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PARITY VIOLATION IN QCD-MOTIVATED HADRONIC MODELS AT EXTREME CONDITIONS

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We investigate the possibility of parity being spontaneously violated in QCD at finite baryon density and temperature. QCD is approximated by a generalized σ model with two isomultiplets of scalars and pseudoscalars. The mechanism of parity violation is based on interplay between lightest and heavier degrees of freedom and it cannot be understood in simple models retaining the pion and nucleon sectors solely. We argue that, in the dense and hot nuclear matter of a few normal densities and moderate temperatures, parity violation may arise due to a second-order phase transition and its occurrence is well compatible with the existence of stable bound state of normal nuclear matter.

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INTRODUCTION

The appearance of parity (P) violation via pseudoscalar condensation for sufficiently large values of temperature and/or chemical potential has been attracting much interest during last decades to search it both in dense nuclear matter (in neutron/quark stars and heavy-ion collisions at intermediate energies) and in strongly interacting quark–gluon matter («quark–gluon plasma» in heavy-ion collisions at very high energies). At finite baryon density it was conjectured by A. Migdal in [1] long ago. One should also mention the possibility of $(C)P$-parity violation in metastable nuclear bubbles created in hot nuclear matter [2]. Finally P violation might conceivably accompany the transitions to open color phases [3] such as CFL (color-flavor locking) or SC (superconducting), but these are phases beyond the range of validity of our analysis. While it was argued in [4] that parity and vector flavor symmetry could not undergo spontaneous symmetry breaking in a

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vector-like theory such as QCD, the conditions under which the results of [4] hold are not valid for nonzero chemical potential.

Parity violation in QCD would lead to rather remarkable experimental signals such as the same in-medium resonance being able to decay into even and odd number of pions, the presence of additional Goldstone bosons (in the exact chiral limit, six right at the phase transition, and five throughout the broken parity phase), changes in the nuclear equation of state, isospin breaking effects in the pion decay constant and substantial modification of the weak decay constant $F_{\pi'}$ for massless charged pions, giving an enhancement of electroweak decays.

In this talk we consider the possibility of spontaneous parity violation employing effective Lagrangian techniques in the range of nuclear densities where the hadron phase persists and quark percolation does not occur yet. Our effective Lagrangian is a realization of the generalized linear σ model, but including the two lowest lying resonances in each channel. This is the minimal model where this interesting possibility can be realized. The use of effective Lagrangians is also crucial to answer the second question of interest, namely, how would parity violation originating from a finite baryon density eventually reflect in hadronic physics. We present the analysis [5] extended to the case of nonzero temperature [6], showing that the parity violation phase persists in some finite domain in the μ–T plane.

We also address the issue of how our model can describe the saturation point and the formation of stable nuclear matter and find that our description turns out to be rather accurate in describing nuclear matter formation avoiding the unacceptable «chiral collapse» [7].

Many techniques have been used to study QCD in extreme conditions: from meson–nucleon [1, 8] or quark–meson [9, 10] Lagrangians for low-dense nuclear matter to models of Nambu–Jona-Lasinio type [11, 12] for high-dense quark matter [3]. However, all hadronic models lack, for one reason or another, some essential ingredient to detect spontaneous parity violation. One should also mention the extensive lattice investigations, plagued with technical difficulties when $\mu \neq 0$ [13]. Let us finally comment that the range of intermediate nuclear densities (from 3 to 10 times the usual nuclear density) where we expect parity violation to occur is of high interest as it may be reached both in heavy-ion collisions [14, 15] and in compact stars [10].

1. A GENERALIZED SIGMA MODEL FOR QCD

The oldest hadronic effective theory is the linear σ model of Gell-Mann and Levy [16], which contains a multiplet of the lightest isoscalar σ and isotriplet pseudoscalar π^a fields. Spontaneous chiral symmetry breaking emerges due to a nonzero value for $\langle\sigma\rangle \sim \langle\bar{q}q\rangle/F_\pi^2$.

