Fluctuations of conserved charges, chiral-spin symmetry, and deconfinement in an $SU(2)_{\text{color}}$ subgroup of $SU(3)_{\text{color}}$ above T_c

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Above a pseudocritical temperature of chiral symmetry restoration T_c , the energy and the pressure are very far from the quark-gluon-plasma limit (i.e., ideal gas of free quarks and gluons). At the same time, very soon above T_c , fluctuations of conserved charges behave as if quarks were free particles. Within the $T_c - 3T_c$ interval, a chiral-spin symmetry emerges in QCD which is not consistent with free quarks and suggests that degrees of freedom are chirally symmetric quarks bound into the color-singlet objects by the chromoelectric field. Here we analyze temporal and spatial correlators in this interval and demonstrate that they indicate simultaneously the chiral-spin symmetry as well as the absence of the interquark interactions in channels constrained by a current conservation. The latter channels are responsible for both fluctuations of conserved charges and for dileptons. Assuming that an $SU(2)_{\text{color}}$ subgroup of $SU(3)_{\text{color}}$ is deconfined soon above T_c but confinement persists in $SU(3)_{\text{color}}/SU(2)_{\text{color}}$ in the interval $T_c - 3T_c$, we are able to reconcile all empirical facts listed above.

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I. INTRODUCTION

It is understood that QCD at low temperatures and vanishing baryon chemical potential is a hadron (meson) gas. At temperatures between 100 and 200 MeV, a very smooth chiral symmetry restoration crossover is observed, and a temperature $T_c \sim 155$ MeV could be approximately considered as a pseudocritical temperature of chiral symmetry restoration [1,2]. There is no definition and an order parameter for deconfinement in QCD with light quarks except that deconfinement should be accompanied by a free motion of colored quarks and gluons within a macroscopical piece of matter. Such a situation should take place at very high temperatures where due to asymptotic freedom the strong interaction coupling vanishes. What the strongly interacting matter is above the chiral symmetry restoration is a big puzzle. It cannot be a quark-gluon-plasma (QGP) that is defined as a gas of free (i.e., deconfined) quarks and gluons. For example, a perfect fluidity of the QCD matter above T_c is not consistent with free quarks and gluons as degrees of freedom. Other evidence that QCD is strongly interacting matter above T_c is its pressure and energy density that are very far from the Stefan-Boltzmann limit, i.e., the limit of free noninteracting quarks and gluons above T_c but below $3T_c$ [3,4].

At the same time, fluctuations of conserved charges [5–8] approach the limit of noninteracting quarks very soon after T_c ; see Fig. 1. This behavior of fluctuations is considered sometimes as a signal of deconfinement. But how to reconcile this with pressure and energy density as well as with perfect fluidity?

It was recently found on the lattice [9–11] that temporal and spatial correlators of $N_F = 2$ QCD reveal above T_c but below $3T_c$ multiplet patterns of chiral-spin $SU(2)_{CS}$ and $SU(2N_F)$ groups [12,13]. These groups are not symmetries of the Dirac action and hence inconsistent with free (deconfined) quarks. At the same time, they are symmetries of the Lorentz-invariant color charge and of the chromoelectric interaction in a given reference frame. Observation of these symmetries suggests that in the $T_c - 3T_c$ interval, the degrees of freedom are chirally symmetric quarks bound into color-singlet objects by the chromoelectric field, and effects of the chromomagnetic field are at least strongly suppressed. The chemical potential term in the QCD action is manifestly chiral-spin symmetric, which means that this symmetry should persist at a finite chemical potential [14]. These symmetries prohibit a finite axial chemical potential, and consequently, an electric current in QCD matter induced by an external magnetic field should be either absent or very small [15]. The experimental search of the chiral magnetic effect [16,17] at the RHIC and LHC hints that indeed the magnetically induced electric current either vanishes or is very small [18]. If confirmed, it would be an experimental verification of emerged chiral-spin symmetry above T_c .

