

INVERSE MAGNETIC CATALYSIS IN THE POLYAKOV–NAMBU–JONA-LASINIO AND ENTANGLED POLYAKOV–NAMBU–JONA-LASINIO MODELS*

CONSTANÇA PROVIDÊNCIA, MÁRCIO FERREIRA, PEDRO COSTA

Centro de Física Computacional, Department of Physics, University of Coimbra
3004-516 Coimbra, Portugal

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We investigate the QCD phase diagram at zero chemical potential and finite temperature in the presence of an external magnetic field within the three flavor Polyakov–Nambu–Jona-Lasinio and entangled Polyakov–Nambu–Jona-Lasinio models looking for the inverse magnetic catalysis. Two scenarios for a scalar coupling parameter dependent on the magnetic field intensity are considered. These dependencies of the coupling allow to reproduce qualitatively lattice QCD results for the quark condensates and for the Polyakov loop: due to the magnetic field, the quark condensates are enhanced at low and high temperatures and suppressed for temperatures close to the transition temperatures, while the Polyakov loop increases with the increase of the magnetic field.

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1. Introduction

Presently, the investigation of magnetized quark matter is attracting the attention of the physics community due to its relevance for different regions of the QCD phase diagram [1]: from the heavy ion collisions at very high energies, to the understanding of the early stages of the Universe and for studies involving compact objects like magnetars. In the presence of an external magnetic field B , the competition between two different mechanisms determine the behavior of quark matter: on the one hand, the increase of low energy contributions leads to an enhancement of the quark condensate; on the other hand, the suppression of the quark condensate due to the partial restoration of chiral symmetry.

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At zero baryonic chemical potential, almost all low-energy effective models, including the Nambu–Jona-Lasinio (NJL)-type models, find an enhancement of the condensate due to the magnetic field, the so-called magnetic catalysis (MC), and no reduction of the pseudocritical chiral transition temperature with the magnetic field [2]. However, the suppression of the quark condensate, also known as inverse magnetic catalysis (IMC), was obtained in lattice QCD (LQCD) calculations with physical quark masses [3–5]. Due to the IMC effect, the pseudocritical chiral transition temperature decreases and the Polyakov loop increases with increasing B . In Ref. [5], it is argued that the IMC may be a consequence of how the gluonic sector reacts to the presence of a magnetic field, and it is shown that the magnetic field drives up the expectation value of the Polyakov field. The distribution of gluon fields changes as a consequence of the distortion of the quark loops in the magnetic field background. Therefore, the backreactions of the quarks on the gauge fields should be incorporated in effective models in order to describe the IMC.

It is also known that in the region of low momenta, relevant for chiral symmetry breaking, there is a strong screening effect of the gluon interactions which suppresses the condensate [5, 6]. In this region, the gluons acquire a mass M_g of the order of $\sqrt{N_f \alpha_s |eB|}$, due to the coupling of the gluon field to a quark–antiquark interacting state. In the presence of a strong enough magnetic field, this mass M_g for gluons becomes larger. This, along with the property that the strong coupling α_s decreases with increasing B ($\alpha_s(eB) \sim [b \ln(|eB|/\Lambda_{\text{QCD}}^2)]^{-1}$ with $b = (11N_c - 2N_f)/12\pi = 27/12\pi$ [6]), leads to an effective weakening of the interaction between the quarks in the presence of an external magnetic field, and damps the chiral condensate. This suggests that the effective interaction between the quarks should include the reaction of the gluon distribution to the magnetic field background. Having this in mind, the present work shows two different approaches of taking into account the influence of the presence of an external magnetic field in the gluonic sector.

We perform our calculations in the framework of the Polyakov–Nambu–Jona-Lasinio (PNJL) model. The Lagrangian in the presence of an external magnetic field is given by

$$\begin{aligned} \mathcal{L} = & \bar{q} [i\gamma_\mu D^\mu - \hat{m}_f] q + G_s \sum_{a=0}^8 [(\bar{q}\lambda_a q)^2 + (\bar{q}i\gamma_5 \lambda_a q)^2] \\ & - K \{ \det [\bar{q}(1+\gamma_5)q] + \det [\bar{q}(1-\gamma_5)q] \} + \mathcal{U}(\Phi, \bar{\Phi}; T) - \frac{1}{4} F_{\mu\nu} F^{\mu\nu}, \quad (1) \end{aligned}$$

where the quarks couple to a (spatially constant) temporal background gauge field, represented in terms of the Polyakov loop. Besides the chiral point-like coupling G_s , that denotes the coupling of the scalar-type four-quark interaction in the NJL sector, in the PNJL model the gluon dynamics is reduced

to the chiral-point coupling between quarks together with a simple static background field representing the Polyakov loop. The Polyakov potential $\mathcal{U}(\Phi, \bar{\Phi}; T)$ is introduced and depends on the critical temperature T_0 , that for pure gauge is 270 MeV. In addition to the PNJL model, we also consider the effective vertex depending on the Polyakov loop [7] (EPNJL model), $G_s(\Phi; \bar{\Phi}) = G_s[1 - \alpha_1 \Phi \bar{\Phi} - \alpha_2(\Phi^3 + \bar{\Phi}^3)]$, that generates an entanglement interaction between the Polyakov loop and the chiral condensate.