The minimal model to explore spontaneous parity breaking (SPB) contains two multiplets of scalar/pseudoscalar fields $H_j = \tilde{\sigma}_j \mathbf{I} + i\hat{\pi}_j$, $j = 1, 2$, with $\hat{\pi}_j \equiv \tilde{\pi}^a_j \tau^a$, where τ^a are Pauli matrices. We require an exact $SU(2)_L \times SU(2)_R$ symmetry in the chiral limit. The effective potential of this generalized σ model is

$$V_{\text{eff}} = \frac{1}{2}\,\mathrm{tr}\left\{ -\sum_{j,k=1}^{2} H_j^\dagger \Delta_{jk} H_k + \lambda_1 (H_1^\dagger H_1)^2 + \lambda_2 (H_2^\dagger H_2)^2 + \lambda_3 H_1^\dagger H_1 H_2^\dagger H_2 + \right.$$
$$+ \frac{1}{2}\lambda_4 (H_1^\dagger H_2 H_1^\dagger H_2 + H_2^\dagger H_1 H_2^\dagger H_1) + \frac{1}{2}\lambda_5 (H_1^\dagger H_2 + H_2^\dagger H_1) H_1^\dagger H_1 +$$
$$\left. + \frac{1}{2}\lambda_6 (H_1^\dagger H_2 + H_2^\dagger H_1) H_2^\dagger H_2 \right\} + \mathcal{O}\left(\frac{|H|^6}{\Lambda^2}\right) \quad (1)$$

and contains nine real constants. QCD bosonization rules imply that they are $\sim N_c$. The neglected terms will be suppressed by inverse powers of the chiral symmetry breaking (CSB) scale $\Lambda \simeq 1.2$ GeV.

This effective potential can be further simplified by making a general linear transformation on the H_j fields

$$\tilde{H}_j = \sum_{k=1,2} L_{jk} H_k. \quad (2)$$

The L_{jk} must be real in order not to mix states of different parities. We shall consider here the case where the eigenvalues of the matrix $-\Delta_{ij}$ in (1) are all negative. As the transformation (2) has four real parameters, it suffices to have $\Delta_{jk} = \Delta\delta_{jk}$ and therefore $\sum_{j,k=1}^{2} H_j^\dagger \Delta_{jk} H_k \Rightarrow \Delta(|H_1|^2 + |H_2|^2)$.

We shall take H_1 as the chiral multiplet coupling locally to the quark fields (see Sec. 2) and the above diagonalization can be implemented without changing this prescription. Using the global invariance of the model, we parameterize,

$$\begin{aligned} H_1(x) &= \sigma_1(x)\xi^2(x) = \sigma_1(x)\exp\left(i\frac{\pi_1^a \tau_a}{F_0}\right), \\ H_2(x) &= \xi(x)\Big(\sigma_2(x) + i\hat{\pi}_2(x)\Big)\xi(x). \end{aligned} \quad (3)$$

The parities of $\sigma_2(x)$ and $\hat{\pi}_2$ are even and odd, respectively (in the absence of SPB). In these variables the corresponding gap equations are

$$\begin{aligned} 2\Delta\sigma_1 &= 4\lambda_1\sigma_1^3 + 3\lambda_5\sigma_1^2\sigma_2 + 2(\lambda_3 + \lambda_4)\sigma_1\sigma_2^2 + \lambda_6\sigma_2^3 + \rho^2(2(\lambda_3 - \lambda_4)\sigma_1 + \lambda_6\sigma_2), \\ 2\Delta\sigma_2 &= \lambda_5\sigma_1^3 + 2(\lambda_3 + \lambda_4)\sigma_1^2\sigma_2 + 3\lambda_6\sigma_1\sigma_2^2 + 4\lambda_2\sigma_2^3 + \rho^2(\lambda_6\sigma_1 + 4\lambda_2\sigma_2), \end{aligned} \quad (4)$$

$$0 = \rho(-\Delta + (\lambda_3 - \lambda_4)\sigma_1^2 + \lambda_6\sigma_1\sigma_2 + 2\lambda_2\sigma_2^2 + 2\lambda_2\rho^2), \quad (5)$$

where the following notation has been introduced in anticipation of neutral pseudoscalar condensate: $\langle\pi_1^a\rangle = \langle\pi^0\rangle\delta^{0a}, \langle\pi_2^a\rangle = \rho\delta^{0a}$.