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FIG. 1. Left: The ratio of fourth and second order cumulants of net-baryon number fluctuations versus the temperature. Right: Same as the left-hand side, but for the ratio of sixth and second order cumulants of net-baryon number fluctuations B . The boxes indicate the transition region $T_c = (154 \pm 9)$ MeV. Gray bands show the continuum estimate. From Ref. [8].

A very strange and self-contradictory picture emerges. On the one hand, there are signals that quarks are deconfined soon above T_c , as it follows from fluctuations of conserved charges and from the absence of ρ -like bound states above T_c seen via dilepton production. At the same time, quarks cannot be deconfined because of very clear chiral-spin symmetry in QCD above T_c and because of the energy density and pressure and perfect fluidity properties. In this paper, we reanalyze existing correlators [9–11] and demonstrate that while correlators show clear patterns of the chiral-spin symmetry, they are simultaneously consistent with the absence of the interquark interactions in channels constrained by a conserved current. Hence, these correlators are compatible with the known results about fluctuations of conserved charges and the absence of bound states as seen by dilepton production. This suggests that we have evidence of confinement of quarks and at the same time their deconfinedlike behavior in channels controlled by a conserved current. This paradox could be explained by the assumption that while deconfinement, i.e., a screening of the color charge happens above T_c in the $SU(2)_{\text{color}}$ subgroup of $SU(3)_{\text{color}}$, it does not happen in the $SU(3)_{\text{color}}/SU(2)_{\text{color}}$. This way, we automatically obtain the chiral-spin symmetry of correlators, and at the same time, the absence of interactions in channels constrained by a conserved current. This assumption allows us to reconcile all empirical facts listed above.

II. CHIRAL-SPIN SYMMETRY AS A SYMMETRY OF THE COLOR CHARGE

The $SU(2)_{CS}$ chiral-spin (CS) transformations [12,13] are defined by

$$\begin{aligned}\psi(x) &\rightarrow \exp\left(\frac{i}{2}\vec{\Sigma}\vec{\epsilon}\right)\psi(x), \\ \bar{\psi}(x) &\rightarrow \bar{\psi}(x)\gamma_4 \exp\left(-\frac{i}{2}\vec{\Sigma}\vec{\epsilon}\right)\gamma_4,\end{aligned}\quad (1)$$

where $\vec{\epsilon}$ are the rotation parameters.¹ For the generators $\vec{\Sigma}$, there are four different choices of $\vec{\Sigma} = \vec{\Sigma}_k$ with $k = 1, 2, 3, 4$. The choice of k for a given observable is constrained by the rotational $O(3)$ symmetry. Namely, the CS transformations should not mix operators with a different angular momentum. The generators are given by

$$\vec{\Sigma}_k = \{\gamma_k, -i\gamma_5\gamma_k, \gamma_5\}, \quad (2)$$

and the $SU(2)$ algebra is satisfied for any choice $k = 1, 2, 3, 4$. Here we use the set γ_μ , $\mu = 1, 2, 3, 4$ of Hermitian Euclidean γ matrices that fulfill the anticommutation relations

$$\gamma_\mu\gamma_\nu + \gamma_\nu\gamma_\mu = 2\delta_{\mu\nu}, \quad \gamma_5 \equiv \gamma_1\gamma_2\gamma_3\gamma_4. \quad (3)$$

The CS symmetry is not a symmetry of the Dirac Lagrangian; hence, it cannot exist for free deconfined quarks. At the same time, it is a symmetry of the Lorentz-invariant color charge in QCD

$$Q_c^a = \int d^3x \psi^\dagger(x) \frac{t^a}{2} \psi(x), \quad a = 1, 2, \dots, N_c^2 - 1, \quad (4)$$

where t^a are the color $SU(N_c)$ generators. The color charge is invariant under a unitary transformation that acts only in the Dirac space. The gluon part of the color charge is automatically CS invariant.

This symmetry of the color charge has important implications. Namely, the chromoelectric field in a given reference frame is defined via its action on the probe color charge. Since the color charge is invariant under the CS transformations, the chromoelectric interaction in a given reference frame is also CS invariant. The chromomagnetic field is defined via its action on the color spatial current

¹This symmetry was reconstructed from a hadron degeneracy observed on the lattice upon artificial truncation of the near-zero modes of the Dirac operator [19,20].

