Phase structure of cold magnetized color superconducting quark matter

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Abstract. The influence of intense magnetic fields on the behavior of color superconducting cold quark matter is investigated using an SU(2)\textsubscript{f} NJL-type model for which a novel regulation scheme is introduced. In such a scheme the contributions which are explicitly dependent on the magnetic field turn out to be finite and, thus, do not require to be regularized. As a result of this, non-physical oscillations that arise from regularizing magnetic field dependent terms are naturally removed, and oscillations that are actually physical can be better appreciated. The phase diagrams in the $\tilde{e}B - \mu$ plane are presented for different values of the diquark coupling.

1. Introduction

At asymptotically large chemical potentials, the fact that cold quark matter behaves as a color superconductor can be shown by using perturbative methods in the context of quantum chromodynamics (QCD)[1]. However, such methods cannot be applied in the range of moderate densities relevant for, amongst others, the astrophysics of strongly magnetized compact stellar objects[2]. Moreover, since the well-known sign problem prevents lattice QCD calculations from being performed at sufficiently low temperatures and finite chemical potential, one has to rely on effective models to analyze the behavior of magnetized quark matter in this region. One particular model that has been extensively used for this purpose is the Nambu-Jona-Lasinio (NJL) model[3]. This is an effective model originally devised to study the dynamics of chiral symmetry breaking, in which gluon degrees of freedom are integrated out and interactions are described by local four-quark interactions. The incorporation of additional diquark interactions into the model allows for the description of color superconducting matter[4].

In this context, the effect of a constant magnetic field has been analyzed by several authors [5, 6, 7, 8, 9, 10]. However, the Landau level (LL) structure acquired by the quark dispersion relations introduces complications in the regularization procedure necessary to treat the divergent vacuum integrals characteristic of the local interactions in NJL-type models. In the superconducting case, the only prescriptions used until now applied regulating functions on each Landau level integral separately [6, 7, 8, 9, 10]. This procedure, however, might introduce non-physical oscillations [11, 12, 13], making the interpretation of the results unclear. On the other hand, another way of regularizing the NJL model under magnetic fields was reported in Ref.[14] in the absence of color superconductivity. In this procedure, a sum over all Landau levels is applied.

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levels in the vacuum term is performed, and the divergence is isolated into a term that has no explicit dependence on the magnetic field, and can be regularized in the standard fashion. The main purpose of this work is to introduce a regularization scheme which can be considered an extension of this procedure to the superconducting case. The results obtained within this “Magnetic field independent regularization” (MFIR) will be compared to the ones obtained with a few regularizations which have been used in the literature, showing that the aforementioned non-physical oscillations are removed. Then, the behavior of the model under magnetic field and varying coupling strength will be studied.

2. Magnetized cold quark matter within the SU(2)_f NJL model in the presence of color pairing interactions

2.1. The thermodynamical potential in the mean field approximation

We consider a NJL-type SU(2)_f Lagrangian density which includes scalar-pseudoscalar and color pairing interactions. In the presence of an external magnetic field and chemical potential it reads:

\[ \mathcal{L} = \bar{\psi} \left[ i \hat{D} - m_c + \mu \gamma^0 \right] \psi + G \left( \bar{\psi} \psi \right)^2 + \left( \bar{\psi} i \gamma_5 \bar{\tau} \psi \right)^2 + H \left[ (i \bar{\psi} \gamma^C \gamma_5 \psi)(i \bar{\psi} \gamma^C \gamma_5 \gamma^C) \right]. \]  

(1)

Here, \( G \) and \( H \) are coupling constants, \( \psi = (u, d)^T \) represents a quark field with two flavors, \( \psi^C = C \bar{\psi}^T \) and \( \bar{\psi}^C = \bar{\psi}^T C \), with \( C = i \gamma^2 \gamma^0 \), are charge-conjugate spinors and \( \bar{\tau} = (\tau_1, \tau_2, \tau_3) \) are Pauli matrices. Moreover, \((\epsilon^C)^{ab} = (\epsilon^3)^{3ab}\) and \((\epsilon^f)^{13}\) are antisymmetric matrices in color and flavor space respectively. Furthermore, \( m_c \) is the (current) quark mass that we take to be the same for both flavors and \( \mu \) is the quark chemical potential. The coupling of the quarks to the electromagnetic field \( \mathcal{A}_\mu \) is implemented through the covariant derivative \( \hat{D}_\mu = \partial_\mu - i e \mathcal{A}_\mu \).