The above effective potential must exhibit the usual chiral symmetry breaking pattern at $\mu = T = 0$. For this to happen, $\langle\sigma_1\rangle$ must acquire a real and positive v.e.v. to agree with current algebra considerations. Note that $\langle\pi^0\rangle$ does not appear at all in the gap equations and hence its value is completely undetermined, but the addition of a small mass for the quarks fixes the phase of the breaking $\langle\pi^0\rangle = 0$.

The previous set of gap equations may have several solutions for σ_1 and σ_2, but since we know that in normal conditions QCD does not break parity, ρ must vanish. Since for the potential to be well defined requires $\lambda_2 > 0$, a *sufficient* condition for the absence of SPB is

$$(\lambda_3 - \lambda_4)\sigma_1^2 + \lambda_6\sigma_1\sigma_2 + 2\lambda_2\sigma_2^2 > \Delta. \tag{6}$$

On the other hand, the mass of the pseudoscalar π_2 is governed by the second variation

$$\frac{1}{2}V^{(2)}_{\pi_2\pi_2} = -\Delta + (\lambda_3 - \lambda_4)\sigma_1^2 + \lambda_6\sigma_1\sigma_2 + 2\lambda_2\sigma_2^2 + 6\lambda_2\rho^2. \tag{7}$$

Positivity of this mass for $\rho = 0$ implies (6). The condition is therefore *necessary* too.

The necessary condition for CSB in normal conditions is to have a minimum of V_{eff} for nonzero σ_j (for vanishing ρ). It can be derived from the condition to get a local maximum (or at least a saddle point) for zero σ_j. This extremum is characterized by the matrix $-\Delta_{ij} \Rightarrow -\Delta\delta_{jk}$ in (1). The sufficient condition for CSB follows from the positivity of the second variation of V_{eff} for a nontrivial solution of the two equations (4) at $\rho = 0$ (see Sec. 3).

2. INCLUSION OF TEMPERATURE AND CHEMICAL POTENTIAL

After bosonization the baryon chemical potential μ is transmitted to the meson sector (in the leading order of chiral expansion) via a local quark–meson coupling. In turn, in the large N_c limit and for moderate temperatures, one can neglect the temperature dependence due to meson collisions and assume that the temperature T is induced with the help of the imaginary time Matsubara formalism for Green functions — Matsubara frequencies for quarks $\omega_n = (2n+1)\pi/\beta$ with $\beta = 1/kT$. In the real world with three colors this is of course an approximation, but nevertheless it should be sufficient to describe qualitatively the interplay between baryon density and temperature, and it is the one consistent with our mean-field approach.

As already mentioned, we take the chiral multiplet to have local couplings with the quark fields as being H_1. The set of coupling constants in (1) allows us to fix the Yukawa coupling constant to unity. Thus, μ and T are transmitted to the boson sector by the term

$$\Delta\mathcal{L} = -(\bar{q}_R H_1 q_L + \bar{q}_L H_1^\dagger q_R) \longrightarrow -\bar{q}\sigma_1 q, \tag{8}$$

where $q_{L,R}$ are assumed to be constituent quarks. We do not include baryon fields explicitly and therefore quark matter and nuclear matter are indistinguishable in our approach.