In Case I, we adopt a running coupling of the chiral invariant quartic quark interaction in the PNJL model with the magnetic field [8, 9]. The damping of the strength of the effective interaction is built phenomenologically: since there is no available LQCD data for $\alpha_s(eB)$, we fit $G_s(eB)$ in order to reproduce the chiral pseudocritical temperature $T_c^X(eB)$ obtained in LQCD calculations [3]. The $G_s(eB)$ coupling, that reproduces $T_c^X(eB)$, is calculated in the NJL model and is shown in the left panel of Fig. 1 (solid black line). Now, using this $G_s(eB)$ coupling in the PNJL model, both the deconfinement transition and chiral transition pseudocritical temperatures are decreasing functions with eB , up to $eB \sim 1 \text{ GeV}^2$. Due to the existing coupling between the Polyakov loop field and quarks, the $G_s(eB)$ does not only affect the chiral transition but also the deconfinement transition.

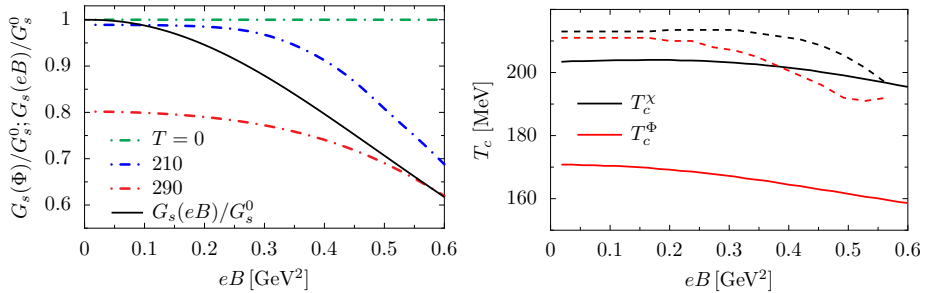


Fig. 1. Comparison between both cases: left panel – the scalar coupling G_s versus the magnetic field; right panel – the chiral and deconfinement temperatures as a function of eB , being the full (dashed) lines for Case I (Case II).

In Case II, we introduce an eB -dependence on the pure-gauge critical temperature T_0 , reproducing the LQCD data for the deconfinement transition [3], in order to mimic the reaction of the gluon sector to the presence of an external magnetic field.

We will use the EPNJL model because within the PNJL it is not possible to implement the above scheme, since the chiral transition temperatures increase strongly with the external magnetic field. In order to bring these temperatures down, it would be necessary to use very small values of T_0 , for which the deconfinement phase transition becomes of first order.

Nevertheless, within EPNJL the chiral condensates and the Polyakov loop are entangled. Thus, the chiral transition temperatures are pulled down to temperatures close to the deconfinement transition temperature. This model, however, at moderate magnetic fields, still predicts a first order transition for both transitions, when a small T_0 is needed. As a consequence, a too small value of T_0 leads to a first order phase transition within both PNJL and EPNJL models, and, therefore, the range of T_0 values we are interested in is limited to the values that maintain the crossover transition. A larger range of validity is obtained if the quark backreactions are not taken into account at $eB = 0$, *i.e.* when $T_0 = 270$ MeV as obtained in pure gauge. This gives $T_c^\Phi = 214$ MeV, 40 MeV higher than the prediction of lattice QCD data in [3]. This parametrization reproduces the referred lattice QCD data for $T_c^\Phi(eB)$, shifted by 40 MeV, for magnetic fields up to 0.61 GeV^2 . Above 0.61 GeV^2 , a first order phase transition occurs. We will use the last scenario in Case II to illustrate our results because larger magnetic fields are achieved.

Moreover, in the EPNJL the coupling G_s depends on the Polyakov loop, thus, in the crossover region, where the Polyakov loop increases with temperature, the coupling G_s becomes weaker. This is shown in Fig. 1 (left panel), where the coupling $G_s[\Phi(T)]$ is plotted for several temperatures (dashed curves) [10]. Within the PNJL model with constant coupling G_s , no IMC effect was obtained even with $T_0(eB)$, because $T_0(eB)$ does not affect the coupling G_s .

In Fig. 1 (right panel), the results for the pseudocritical temperatures for both cases are compared. In Case II, the pseudocritical temperatures have a much flatter behavior at small values of eB than in Case I, reflecting the softer decrease of the coupling G_s at small magnetic field values as shown in Fig. 1 (left panel). Also, within the EPNJL with $T_0(eB)$, the difference between the pseudocritical temperatures T_c^χ and T_c^Φ is much smaller, due to the strong coupling between the Polyakov loop and the quark condensates. For $eB = 0$, these temperatures are almost coincident, but a finite strong magnetic field destroys this coincidence. The PNJL model with $G_s(eB)$ does not have this feature and different temperatures for T_c^χ and T_c^Φ are predicted.