Note that here we are dealing with “rotated” fields. In fact, as is well known, in the presence of a non-vanishing superconducting gap \( \Delta \), the photon acquires a finite mass. However, as shown in Ref.[15], there is a linear combination of the photon and the eighth component of the gluon field that leads to a massless rotated \( U(1) \) field. The associated rotated charge matrix \( \bar{Q} \) is given by

\[ \bar{Q} = Q_f \otimes \frac{\lambda_8}{2\sqrt{3}} \]

(2)

where \( Q_f = \text{diag}(2/3, -1/3) \) and \( \lambda_8 = \text{diag}(1, 1, -2)/\sqrt{3} \). Then, in a six dimensional flavor-color representation \((u_r, u_g, u_b, d_r, d_g, d_b)\), the rotated \( \bar{q} \) for different quarks are: \( u_r = 1/2, u_g = 1/2, u_b = 1, d_r = -1/2, d_g = -1/2, d_b = 0 \). The rotated unit charge \( \bar{e} \) is given by \( \bar{e} = e \cos \theta \), where \( \theta \) is the mixing angle which is estimated to be \( \approx 1/20\).[17]. In the present work we consider a static and constant magnetic field in the 3-direction, \( \mathcal{A}_\mu = \delta_{\mu 3} x_1 B \), which in fact is also a mixture of the electromagnetic field and color fields.

In what follows we work in the mean field approximation (MFA), assuming that the only non-vanishing expectation values are \( < \bar{\psi} \psi > = (M - m_c)/2G \) and \( < i \bar{\psi} \gamma^C \gamma_5 \psi > = -\Delta/2H \), which can be chosen to be real. Here, \( M \) and \( \Delta \) are the so-called dressed quark mass and superconducting gap, respectively. The resulting MFA thermodynamic potential at vanishing temperature reads

\[ \Omega_{\text{MFA}} = \frac{(M - m_c)^2}{4G} + \frac{\Delta^2}{4H} - \sum_{|\bar{q}|=0,1} \frac{P_{|\bar{q}|}}{4}. \]  

(3)

where

\[ P_{|\bar{q}|=0} = \int \frac{d^3p}{(2\pi)^3} \left[ E_0^+ + |E_0^-| \right]. \]

(4)
\[ P_{|\vec{q}|=1} = \frac{e\vec{B}}{8\pi^2} \sum_{k=0}^{\infty} \alpha_k \int_{-\infty}^{\infty} dp_z \left[ E_1^+ + E_1^- \right], \]
\[ P_{|\vec{q}|=1/2} = \frac{e\vec{B}}{4\pi^2} \sum_{k=0}^{\infty} \alpha_k \int_{-\infty}^{\infty} dp_z \left[ E_{1/2}^+ + E_{1/2}^- \right]. \]

Here, we have introduced \( \alpha_k = 2 - \delta_{k0} \) and

\[ E_0^+ = \sqrt{p^2 + M^2} \pm \mu \]
\[ E_1^+ = \sqrt{p_z^2 + 2k\vec{e}B + M^2} \pm \mu \]
\[ E_{1/2}^+ = \sqrt{\left( \sqrt{p_z^2 + k\vec{e}B + M^2} \pm \mu \right)^2 + \Delta^2}. \]

Clearly, Eqs.(4-6) are divergent and, thus, require to be regularized. Some alternative schemes to achieve this will be discussed in the following subsection. Given the corresponding regularized form \( \Omega_{MFA}^{reg} \), the associated gap equations for \( M \) and \( \Delta \) then read

\[ \frac{\partial \Omega_{MFA}^{reg}}{\partial (M, \Delta)} = 0. \] 

For each value of \( \mu \) and \( \vec{e}B \), several solutions of these equations will generally exist, corresponding to different possible phases, and the most stable solution is that associated to the absolute minimum of the thermodynamic potential.

