Nonlocality effects on Color Spin Locking condensates

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We consider the color spin locking (CSL) phase of two-flavor quark matter at zero temperature for nonlocal instantaneous, separable interactions. We employ a Lorentzian-type form factor allowing a parametric interpolation between the sharp (Nambu-Jona-Lasinio (NJL) model) and very smooth (e.g., Gaussian) cut-off models for systematic studies of the influence on the CSL condensate the deviation from the NJL model entails. This smoothing of the NJL model form factor shows advantageous features for the phenomenology of compact stars: (i) a lowering of the critical chemical potential for the onset of the chiral phase transition as a prerequisite for stability of hybrid stars with extended quark matter cores and (ii) a reduction of the smallest pairing gap to the order of 100 keV, being in the range of values interesting for phenomenological studies of hybrid star cooling evolution.

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I. INTRODUCTION

Recently, the investigation of color superconducting phases in cold dense quark matter has received much attention $^{1, 2, 3, 4, 5, 6}$, in particular due to the possible consequences for the physics of compact stars $^{7, 8}$. From the point of view of observational constraints on quark matter and color superconductivity in compact stars the cooling characteristics play a central role. It has been shown that the occurrence of a normal quark matter core would lead to a conflict with observations since the direct Urca (DU) process in normal quark matter would lead to enhanced cooling in disagreement with the data $^{8, 10}$. The DU conflict would be solved provided no ungapped quark modes occur in the quark core. This has been demonstrated on the example of a hypothetical pairing channel (X-gap) for the quark color which is ungapped in the 2SC phase (2SC+X phase) $^{10, 11}$. However, the microscopic origin of the X-gap could not yet be specified. A microscopically well-defined pairing pattern which could solve the quark DU problem would be the CSL phase $^{12}$ corresponding to a spin-one condensate $^{12, 13, 14, 15}$. A prerequisite for the realization of this pairing pattern in quark matter would be a sufficient flavor asymmetry to prevent the u-d pairing in the otherwise dominant scalar diquark channel of the 2SC phase. It has been demonstrated that under neutron star conditions the 2SC phase is indeed rather fragile and may not be realized for moderate coupling strengths $^{10}$. Thus the CSL phase becomes particularly interesting for the solution of the quark DU cooling problem, and corresponding simulations will be performed as soon as the cooling regulators such as emissivities, specific heat and thermal conductivity will be provided. First steps in this direction have been made recently $^{17, 18}$.

Most of the calculations of QCD superconducting phases have been done using the sharp cut-off NJL model (see Ref. $^1$ and references therein). However, lattice QCD calculations $^{19}$ indicate that quark interactions should act over a certain range in the momentum space, and various approaches to include nonlocality effects beyond the NJL model have been suggested $^{20}$. We refer to nonlocal separable interaction models as introduced, e.g., in the works $^{21, 22, 23, 24, 25, 26}$ and references therein, where it has been concluded that smoothing the cutoff leads to a reduction of the chiral condensate and a lowering of the critical temperature for the chiral phase transition. The question arises for the effects of nonlocality on the spin-one gaps, to be explored by varying the form factor of the quark interaction from a sharp cutoff in the NJL model to smoothly decreasing form such as a Gaussian. First exploratory calculations reported in $^{27}$ have shown that the nonlocality could lead to a sizeable
decrease of the energy gaps.

In this paper, we investigate the robustness of CSL pairing against a modification of the sharp cut-off (NJL) in a systematic way by employing a separable, instantaneous interaction with a Lorentzian-type interaction which allows to interpolate between the NJL case and very smooth interaction form factors of, e.g., the Gaussian type. This investigation is performed on the basis of recently developed parameterizations for the instantaneous three flavor case [23].

II. NONLOCAL CHIRAL QUARK MODEL FOR THE COLOR-SPIN LOCKING (CSL) PHASE

We investigate a nonlocal chiral quark model in which the quark interaction is represented in a separable way by introducing form factor functions \(g(p)\) in the bilinear of the current-current interaction terms in the Lagrangian [16, 21, 29]. It is assumed that this four-fermion interaction is instantaneous and therefore the form factors do not depend on the energy but only on the modulus of the three momentum \(p = |p|\). The ansatz for the s-wave, single flavor diquark condensate characterizing the CSL phase as introduced in Ref. [12] is a scalar product (locking) of the three-vector of antisymmetric color matrices \((\lambda_2, \lambda_5, \lambda_7)\) with the three-vector of Dirac spin matrices \((\gamma_3, \gamma_2, \gamma_1)\). Thus, the corresponding gap matrix \(\hat{\Delta}\) for the CSL phase reads

\[
\hat{\Delta} = \Delta(\gamma_3\lambda_2 + \gamma_2\lambda_5 + \gamma_1\lambda_7) .
\]