Next, we discuss the effect of the magnetic field on the quark condensates and on the Polyakov loop, for both cases.

According to [4], we define the change of the light quark condensate due to the magnetic field as $\Delta\Sigma_f(B, T) = \Sigma_f(B, T) - \Sigma_f(0, T)$, with $\Sigma_f(B, T) = \frac{2M_f}{m_\pi^2 f_\pi^2} [\langle \bar{q}_f q_f \rangle(B, T) - \langle \bar{q}_f q_f \rangle(0, 0)] + 1$, where the factor $m_\pi^2 f_\pi^2$ in the denominator contains the pion mass in the vacuum ($m_\pi = 135$ MeV) and (the chiral limit of the) pion decay constant ($f_\pi = 87.9$ MeV) in NJL model.

In Fig. 2, the light chiral condensate $\Delta(\Sigma_u + \Sigma_d)/2$ is plotted as a function of the magnetic field, for $eB < 1$ ($eB < 0.61$) GeV^2 in Case I (Case II), at temperatures close to the respective $T_c^\Phi(eB = 0)$. The main conclusions are: (i) for both cases the qualitative behavior shown in Fig. 2 of Ref. [3] and in Fig. 6 of Ref. [5] is reproduced, that is, the non-monotonic behavior of the condensates as a function of the magnetic field, having the $T = 0$ curves the highest $\Delta(\Sigma_u + \Sigma_d)/2$; (ii) for temperatures close $T_c^\chi(eB = 0)$ the strong interplay between the partial restoration of chiral symmetry and the condensate enhancement due to the magnetic field gives rise to curves that increase, for small values of eB , and as soon as the partial restoration of chiral symmetry becomes dominant the curve starts to decrease.

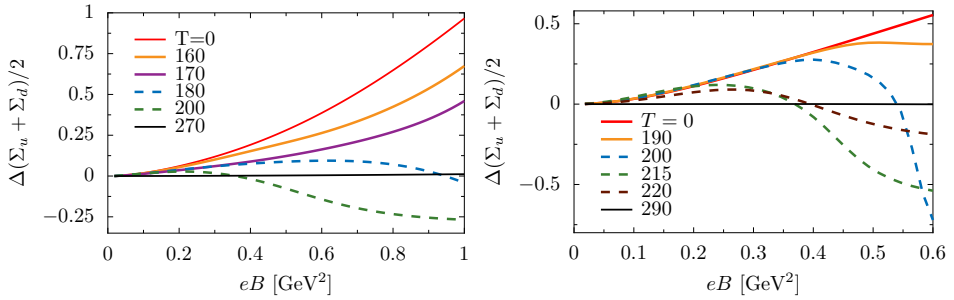


Fig. 2. The light chiral condensate, $\Delta(\Sigma_u + \Sigma_d)/2$, as a function of eB for several values of T in MeV: left panel — Case I; right panel — Case II.

Finally, the effect of the magnetic field on the Polyakov loop is seen in Fig. 3, where the Polyakov loop value, Φ , is plotted as a function of the temperature, for several magnetic field strengths. As can be seen, for both cases, the Polyakov loop increases sharply with the magnetic field around

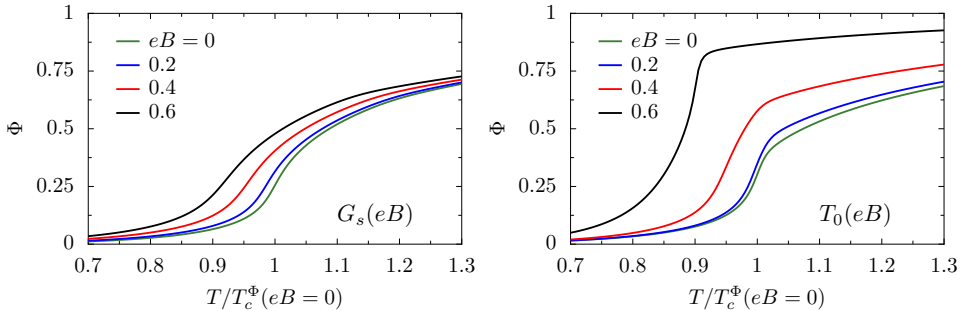


Fig. 3. The Polyakov loop as a function of T for different values of eB (in GeV^2) renormalized by the deconfinement pseudocritical temperature at $eB = 0$: $T_c^\Phi = 171$ MeV for Case I (left panel) and $T_c^\Phi = 214$ MeV for Case II (right panel).

the transition temperature, and the transition temperature decreases with the magnetic field with B , in close agreement with the LQCD results [5]. Indeed, the suppression of the condensates achieved by the magnetic field dependence of the coupling parameter results in an increase of the Polyakov loop, with this effect being stronger precisely for temperatures close to the transition temperature.

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