After integrating out the constituent quarks the full temperature and chemical potential dependence, to the leading orders in chiral expansion, find their way into the gap equations. Namely, the first Eq. (4) is modified to

$$2\Delta\sigma_1 = 4\lambda_1\sigma_1^3 + 3\lambda_5\sigma_1^2\sigma_2 + 2(\lambda_3+\lambda_4)\sigma_1\sigma_2^2 + \lambda_6\sigma_2^3 + \rho^2(2(\lambda_3-\lambda_4)\sigma_1 + \lambda_6\sigma_2) + \\ + 2\mathcal{N}\sigma_1\mathcal{A}(\sigma_1,\mu,\beta), \quad \mathcal{N} \equiv \frac{N_c N_f}{4\pi^2}, \tag{9}$$

$$\mathcal{A}(\sigma_1,\mu,\beta) = 2\int_{\sigma_1}^{\infty} dE\,\sqrt{E^2-\sigma_1^2}\,\frac{\cosh(\beta\mu) + \exp(-\beta E)}{\cosh(\beta\mu) + \cosh(\beta E)}, \tag{10}$$

where the Fermi distribution has been introduced. All the dependence on the environment is in the function $\mathcal{A}$ which originates from the one-loop contribution [6] to V_{eff}. At $T = 0$

$$\Delta V_{\text{eff}}(\mu) = \frac{\mathcal{N}}{2}\Theta(\mu-\sigma_1)\times \\ \times\left[\mu\sigma_1^2\sqrt{\mu^2-\sigma_1^2} - \frac{2\mu}{3}(\mu^2-\sigma_1^2)^{3/2} - \sigma_1^4\ln\frac{\mu+\sqrt{\mu^2-\sigma_1^2}}{\sigma_1}\right] \tag{11}$$

and

$$\mathcal{A}(\sigma_1,\mu,\beta=\infty) = \\ = 2\theta(\mu-\sigma_1)\int_{\sigma_1}^{\mu} dE\,\sqrt{E^2-\sigma_1^2} = \mu\sqrt{\mu^2-\sigma_1^2} - \sigma_1^2\ln\frac{\mu+\sqrt{\mu^2-\sigma_1^2}}{\sigma_1}. \tag{12}$$

By using the gap equations (4), (5) and (9), the value of the effective potential at its minima is given by the compact expression

$$V_{\text{eff}}(\mu) = \\ = -\frac{1}{2}\Delta\sum_{j=1}^{2}(\sigma_j(\mu))^2 - \frac{1}{2}\Delta\rho^2(\mu) - \frac{\mathcal{N}}{3}\mu(\mu^2-\sigma_1(\mu)^2)^{3/2}\theta(\mu-\sigma_1(\mu)). \tag{13}$$

It should be emphasized that all the above results have corrections of $\mathcal{O}(\mu^2/\Lambda^2, \sigma_1^2/\Lambda^2)$.

The stabilization of nuclear matter requires not only attractive scalar forces (scalars) but also repulsive ones (vector-mediated) [17]. Conventionally, the latter ones are associated to the interactions mediated by the iso-singlet vector ω meson. Let us supplement our action with

$$\Delta\mathcal{L}_\omega = -\frac{1}{4}\omega_{\mu\nu}\omega^{\mu\nu} + \frac{1}{2}m_\omega^2\omega_\mu\omega^\mu - g_{\omega\bar{q}q}\bar{q}\gamma_\mu\omega^\mu q, \tag{14}$$

with a coupling constant $g_{\omega\bar{q}q} \sim \mathcal{O}(1/\sqrt{N_c})$. In the quark sector the time component ω_0 interplays with the chemical potential [6]. A constant v.e.v. for this component $g_{\omega\bar{q}q}\langle\omega_0\rangle \equiv \bar{\omega}$. Then one needs to compute the modification of the effective potential due to the replacement $\mu \to \mu + \bar{\omega} \equiv \bar{\mu}$,

$$\Delta V_\omega = -\frac{1}{2}m_\omega^2\langle\omega_0^2\rangle = -\frac{1}{2}\frac{(\bar{\mu}-\mu)^2}{G_\omega}, \quad G_\omega \equiv \frac{g_{\omega\bar{q}q}^2}{m_\omega^2} \simeq \mathcal{O}\left(\frac{1}{N_c}\right). \tag{15}$$

$\bar{\mu}$ can be determined via the variation of $\bar{V}_{\text{eff}}$

$$\frac{\bar{\mu}-\mu}{G_\omega} = -N_c\varrho_B(\mu) = -\frac{N_cN_f}{3\pi^2}(\bar{\mu}^2 - \sigma_1(\bar{\mu})^2)^{3/2}. \tag{16}$$