Since the two flavor channels decouple, the quark thermodynamical potential can be decomposed into single-flavor components

\[
\Omega_q(T, \{\mu_f\}) = \sum_{f=u,d} \Omega(T, \mu_f) ,
\]  
and it is sufficient to consider in the following the contribution of a single flavor only, which in the mean field approximation is given by

\[
\begin{align*}
\Omega(T, \mu) &= \frac{\phi^2}{8G} + 3\frac{\Delta^2}{8H_v} \\
&- T \sum_n \int \frac{d^3p}{(2\pi)^3} \frac{1}{2} \text{Tr} \left( \frac{1}{T} S^{-1}(i\omega_n, \bar{p}) \right) .
\end{align*}
\]

where \(\mu\) stands for the chemical potential of that flavor. The first two terms are quadratic contributions of the mean field values \(\phi\) and \(\Delta\) of the order parameter fields that signal chiral symmetry breaking and CSL superconductivity, respectively. Their denominators contain the coupling constants \(G\) and \(H_v\) in the corresponding channel. In the sum is over fermionic Matsubara frequencies \(\omega_n = (2n + 1)\pi T\), and the trace is over Dirac, color and Nambu-Gorkov indices.

In our nonlocal extension, the inverse fermion propagator differs from the NJL model case by momentum dependent form factors \(g(p)\) modifying the mesonic and diquark mean fields

\[
S^{-1}(p) = \begin{pmatrix} \rho + \mu^2 & 0 & g(p)\hat{\Delta} \\
0 & -g(p)\hat{\Delta} & 0 \\
-\rho - \mu^2 & 0 & M(p) \end{pmatrix}
\]  

where \(M(p)\) is the dynamical quark mass function

\[
M(p) = m + g(p)\phi .
\]

Note that although the first two terms in Eq. (3) do not have any explicit dependence on the form factors, the quantities \(\phi\) and \(\Delta\) do depend implicitly on them through the gap equations, see Eq. (14) below.

After evaluation of the trace [12] and Matsubara summation the thermodynamical potential takes the form

\[
\Omega(T, \mu) = \frac{\phi^2}{8G} + 3\frac{\Delta^2}{8H_v} \\
- \frac{6}{T} \sum_k \int \frac{d^3p}{(2\pi)^3} \left[ E_k(p) + 2T \ln(1 + e^{-E_k(p)/T}) \right] ,
\]

where \(E_k(p)\) denote the excitation energies for the modes \(k = 1 \ldots 6\). The odd (even) indices denote particle (antiparticle) excitations corresponding each to a triplet of spin-one eigenstates. All modes have a gap in the excitation spectrum and can be brought into a standard form, which for \(E_1(p)\) reads

\[
E_1^2(p) = (\varepsilon_{\text{eff}}(p) - \mu_{\text{eff}}(p))^2 + \Delta_{\text{eff}}^2(p) ,
\]

with the effective quantities

\[
\begin{align*}
\varepsilon_{\text{eff}}(p) &= \sqrt{p^2 + M_{\text{eff}}^2(p)} , \\
M_{\text{eff}}(p) &= \mu_{\text{eff}}(p) M(p) , \\
\mu_{\text{eff}}(p) &= \mu \sqrt{1 + \Delta^2 g^2(p)/\mu^2} , \\
\Delta_{\text{eff}}(p) &= M(p) \mu_{\text{eff}}(p) \Delta g(p) .
\end{align*}
\]

and for \(E_{3,5}(p)\) is given by

\[
E_{3,5}^2(p) = (\varepsilon(p) - \mu)^2 + a_{3,5}(p) \Delta g^2(p) ,
\]

with the momentum-dependent coefficients

\[
a_{3,5}(p) = \frac{1}{2} \left[ 5 - \frac{p^2}{\varepsilon(p)\mu} \pm \sqrt{\left( 1 - \frac{p^2}{\varepsilon(p)\mu} \right)^2 + 8\frac{M^2(p)}{\varepsilon^2(p)}} \right] .
\]
where \(\varepsilon(p) = \sqrt{p^2 + M^2(p)}\). The remaining modes \(E_{2,4,6}(p)\) are obtained from \(E_{1,3,5}(p)\) by changing \(\mu \to -\mu\) in Eqs. (7)-(13). Note that the modes \(E_{1,2}(p)\) correspond to the vanishing z-projection of the spin, \(S_z = 0\), thus being inert against an external B-field. The remaining modes corresponding to \(S_z = \pm 1\) are expected to get shifted (Zeeman effect).