To regularize the integrals, several prescriptions were introduced in which some cutoff function \( h_{\Lambda}(q) \) is introduced, with \( q = p \) in the case of Eq.(4) and \( q = \omega_{p_z,k} \equiv \sqrt{p_z^2 + 2k|\vec{q}|\vec{B}} \) for Eqs.(5,6). The obvious and simplest choice would be to take \( h_{\Lambda}(q) = \Theta(\Lambda - q) \), scheme which will be referred to as “sharp function regularization” (ShFR). In this case, the integral in Eq.(4), which is magnetic field independent, is cut off when \( p = \Lambda \). The contributions coming from Eqs.(5-6) include a sum over Landau levels, and the integral in each of these is cut off for the \( p_z \) that satisfies \( \Lambda = \sqrt{p_z^2 + 2k|\vec{q}|\vec{B}} \), that is, such that the momentum and magnetic field contribution to the quark dispersion relation does not exceed the value of the cut-off. This would seem like a natural way of extending the 3D sharp cutoff zero magnetic field regularization to the finite \( \vec{e}B \) case. However, the magnetic field dependence of the prescription brings in strong unphysical oscillations. In an attempt to minimize the effects of this magnetic field dependence, the Heaviside function can be replaced with a smooth regulator, prescription that will be referred to as “smooth function regularization” (SmFR). For definiteness in this work, we consider the function \( h_{\Lambda}(q) = 1/(1 + \exp[(q/\Lambda - 1)/\alpha]) \). We have verified that other possible choices lead to similar results. To choose the value of the constant \( \alpha \) that determines the regulator smoothness one is limited by the fact that a too steep function does not improve over the ShFR results and that a too smooth function leads to values of the quark condensate in absence of the magnetic field which are quite above the phenomenological range. Here, we follow Refs.[8, 9] and consider \( \alpha = 0.05 \). We will see that this procedure softens the oscillations but does not remove them.

To fully get rid of the above mentioned regularization artifacts, we introduce in what follows an alternative scheme in which the contributions that are explicitly dependent on the magnetic field turn out to be finite and thus do not need to be regularized. We will refer to this regularization scheme as the “magnetic field independent regularization” (MFIR). We start by considering the \( P_{|\vec{q}|=0} \) contribution. Since it is independent of the magnetic field, it can be treated in the usual way [3]. Introducing a sharp 3D cutoff we get
\[ P_{|q|=0} = \frac{1}{\pi^2} \int_0^{\Lambda} dp \ p^2 \sqrt{p^2 + M^2} + \frac{\Theta(\mu - M)}{\pi^2} \left[ \frac{\mu(\mu^2 - M^2)^{3/2}}{3} - \frac{(\mu^2 - M^2)^2}{8} h \left( \frac{M}{\sqrt{\mu^2 - M^2}} \right) \right], \tag{9} \]

where \( h(z) = (2 + z^2)\sqrt{1 + z^2} + z^4 \ln[z/\sqrt{1 + z^2}] \). In the case of \( P_{|q|=1} \) we note that, except for the value of the quark charge, the expression coincides with that analyzed in Ref.[14] where no color pairing interactions were considered. Following the steps in that reference, we get

\[ P_{|q|=1} = \frac{1}{\pi^2} \int_0^{\Lambda} dp \ p^2 \sqrt{p^2 + M^2} + \frac{\tilde{e}B}{4\pi^2} k_{\text{max}} \sum_{k=0}^{k_{\text{max}}} \alpha_k \left[ \mu\sqrt{\mu^2 - s_k^2} - s_k^2 \ln \left( \frac{\mu + \sqrt{\mu^2 - s_k^2}}{s_k} \right) \right] \]

\[ + \frac{\tilde{e}B^2}{2\pi^2} \left[ \xi'(-1, x) + \frac{x - x^2}{2} \ln x + \frac{x^2}{4} \right] \tag{10} \]