The first-order phase transition from nuclear vapor to nuclear liquid — stable nuclear matter, occurs at zero pressure when $\bar{\mu}^* < \sigma_1^0$, $\sigma_j^* \equiv \sigma_j(\bar{\mu}^*)$. The energy crossing condition can be written, taking into account (13) and (15), as

$$\sum_{j,k=1}^{2}(\sigma_j^0\Delta_{jk}\sigma_k^0 - \sigma_j^*\Delta_{jk}\sigma_k^*) = \frac{N_cN_f}{6\pi^2}\bar{\mu}^*p_F^3(\bar{\mu}^*) + G_\omega\frac{N_c^2N_f^2}{9\pi^4}p_F^6(\bar{\mu}^*) = $$
$$= \frac{N_c}{2}\bar{\mu}^*\varrho_B(\mu^*) + G_\omega N_c^2\varrho_B^2(\mu^*), \tag{17}$$

where $\bar{\mu}^*$ is related to the physical value of μ^* by Eq. (16). This relation represents the condition for the existence of symmetric nuclear matter. It can always be fulfilled by an appropriate choice of G_ω.

3. THE SPB PHASE TRANSITION

We shall consider from now on the solution corresponding to the most stable minima for $\mu > \mu^*$.

The possibility of SPB is controlled by the inequality (6). In order to approach an SPB phase transition when the chemical potential is increasing, we have to diminish the l.h.s. of the inequality (6) and therefore we need to have

$$\partial_\mu((\lambda_3-\lambda_4)\sigma_1^2+\lambda_6\sigma_1\sigma_2+2\lambda_2\sigma_2^2)<0\Rightarrow$$
$$\Rightarrow(\lambda_6\sigma_1+4\lambda_2\sigma_2)V^{(2)}_{\sigma_1\sigma_2}<(2(\lambda_3-\lambda_4)\sigma_1+\lambda_6\sigma_2)V^{(2)}_{\sigma_2\sigma_2}. \quad (18)$$

This last inequality is a *necessary* condition that has to be satisfied by the model for it to be potentially capable of yielding SPB at large densities.

Let us examine the possible existence of a region of μ where $\rho\neq 0$. Then

$$(\lambda_3-\lambda_4)\sigma_1^2+\lambda_6\sigma_1\sigma_2+2\lambda_2(\sigma_2^2+\rho^2)=\Delta. \quad (19)$$

After substituting Δ from (19) into the second Eq. (4), one finds that

$$\lambda_5\sigma_1^2+4\lambda_4\sigma_1\sigma_2+\lambda_6(\sigma_2^2+\rho^2)=0, \quad (20)$$

where we have taken into account that $\sigma_1\neq 0$. Together with (19), this completely fixes the relation between the two v.e.v.'s of the scalar fields $\sigma_{1,2}$ throughout the SPB phase independently of μ and ρ. If $\lambda_2\lambda_6\neq 0$, (19) and (20) allow us to get rid of the v.e.v. ρ for $\mu>\mu_{\rm crit}$ and

$$\sigma_2=A\sigma_1+\frac{B}{\sigma_1},\quad A\equiv\frac{2\lambda_5\lambda_2+\lambda_6(\lambda_4-\lambda_3)}{\lambda_6^2-8\lambda_2\lambda_4},\quad B\equiv\frac{\lambda_6\Delta}{\lambda_6^2-8\lambda_2\lambda_4}. \quad (21)$$

Let us now determine the critical value of the chemical potential, namely, the value $\mu_{\rm crit}$, where $\rho(\mu_{\rm crit})=0$, but Eqs. (19), (20), and (21) hold:

$$\lambda_6x^2+4\lambda_4x+\lambda_5=0,\quad x=\frac{\sigma_2}{\sigma_1}. \quad (22)$$

In order for an SPB phase to exist, this equation has to possess real solutions, i.e., $4\lambda_4^2\geqslant\lambda_5\lambda_6$. We stress that Eqs. (21) and (22) contain only the constants of the potential and do not depend on temperature and chemical potential manifestly.