For given values of \(T\) and \(\mu\), the global minimum of \(\Omega(T, \mu)\) in the space of the order parameters \(\phi\) and \(\Delta\) corresponds to the thermodynamical equilibrium state. We obtain this state by comparing solutions of the gap equations

\[
\frac{\delta \Omega(T, \mu)}{\delta \phi} = \frac{\delta \Omega(T, \mu)}{\delta \Delta} = 0.
\]

We present results for the case of vanishing temperature and finite chemical potential in the next section.

III. MODEL CALCULATIONS

A. Form factors and their parameters

In [4] and [5] we have introduced the same form factors \(g(p)\) to represent the nonlocality of the interaction in the meson (qq) and diquark (qq) channels. In our calculations we use the sharp cutoff (NJL), Lorentzian with integer parameter \(\alpha\) (L\(\alpha\)) and Gaussian (G), form factors defined as

\[
g_{\text{NJL}}(p) = \theta(1 - p/\Lambda) , \tag{15}
\]

\[
g_{\text{L}\alpha}(p) = [1 + (p/\Lambda)^{2\alpha}]^{-1}, \quad \alpha \geq 2 , \tag{16}
\]

\[
g_{\text{G}}(p) = \exp(-p^2/\Lambda^2) . \tag{17}
\]

where \(\Lambda\) is a cut-off parameter. These form factors are plotted in Fig. 1. We achieve deviations from the NJL case (step function) by using the Lorentzian form factor with decreasing the \(\alpha\) parameter. The Gaussian form factor appears on the other limit having a very soft momentum dependence.

To perform numerical calculations one has to specify, for each form factor, the following set of parameters: the light quark current mass \((m)\), the coupling strength \((G)\) and the range of the interaction \((\Lambda)\). The diquark coupling constant \(H_v\) is fixed to the ratio \(H_v/G = 8/3\) in accordance with the result of the Fierz transformation for a one-gluon exchange interaction. In this work we use the parameterizations recently given in Ref. [28] and listed in Table I. They have been obtained by fitting the vacuum properties of the pion \((f_\pi = 92.4\ \text{MeV}, \ M_\pi = 135\ \text{MeV})\) and the vacuum constituent quark mass at zero momentum, \(M = m + \phi\). For the latter phenomenologically reasonable values \(M = 330\) and 400 MeV are used.

Note that the results to be presented below do not depend on the choice of the Lorentzian-type function as interpolating form factor. In fact, similar results have been obtained using other interpolating functions as, e.g., the Woods-Saxon form factors the parameterization of which is given in [28].

B. Quark mass and CSL pairing gap

First, we analyze form factors which do not deviate strongly from the NJL case, i.e. \(L\alpha\) for \(\alpha \geq 3\), shown as the grey area in Fig. 1. In Fig. 2 we compare the solutions obtained for the mass and the CSL gaps for two different sets of regularizations for fixed constituent mass: \(M = 330\ \text{MeV}\) (left) and \(M = 400\ \text{MeV}\) (right). For the parameterizations with a larger constituent mass

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**TABLE I: Parameter sets for the nonlocal chiral quark model**

| \(M\) [MeV] | Form factor | \(\Delta\) [MeV] | \(G\) | \(\Lambda\) | \(m\) [MeV] |
|---|---|---|---|---|---|
| 330 | NJL | 629.5 | 2.17 | 5.28 |
|   | L10 | 649.2 | 2.36 | 4.71 |
|   | L5  | 666.5 | 2.49 | 4.09 |
|   | L3  | 685.8 | 2.59 | 3.25 |
|   | L2  | 703.4 | 2.58 | 2.37 |
|   | L5  | 735.4 | 2.58 | 2.37 |
|   | L3  | 756.1 | 2.58 | 2.37 |
| 400 | NJL | 587.9 | 2.44 | 5.58 |
|   | L10 | 600.3 | 2.64 | 5.01 |
|   | L5  | 609.3 | 2.78 | 4.39 |
|   | L3  | 616.2 | 2.87 | 3.55 |
|   | L2  | 617.8 | 2.83 | 2.65 |
|   | L5  | 649.2 | 2.58 | 2.37 |
|   | L3  | 756.1 | 2.58 | 2.37 |

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in vacuum, one obtains a larger critical chemical potential $\mu_c$ for the phase transition from the chirally broken phase to the restored one, where the CSL pairing can occur. The gaps at the onset, $\Delta(\mu_c)$, are larger whereas the mass gaps after the chiral transition are smaller.

