Quark and gluon condensates in the quark-meson coupling model

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Abstract

Using the quark-meson coupling (QMC) model, we study the density dependence of the quark and gluon condensates in nuclear matter. We show that the change of the quark condensate is mainly driven by the scalar field in the medium and that the reduction of the quark condensate is suppressed at high density, even in the mean-field approximation. The gluon condensate decreases by 4 - 6 % at nuclear saturation density. We also give a simple relationship between the change of the quark condensate and that of a hadron mass in the medium.

PACS numbers: 24.85.+p, 21.65.+f, 24.10.Jv, 12.39Ba
Keywords: quark and gluon condensates, infinite nuclear matter, relativistic mean-field theory, quark degrees of freedom

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The QCD ground state is highly non-trivial, and the strong condensates of scalar quark-antiquark pairs $\langle \bar{q}q \rangle$ and gluon fields $\langle G^a_{\mu\nu}G^{a\mu\nu} \rangle$ may play important roles in a wide range of low-energy hadronic phenomena [1, 2, 3, 4]. Therefore, it is quite interesting to study the density dependence of the condensates in nuclear matter. The vacuum values of the lowest-dimensional quark and gluon condensates are typically given by [2]

\[ Q_0 \equiv \langle \bar{q}q \rangle_0 \simeq -(225 \pm 25\text{MeV})^3, \]  
\[ G_0 \equiv \langle G^a_{\mu\nu}G^{a\mu\nu} \rangle_0 \simeq (360 \pm 20\text{MeV})^4. \]  

Drukarev et al. [3], Cohen et al. [5] and Lutz et al. [6] have shown that the leading dependence on the nuclear density, $\rho_B$, of the quark condensate in nuclear matter, $Q(\rho_B)$, is given by the model-independent form:

\[ \frac{Q(\rho_B)}{Q_0} \simeq 1 - \frac{\sigma_N}{f^2\pi m^2} \rho_B, \]  

where $\sigma_N$ is the pion-nucleon sigma term (empirically $\sigma_N \simeq 45 \text{MeV}$ [7]), $m_\pi$ is the pion mass (138 MeV) and $f_\pi \simeq 93 \text{MeV}$, the pion decay constant. Further, the strange quark content in the nucleon at finite density (and temperature) was studied in Ref. [8] using the Nambu–Jona-Lasinio (NJL) model, supplemented by an instanton induced interaction involving the in-medium quark condensates. The gluon condensate at finite density, $G(\rho_B)$, has also been discussed in Ref. [5].

Several years ago M. Ericson [9] suggested that a “distortion factor”, coming from rescattering of soft pions in the nuclear medium, would tend to reduce the amount of the chiral symmetry restoration (i.e., to oppose the reduction of the quark condensate). However, Birse et al. [10] pointed out that there was an incompleteness in the treatment of the rescattering of soft pions using the simple linear $\sigma$ model, and showed that the full amplitude for soft-pion scattering from two nucleons leads to an enhancement of the chiral symmetry restoration at finite density.

Recently the in-medium quark condensate has been calculated in several, more elaborate ways [11, 12]. In particular, using the Dirac-Brueckner-Hartree-Fock (DBHF) approach, Li and Ko and Brockmann and Weise have shown that higher-order contributions...
become increasingly important at high density, and tend to hinder the restoration of chiral symmetry [11].

We study the density dependence of the quark and gluon condensates in nuclear matter within the framework of the quark-meson coupling (QMC) model [13, 14, 15, 16, 17]. The QMC model may be viewed as an extension of QHD [18] in which the nucleons still interact through the exchange of scalar ($\sigma$) and vector ($\omega$ and $\rho$) mesons. However, the mesons couple not to point-like nucleons but to confined quarks (in the nucleon bag). In studies of infinite nuclear matter it was found that the extra degrees of freedom provided by the internal structure of the nucleon give an acceptable value for the incompressibility once the coupling constants are chosen to reproduce the correct saturation energy and density for symmetric nuclear matter. This is a significant improvement on QHD at the same level of sophistication (see also Ref. [19]). Furthermore, the model has been successfully applied to finite nuclei within the Born-Oppenheimer approximation [15, 16, 20, 21]. It has been found that the QMC model can reproduce the properties of finite, closed-shell nuclei (from $^{12}$C to $^{208}$Pb) quite well.

As shown in Refs. [15, 16], the basic result in the QMC model with mean-field approximation (MFA) is that, in the scalar and vector meson fields, the nucleon behaves as if it had an effective mass $M_N^\star$. The latter can be calculated using a relativistic quark model for the nucleon (e.g., the MIT bag model) and depends on the nuclear density only through the $\sigma$ field.