$\bar{\psi}\gamma^i \frac{t^a}{2}\psi$. The latter current is not CS invariant. Hence, the chromomagnetic and chromoelectric interactions in a given reference frame can be distinguished by the CS symmetry. Obviously, to discuss the chromoelectric and chromomagnetic parts of the gluonic field, one needs to fix a reference frame since they transform through each other upon a Lorentz transformation.

This can be made more explicit. In Minkowski space in a given reference frame the chromoelectric and chromomagnetic fields are different fields. The quark-gluon interaction Lagrangian can be split into temporal and spatial parts:

$$\bar{\psi}\gamma^\mu D_\mu\psi = \bar{\psi}\gamma^0 D_0\psi + \bar{\psi}\gamma^i D_i\psi, \quad (5)$$

where

$$D_\mu\psi = \left(\partial_\mu - ig\frac{t \cdot \mathbf{A}_\mu}{2}\right)\psi. \quad (6)$$

The temporal term includes the interaction of the color-octet charge density

$$\bar{\psi}(x)\gamma^0 \frac{t^a}{2}\psi(x) = \psi(x)^\dagger \frac{t^a}{2}\psi(x) \quad (7)$$

with the chromoelectric component A_0^a of the gluonic field. It is invariant under $SU(2)_{CS}$. We stress that the $SU(2)_{CS}$ transformations defined in Eq. (1) via the Euclidean Dirac matrices can be identically applied to Minkowski Dirac spinors without any modification of the generators. The spatial part contains the quark kinetic term and the interaction with the chromomagnetic field. This term breaks $SU(2)_{CS}$. In other words, the $SU(2)_{CS}$ symmetry distinguishes between quarks interacting with the chromoelectric and chromomagnetic components of the gauge field. We note that discussing “electric” and “magnetic” components can be done only in Minkowski space, and in addition, one needs to fix a reference frame. At high

temperatures, Lorentz invariance is broken and a natural frame to discuss the physics is the rest frame of the medium.

The group $SU(2)_{CS} \otimes SU(N)_F$ can be extended to $SU(2N_F)$ with generators

$$\{(\vec{\tau} \otimes \mathbb{1}_D), (\mathbb{1}_F \otimes \vec{\Sigma}_k), (\vec{\tau} \otimes \vec{\Sigma}_k)\}. \quad (8)$$

The corresponding transformations are a straightforward generalization of Eq. (1) obtained by replacing the generators $\vec{\Sigma}$ by those listed in (8). Also, the group $SU(2N_F)$ is a symmetry of the color charge (4) as well as of the quark-chromoelectric interaction terms of the QCD Lagrangian.

III. REVIEW OF THE TEMPORAL CORRELATORS ABOVE T_c

The $SU(2)_L \times SU(2)_R$ and $U(1)_A$ transformation properties of the $J = 1$ operators relevant for temporal correlators are given in the left panel of Fig. 2, while their $SU(2)_{CS}$ and $SU(4)$ multiplets are presented in the right panel of this figure [13].

On the rhs of Fig. 3, we show temporal correlators

$$C_\Gamma(t) = \sum_{x,y,z} \langle \mathcal{O}_\Gamma(x, y, z, t) \mathcal{O}_\Gamma(\mathbf{0}, 0)^\dagger \rangle \quad (9)$$

at a temperature $T = 220$ MeV ($1.2T_c$) calculated in $N_F = 2$ QCD with a chirally symmetric domain wall Dirac operator with physical quark masses [11]. Here, $\mathcal{O}_\Gamma(x, y, z, t)$ is an operator that creates a quark-antiquark pair with fixed quantum numbers. Summation over x, y, z projects out the rest frame.

