Chiral density wave in nuclear matter

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Inspired by recent work on inhomogeneous chiral condensation in cold, dense quark matter within models featuring quark degrees of freedom, we investigate the chiral density-wave solution in nuclear matter at zero temperature and nonvanishing baryon number density in the framework of the so-called extended linear sigma model (eLSM). The eLSM is an effective model for the strong interaction based on the global chiral symmetry of quantum chromodynamics (QCD). It contains scalar, pseudoscalar, vector, and axial-vector mesons as well as baryons. In the latter sector, the nucleon and its chiral partner are introduced as parity doublets in the mirror assignment. The eLSM simultaneously provides a good description of hadrons in vacuum as well as nuclear matter ground-state properties. We find that an inhomogeneous phase in the form of a chiral density wave is realized, but only for densities larger than $2.4 \rho_0$, where $\rho_0$ is the nuclear matter ground-state density.

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Introduction: The spontaneous breaking of chiral symmetry in the QCD vacuum is a nonperturbative phenomenon which has to be reflected in low-energy hadronic theories, see e.g. Refs. [1–3]. The order parameter of chiral symmetry breaking is the chiral condensate, denoted as $\langle \bar{q}q \rangle \sim \langle \sigma \rangle$, which contributes to hadronic masses and is responsible for the mass splitting of so-called chiral partners, i.e., hadrons with the same quantum numbers except for parity and G-parity.

At sufficiently large temperature and density, it is expected that the spontaneously broken chiral symmetry is (at least partially) restored. Lattice-QCD calculations [4,5] show that, for values of the quark masses realized in nature, this so-called chiral transition is cross-over along the temperature axis of the QCD phase diagram. Along the density axis, lattice-QCD calculations are not yet available (for realistic quark masses), but phenomenological models [6,7] indicate that chiral symmetry restoration may occur through a first-order phase transition.

An interesting possibility is that the effective potential is minimized by an order parameter which varies as a function of spatial coordinate. Such inhomogeneous phases were already suggested in the pioneering works of Ref. [8–19]. In particular, the chiral condensate may assume the form of the so-called chiral-density wave where not only the chiral condensate $\langle \sigma \rangle$ but also the expectation value of the neutral pion field is nonvanishing, $\langle \pi^0 \rangle \neq 0$. However, the problem of the aforementioned approaches was that, without nucleon-nucleon tensor forces, inhomogeneous chiral condensation took place already in the nuclear matter ground state, in contradiction to experimental findings.

More recently, inhomogeneous phases were studied in the framework of the (1+1)-dimensional Gross-Niveu model [20, 21], where it was indeed found that a spatially varying order parameter minimizes the effective potential at high density. In Ref. [22, 23], the authors coined the phrase “quarkyonic matter” for an inhomogeneous phase at high density, where the chiral density-wave solution is realized within QCD in the large–$N_c$ limit. Inhomogeneous phases were also investigated in Refs. [24–27] in the framework of the Nambu–Jona-Lasinio model as well as the quark-meson model.

In this work, we re-investigate the question whether inhomogeneous condensation at nonzero density occurs in a model solely based on the degrees of freedom of the QCD vacuum, i.e., hadrons. We employ dilatation invariance and chiral symmetry of QCD, and include scalar, pseudoscalar, vector, and axial-vector mesons, as well as baryons and their chiral partners. This approach, developed in Refs. [28–30] and denoted as extended Linear Sigma Model (eLSM), successfully describes hadron vacuum phenomenology both in the meson [28, 29] and baryon [30] sectors. In the latter, the nucleon and its chiral partner are treated in the mirror assignment [31, 32], in which a chirally invariant mass term exists [see also Refs. [33–37] and refs. therein].

The chiral condensate $\langle \sigma \rangle$ is the expectation value of the $\sigma$ field which is the chiral partner of the pion $\pi$. In the framework of the eLSM the resonance corresponding to the $\sigma$ field is not the lightest scalar resonance $f_0(500)$ (as proposed in older versions of the $\sigma$ model), but the heavier state $f_0(1370)$. This result is in agreement with a variety of studies of low-energy QCD, e.g. Refs. [38, 39] and refs. therein. In Ref. [40], $f_0(500)$ was coupled in a chiral linear sigma model, both resonances $f_0(500)$ and $f_0(1370)$ should be taken into account, the former being the lightest scalar state and the latter being an excitation of the chiral condensate. This was, for instance, done in Ref. [36] where, in the framework of the eLSM, $f_0(500)$ was coupled in a chirally invariant manner to nucleons and their chiral partners. An important result of this study was that the nuclear matter ground-state properties (i.e., density, bind-
ing energy, and compressibility) could be success-fully described. In the mean-field approximation, and assuming homogeneous condensates, Ref. 30 reports the onset of a first-order phase transition at a density of about 2.5\rho_0. The important role of both aforementioned scalar resonances has been also investigated at nonzero temperature in the framework of a simplified version of the eLSM in Ref. 44. Interestingly, the necessity to include both scalar-isoscalar states f_0(500) and f_0(1370) has been also shown in the framework of the Bonn nucleon-nucleon po-tential 45.

