\epsilon_{\text{vac}}$. The variation of the bag radius for hot nuclear matter at $T = 200 \text{ MeV}$ is displayed in Fig. 7(b). It is seen that the bag radius increases with both $\rho_B$ and $\epsilon_{\text{vac}}$ except for the smallest value $\epsilon_{\text{vac}} = (200\text{MeV})^4$ where it is found that the bag radius $R$ at first increases weakly with $\rho_B$ and then starts to decrease.

The scalar field $\sigma$ as a function of the baryon density $\rho_B$ is displayed in Fig. 8. For cold nuclear matter, Fig. 8(a), it is seen that $\sigma$ increases with increasing $\rho_B$ and is almost independent of $\epsilon_{\text{vac}}$ for $\rho_B \leq 0.24 \text{ fm}^{-3}$. For $\rho_B > 0.24 \text{ fm}^{-3}$, it decreases weakly with $\epsilon_{\text{vac}}$. On the other hand for the hot nuclear matter case at $T = 200 \text{ MeV}$, it is seen in Fig. 8(b) that the scalar field $\sigma$ increases with both $\rho_B$ and $\epsilon_{\text{vac}}$.

We also examined the saturation properties of nuclear matter at $T=0$. In Fig. 9, we display $E_{\text{tot}} - M_N$ as a function of the baryon density for several values of $\epsilon_{\text{vac}}$. It is seen that as $\epsilon_{\text{vac}}$ decreases, the equation of state becomes stiffer and the compressibility, $K$, increases [12,13]. The compressibility was found to have the reasonable value $K=300 \text{ MeV}$ for $\epsilon_{\text{vac}} = (800\text{MeV})^4$. However, for smaller values of $\epsilon_{\text{vac}}$, the compressibility increases. For example, $K$ is increased by about a factor of two if $\epsilon_{\text{vac}}$ is changed from $(800\text{MeV})^4$ to $(200\text{MeV})^4$.

**ACKNOWLEDGMENTS**

Financial support by the Deutsche Forschungsgemeinschaft through the grant GR 243/51-1 is gratefully acknowledged.
REFERENCES

[1] J. D. Walecka, Ann. of Phys.83, 491 (1974); Phys. Lett. B 59, 109 (1975).

[2] B. D. Serot and J. D. Walecka, Advan. Nucl. Phys. vol 16, 1 (1986).

[3] EMC Collaboration, J. J. Aubert et al., Phys. lett. B 123, 275 (1983); M. Arneodo, Phys. Rep. 240, 301 (1994).

[4] P. A. M. Guichon, Phys. Lett. B 200 235 (1988).

[5] K. Saito and A. W. Thomas, Phys. Lett. B 327, 9 (1994).

[6] X. Jin and B. K. Jennings, Phys. Rev. C 54, 1427 (1996); ibid Phys. Lett. B 374, 13 (1996).

[7] P.K. Panda, A. Mishra, J. M. Eisenberg and W. Greiner; Phys. Rev. C 56, 3134 (1997).

[8] I. Zakout and H. R. Jaqaman, nucl-th 9705034, submitted to Phys. Rev. C.

[9] J. Schechter, Phys. Rev. D 21, 3393 (1980); A. A. Migdal and M. A. Shifman, Phys. lett. B 114, 445 (1982).

[10] R. G. Rodriguez and J. I. Kapusta, Phys. Rev. C 44 870 (1991).

[11] I. Mishustin, J. Bondorf and M. Rho, Nucl. Phys. A 555 215 (1993).

[12] E. K. Heide, S. Rudaz and P. Ellis, Phys. lett. B 293, 259 (1992).

[13] E. K. Heide, S. Rudaz and P. J. Ellis, Nucl. Phys. A571, 713 (1994).

[14] G. W. Carter and P. J. Ellis Nucl Phys. A628, 325 (1998).

[15] G. Kälbermann, J. M. Eisenberg and B. Svetitsky, Nucl. Phys. A600, 436 (1996).

[16] R. J. Furnstahl and B. D. Serot, Phys. Rev. C47 2338 (1993).

[17] S. Fleck, W. Bentz, K. Shimizu and K. Yazaki, Nucl. Phys. A 510, 731 (1990).

[18] R. J. Furnstahl and B. D. Serot, Phys. Rev. C 41, 262 (1990).
FIGURES

FIG. 1. The dilaton scale $\chi = \frac{\phi}{\phi_0}$ for nuclear matter as a function of the baryon density $\rho_B$ with a dilaton potential constant $\epsilon_{\text{vac}} = (250 \text{MeV})^4$ for various values of temperature.