In an earlier version of the QMC model [13, 20], we considered the effect of the nuclear medium on the structure of the nucleon alone and froze the quark degrees of freedom in the mesons. We call this version QMC-I. We have calculated the quark condensate in nuclear matter using this version [14, 22, 23]. (For a recent study, see also Ref. [24].) However, strictly speaking, those calculations (as well as the DBHF calculations [11]) were not complete because the meson structure effects were not treated consistently. The mesons themselves are built of quarks and anti-quarks, and their structure may also change in matter [16, 21]. An additional, technical difference from Refs. [14, 22, 23] is that for
reasons explained in Ref. [15] – see especially Appendix – the c.m. correction to the bag energy is now treated as being independent of the applied scalar field.

To incorporate the effect of meson structure in the QMC model in MFA, we suppose that the vector mesons are again described by the MIT bag model with common scalar and vector mean-fields (like the nucleon in QMC-I). In this case the effective vector-meson mass in matter, \( m_v^* \) (\( v = \omega, \rho \)), will also depend on the \( \sigma \) mean-field. The \( \sigma \) meson itself is, however, not so readily represented by a simple quark model (like a bag), because it couples strongly to two pions and a direct treatment of chiral symmetry in medium is important [10, 25]. On the other hand, many approaches, including the NJL model [4], the Walecka model [26] and Brown-Rho scaling [27] suggest that the \( \sigma \)-meson mass in medium, \( m_\sigma^* \), should be less than the free value, \( m_\sigma \) (= 550 MeV). It has therefore been parametrized as a quadratic function of the scalar field [16]:

\[
\left( \frac{m_\sigma^*}{m_\sigma} \right) = 1 - a_\sigma(g_\sigma\bar{\sigma}) + b_\sigma(g_\sigma\bar{\sigma})^2,
\]

with \( g_\sigma\bar{\sigma} \) in MeV. Here \( \bar{\sigma} \) is the mean-field value of the \( \sigma \) field and \( g_\sigma \) is the \( \sigma \)-nucleon coupling constant (in free space). Three parameter sets: \( (a_\sigma ; b_\sigma) = (3.0, 5.0 \text{ and } 7.5 \times 10^{-4} \text{ MeV}^{-1} ; 10, 5 \text{ and } 10 \times 10^{-7} \text{ MeV}^{-2}) \), called A, B and C, respectively, were determined in Ref. [16] so as to reduce the mass of the \( \sigma \)-meson by about 2%, 7% and 10% (respectively) at saturation density \( \rho_0 \) (= 0.15 fm\(^{-3}\)). This model, involving the structure effects of both the nucleon and the mesons, was called QMC-II [16].

Within QMC-II, the total energy per nucleon, \( E_{tot} \), can be written as [16]

\[
E_{tot} = \frac{2}{\rho_B(2\pi)^3} \sum_{i=p,n} \int^{k_{F_i}} d\vec{k} \sqrt{\vec{M}_i^*} + \vec{k}^2 + \frac{m_{\omega}^*}{2\rho_B} \vec{\omega}^2 + \frac{m_{\rho}^*}{2\rho_B} \vec{\rho}^2 + \frac{g_\omega^2}{8m_{\omega}^*} \rho_B + \frac{g_\rho^2}{8m_{\rho}^2} \rho_B^2.
\]

where \( g_v \) (\( v = \omega, \rho \)) is the \( v \)-nucleon coupling constant and \( k_{F_i} \) (\( i = \text{proton or neutron} \)) is the Fermi momentum. The density of protons (neutrons) \( \rho_p \) (\( \rho_n \)) is defined by \( \rho_i = k_{F_i}^3/(3\pi^2) \), and then \( \rho_B = \rho_p + \rho_n \) and \( \rho_3 = \rho_p - \rho_n \). Detailed values of the coupling constants and properties of nuclear matter in QMC-II can be found in Ref. [16].

The density-dependent quark condensate, \( Q(\rho_B) \), is formally derived by applying the Hellmann–Feynman theorem to the chiral-symmetry-breaking quark mass term of the
total Hamiltonian. One finds the relation for the quark condensate in nuclear matter at
the baryon density \(\rho_B\):

\[
m_q(Q(\rho_B) - Q_0) = m_q \frac{d}{dm_q} \mathcal{E}(\rho_B),
\]

where \(\mathcal{E}(\rho_B) = \rho_B E_{tot}\) and \(m_q\) is the average, current quark mass of the u and d quarks.

