The connection between the nuclear matter mean-field equation of state and the quark and gluon condensates at high density.

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Abstract: It is known now that chiral symmetry restoration requires the meson-nucleon coupling to be density dependent in nuclear-matter mean-field models. We further show that quite generally, the quark and gluon condensates in medium are related to the trace of energy-momentum tensor of nuclear matter and in these models the incompressibility, K, must be less than thrice the chemical potential, \( \mu \). In the critical density \( \rho_c \), the gluon condensate is only reduced by 20% indicating a larger effective nucleon mass.

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Most of the models used in nuclear physics do not have explicit chiral invariance, being based on effective Lagrangians. However we have shown [1] that some of these so-called QHD models by virtue of the density-dependent coupling constants produce $\langle \bar{q}q \rangle_\rho \to 0$ at high $\rho$ and $T$. These are the Zimanyi-Moskowski models [2].

We inspect the density variation of the condensate in hadron models. These models yield different effective nucleon mass - its variation with density is also model dependent. The pivotal question is the restoration of chiral symmetry at high density: is there any correlation between the nuclear matter equation of state (EOS) and the change of the condensates in the medium?

In [1] we have shown that in the linear Walecka model [3], the ratio of the quark condensate of the nucleon in the medium, to that of the vacuum,

$$R_{\bar{q}q} = \frac{\langle N|\bar{q}q|N\rangle_\rho}{\langle 0|\bar{q}q|0\rangle_0}$$

(1)

goes down linearly upto 1.5 times the nuclear matter density but then tends to increase (see Fig.1), contrary to expectations based on ideas of chiral symmetry restoration. It could be argued that one should not push the Walecka model beyond its range of validity and explore it at high density. However, the same kind of result is obtained by Li and Ko [4] very recently from Bonn potential in the relativistic Dirac-Brueckner approach. The problem is not with the mean field approach which should work better at high density. Li and Ko also found the results unpalatable and conclude that effects like dependence of meson - nucleon coupling constants on the current quark mass may be crucial in obtaining a reliable result for the density dependence of the quark condensate. Indeed we found [1] contrary result for variants of the Walecka model given by Zimanyi and Moszkowski (see [2]). These models are labelled by ZM and ZM3 where the scalar and the scalar-vector mesons interact non-linearly and couplings are $\rho$ - dependent. In ZM and ZM3, the effective nucleon remains massive but $\langle \bar{q}q \rangle$ goes to zero at high $\rho$ (in Fig. 1 we only show Walecka and ZM3 models. The results are similar for ZM [1]). Thus we were forced to conclude that the increase of $|\langle \bar{q}q \rangle|$ in the Walecka model is not a fault of the QHD approach or the mean field approximation but due to neglect of non-linear coupling [1].

The present paper generalizes our previous work and gives the following important results: (1) We show that it is possible to relate the condensate ratio, $R_{\bar{q}q}$ (eq.1), to the trace of the energy momentum tensor of nuclear matter, $\mathcal{E} - 3P$, where $\mathcal{E}$ is the energy density and $P$ the pressure of the fluid. We remind the reader that for a non-interacting fluid $\mathcal{E} > 3P$ [5] but in presence of non-electromagnetic interaction this is not necessarily true; for example in the Walecka model the violation is due to strong vector meson interaction. The condition, $\mathcal{E} > 3P$, however must be satisfied if $R_{\bar{q}q}$ is to decrease, rather than increase. Of course we subscribe to the common belief that increase of the condensate ratio with density is unphysical. (2) In interacting nuclear models we
must not only have \( \mathcal{E} > 3P \) but to ensure that the ratio of the condensate \( R_{\bar{q} q} \) should fall monotonically we must have a relation satisfied between their derivatives, \( \partial \mathcal{E} / \partial \rho \equiv \mu \) and \( \partial P / \partial \rho \equiv K/9 \); namely \( K < 3 \mu \), \( K \) being the nuclear matter incompressibility. (3) The relation that we find to relate \( R_{\bar{q} q} \) and \( \mathcal{E} - 3P \) is a quite general expression for mean-field relativistic nuclear matter models. This is also valid for the non-linear Walecka Model and still for chiral symmetric models as the chiral linear \( \sigma \)-\( \omega \) model. (4) This relation, at the critical density (when the quark condensate vanishes meaning a start of the quark-gluon plasma phase), yields to the right bag constant found from phenomenology [6]. (5) We obtain expressions for gluon condensates in medium, relating this to changes of the nucleon mass at high densities.

