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Restoration of dynamically broken chiral and color symmetries for an accelerated observer

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Abstract

We study the behavior of quark and diquark condensates at finite Unruh temperature as seen by an accelerated observer. The gap equations for these condensates have been obtained with consideration of a finite chemical potential. Critical values of the acceleration for the restoration of chiral and color symmetries have been estimated.
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1. Introduction

It is well known that, according to the Hawking–Unruh effect [1,2], an observer accelerated uniformly in the QCD vacuum behaves as if he were in a thermal bath at the Unruh temperature $T_U = a/2\pi$ (a is an acceleration constant). For interacting field theories this is demonstrated by the fact that Euclidean Green's functions written in terms of the Rindler coordinates are periodic in time, and therefore they may be interpreted as thermal. Several problems have been studied in relation to the Hawking–Unruh effect: see, for instance, [3] concerning the discussion of a uniformly accelerated oscillator, or [4] with the study of pair creation by a homogeneous electric field from the point of view of an accelerated observer. At the same time, in a recent paper [5], it was argued that the phase transition from a color glass condensate to a quark–gluon plasma through the mechanism of the Hawking–Unruh thermalization can become experimentally observable in relativistic heavy ion collisions.

It was proposed more than twenty years ago [6–8] that at high baryon densities a colored diquark condensate $\langle qq\rangle$ might appear. In analogy with the ordinary superconductivity, this effect was called color superconductivity (CSC). In particular, the CSC phenomenon was investigated in the framework of the one-gluon exchange approximation in QCD [9], where the colored Cooper pair formation is predicted selfconsistently at extremely high values of the chemical potential $\mu \gtrsim 10^8$ MeV [10]. Unfortunately, such baryon densities are not observable in nature and not accessible in experiments (the typical densities inside the neutron stars or in the future heavy ion experiments correspond to $\mu \sim 500$ MeV). In order to study the problem at lower values of μ, various effective theories for low energy QCD, such as the instanton model [11] and the Nambu–Jona-Lasinio (NJL) model [12] can be employed.

It is well known that effective field theories with four-fermion interaction (the so-called Nambu–Jona-Lasinio (NJL) models), which incorporate the phenomenon of dynamical chiral symmetry breaking, are quite useful in describing low-energy hadronic processes (see, e.g., [13,14] and references therein). Since the NJL model displays the same symmetries as QCD, it can be successfully used for simulating some of the QCD ground state properties under the influence of external conditions such as temperature, baryonic chemical potential, or even curved spacetime [14–18]. In particular, the role of the NJL approach increases, when detailed numerical lattice cal-

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culations are not yet admissible in QCD with nonzero baryon density and in the presence of external gauge fields [19–21].

The possibility for the existency of the CSC phase in the region of moderate densities was recently proved (see, e.g., the papers [11,17,22,23] as well as the review articles [24] and references therein). In these papers it was shown that the diquark condensate $\langle qq\rangle$ can appear already at a rather moderate baryon density ($\mu \sim 400$ MeV). The conditions favorable for this condensate to be formed can possibly exist in the cores of cold neutron stars. Since quark Cooper pairing occurs in the color antitriplet channel, a nonzero value of $\langle qq\rangle$ means that, apart from the electromagnetic $U(1)$ symmetry, the color $SU_c(3)$ symmetry should be spontaneously broken inside the CSC phase as well. In the framework of NJL models the CSC phase formation has generally been considered as a dynamical competition between diquark $\langle qq\rangle$ and usual quark–antiquark condensation $\langle \bar{q}q\rangle$.

Recently, the dynamical chiral symmetry breaking and its restoration for a uniformly accelerated observer due to the thermalization effect of acceleration was studied in [25] at zero chemical potential. Further investigation of the possible influence of the Unruh temperature on the phase transitions in dense quark matter with a finite chemical potential, and especially on the restoration of the broken color symmetry in CSC is thus especially interesting. Related problems have also been studied for chiral symmetry breaking in curved spacetime [16,26], which may be useful for the investigation of compact stars, where the gravitational field is strong and its effect cannot be neglected. Obviously, the results of these studies might have some relevance to the physics of black holes, whose surface gravity causes the finite temperature effects (the Hawking effect). The Rindler metric may be regarded here as an approximation of the situation near the event horizon of a black hole. Thus, the study of phase transitions in quark matter with quark and diquark condensates under the influence of acceleration or, equivalently, strong gravitational fields is of great interest.

In this Letter, we study quark and diquark condensates as functions of the Unruh temperature and finite chemical potential by using a NJL-type model formulated in Rindler coordinates. As our main result, from the gap equations for these condensates written in Rindler coordinates, the critical values of acceleration (the critical Hawking–Unruh temperatures) for the restoration of the broken chiral and color symmetries were obtained. The results exactly coincide with the usual temperature restoration of these symmetries, when the corresponding relation between the acceleration and the Unruh temperature is taken into account.