The resulting value for \(Q/Q_0\) as a function of density is shown in Fig. 1 – dashed line.
(We have chosen the quark mass to be \(m_q = 5\) MeV and the bag radius of the free nucleon \(R_N = 0.8\) fm, but the result is quite insensitive to these choices.)

Using Eq.(6), the Gell-Mann–Oakes–Renner relation \[28\] and the explicit expression
for the self-consistency condition of the \(\sigma\) field in nuclear matter \[16\], Eq.(5) leads to the
following explicit relation for the ratio of \(Q(\rho_B)\) to \(Q_0\):

\[
\frac{Q(\rho_B)}{Q_0} = 1 - \left( \frac{\sigma_N}{m^2_\pi f^2_\pi} \right) \frac{g^2_\sigma}{g^2_\pi} \left( \frac{m^*_\sigma}{g^2_\pi} \right) \left[ 1 - \frac{1}{g^2_\pi} \left( \frac{dg^2_\pi}{dm^*_\pi} \right) \frac{(g^2_\sigma)}{g^2_\pi} \left( g^2_\sigma \right) \right] \\
- \left( \frac{\sigma_N \rho_0}{6S_N(0)m^2_\pi f^2_\pi} \right) \rho^2 \left( \frac{\rho_0}{m^*_\pi} \frac{dg^2_\pi}{dm^*_\pi} + \frac{\rho_0}{4m^*_\pi} (2f_p - 1)^2 \frac{dg^2_\pi}{dm^*_\pi} \right),
\]

(7)

where \(\rho_r = \rho_B/\rho_0\), \(f_p = \rho_p/\rho_B\), \(S_N(0)\) is the quark scalar charge of the free nucleon
\((= \int_{bag} d\vec{r} \bar{\psi}_q \psi_q)\) and \(g^2_\sigma = g^2_\sigma/(3S_N(0))\) \[13\], \[16\]. The vector-meson mass is calculated
using the bag model, while the \(\sigma\)-meson mass is given by Eq.(4). Because \(m_q\) enters only
in the combination \(m_q - g^2_\sigma \bar{\sigma}\), which is generally regarded as the chiral-symmetry-breaking
term in nuclear medium, we were able to evaluate \(\left( \frac{dm^*_\pi}{dm^*_\pi} \right)\) in terms of the derivative of \(m^*_\sigma\)
with respect to the applied scalar field \(\bar{\sigma}\).

In Eq.(7) we have followed the usual convention of identifying \(3m_q S_N(0)\), which is the
sigma commutator in the free MIT bag, as the experimental pion-nucleon sigma term,
\(\sigma_N\) \[14\]. It is well known that the meson cloud of the nucleon (mainly the pions), as well
as its strange quark content contribute significantly to \(\sigma_N\) \[29\]. However, because we are
concerned primarily with the variation of \(Q\) in matter from its free value, \(Q_0\), it should
be reasonable to replace \(\sigma_N\) in Eq.(7) by its empirical value. (We note that the main
variation of \(Q\) in medium is generated by the \(\sigma\) mean-field.)

Clearly, from Eq.(7), the leading dependence of the quark condensate on the density
is given by the scalar field:

\[
\frac{Q(\rho_B)}{Q_0} \simeq 1 - \frac{\sigma_N}{m^2_{\pi}} \left( \frac{m_\sigma}{g_\sigma} \right)^2 (g_\sigma \bar{\sigma}).
\]  

(8)

One can easily show that Eq.(8) reduces to the model-independent result, Eq.(3), to leading order in the density, so that for small \( \rho_r \) one has (for the set B) [21] :

\[
\frac{Q(\rho_B)}{Q_0} \simeq 1 - 0.357 \rho_r.
\]  

(9)

This is shown as the dotted line in Fig. 1.

Equation (7) also involves deviations of the quark-meson coupling constants with respect to \( m_q \). In principle, if one could derive these coupling constants from QCD, their dependence on \( m_q \) would be given. Within the present model there is no reason to believe that the couplings should vary with \( m_q \). This is especially so for the vector couplings since they involve conserved vector currents. On the other hand, we require that our model reproduces the correct saturation energy and density of nuclear matter whatever parameters are chosen for the free nucleon. As a consequence, the coupling constants depend on \( m_q \) in a way that has nothing to do with chiral symmetry breaking. (For example, for set B, we find \( g^2_\sigma = 4.891 - 0.005880m_q + 1.200 \times 10^{-5}m_q^2 \), \( g_\omega^2 = 39.59 + 0.03828m_q + 1.144 \times 10^{-3}m_q^2 \) and \( g_\rho^2 = 66.3 - 0.02m_q \), with \( m_q \) in MeV.)