where \( x = M^2/(2\tilde{e}B), k_{\text{max}} = \text{Floor}[\mu^2 - M^2]/(2\tilde{e}B)] \) and \( s_k = \sqrt{M^2 + 2k\tilde{e}B} \). In Eqs.(9-10), the first term is a vacuum contribution which does not explicitly depend on the magnetic field and the second term is the matter contribution. The last term in Eq.(10) is the explicit magnetic field contribution to the vacuum, which has been isolated into a finite term. The case of \(|q|=1/2 \) is more involved. However, it can be cast into the form (see Appendix of [18] for details)

\[ P_{|q|=1/2} = \frac{2}{\pi^2} \int_0^{\Lambda} dp \ p^2 \left( E_\Delta^+ + E_\Delta^- \right) + \frac{\tilde{e}B^2}{2\pi^2} \left[ \xi'(-1, y) + \frac{y - y^2}{2} \ln y + \frac{y^2}{4} \right] \]

\[ + \frac{\tilde{e}B^2}{2\pi^2} \int_0^{\infty} dp \left[ \sum_{k=0}^{\infty} \alpha_k \ f(p^2 + k) - 2 \int_0^{\infty} dx \ f(p^2 + x) \right]. \tag{11} \]

where

\[ E_\Delta^\pm = \sqrt{(\sqrt{p^2 + M^2} \pm \mu)^2 + \Delta^2} \tag{12} \]

\[ y = (M^2 + \Delta^2)/(\tilde{e}B) \tag{13} \]

and

\[ f(z) = \sum_{s=\pm 1} \left\{ \sqrt{(\sqrt{z + 2s} + \mu/\sqrt{\tilde{e}B})^2 + y - 2s - \sqrt{z + y}} \right\}. \tag{14} \]

In this expression a 3D sharp cutoff has been introduced to regularize the first term, i.e. the one that contains contributions from the vacuum and matter which do not explicitly depend on the magnetic field. Note that, as in the case of vanishing \( \tilde{e}B \) discussed in e.g. Ref.[19], these cannot be disentangled into two terms unless \( \Delta = 0 \). The second term is the vacuum magnetic contribution analogous to the \( |q|=1 \) case. Finally, the third term is an additional explicitly magnetic field dependent matter contribution which turns out to be finite.

In this work, the parameters will be set to the following values: \( m_c = 5.595 \text{ MeV}, \)

\( \Lambda = 620.9 \text{ MeV}^2 \) and a coupling constant \( G \) such that \( GA^2 = 2.212 \). This parameter choice leads \( M_0 = 340 \text{ MeV} \) within the MFIR regularization, where \( M_0 \) represents the vacuum quark effective mass in the absence of external magnetic fields.
3. Numerical results

3.1. Comparison between regularization schemes

The behavior of the dressed mass as a function of magnetic field is displayed in Fig. 1, for the three regularizations and for \( \mu = 0 \). The results for ShFR and SmFR exhibit non-physical oscillations, whose origin lies in the magnetic field dependence of the regularization in the vacuum term, that causes the contribution of a given LL to be larger for lower magnetic fields. In the ShFR, which is the most extreme case, the only LLs participating in the sum are those for which \( \Lambda^2 \geq 2k|\tilde{q}|B \). Therefore, depending on the magnetic field, more or less terms appear and each time the relation is satisfied for a given \( k \), there will be a discontinuity in the derivative of the thermodynamic potential. This singularity, hence, does not correspond to a phase transition. The soft regulator, which could be regarded as a way to handle this problem and remove sharp oscillations, still contains this pathology, because in this case the contribution of a given LL also depends on the magnetic field through the Fermi-type regulator function. So, even though the smooth integrals partly conceal this problem, the oscillations are still present and we can actually see that for ShFR and SmFR they are in phase. On the other hand, in the MFIR scheme the mass increases steadily with the magnetic field displaying the usual “magnetic catalysis effect” as in e.g. Refs. [14, 20, 21, 22].