2. Nambu–Jona-Lasinio model in Rindler coordinates

The required effective NJL field theory in Rindler coordinates will be obtained by a suitable transformation of a NJL-type model in flat spacetime with Minkowski coordinates $(x^0, x^1, \vec{x}_\perp)$. For this aim, let us first quote some useful definitions and formulas. The physics for an accelerated observer can be described by transforming to the Rindler coordinates $(\eta, \rho, \vec{x}_\perp)$ by means of the following coordinate transformation:

$$x^0 = \rho \sinh a\eta, \qquad x^1 = \rho \cosh a\eta, \qquad x^i = x^i \quad (i = 2, 3),$$

defined on the right Rindler wedge,

$$0 < \rho < +\infty, \qquad -\infty < \eta < +\infty,$$

and on the left Rindler wedge,

$$-\infty < \rho < 0, \qquad -\infty < \eta < +\infty,$$

where η is the time variable in Rindler coordinates.

The gamma-matrices γ_μ, the metric $g_{\mu\nu}$ and the vierbein $e^\mu_{\hat{a}}$, as well as the definitions of the covariant derivative ∇_ν and spin connection $\omega^{\hat{a}\hat{b}}_\nu$ are given by the following relations [27]:

$$\begin{aligned}
&\{\gamma_\mu(x), \gamma_\nu(x)\} = 2g_{\mu\nu}(x),\\
&\{\gamma_{\hat{a}}, \gamma_{\hat{b}}\} = 2\eta_{\hat{a}\hat{b}}, \quad \eta_{\hat{a}\hat{b}} = \mathrm{diag}(1, -1, -1, -1),\\
&g_{\mu\nu} g^{\nu\rho} = \delta^\rho_\mu, \qquad g^{\mu\nu}(x) = e^\mu_{\hat{a}}(x) e^{\nu\hat{a}}(x),\\
&\gamma_\mu(x) = e^{\hat{a}}_\mu(x)\gamma_{\hat{a}}.
\end{aligned} \tag{1}$$

$$\begin{aligned}
&\nabla_\nu \equiv \partial_\nu + \frac{1}{2}\sigma_{\hat{a}\hat{b}}\omega^{\hat{a}\hat{b}}_\nu, \quad \sigma_{\hat{a}\hat{b}} \equiv \frac{1}{4}[\gamma_{\hat{a}}, \gamma_{\hat{b}}],\\
&\omega^{\hat{a}\hat{b}}_\mu \equiv \frac{1}{2} e^{\hat{a}\lambda} e^{\hat{b}\rho}[C_{\lambda\rho\mu} - C_{\rho\lambda\mu} - C_{\mu\lambda\rho}],\\
&C_{\lambda\rho\mu} \equiv e^{\hat{a}}_\lambda \partial_{[\rho} e_{\mu]\hat{a}}.
\end{aligned} \tag{2}$$

Here, the index $\hat{a}$ refers to the flat tangent space defined by the vierbein at spacetime point x, and the $\gamma^{\hat{a}}$ $(a = 0, 1, 2, 3)$ are the usual Dirac gamma-matrices of Minkowski spacetime.

The line element

$$ds^2 = \eta_{\hat{a}\hat{b}} e^{\hat{a}}_\mu e^{\hat{b}}_\nu \, dx^\mu \, dx^\nu$$

in these coordinates with the vierbeins

$$e^{\hat{0}}_0 = a\rho, \qquad e^{\hat{1}}_1 = \cdots = 1$$

is given by the relation

$$ds^2 = a^2\rho^2 \, d\eta^2 - d\rho^2 - d\vec{x}^2_\perp$$

with the metric tensor

$$g_{\mu\nu} = \left(a^2\rho^2, -1, -1, -1\right). \tag{3}$$

In what follows, we shall limit our consideration to the right Rindler wedge. An observer at fixed $\rho, \vec{x}_\perp$ measures a proper time $d\tau = a\rho\, d\eta$ and has a proper acceleration $1/\rho$. The observer at $\rho = 1/a$ measures $d\tau = d\eta$ and has a proper acceleration a. The world line of the observer in Rindler coordinates is thus given as

$$\eta(\tau) = \tau, \qquad \rho(\tau) = 1/a, \qquad \vec{x}_\perp(\tau) = \mathrm{const.} \tag{4}$$

For an accelerated observer moving with constant acceleration according to (4), the gamma matrices in Rindler coordinates are obtained by the definition given in (1)

$$\gamma_0 = a^2\rho^2\gamma^0 = a\rho\gamma_{\hat{0}}, \qquad \gamma_\mu = g_{\mu\nu}\gamma^\nu, \tag{5}$$

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