Above the chiral restoration crossover, we *a priori* expect in observables the chiral $SU(2)_L \times SU(2)_R$ symmetry. This symmetry is evidenced by degeneracy of correlators of the operators $\rho_{(0,1)+(1,0)}$ and a_1 connected by the $SU(2)_L \times SU(2)_R$ transformation. While the axial

FIG. 2. Transformations between $J = 1$ operators, $i = 1, 2, 3$. The left columns indicate the $SU(2)_L \times SU(2)_R$ representation for every operator. Red and blue arrows connect operators which transform into each other under $SU(2)_L \times SU(2)_R$ and $U(1)_A$, respectively. Green arrows connect operators that belong to $SU(2)_{CS}$, $k = 4$ triplets. Purple arrow shows the $SU(4)$ 15-plet. The f_1 operator is a singlet of $SU(4)$. From Ref. [13].

FIG. 3. Temporal correlation functions for 12×48^3 lattices. The lhs shows correlators calculated with free noninteracting quarks with manifest $U(1)_A$ and $SU(2)_L \times SU(2)_R$ symmetries. The rhs presents full QCD results at a temperature $T = 220$ MeV ($1.2T_c$), which shows multiplets of all $U(1)_A$, $SU(2)_L \times SU(2)_R$, $SU(2)_{CS}$, and $SU(4)$ groups. From Ref. [11].

anomaly is a pertinent property of QCD, its effect is determined by the topological charge density. There is strong evidence from the lattice that the $U(1)_A$ symmetry is also effectively restored above T_c [21]. The $U(1)_A$ restoration is signaled by degeneracy of the correlators of the isovector scalar (S) and isovector pseudoscalar (PS) operators connected by the $U(1)_A$ transformation. It is also evidenced by the degeneracy of the correlators with b_1 and $\rho_{(1/2,1/2)_b}$ operators.

An approximate degeneracy of the b_1 , $\rho_{(1,0)+(0,1)}$ and $\rho_{(1/2,1/2)_b}$ correlators indicates emergent $SU(2)_{CS}$ symmetry since these three operators form a CS triplet. Its breaking is estimated at a level of less than 5%. A degeneracy of all four $J = 1$ correlators b_1 , $\rho_{(1,0)+(0,1)}$, $\rho_{(1/2,1/2)_b}$, and a_1 on the rhs of Fig. 3 evidences the emergence of the $SU(4)$ symmetry.

On the lhs of Fig. 3, we present correlators calculated with noninteracting quarks on the same lattice. They represent a QGP at a very high temperature. Free quarks are governed by the Dirac equation and only $U(1)_A$ and $SU(2)_L \times SU(2)_R$ chiral symmetries exist. Indeed, here we observe exact degeneracy of all correlators connected by $U(1)_A$ and $SU(2)_L \times SU(2)_R$ transformations.

A qualitative difference between the pattern on the lhs and the pattern on the rhs of Fig. 3 is obvious. While for free quarks (i.e., deconfined quarks) on the lhs of Fig. 3 we observe multiplets of $U(1)_A$ and $SU(2)_L \times SU(2)_R$ groups, on the rhs of the same figure we clearly see multiplets of all $U(1)_A$, $SU(2)_L \times SU(2)_R$, $SU(2)_{CS}$, and $SU(4)$ groups.

The $SU(2)_{CS}$ and $SU(4)$ groups are not symmetries of the Dirac action, and hence, they are incompatible with free deconfined quarks. Since these are symmetries of the color charge in QCD and of the chromoelectric interaction, while the chromomagnetic interaction breaks them, we conclude that elementary objects are the chirally symmetric quarks bound by the chromoelectric field into the color-singlet

objects, and a contribution of the chromomagnetic field is at least very strongly suppressed.

All these properties have been discussed in Ref. [11]. Now we add a very important observation. While correlators of most operators are very different as compared to the noninteracting quark system, this is not true for the correlators of the $\rho_{(1,0)+(0,1)}$ operator. The correlators of the latter operator on the rhs and on the lhs are very close. The vector $\rho_{(1,0)+(0,1)}$ operator is constrained by the current conservation:

$$\partial_\mu \left(\bar{\psi} \frac{\tau^a}{2} \gamma^\mu \psi \right) = 0. \quad (10)$$

The current conservation is a consequence of a global $SU(2)_V$ isospin symmetry. While in general, the axial vector current a_1 is not constrained by the current conservation, it is in the chirally symmetric regime and the correlators of the vector and axial vector operators are identical.