![Figure 1](shfr_smfr_mfir_mass.png)  
**Figure 1.** \( M \) vs \( \tilde{e}B \) in the B phase, for the three regularization schemes considered and \( H/G = 0.75 \). Note that \( \Delta = 0 \) in this phase.

![Figure 2](phasediagram.png)  
**Figure 2.** Phase diagram in the \( \tilde{e}B - \mu \) plane, for the three regularization schemes considered and \( H/G = 0.75 \). All transitions seen in this diagram are first order.

The phase diagrams in the \( \tilde{e}B - \mu \) plane are present in Fig.2. Despite the differences owing to the regularization, a few general properties are common to all prescriptions. For low enough chemical potential, there is always a vacuum phase (B) in which chiral symmetry is broken and the superconducting gap is zero. When chemical potential is raised, there is a first order transition to a region in which the quark population is non-zero, the superconducting gap is finite and chiral symmetry is restored (strictly speaking, restoration is only approximate due to the small but non-zero current quark mass). The near vertical lines, which are known as vAdH transitions, are first order transitions that separate the superconducting region into several phases which are denoted as \( A_i \). In a given \( A_i \) phase, Landau levels from zero to \( i \) are populated, and through any given vAdH transition, the value of \( i \) changes in one unit. The vAdH transitions coincide almost exactly for all three regularizations. We should also emphasize that these transitions are effectively physical because quark density is discontinuous at these, and that they in turn induce physical oscillations in the order parameters as a function of the magnetic field. Despite these general similarities in the phase diagrams for the three
regularizations, we can once again see that the critical chemical potential for the transition from the B phase to the A-type phases exhibits strong oscillations in the ShFR case, and that these continue to exist in SmFR, even though they are slightly suppressed. These are small for low magnetic fields, but become larger in the intermediate $\tilde{e}B$ range and once again make the phase diagram hard to interpret. In the MFIR scheme, we note that the critical chemical potential is approximately independent of $\tilde{e}B$ for values below 0.07 GeV$^2$, then it decreases until it reaches a minimum near $\tilde{e}B = 0.2$ GeV$^2$ and after this value, it increases back again, giving rise to the usual well-shaped curve related to the “inverse magnetic catalysis effect” [23]. Due to the regulation artifacts, this feature is much less evident in the ShFR and SmFR.

3.2. MFIR results for different coupling ratios

Given that the MFIR is seen to make an improvement over the other displayed regularizations, we will now concentrate on analyzing within this scheme the phase diagrams and their dependence on the coupling constant. In the previous section, only $H/G = 0.75$ was considered. However, given that the value of this ratio is subject to certain degree of uncertainty, it is worthwhile to explore the consequences of varying it within a reasonable range. Thus, in what follows, the representative values $H/G = 0.5, 0.75$ and 1 will be considered. Note that values $H/G > 1$ are quite unlikely to be realized in QCD.

The phase diagrams in the $\tilde{e}B - \mu$ plane for the three coupling ratios considered are displayed in Fig.3. As was seen in the previous section for $H/G = 0.75$, the system is in the B phase for low $\mu$, where chiral symmetry is broken, and in one of the possible A-type phases for a high enough $\mu$ value, where chiral symmetry is approximately restored. However, if the coupling ratio is changed, other phases may appear. The first transition encountered if the phase diagram is traversed in the direction of increasing $\mu$ will be referred to as the “main transition”. It has approximately the same shape in all displayed phase diagrams and it connects the B phase to populated phases in general. As $H/G$ increases, the main transition is displaced downwards in its entirety, and the depth of the “inverse magnetic catalysis well” diminishes. Also, the superconducting gap in the A-type phases becomes larger. For $H/G = 0.5$ (left panel), an intermediate phase that will be referred to as a C-type phase exists. It is enclosed by the main transition from below and by another first order transitions from above. Both transitions are horizontal in a large magnetic field range so that the phase extends in an approximately horizontal band up to $\tilde{e}B \simeq 0.11$ GeV$^2$, where it is bounded by a crossover type transition, signalled by the peak of the chiral susceptibility. In the C-type phase, chiral symmetry is partially restored: the mass is lower than in the vacuum phase but higher than in the chirally restored phases, and density is finite for the $|\tilde{q}| = 0$ and $|\tilde{q}| = 1$ quarks (only the lowest LL being occupied for the latter species). It is important to point out that the existence of the C-type phase is not a consequence of the diquark pairing channels. Actually, superconducting effects are still small for $H/G = 0.5$ and the phase diagram remains basically unchanged with respect to the $H/G = 0$ case. In particular, the superconducting gap is always below 1 MeV in the C-phase.