In order to extract a physically meaningful result for \( Q/Q_0 \) we should therefore remove the spurious contributions associated with \( \frac{dg_M}{dm_q} (M = \sigma, \rho, \omega) \) in Eq.(7). In fact, the variation of \( g_M^2 \) with \( m_q \) is extremely small so we need only correct the \( \omega \) and \( \rho \) contributions. The final, corrected result is shown as the solid line in Fig.1. Even in the mean-field approximation, our calculations show that the higher-order contributions in the nuclear density become very important and that they weaken the chiral symmetry restoration at high density (c.f. Ref. [11]). In QMC-II, the \( \sigma \) field in nuclear matter is suppressed at high density (for example, \( g_\sigma \bar{\sigma} \simeq 200 \) (300) MeV at \( \rho_0 \) (3\( \rho_0 \))) because the quark scalar charge, \( S_N \), decreases significantly as the density rises, as a result of the change in the quark structure of the bound nucleon [15, 16, 30]. Since the reduction of the quark condensate is mainly controlled by the scalar field, it is much smaller than in the
Figure 1: Quark condensate at finite density using parameter set B. The dashed and dotted curves are respectively for the full calculation in symmetric nuclear matter \( f_p = 0.5 \) and the linear approximation, Eq.(9). The solid curve is the result corrected by removing the spurious \( \omega \) and \( \rho \) contributions.

simple, linear approximation, Eq.(9). From the difference between the solid and dashed curves we see that the correction for the dependence of the coupling constants on \( m_q \) is significant and this should be born in mind in any phenomenological treatment.

We should note here that, from extensive studies of chiral perturbation theory for nuclear matter, especially the recent work of Birse [10, 25], a reduction of the quark condensate from its vacuum value may not be enough to conclude that the chiral symmetry has been partially restored – especially if part of the change in \( \langle \bar{q}q \rangle \) arises from low-momentum pions. We note also that higher-order condensates may play an increasingly important role as the quark condensate tends to zero.

Next let us consider the in-medium gluon condensate. Cohen et al. [5] also developed a model-independent prediction of the gluon condensate that is valid to first order in the nuclear density through an application of the trace anomaly and the Hellmann–Feynman theorem.
Following their approach, the ratio of the gluon condensate in nuclear matter, \( G(\rho_B) \), to that in vacuum \( (G_0) \) is given by

\[
G(\rho_B)/G_0 \simeq 1 - \left( \frac{8}{9G_0} \right) \left[ \mathcal{E}(\rho_B) - 2m_q(Q(\rho_B) - Q_0) - m_s(Q_s(\rho_B) - Q_{s0}) \right],
\]

where \( m_s \) is the strange-quark mass and \( Q_s(\rho_B) \) (\( Q_{s0} \)) is the strange-quark condensate in nuclear matter (in vacuum). Up to first order in the density, the change of the strange-quark condensate may be written in terms of the strange quark content of the nucleon in free space, \( S \):

\[
m_s(Q_s(\rho_B) - Q_{s0}) = S\rho_B + \mathcal{O}(\rho_B^2).
\]

The strange quark content is commonly specified by the dimensionless quantity, \( y \), defined by

\[
y \equiv \frac{2\langle \bar{s}s \rangle_N}{\langle \bar{u}u + \bar{d}d \rangle_N},
\]

which leads to \( S = (m_s/2m_q)\sigma_N y \). Roughly speaking, \( y \) represents the probability to find \( s \) or \( \bar{s} \) in the nucleon and is a measure of the OZI-rule violation. If \( m_s/m_q \simeq 25 \) [28] and \( y \simeq 0.45 \) [4], we get \( S \simeq 250 \) MeV. We note, however, that \( y \simeq 0.45 \) is an extreme value, and it has recently been suggested that it may be compatible with zero [31]. In the analysis below we shall take care to examine the sensitivity to the full range of variation of \( y \).