QCD, has a non-trivial vacuum with large non-perturbative condensates of quarks and gluons. The non-zero value of the quark condensate is due to the breaking of approximate chiral symmetry by the vacuum, otherwise enjoyed by the QCD Hamiltonian, by virtue of the smallness of the quark mass \( m_q \):

\[
H_{\text{QCD}} = H_0 + 2m_q \bar{q}q
\]

the major part of the above being the chirally symmetric \( H_0 \). The Hellmann-Feynman theorem relates the shift of the quark condensate from its vacuum value to the nucleon sigma term and up and down quark masses [4]:

\[
2m_q \int d^3x (\langle N|\bar{q}q|N\rangle_N - \langle 0|\bar{q}q|0\rangle_0) = 2m_q \int d^3x (\Delta q)_N = \sigma_N = m_q \frac{dM_N}{dm_q}.
\]

where \( |N\rangle \) is the free nucleon at rest and \( (\Delta q)_N = \langle N|\bar{q}q|N\rangle_N - \langle 0|\bar{q}q|0\rangle \). This term can be looked upon as the "mass defect" of the nucleon due to \( m_q \neq 0 \) : \( \sigma_N = M_N - M_N^0 \), \( M_N^0 \) being the contribution from \( H_0 \) towards \( M_N \).

Using also the Hellmann-Feynman theorem and the relation of Gell-Mann, Oakes and Renner, one can show [4] that the condensate ratio in medium to its vacuum value is

\[
R_{\bar{q} q} = 1 - \rho \frac{\sigma_{\text{eff}}}{m^2_\pi f^2_\pi}
\]

with \( f_\pi = 93 \text{ MeV} \). And

\[
\sigma_{\text{eff}} = \frac{m_q}{\rho} \frac{d\mathcal{E}}{dm_q} = \frac{\sigma_N}{\rho} \frac{d\mathcal{E}}{dM_N}
\]

is the effective \( \sigma \)-commutator for a nucleon in the nuclear medium. The assumption is that the \( m_q \) dependence of \( \mathcal{E} \) comes through \( M_N \). If we neglect \( \delta \mathcal{E} \), the nucleon kinetic and interaction energy density in \( \mathcal{E} = \rho M_N + \delta \mathcal{E} \) we get \( \sigma_{\text{eff}} = \sigma_N \).

Using the expression for the energy density in the mean field approximation (MFA) and calculating the derivatives in eq.(5), we obtain a unified expression:
\[ R_{qq} = 1 - \frac{\sigma_N}{m_{\pi}^2 f_{\pi}^2} \left[ \frac{m_{\sigma}^2}{g_{\sigma}^2 M} (M_N - M_N^*) + (1 + \alpha) \frac{m_{\sigma}^2}{g_{\sigma}^2 M} (M_N - M_N^*)^2 - (1 + \beta) \frac{g_{\omega}^2}{m_{\omega}^2 M} \rho^2 \right], \]  
(6)

where the models in terms of \( \alpha \) and \( \beta \) are described as: Walecka \((\alpha = \beta = 0)\), ZM \((\alpha = 1 \text{ and } \beta = 0)\) and ZM3 \((\alpha = \beta = 1)\) [1], [2].