Once we find $x_{\rm crit}=x_\pm(\Delta,\lambda_2,\ldots,\lambda_6)$, one can immediately calculate

$$\sigma_1^\pm(\Delta,\lambda_j)=\sqrt{\frac{B}{x_\pm-A}},\quad \sigma_2^\pm(\Delta,\lambda_j)=x_\pm\sigma_1^\pm. \quad (23)$$

After substituting these values into Eq. (9), one derives the boundary of the P-violation phase

$$\mathcal{N}\mathcal{A}(\sigma_1^\pm,\mu,\beta)=\Delta-2\lambda_1(\sigma_1^\pm)^2-\lambda_5\sigma_1^\pm\sigma_2^\pm-(\lambda_3-\lambda_4)(\sigma_2^\pm)^2, \quad (24)$$

which is a positive combination. The relation (24) defines a P-breaking divide line in the T–μ plane. From (10) one can obtain that $\mathcal{A}>0$ and $\mathcal{A}\to\infty$ when

$T, \mu \to \infty$. It means that for any nontrivial solution $x_\pm, \sigma_1^\pm, \sigma_2^\pm$ $\mathcal{NA}(\sigma_1^\pm, \mu, \beta) > 0$, the P-breaking phase boundary exists. If the phenomenon of P violation is realized for zero temperature, it will take place in a domain involving lower chemical potentials but higher temperatures.

Once a condensate for π_2^0 appears spontaneously, the vector $SU(2)$ symmetry is broken to $U(1)$ and two charged excited π' mesons are expected to possess zero masses.

Quantitatively, the mass spectrum can be obtained only after kinetic terms are normalized. We just note that in the SPB phase the situation is rather peculiar: pseudoscalar states mix with scalar ones. In particular, the diagonalization of kinetic terms is different for neutral and charged pions because the vector isospin symmetry is broken: $SU(2)_V \to U(1)$. This triggers a rather exotic mechanism of isospin breaking via different decay constants. Even in the massless pion sector the isospin breaking $SU(2)_V \to U(1)$ occurs: neutral pions become less stable with a larger decay constant. We refer the reader to [5] and [6] for details.

SPB also induces mixing of both massless and heavy neutral pions with scalars. In fact, in the SPB phase, parity is no longer a conserved quantity in strong interactions, so the distinction between scalars and pseudoscalars is immaterial. This is why while the global broken symmetry at the point of transition to the SPB is a vector one, the two Goldstone bosons are apparently pseudoscalars, but, as emphasized, the distinction is purely semantic once parity is broken.

CONCLUSIONS

Let us summarize here our main findings. Parity violation seems to be quite a realistic possibility in nuclear matter at moderate densities. We have arrived at this conclusion by using an effective Lagrangian for low-energy QCD that retains the two lowest lying states in the scalar and pseudoscalar sectors. We include a chemical potential for the quarks that corresponds to a finite density of baryons, implement the bound state of normal nuclear matter and investigate the pattern of symmetry violation in its presence. We have found the necessary and sufficient conditions for a phase where parity is spontaneously broken to exist.

Salient characteristics of this phase would be the spontaneous violation of isospin and the generation of two additional massless charged pseudoscalar mesons. We have also examined departures from the chiral limit, i.e., allowing for nonzero quark masses. This leads to rather interesting results as in this case the usual pions are not exactly massless, but the new Goldstone bosons appearing at the transition point to the parity-violating phase are. Strong interaction phenomenology becomes indeed very unfamiliar at that point.

We also find a strong mixing between scalar and pseudoscalar states that translate spontaneous parity violation into meson decays. The mass eigenstates will decay in both odd and even number of pions simultaneously. Isospin violation can also be visible in decay constants.

So far there is no sufficient experimental information to be able to fully determine the value of the nine low energy constants appearing in $V_{\rm eff}$ and one could hope that lattice methods [13] may shed some light on this issue and confirm or falsify the existence of this interesting phase in dense nuclear matter.

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