FIG. 2. The effective nucleon mass $M^*_N$ for nuclear matter as a function of the baryon density $\rho_B$ with a dilaton potential constant $\epsilon_{\text{vac}} = (250 \text{MeV})^4$ for various values of temperature.

FIG. 3. The bag radius $R$ for nuclear matter as a function of the baryon density $\rho_B$ with a dilaton potential constant $\epsilon_{\text{vac}} = (250 \text{MeV})^4$ for various values of temperature.

FIG. 4. The scalar mean field $\sigma$ for nuclear matter as a function of the baryon density $\rho_B$ with a dilaton potential constant $\epsilon_{\text{vac}} = (250 \text{MeV})^4$ at various temperatures.

FIG. 5. The dilaton scale $\chi = \frac{\phi}{\phi_0}$ for a nuclear matter as a function of the baryon density $\rho_B$ for several values of the dilaton potential constant $\epsilon_{\text{vac}}$, (a) for cold nuclear matter, (b) for hot nuclear matter at $T = 200 \text{MeV}$.

FIG. 6. The effective nucleon mass $M^*_N$ for nuclear matter as a function of the baryon density $\rho_B$ for several values of $\epsilon_{\text{vac}}$, (a) for cold nuclear matter, (b) for hot nuclear matter at $T = 200 \text{MeV}$.

FIG. 7. The bag radius, $R$, as a function of the baryon density $\rho_B$ for several values of $\epsilon_{\text{vac}}$, (a) for cold nuclear matter, (b) for hot nuclear matter at $T = 200 \text{MeV}$.

FIG. 8. The scalar mean field $\sigma$ as a function of baryon density $\rho_B$ for several values of $\epsilon_{\text{vac}}$, (a) for cold nuclear matter, (b) for hot nuclear matter at $T = 200 \text{MeV}$.

FIG. 9. The energy per nucleon for cold nuclear matter as a function of baryon density $\rho_B$ with different values for $\epsilon_{\text{vac}}$. 
\[ \chi = \left( \frac{\phi}{\phi_0} \right)^{0.5} \]

\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]

**Figure 1**

Cold nuclear matter
- 50 MeV
- 150 MeV
- 200 MeV
- 240 MeV

**Figure 2**

Cold nuclear matter
- 50 MeV
- 150 MeV
- 200 MeV
- 240 MeV

\[ M_N^* (\text{MeV}) \]

\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]
Figure 3

\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]

Cold nuclear matter
50 MeV
150 MeV
200 MeV
240 MeV

Figure 4

\[ \sigma = (250\text{MeV})^4 \]

Cold nuclear matter
50 MeV
150 MeV
200 MeV
240 MeV
\[ \chi = \frac{\phi}{\phi_0} \]

Cold nuclear matter

\[ \epsilon_{\text{vac}} = (200 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800 \text{MeV})^4 \]

Figure 5(a)

\[ \rho_B \text{ (fm}^{-3}) \]

\[ \chi = \left( \frac{\phi}{\phi_0} \right) \]

Temperature 200 MeV

\[ \epsilon_{\text{vac}} = (200 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800 \text{MeV})^4 \]

Figure 5(b)
Cold nuclear matter

\[ \epsilon_{\text{vac}} = (200\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800\text{MeV})^4 \]

![Fig. 6(a)](image)

Temperature 200 MeV$^4$

\[ \epsilon_{\text{vac}} = (200\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350\text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800\text{MeV})^4 \]

![Fig. 6(b)](image)
Fig. 7(a)

Cold nuclear matter

\[ \epsilon_{\text{vac}} = (200 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800 \text{ MeV})^4 \]

Fig. 7(b)

Temperature 200 MeV

\[ \epsilon_{\text{vac}} = (200 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350 \text{ MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800 \text{ MeV})^4 \]
Cold nuclear matter

\[ \epsilon_{\text{vac}} = (200\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (300\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (350\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (800\text{MeV})^4 \]

Temperature 200 MeV

\[ \epsilon_{\text{vac}} = (200\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (250\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (300\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (350\text{MeV})^4 \]

\[ \epsilon_{\text{vac}} = (800\text{MeV})^4 \]
Fig. 9

Cold nuclear matter

\[ \epsilon_{\text{vac}} = (200 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (250 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (300 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (350 \text{MeV})^4 \]
\[ \epsilon_{\text{vac}} = (800 \text{MeV})^4 \]

\[ E_{\text{tot}} - M_N \]
\[ (\text{MeV})^4 \]

\[ \rho_B \text{ (fm}^{-3}\text{)} \]

19