The main question of the present work is whether in-homogeneous condensation takes place within the eLSM. For simplicity, we restrict ourselves to a spatial dependence of the condensate, which is of the type of the chiral-density wave: \langle \sigma \rangle \sim \cos(2fx), \langle \pi^3 \rangle \sim \sin(2fx). We will show, by using the parameters determined in Ref. 36, to describe nuclear matter ground-state properties (and thus having no additional free parameters), that nuclear matter in the ground state is still a (homogeneous) liq-uid. Chiral condensation remains homogeneous up to a chemical potential of 973 MeV (corresponding to a den-sity of 2.4\rho_0), followed by an inhomogeneous phase with a chiral density wavelength \pi/f of approximately 1.5 fm.

The model: In the two-flavor case, N_f = 2, the scalar and pseudoscalar sectors are described by the matrix

\Phi = (\sigma + i\eta_N)\psi_0 + (\bar{\psi}_0 + i\bar{\phi}) \cdot \vec{t},

where \vec{t} = \vec{\tau}/2, with the vector of Pauli matrices \vec{\tau}, and \psi_0 = 1_2/2. The vector and axial-vector mesons enter via the matrices

V^\mu = \omega^\mu \psi_0 + \vec{\rho}^\mu \cdot \vec{t}, \quad A^\mu = f_0^\mu \psi_0 + \vec{a}_0^\mu \cdot \vec{t},

from which the left-handed and right-handed vector fields are defined as R^\mu \equiv V^\mu - A^\mu, L^\mu \equiv V^\mu + A^\mu. Under the chiral group SU(2)_R \times SU(2)_L the fields transform as \Phi \rightarrow U_l \Phi U^\dagger_r, R^\mu \rightarrow U_l R^\mu U^\dagger_r, and L^\mu \rightarrow U_l L^\mu U^\dagger_r. The identification of mesons with particles listed in Ref. 16 is as follows: the fields \vec{t} and \eta_N correspond to the pion and the nonstrange part of the \eta meson, \eta_N \equiv (\bar{\eta}u + \bar{\eta}d)/\sqrt{2}. The fields \omega^\mu and \vec{\rho}^\mu represent the vector mesons \omega(782) and \rho(770), and the fields \vec{f}_0^\mu and \vec{a}_0^\mu the axial-vector mesons f_1(1285) and a1(1260). The scalar fields \sigma and \bar{\sigma}_0 fields are identified with f_0(1370) and a0(1450), re-spectively. The chiral condensate \phi = \langle \sigma \rangle = Zf_\pi emerges upon spontaneous chiral symmetry breaking in the mesonic sector, where f_\pi \approx 92.4 MeV is the pion decay constant and Z \approx 1.67 is the wave-function renormalization constant of the pseudoscalar fields 28.

In the present work, besides ordinary quarkonium mesons, also the lightest resonance f_0(500) is introduced 36. For N_f = 2, it does not matter whether we interpret f_0(500) as (predominantly) a tetraquark field or as a pion-pion resonance; differences in the coupling to other fields occur, however, for N_f \geq 3. The bare f_0(500) field is denoted as \chi. It is a singlet under chiral transformation and is coupled to mesons following Refs. 43, 44.

We first consider the mesonic part of the eLSM Lagrangian and keep only terms involving fields which will eventually condense, i.e., \sigma, \pi = \pi^3, \omega_\mu, and \chi [for the full Lagrangian, see Ref. 28];

\[ L_{\text{mes}} = \frac{1}{2} \partial{}_{\mu}\sigma \partial{}^{\mu}\sigma + \frac{1}{2} \partial{}_{\mu}\pi \partial{}^{\mu}\pi - \frac{1}{4} (\partial{}_{\mu}\omega_\nu - \partial{}_{\nu}\omega_\mu)^2 + \frac{1}{2} \partial{}_{\mu}\chi \partial{}^{\mu}\chi + \frac{m^2}{2} (\sigma^2 + \pi^2) - \frac{\lambda}{4} (\sigma^2 + \pi^2)^2 + \varepsilon \sigma + \frac{1}{2} m_\omega^2 \omega_0^2 - \frac{1}{2} m_\chi^2 \chi^2 + g \chi (\sigma^2 + \pi^2). \] (1)

The following numerical values are used 28, 43: \[ m^2 = (896.3 \text{ MeV})^2, \lambda = 35.05, \varepsilon = 1.054 \cdot 10^6 \text{ MeV}^3, g = 438 \text{ MeV}, m_\omega = 782 \text{ MeV}, \text{and} m_\chi = 611 \text{ MeV}. As a consequence of this choice of parameters, m_\sigma = 1295 \text{ MeV} and \sigma is identified with f_0(1370).