To our surprise we find that the expression between parenthesis in eq.(6) is equal to \((E - 3P)/M_N\) which (combining eqs. 4 and 5) can be written as

\[ \frac{dE}{dM_N} = \frac{(E - 3P)}{M_N}. \]  
(7)

Therefore, a general expression for the quark condensate in the medium arises:

\[ \rho \sigma_{\text{eff}} = 2m_q (\langle N | \bar{q}q | N \rangle - \langle 0 | \bar{q}q | 0 \rangle) = 2m_q (\Delta q)_\rho = \frac{\sigma_N}{M_N} (E - 3P) \]  
(8)

or

\[ R_{qq} = 1 - \frac{\sigma_N}{m_{\pi}^2 f_{\pi}^2} \frac{(E - 3P)}{M_N}. \]  
(9)

Further, to connect the change of \( R_{qq} \) with \( \rho \) we obtain

\[ \frac{\partial R_{\bar{q}q}}{\partial \rho} = \frac{1}{3} \frac{\sigma_N}{m_{\pi}^2 f_{\pi}^2} \frac{(K - 3\mu)}{M_N}. \]  
(10)

Thus the behaviour of the condensate is controlled not by the effective nucleon mass but by the EOS. For a softer EOS there is cancellation of the vector and the scalar fields. This makes \( E > 3P \) and the condensate ratio decreases. Monotonic decrease, however, depends on the derivative of the EOS. From eq.(10) we find that \( K < 3\mu \equiv 3(E + P)/\rho \) makes the decrease monotonic. Otherwise the medium enhances the chiral breaking yielding models in which this happens to become unphysical. In other words, this means that the coupling constants must depend on \( \rho \) to give a meaningful and soft EOS valid at high \( \rho \). Since \( \mu \) is directly determined from \( \rho \) the condition \( K < 3\mu \) is the concrete criterion of softness for the EOS. This is one of the central results of this paper.

Stiff EOS is not desirable in the context of phenomenology of neutron stars. Also it leads to very low effective nucleon mass and high nuclear matter incompressibility. We find the absence of increasing the condensate ratio as another criterion to rule out stiff EOS.

The models we used so far do not have chiral symmetry to start with and this is not very satisfactory. To test whether chiral symmetric models may reveal something different, we calculate the condensate in the chiral linear \( \sigma - \omega \) model [3]. Again we find...
\[
\frac{d\mathcal{E}}{dM_N} = \frac{(\mathcal{E} - 3P)}{M_N} = \left[ \frac{m_\sigma^2 (M_N^2 - M_N^*^2)}{g_\sigma^2 2M_N} - \frac{g_\omega^2}{m_\omega^2 M_N} \rho^2 \right].
\] (11)

The vector field, responsible for the \(\rho^2\) term above, controls the condensate at high \(\rho\). Above the density, when \(M^* \to 0\), the condensate ratio goes up! It seems a common feature that in any model, whether chirally symmetric or not, the vector field (if not reduced by the density dependence of the coupling constant) makes the EOS stiff and the condensate ratio increases.

Let us analyse the eq. (8), connecting the trace of energy momentum tensor, \(T_\mu^\nu\), via trace anomaly of QCD and a hydrodynamical model of nuclear matter. Consider the latter first. Embedding the nucleon in the infinite medium and treating the resulting nuclear matter as perfect relativistic hydrodynamical fluid, the energy-momentum tensor \(T^\mu_\nu\) (as we are dealing with a uniform and isolated system with no heat flow-vector) is given by:

\[
T^\mu_\nu = (\mathcal{E} + P) u^\mu u^\nu - Pg^\mu_\nu.
\] (12)

where \(\mathcal{E}\) is the proper energy density, \(P\) the hydrostatic pressure and \(u^\mu\) the four-velocity of the fluid [9].