These results suggest that while the interquark interaction is strong and the system in QCD above T_c is not a system of free quarks (the latter would require that correlators in QCD and for free quarks were identical for all possible quantum numbers), the interquark interaction in channels that are constrained by the current conservation is absent or nearly absent. How could it be?

Needless to add that it is a conserved vector current that drives a dilepton production. In other words, the absence of a structure seen in a dilepton production is consistent with the chiral-spin symmetry.

IV. REVIEW OF THE SPATIAL CORRELATORS ABOVE T_c

The spatial correlators of the isovector operators $\mathcal{O}_T(0,0)$,

TABLE I. Fermion isovector bilinears and their $U(1)_A$ and $SU(2)_L \times SU(2)_R$ transformation properties (last column). This classification assumes propagation in the z direction. The open vector index k here runs over the components 1,2,4, i.e., x , y , and t .

Name	Dirac structure	Abbreviation	
Pseudoscalar	γ_5	PS	] $U(1)_A$
Scalar	$\mathbb{1}$	S	
Axial vector	$\gamma_k \gamma_5$	$\mathbf{A}$	] $SU(2)_A$
Vector	γ_k	$\mathbf{V}$	
Tensor vector	$\gamma_k \gamma_3$	$\mathbf{T}$	] $U(1)_A$
Axial-tensor vector	$\gamma_k \gamma_3 \gamma_5$	$\mathbf{X}$	

$$C_\Gamma(z) = \sum_{x,y,t} \langle \mathcal{O}_\Gamma(x, y, z, t) \mathcal{O}_\Gamma(\mathbf{0}, 0)^\dagger \rangle \quad (11)$$

in $N_F = 2$ QCD with the domain wall Dirac operator at physical quark masses have been calculated in Refs. [9,10]. The isovector fermion bilinears are named according to Table I. The change in the names of bilinears as compared to the temporal correlators is because a propagation along the t axis and along a spatial axis requires different gamma structures of operators.

At finite temperature, the rotational $SO(2)$ symmetry connects the correlators of the spatial components $V_x \leftrightarrow V_y$, $A_x \leftrightarrow A_y$, etc. The CS transformations (1) with $k = 1, 2$ together with the $x \leftrightarrow y$ symmetry generate the following multiplets:

$$(V_x, V_y); (A_x, A_y, T_t, X_t), \quad (12)$$

$$(V_t); (A_t, T_x, T_y, X_x, X_y). \quad (13)$$

Extending the $SU(2)_{CS}$ group to $SU(4)$, one obtains larger multiplets of the isovector operators:

$$(V_x, V_y, A_x, A_y, T_t, X_t), \quad (14)$$

$$(V_t, A_t, T_x, T_y, X_x, X_y). \quad (15)$$

Complete $SU(4)$ multiplets include also the isoscalar partners of the operators A_x, A_y, T_t, X_t in Eq. (14) as well as isoscalar partners of the A_t, T_x, T_y, X_x, X_y operators in Eq. (15).

In Fig. 4, we show correlators (full lines) calculated in $N_F = 2$ QCD at $T = 380$ MeV ($\sim 2.2T_c$) normalized to 1 at $n_z = 1$ where n_z is the dimensionless distance in lattice units [9]. The dashed lines represent the calculation with free noninteracting quarks (i.e., they represent a QGP at an asymptotically high temperature).

Degeneracy of the S and PS correlators indicates the $U(1)_A$ symmetry. The same is true for degenerate T_t and X_t correlators that are connected by $U(1)_A$. Equality of the V_x and A_x correlators is because of the restored

FIG. 4. Correlators of the $PS, S, V_x, A_x, T_t, X_t$ operators in full QCD at $T = 380$ MeV ($\sim 2.2T_c$) for an 8×32^3 lattice (abbreviated as *dressed*) and with noninteracting quarks (*free*) on the same lattice. From Ref. [9].