As the coupling constant ratio is increased from the value $H/G = 0.5$, the upper phase transition is displaced downwards. This causes the C-type phase to shrink until it eventually disappears around $H/G \sim 0.65$, so that a relatively simple phase structure results, in which a single phase transition connects the B phase to the A-type phases. For an even larger value of $H/G$ ($\sim 0.94$), the phase transition splits into two once again, so that for $H/G = 1$ (right panel) there is an intermediate phase, which will be referred to as a D phase. This phase is qualitatively different from the one found in $H/G = 0.5$. To begin with, the transition from vacuum to this phase is second order. Also, both the dressed mass and the superconducting gap are appreciably large. This type of phase is sometimes referred to as a “mixed phase” [26], even though other meanings may be found in the literature for this term [27]. There is no quark
Figure 3. Phase diagrams within the MFIR scheme for different $H/G$ ratios: (a) $H/G = 0.5$, (b) $H/G = 0.75$, (c) $H/G = 1$. Full black lines correspond to first order transitions and dotted black lines to second order transitions. The dash-dotted red lines represent chiral susceptibility crossovers. Blue dashed lines represent quark number susceptibility crossovers related to $|\tilde{q}| = 1/2$ quark which are too weak to be actually considered to separate distinct phases.

population for the $|\tilde{q}| = 0$ and $|\tilde{q}| = 1$ species, but the finiteness of $\Delta$ induces a non-zero density for $|\tilde{q}| = 1/2$ quarks. The transition leading to the A-type phases is first order as in the previous cases. The existence of this phase is a consequence of the diquark pairing alone and hence it exists for $\tilde{e}B = 0$. It is interesting to note, however, that it extends along the magnetic field axis of the phase diagram for an appreciable range.

4. Concluding Remarks
In the present work we explored the effects of magnetic field on cold color superconducting quark matter in the framework of the NJL-type model, using a regularization scheme in which the contributions which are explicitly dependent on the magnetic field turn out to be finite and, thus, do not require to be regularized. Such a “magnetic field independent regularization” (MFIR) scheme can be considered an extension of the method described in e.g. Ref.[14] to the case in which color pairing interactions are present. We compared the corresponding results with those obtained through the regularization methods used in previous works [6, 7, 8, 9, 10] and concluded that the MFIR scheme removes the unphysical oscillations originating in the regularization procedure. Taking this into account, we studied the phase structure in the $\tilde{e}B - \mu$ plane using three coupling ratio values which lead to qualitatively different results. We found that up to $H/G = 0.5$ superconducting effects are still small, and the phase diagram remains basically unchanged. In particular, there is an intermediate phase in which quark population exists but chiral symmetry is still strongly broken. The diquark gap in this phase is finite but extremely small. It is connected to other phases by first order transitions and a chiral crossover.

As the coupling ratio is increased, the two transitions surrounding this phase move close to each other, causing it to disappear at $H/G \sim 0.65$. A relatively simple phase diagram hence exists for a narrow range around the standard value $H/G = 0.75$. If the coupling ratio is increased beyond 0.94, a mixed phase is present in both parameter sets, where both condensates are large. Finally, we found that the inverse catalysis phenomenon is present, although it becomes less pronounced as $H/G$ increases. Even though charge neutrality and $\beta$ equilibrium conditions have not been taken into account in this work, they are certainly necessary for a correct description of the physics of compact stars. We expect to report on this in the near future.
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