Write the fluid four-velocity vector \(u^\mu\) in terms of the three-velocity as

\[
u^\mu = \gamma (1, \vec{v}), \gamma = \left(1 - \vec{v}^2\right)^{-1/2}
\] (13)

with \(u_\mu u^\mu = 1\). In the frame where nuclear matter is at rest (comoving frame) \(u^\mu = (1, 0)\) the energy-momentum tensor is diagonal

\[
\langle T^{\rho\rho} \rangle_\rho = \mathcal{E}, \quad \langle T^{ij} \rangle = P \delta_{ij},
\] (14)

the expectation value of the trace of the energy-momentum tensor \(T^\mu_\rho\) in nuclear matter from its vacuum value becomes

\[
\langle T^\mu_\rho \rangle_\rho - \langle T^\mu_\rho \rangle_0 = (\mathcal{E} - 3P).
\] (15)

To connect with the microscopic picture we turn our attention to a recent paper by Ji [11], the "QCD analysis of the Mass Structure of the Nucleon". The energy-momentum tensor is separated into a traceless \(\tilde{T}^{\mu\nu}\) and a trace \(\hat{T}^{\mu\nu}\) part contributing to \(M_N\) as \((3/4)M_N\) and \((1/4)M_N\) respectively. The trace part is a well known expression (in the leading order and neglecting the anomalous dimension of the mass term):

\[
\hat{T}^{\mu\nu} = \frac{1}{4} g^{\mu\nu} [\bar{\psi} \gamma^\rho \psi] - \frac{9\alpha_s}{8\pi} G^a_\mu \cdot G^{a\rho\nu}.
\] (16)
where the second term is the anomaly term. Thus we have

$$\frac{M_N}{4} = \langle N | H_a | N \rangle + (\sigma_N + S)/4 \,, \quad (17)$$

where $H_a$ is the anomaly Hamiltonian and $S$ is the strangeness content of the nucleon, defined by $S = \int d^3x \ m_s (\langle N | \bar{s}s | N \rangle - \langle 0 | \bar{s}s | 0 \rangle) \equiv \int d^3x \ m_s (\Delta s)_N$.

The divergence of the dilatation current, $\partial_\mu J^\mu_{\text{dil}}$, is equal to the trace of $\hat{T}^{\mu\nu}$ (eq. 16) which in the light sector becomes:

$$T^\mu = -\frac{9\alpha_s}{8\pi} G^\mu G^{\alpha\nu} + m_u \bar{u}u + m_d \bar{d}d + m_s \bar{s}s \,. \quad (18)$$

It vanishes in the classical level for QCD in the chiral limit. Its space part vanishes for a localised physical state such as $|N>$:

$$\int d^3x \langle N | T^\mu_{\text{dil}} | N \rangle = \int d^3x \langle N | \partial_0 J^0_{\text{dil}} - \partial_i j^i_{\text{dil}} | N \rangle = \int d^3x \langle N | \partial_0 j^0_{\text{dil}} | N \rangle \quad (19)$$

so that we get back $M_N$,

$$\int d^3x \langle N | T^\mu | N \rangle = \int d^3x \langle N | T^\sigma | N \rangle = M_N. \quad (20)$$

For a nucleon in the matter, this space part of $\partial_\mu J^\mu_{\text{dil}}$ does not vanish because of its infinite extent and we have the full trace on both sides. Considering the translational invariance the volume of a single nucleon in the medium is $1/\rho$ which appears on both sides and we have

$$(\mathcal{E} - 3P) = (-9/8)(\Delta G)_\rho + 2m_q(\Delta q)_\rho + m_s(\Delta s)_\rho \quad (21)$$

where $(\Delta G)_\rho \equiv \langle N | \frac{\alpha_s}{\pi} G^\mu G^{\alpha\nu} | N \rangle_\rho - \langle 0 | \frac{\alpha_s}{\pi} G^\mu G^{\alpha\nu} | 0 \rangle$. Similar to the fraction of the up and down quark mass term to $M_N$, we have

$$b_1 = \frac{\sigma_N}{M_N}, \quad (22)$$

if we define for the nucleon in the matter a fraction $b_2$,

$$b_2 = \frac{\sigma_{\text{eff}}}{(\mathcal{E} - 3P)/\rho} = \frac{2m_q(\Delta q)_\rho}{(\mathcal{E} - 3P)} \,. \quad (23)$$

and impose the restriction, $b_2 \equiv b_1$, so that

$$\rho \sigma_{\text{eff}} = 2m_q(\Delta q)_\rho = \frac{\sigma_N}{M_N}(\mathcal{E} - 3P) \quad (24)$$
we arrive at the general expression given by (eq. 8). The mean-field hadronic model calculations imply above relation factoring out the medium dependence only in the trace energy of the medium.