At very low \( \rho_B \), \( \mathcal{E}(\rho_B) \) can be expanded as [18]

\[
\mathcal{E}(\rho_B) = M_N \rho_B \left[ 1 + \frac{3}{10M_N^2} \left( \frac{3\pi^2}{2} \right)^{2/3} \rho_B^{2/3} \right] + \mathcal{O}(\rho_B^2),
\]

where the second term in the bracket is the nonrelativistic Fermi-gas energy. Using the approximate form, \( g_\sigma \bar{\sigma} \simeq 214 \) (MeV) \( \times \rho_r \) (for the parameter set B in \( m_\sigma^* \)) [21], we find

\[
2m_q(Q(\rho_B) - Q_0) = 214(\text{MeV}) \times \sigma_N \left( \frac{m_\sigma}{g_\sigma} \right)^2 \rho_r + \mathcal{O}(\rho_r^3).
\]

Choosing the central value of \( G_0 \) in Eq.(2), we then get the in-medium gluon condensate at low \( \rho_B \) (for the set B):

\[
G(\rho_B)/G_0 = 1 - (0.03892\rho_r + 0.001292\rho_r^{5/3}) + \mathcal{O}(\rho_r^2).
\]
Figure 2: Gluon condensate at finite density using parameter set B. The solid and dotted curves are respectively for the full calculation in symmetric nuclear matter ($f_p = 0.5$) and the approximation, Eq.(15).

Our numerical results for the full calculation as well as the approximate calculation with Eq.(15) are shown in Fig. 2 (for $m_q = 5$ MeV and $R_N = 0.8$ fm). The reduction of the gluon condensate at finite density is not large; for example, it is reduced by only 4% at $\rho_0$, which is consistent with the results of Cohen et al. [5]. The approximation of Eq.(15) works very well for a wide range of the nuclear density, which may imply that the effect of higher-order contributions (in powers of the density) is small for the gluon condensate. However, one should keep in mind that the in-medium gluon condensate evaluated here contains a large uncertainty, originating from the uncertainty in the value for the strange quark content of the nucleon in free space. We note that if we assume a vanishing strange quark content of the nucleon in free space ($S = 0$ or $y = 0$) [31], the gluon condensate would be reduced by about 6% at $\rho_0$.

Finally, we relate the quark condensate to the variation of the hadron mass in nuclear matter. In QMC-II, the hadron mass at low $\rho_B$ is simply given in terms of the scalar
field \([\sigma]\):

\[
M_j^* \simeq M_j - \frac{n_0}{3} (g_\sigma \sigma),
\]

(16)

where \(n_0\) is the number of non-strange quarks in the hadron \(j(= N, \omega, \rho, \Lambda, \text{etc.})\). Since the quark condensate at low \(\rho_B\) is also determined by the scalar field (see Eq.(3)), we find a simple relation between the variations of the hadron mass and the quark condensate:

\[
\delta M_j^* \simeq \left( \frac{m_\pi^2 f_\pi^2}{3 \sigma_N} \right) \left( \frac{g_\sigma}{m_\sigma} \right)^2 n_0 \left( 1 - \frac{Q(\rho_B)}{Q_0} \right) \approx 200(\text{MeV}) \times n_0 \left( 1 - \frac{Q(\rho_B)}{Q_0} \right),
\]

(17)

where \(\delta M_j^* = M_j - M_j^*\) (cf. Ref. [27]).

However, as shown in Ref. [25], we know that the nucleon mass in matter cannot depend in any simple way on the quark condensate alone because the leading non-analytic contribution (LNAC) to the pion-nucleon sigma term – the term of order \(m_\pi^3\) – should not appear in the nucleon-nucleon interaction [32]. To discuss this problem further, we have to include pions self-consistently in the QMC model, which is beyond the scope of the present work.

In summary, we have calculated the quark and gluon condensates in nuclear matter. In the QMC-II model, the quark condensate at finite density is given in terms of the scalar field in the medium and the variation of the coupling constants with respect to the quark mass. We have shown that the reduction of the quark condensate at high density is much less than that suggested by the (model-independent) leading-order prediction, Eq.(3), even in the mean-field approximation. We also point out that the need to correct a naive use of the Hellmann–Feynman theorem to calculate \(Q(\rho_B)\) for any purely phenomenological dependence of the quark-meson coupling constants on \(m_q\).

In comparison with the quark condensate, the gluon condensate does not decrease much in nuclear matter. We have also provided a simple relationship between the change of the quark condensate and that of the hadron mass in nuclear matter. We should notice here that the validity of our model is limited to low and moderate density (probably less than \(\sim 3 \rho_0\)), because the short-range correlations between quarks in overlapping hadron bags have been ignored. The effect of the pion cloud of the hadrons [10, 25, 33] should
also be considered explicitly in any truly quantitative study of the condensate properties in the medium.

We would like to thank M. Birse, P. Guichon and M. Ericson for helpful comments on the issues discussed here during the Workshop on Hadrons in Dense Matter held at the CSSM. This work was supported by the Australian Research Council and the Japan Society for the Promotion of Science.
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