We now make the following Ansatz for the condensates, which is of the form of a chiral density wave:

\[ \langle \sigma \rangle = \phi \cos(2fx), \langle \pi \rangle = \phi \sin(2fx), \] (2)

In the limit f \rightarrow 0 we obtain the usual homogeneous condensation which is realized in the vacuum and, as we shall see, for low densities. Besides these inhomogeneous condensates, also \chi and \omega_0 develop nonvanishing con-densates which are, however, homogeneous: \langle \chi \rangle = \bar{\chi} and \langle \omega_0 \rangle = \bar{\omega}_0.

In mean-field approximation, we insert the Ansatz 24 into the mesonic part of the Lagrangian 11 and obtain the tree-level potential

\[ U_{\text{mean-field}}^{\text{mes}} = 2f^2 \phi^2 + \frac{\lambda}{4} \phi^4 - \frac{1}{2} m^2 \phi^2 - \varepsilon \phi \cos(2fx) - \frac{1}{2} \frac{m_\omega^2 \omega_0^2 + 1}{2} m_\chi^2 \chi^2 - g \chi \phi^2. \] (3)

Note that, in a spatial volume V > (\pi/f)^3, the spatially dependent term \sim \cos(2fx) averages to zero for any nonzero value of f.

We now turn to the baryonic sector where we introduce two doublets \Psi_1 and \Psi_2 transforming according to the mirror assignment:

\[ \Psi_{1,R} \rightarrow U_R \Psi_{1,R}, \quad \Psi_{1,L} \rightarrow U_L \Psi_{1,L}, \] (4)
\[ \Psi_{2,R} \rightarrow U_R \Psi_{2,R}, \quad \Psi_{2,L} \rightarrow U_L \Psi_{2,L}. \] (5)

The mirror assignment has the consequence that the following mass term is allowed by chiral symmetry 30–32, 37:

\[ m_0 (\bar{\Psi}_{1,L} \Psi_{2,R} - \bar{\Psi}_{1,R} \Psi_{2,L} - \bar{\Psi}_{2,L} \Psi_{1,R} + \bar{\Psi}_{2,R} \Psi_{1,L}). \] (6)

However, in order to preserve dilatation invariance for the baryon sector, the constant m_0 should emerge upon condensation of scalar fields. These fields can be assigned as a scalar-isoscalar glueball or a tetraquark (molecular) state 30, 36; here we only consider for simplicity the
latter possibility, i.e., the coupling of baryons to the field \( \chi \). The resulting baryonic Lagrangian in which only the mesons \( \sigma, \omega, \) and \( \chi \) are retained reads:

\[
\mathcal{L}_{\text{bar}} = \nabla_\mu \bar{\Psi}_1 \gamma^\mu \Psi_1 + \nabla_\mu \bar{\Psi}_2 \gamma^\mu \Psi_2 - \frac{\hat{g}_1}{2} \nabla_\mu \bar{\Psi}_1 \left( \sigma + i \gamma_5 \tau^3 \pi \right) \Psi_1 - \frac{\hat{g}_2}{2} \nabla_\mu \bar{\Psi}_2 \left( \sigma - i \gamma_5 \tau^3 \pi \right) \Psi_2 - g_\omega \bar{\Psi}_1 \gamma^\mu \omega \Psi_1 - g_\omega \bar{\Psi}_2 \gamma^\mu \omega \Psi_2 - a \chi \left( \bar{\Psi}_2 \gamma_5 \Psi_1 - \bar{\Psi}_1 \gamma_5 \Psi_2 \right),
\]

(7)

where the parameters \( \hat{g}_1, \hat{g}_2, g_\omega, \) and \( a \) are dimensionless, in accordance with dilatation invariance. The mass term \( m_0 \) of Eq. (6) emerges upon condensation of the tetraquark field: \( m_0 = a \chi \). The choice of parameters which reproduces nuclear matter ground-state properties is [30]: \( \hat{g}_1 = 10.80, \hat{g}_2 = 18.53, g_\omega = 5.04, \) and \( a = 17.92 \) (which implies \( m_0 = 500 \text{ MeV} \)). For these values the compressibility is \( K = 226.9 \text{ MeV} \). This is consistent with the range \( 200 - 300 \text{ MeV} \) quoted in Refs. 47, 48. In particular, our result is in good agreement with the value \( K = 235 \pm 14 \text{ MeV} \) obtained in Ref. 49 by studying inelastic scattering of \( \alpha \) particles off nuclei.