Using this trace anomaly , one can relate the shift of the gluon condensate in medium (from its vacuum value)[10], [7]. Combining the eqs.(17 and 20) we obtain for the free nucleon at rest:

\[ M_N = \int d^3x (-\frac{9}{8})(\Delta G)_N + \sigma_N + S = 4 \langle N|H_a|N \rangle + \sigma_N + S \]  

(25)

where \((\Delta G)_N\) is the shift of the gluon condensate in free nucleon state relative to vacuum.

Analogously, defining a ”gluon sigma” term \(G_N\) given by

\[ G_N = \frac{8}{9}(M_N - \sigma_N - S) = -\frac{32}{9} \langle N|H_a|N \rangle \]  

(26)

we find that in the two limits of \(m_s\) (150 MeV and larger ), studied in [11], where \(S + \sigma_N = 110\) or 160 MeV respectively, the corresponding values for \(G_N\) are \(-737\) or \(-693\) MeV. For \(S = 0\), the \(G_N = -795\) MeV.

Following the same factorization as that of the light quarks we can write the shift in strange quark condensate in the medium as

\[ \rho S_{\text{eff}} = m_s(\Delta \rho)_s = \frac{S}{M_N}(\mathcal{E} - 3P) \]  

(27)

where the density dependence comes again by the trace energy and \(S\) is the strangeness content of the nucleon, the precise value of which is a subject of some controversy [11]. We obtain only a range for this change depending on the estimation of this quantity, and in our calculations we consider \(S = 0\) for which the shift in the gluon condensate is maximum.

The expressions for the quark condensates and the eqs.( 21 and 26 ) can be used to extract the ”effective gluon sigma term” , \(G_{\text{eff}}\), for the nucleon embedded in the medium :

\[ \rho G_{\text{eff}} = \frac{G_N}{M_N}(\mathcal{E} - 3P) \]  

(28)

The eq.(4) shows that at a certain critical density \(\rho_c\), it may happen that \(R_{qq} \to 0\). This corresponds \(\mathcal{E} - 3P = 4(171MeV)^4\). It is to be noted that even at this density the gluon condensate in the nucleon is substantial :

\[ \langle N|\bar{\alpha}s\pi^aG_{\mu\nu}G^{a\mu\nu}|N \rangle \rho = (350)^4 - 3.4 (171)^4 \sim (332MeV)^4, \]  

(29)

which is only a reduction of 20% , as we can see from fig.2. In normal density this reduction is even smaller and about 5% using the vacuum value \(\langle 0|\bar{\alpha}s\pi^aG_{\mu\nu}G^{a\mu\nu}|0 \rangle = (350MeV)^4\).
Gluon contribution to the nucleon mass is more than the quark contribution (510 MeV, as against 430 MeV, \[^{[11]}\]) and gluons seem to persist even at high density. Quark part of nucleon mass in the medium might change but the gluonic part remains. This explains why even at high densities the nucleon mass can not be small and ZM3 model "effectively" reproduces this (this result is reproduced in the non-linear Walecka model which also gives a good value for the nuclear matter incompressibility).

Although the derivation of the eq.\(^{(24)}\) is very general- uses only the trace of the energy momentum tensor and the mean field approach which should work better at high density, when applied to models we find that the condensate can behave very differently. However, the following comments can be made with regard to some model independent results:

(1) At the nuclear saturation density \(\rho_0\), in any hadron model, we shall have the same value of the condensate. The reason is obvious. At this \(\rho_0\) the \(P = 0\) and \(E\) is made to be the same. Actually, with the binding energy of -15.75 MeV and \(M_N = 939\) MeV the condensate ratios become, for example, \(R_{\bar{q}q} = 0.69\) and \(R_{\bar{g}g} = 0.95\) for \(\sigma_N = 45\) Mev and S=0.