Like in the NJL case [12], the CSL gaps are strongly increasing functions of $\mu$ in the range that is relevant for compact stars, $\mu_c < \mu < 500$ MeV, where the upper limit is due to the threshold for the occurrence of strange quarks which allow pairing patterns like the CFL phase, more favorable than CSL (for recent phase diagrams for neutral matter in the three flavor case see [30, 31, 32]).

On the other hand, Fig. 2 clearly shows that the pairing gaps in this nonlocal extension are reduced relative to the NJL ones: the smoother the form factor, the smaller the gap. The reduction could be up to a factor three in the case of Lorentzian with $\alpha = 3$.

However, the qualitative behaviour of the chiral and CSL gaps is not affected by the choice of the form factors.

To obtain the above results, we have kept fixed the ratio $\frac{H_v}{G}$ at the standard value obtained from Fierz transforming the one-gluon exchange interaction. However, in order to estimate the effect of possible uncertainties in this value, we have considered also the situation in which this ratio is taken to be twice its Fierz value. The corresponding results for $2H_v$ and all Lorentzian-type form factors under consideration are shown in Fig. 4. It is worth noticing that while the increase of $H_v$ by a factor 2 increases the CSL gaps by one ($M = 400$ MeV) or two ($M = 330$ MeV) orders of magnitude, the qualitative
we show the quasiparticle excitation spectrum in the diquark channel is doubled with respect to the usual value coming from Fierz transformed one gluon exchange interaction.

C. Quasiparticle excitation spectrum

In Fig. 5 we show the quasiparticle excitation spectrum at the critical chemical potential for both parameterizations, respectively. We observe that the quasiparticle mode with the lowest energy band corresponds is $E_1(p)$ with a minimum

$$E_{1\text{,min}} = \min_p \{E_1(p)\},$$

being the most relevant quantity for possible applications of the CSL phase of quark matter to compact star cooling phenomenology. This minimum occurs at the Fermi momentum $p = p_F$. In a very good approximation $p_F$ can be represented by the lowest orders of a series expansion in the parameter $s = p_F g_T^{\alpha\lambda}(p_F)/g_{\alpha\lambda}(p_F)$, which is a measure for the influence of the form factor

$$p_F^2 = \mu_{\text{eff}}^2(p_F) - M_{\text{eff}}^2(p_F)$$

$$+ 2\Delta_{\text{eff}}^2(p_F) \frac{M(p_F)[M(p_F) - m]}{M^2(p_F) - M_{\text{eff}}^2(p_F)/\mu^2} s$$

$$+ O(s^2).$$

In the same order of the expansion in $s$, we obtain for the minimal excitation energy

$$E_{1\text{,min}} = \Delta_{\text{eff}}(p_F) + O(s^2).$$

Although $E_{1\text{,min}}$ might be quite small (see below) it never vanishes. In fact, as in the NJL case [12], in the present class of models none of the dispersion relations lead to gapless modes. It is interesting to note that this lowest energy mode $E_1(p)$, relevant for compact star cooling phenomenology, is the one which is inert against the influence of a strong external magnetic field typical for neutron stars since it belongs to vanishing spin projection, $S_z = 0$.

In Fig. 4 and Fig. 7 we plot $E_{1\text{,min}}$ as a function of $\mu$ for different models from the NJL-like and the smooth form factor groups, respectively. The calculations are made for both sets of parameterizations of Table 1, corresponding to constituent masses of $M = 330$ MeV and $M = 400$ MeV. According to the analytical approximative result of Eq. (20), the behavior of the minimal excitation energies can be understood as a product of the increasing $\mu$-dependence of the CSL pairing gaps and the decreasing one of the other factors in Eq. (11). For the parameterizations with $M = 400$ MeV on the right panel we obtain that $E_{1\text{,min}}$ is a decreasing function of $\mu$ since the increase in $\Delta(\mu)$ cannot overcompensate the decrease in $M(p_F)g(p_F)/\mu_{\text{eff}}(p_F)$. A similar effect has been reported for NJL models [12]. On the other hand, for parameterizations with $M = 330$ MeV the interplay between $\Delta(\mu)$
and the other factors in Eq. (11) is density dependent: a slightly increasing behavior of $E_{1,\text{min}}$ at low densities is followed by a tendency to a saturation or even decreasing behavior at high densities. For the group of NJL-like form factors, our results for $E_{1,\text{min}}$ lie in the range of 50 – 500 keV and for the case of smooth form factors they are between 1 and 100 keV.