When studying the chiral density-wave Ansatz of Eq. [2], it is useful to make the following field redefinitions of the baryon fields [27]:

\[
\Psi_1 \rightarrow \Psi_1 \exp \left( -i \gamma_5 \tau^3 f x \right), \quad \Psi_2 \rightarrow \Psi_2 \exp \left( +i \gamma_5 \tau^3 f x \right),
\]

thanks to which the explicit spatial coordinate dependence transforms into a momentum dependence. In the effective potential, we treat the meson fields in the mean-field approximation, while the fermions are integrated out. In the no-sea approximation for the fermions, the effective potential then reads

\[
U_{\text{eff}}(\phi, \bar{\chi}, \bar{\omega}_0, f) = \sum_{k=1}^{4} \int \frac{2d^3p}{(2\pi)^3} \left[ E_k(p) - \mu^* \Theta[\mu^* - E_k(p)] + U_{\text{mean-field}} \right],
\]

(8)

where \( \mu^* = \mu - g_\omega \bar{\omega}_0 \) and \( E_k(p) = \sqrt{p^2 + m_k(p_x)^2} \). In the case of homogeneous condensation there are only two (two-fold degenerate) energy eigenstates, corresponding to the nucleon and its chiral partner, while for inhomogeneous condensation the degeneracy is lifted and four different energy eigenstates emerge. The values for \( E_k(p) \) are calculated numerically as solutions of characteristic polynomials.

**Results:** The effective potential is numerically minimized with respect to \( \phi, \bar{\chi}, \) and \( f \) and maximized with respect to \( \bar{\omega}_0 \), respectively. In Fig. 1 we show the effective potential as a function of \( \phi \) (at the extrema of \( \bar{\chi}, \bar{\omega}_0, \) and \( f \)) for the chemical potential \( \mu = 923 \text{ MeV} \), which corresponds to the nuclear matter ground state. There are two degenerate global minima, one for \( \phi = 154.3 \text{ MeV} \) corresponding to the vacuum, and one for \( \phi = 149.5 \text{ MeV} \) corresponding to the nuclear matter ground state, respectively. The decrease of the chiral condensate as compared to the vacuum is very small, which is a consequence of the pseudoscalar wave-function renormalization \( Z = 1.67 > 1 \) [and thus, indirectly, of the presence of (axial-)vector mesons; for results where (axial-)vector mesons were not taken into account, see Ref. 50]. Moreover, we also notice the presence of a local minimum at \( \phi = 38.3 \text{ MeV} \), which corresponds to inhomogeneous condensation. For increasing \( \mu \) the position of this minimum changes only slightly with \( \phi \), but it eventually becomes the global minimum and thus the thermodynamically realized state.

In Fig. 2 the condensates \( \phi \) and \( \bar{\chi} \) are shown as functions of \( \mu \). For \( \mu = 923 \text{ MeV} \) a first-order phase transition to the nuclear matter ground state takes place.
Both condensates drop and then further decrease slowly for increasing \( \mu \). At \( \mu = 973\text{ MeV} \) a transition to the inhomogeneous phase occurs. The condensate \( \overline{\chi} \) drops to (almost) zero and the chiral condensate \( \phi \) to the value \( \phi = 37.6\text{ MeV} \). For larger \( \mu \) the condensates \( \overline{\chi} \) and \( \phi \) change very slowly. Note that \( \phi \) does not vanish, thus chiral symmetry is not completely restored. In terms of density, the onset of inhomogeneous condensation is at \( 2.4\rho_0 \); this density is still small enough such that a hadronic description of the system is valid. Then, a mixed phase is realized between \( 2.4\rho_0 \) to \( 10\rho_0 \). This latter value is too large to trust in a hadronic description, thus we are led to believe that somewhere in the mixed phase the deconfinement phase transition should occur.

The onset of the chiral density-wave phase can be best shown plotting the behavior of the parameter \( f \), see Fig. 3: \( f \) vanishes for small \( \mu \), but jumps to a nonzero value \( (f = 389.5\text{ MeV}) \) at the critical value \( \mu = 973\text{ MeV} \) and then slightly increases for increasing \( \mu \). This is the chiral density-wave phase, with a one-dimensional harmonic modulation with a wavelength of about 1.5 fm.

**Conclusions:** We have shown that in the framework of a chiral hadronic approach (the eLSM) an inhomogeneous chiral condensate becomes favoured at large density. Using a set of parameters which allows for a correct description of the nuclear matter ground-state properties, we find that the onset of the inhomogeneous phase occurs at \( 2.4\rho_0 \). To our knowledge, our result is the first demonstration of inhomogeneous chiral condensation in an approach capable of describing vacuum as well as nuclear matter properties. In future studies one should go beyond the chiral density-wave Ansatz and perform a more detailed analysis of the phases realized in the system. At large densities, one cannot avoid to include strange degrees of freedom.

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