(2) There exists a critical density \(\rho_c\) for which the condensate goes to zero. At this density \(E - 3P = m_\pi^2 f_\pi^2 M_N / \sigma_N\). This will imply a bag constant \(B^{1/4} = 171 MeV\). Notice this is the optimum bag constant if one fits all the known hadrons in the manner of Aerts and Rafelski \[^{[6]}\]. After \(\rho_c\) it is not realistic to picture nucleons as point particles. The quark structure is almost inevitable at this density.

(3) If there is scaling for the condensate with the effective nucleon mass, \(M^*_N\) vanishes when the condensate vanishes. This is awkward for nuclear physicists since the meaning of nuclei or nuclear matter composed of zero mass particles is ill-defined. It is known that for a relativistic gas composed of non interacting particles with rest mass \(M_N\), \(E - 3P = \rho M_N (1 - \frac{v^2}{c^2})^{1/2}\) where \(v\) is the mean velocity of the particles \[^{[1]}\]. This means that for small densities \(E - 3P \sim \rho M_N\) and we obtain the usual relations for the quark and gluon condensates in the leading order approximation \[^{[11]}\] and \[^{[7]}\]. However, when the density and pressure increase \((E - 3P) < \rho M_N\). If we define an effective mass \(M^*_N = M_N (1 - \frac{v^2}{c^2})^{1/2}\), \(M^*_N\) becomes very light when \(E \sim 3P\). But this leads to \(R_{\bar{q}q} \to 1\) and \(R_{\bar{g}g} \to 1\) i.e. the condensates resume their vacuum values which is opposite to our notion of chiral symmetry restoration. Thus we stress that light effective nucleon mass, rather than restoring the chiral symmetry, will enhance the symmetry breaking.

(4) Putting all the eqs. \[^{[24]}\], \[^{[27]}\] and \[^{[28]}\), we can write the effective quantities for the nucleon in the medium as

\[
\rho \xi_{eff} = \frac{\xi_N}{M_N} (E - 3P)
\]

where \(\xi_{eff}\) stands for the \(\sigma_{eff}\), \(G_{eff}\) or \(S_{eff}\). As long as the \((E - 3P)/M_N < \rho\) and positive, these effective quantities will decrease in the medium. In Fig.3 we show that they are
not far from the leading order in the ZM3 model but in the Walecka model they become unphysical. This is another illustration of the breakdown of the Walecka model, where the $M_N^*$ is very light and the condensates increase at high density.

(5) The condition $K < 3\mu$ allows us to conclude that the pressure can not increase fast with the density (small incompressibility $K$ at high $\rho$), which manifests that a soft EOS at high density is needed to expect that the medium will enhance chiral symmetry and the restoration occurs (the case of ZM3 model). This conclusion is also satisfied if we include the temperature effects, where now $\mu \equiv \partial \Omega / \partial \rho$ where $\Omega$ stands for the grand-canonical potential.

It is important to point out that in this density analysis of the condensates our hadronic models do not have pions in 2nd order. In fact the pionic corrections will be there - but as commented and observed in the reference [12] - a calculation with linear sigma model - their net effect should be small.

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Figure Captions

Fig. 1: Ratio of the up-and-down quark condensate to the vacuum value and the effective nucleon mass ratio $M^*/M$ as a function of density for Walecka and ZM3 models.
Fig. 2: Ratio of the gluon condensate to the vacuum value for the models with $S=0$.
Fig. 3: The trace energy of nuclear matter divided by the nucleon rest mass as a function of the density is plotted and compared with the leading order approximation given by $\rho/\rho_o$ (full line).
$\sigma_N = 45$

Walecka model

ZM3 model

$R_{\bar{q}q}$ vs $\frac{\rho}{\rho_0}$
Ratio of Gluon Condensate (S=0)

- Walecka
- ZM3