![FIG. 6](image)

FIG. 6: Minimal excitation energies $E_{1,\text{min}}$ in the CSL phase as functions of the chemical potential $\mu$ for different NJL-like form factors. Parameterizations correspond to fixed $M = 330$ MeV on the left and $M = 400$ MeV on the right.

Our main results are summarized in Fig. 8. In the upper panel we show the scaling of the critical chemical potential $\mu^L_{c}$ for the onset of the CSL phase for the Lorentztian-type model $L_\alpha$ with the smoothness parameter $1/\alpha$ normalized to the NJL limit case, $\mu^\text{NJL}_{c}$. In the lower panel we show a comparative plot of the minimal excitation energies $E_{1,\text{min}}^L$ in units of the corresponding NJL counterpart $E_{1,\text{min}}^{\text{NJL}}$, evaluated at $\mu^L_{c}$ and $\mu^\text{NJL}_{c}$, respectively. The corresponding results for $2H_v$ are also shown in Fig. 8 as open symbols. As we see the qualitative behaviour of both $\mu^L_{c}/\mu^\text{NJL}_{c}$ and $E_{1,\text{min}}^L/E_{1,\text{min}}^{\text{NJL}}$ as a function of $1/\alpha$ remains unchanged.

It is remarkable that, as it would be expected from an expansion to lowest order in $s$, both quantities scale almost linearly with $1/\alpha$. In fact, a very good approximation to our numerical results is obtained with

$$\mu^L_{c} \approx \left(1 - \frac{\xi'}{\alpha}\right)\mu^\text{NJL}_{c},$$

(21)

$$E_{1,\text{min}}^L \approx \left(1 - \frac{\xi}{\alpha}\right)E_{1,\text{min}}^{\text{NJL}},$$

(22)

for $\alpha$ down to 2, where the slope parameters $\xi$ and $\xi'$ do only moderately depend on the model parameterization $(M)$ and the CSL coupling strength $(H_v)$. For $M = 330$ MeV we get $\xi = 1.8$ (1.5) and $\xi' = 0.13$ (0.12), while for $M = 400$ MeV the corresponding values are $\xi = 1.6$ (1.2), $\xi' = 0.17$ (0.16). The numbers in parentheses are obtained by doubling $H_v$.

**IV. CONCLUSION**

We have studied the effect of instantaneous nonlocal interactions in the color spin locking (CSL) phase of quark matter. We have introduced momentum dependent form factors to model the nonlocality and compared systematically with the local NJL counterpart.

We have shown that there is a systematic lowering of the critical chemical potential for the onset of the CSL phase as well as for the minimal excitation energy (effective CSL gap) as a function of the nonlocality which can be represented as a linear dependence on the smoothness parameter $1/\alpha$ of the Lorentz-type form factor. These qualitative effects are shown to be robust under changes in the coupling constant used to represent the CSL interaction.

It has been found that hybrid star cooling requires all quark modes to be paired with a minimal pairing of the order of $10 – 100$ keV to suppress the direct Urca process in quark matter. The present model for the CSL phase meets this requirement and calls for a more detailed analysis of the cooling phenomenology based on this microscopically justified pairing pattern.

The smallest gap which governs the cooling phenomenology corresponds to the zero $z$-projection of the
spin and thus remains unaffected by the external magnetic field of a compact star. Moreover, the CSL pairing pattern is a flavor singlet and insensitive to the flavor asymmetry in a compact star under β-equilibrium.

Therefore, the CSL phase with nonlocal instantaneous interactions is particularly interesting for applications in compact stars and allows to achieve a suitable description of quark matter properties by choosing the appropriate form factor models. Although it remains to be shown that hybrid star configurations with the CSL quark matter phase could be stable, due to the small gaps, we expect to have results reproducing those of the normal quark matter case, where stable quark matter cores in nonlocal models have been found.

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