$SU(2)_L \times SU(2)_R$ symmetry. An approximate degeneracy of the (A_x, T_t, X_t) correlators evidences the $SU(2)_{CS}$ symmetry, while an approximate degeneracy of all four correlators (V_x, A_x, T_t, X_t) reflects the $SU(4)$ symmetry. This multiplet structure is very different from the pattern of free correlators that demonstrates only $U(1)_A$ and $SU(2)_L \times SU(2)_R$ symmetries. All these features have been discussed in detail in Refs. [9,10].

In Fig. 5, we show correlators normalized to 1 at $n_z = 1$ built with the V_t, A_t, T_x, X_x operators. A degeneracy according to the multiplets (13) and (15) is obvious. Now comes an important point. The correlators of the V_t and A_t operators coincide with their free counterparts. This should not be true for operators T_x and X_x . Their free correlators are different from the free correlators V_t and A_t exactly by a factor of 2 as it follows from analytical continuum calculations [10]. The free correlators T_x, X_x look degenerate with the free correlators V_t, A_t only because of the normalization of all correlators to 1 at $n_z = 1$ in Fig. 5.²

What distinguishes the V_t (and A_t) operator from the other ones in Figs. 4 and 5? This operator is constrained by the current conservation (10).³

²Note that while both dashed curves on Fig. 5 must in reality exactly coincide after normalization to 1, they are slightly split only because of discretization errors on the lattice. Also, their bending down around $n_z = 12$ is a finite volume effect on the lattice.

³The $V_x(A_x)$ and $V_y(A_y)$ operators are transverse to the propagation direction z and hence are not constrained by (10).

FIG. 5. Correlators of the V_t , A_t , T_x , X_x operators in full QCD at $T = 380$ MeV ($\sim 2.2T_c$) for an 8×32^3 lattice (abbreviated as *dressed*) and with noninteracting quarks (*free*) on the same lattice. From Ref. [9].

Given the results depicted in Figs. 4 and 5, as well as the discussion above, we conclude that similar to the temporal correlators, one observes multiplets of the chiral spin and $SU(4)$ groups, and at the same time, correlators constructed with operators that are constrained by the $SU(2)_V$ current conservation show an absence of the interquark interaction.

The chiral-spin and $SU(4)$ symmetries are not symmetries of the Dirac action and are incompatible with free deconfined quarks. Indeed, they are symmetries of the color charge (4) and consequently indicate confinement. Deconfinement can happen only when the color charge is screened, and consequently, the CS and $SU(4)$ symmetries would disappear. But what does it mean when at the same time in channels that are constrained by a current conservation, the temporal and spatial correlators demonstrate a free-like behavior? We will answer this question in the next section.

Now we will address the fluctuations of conserved charges that have been mentioned in the Introduction. These fluctuations are typically obtained via chemical potential derivatives of pressure [7,8]. However, they can also be obtained in an alternative way.⁴

The Lorentz-invariant conserved quark (baryon) $U(1)_V$ charge Q and flavor (isospin) $SU(N_F)_V$ charge Q_F^a are defined as spatial integrals of the corresponding charge densities:

$$Q = \int dx dy dz \bar{\psi}(x, y, z, t) \gamma^4 \psi(x, y, z, t), \quad (16)$$

$$Q_F^a = \int dx dy dz \bar{\psi}(x, y, z, t) \gamma^4 \frac{\tau^a}{2} \psi(x, y, z, t), \quad (17)$$

$$a = 1, 2, \dots, N_F^2 - 1,$$

where τ^a are the isospin (flavor) generators. The fluctuations of the conserved charges are completely determined by the spatial equal time correlators of the charge density operator V_t for both flavor singlet and nonsinglet operators. This is because the quark (baryon) charge Q (16) and the flavor charge Q_F^a (17) are spatial integrals of the corresponding charge densities. Consequently, the spatial equal time correlators of the quark (baryon) charge density and of the flavor charge density should contain complete information about the fluctuations of conserved charges.

The spatial correlator V_t of conserved charge shown in Fig. 5 (the flavor singlet correlator of the quark charge density that is not shown in Fig. 5 should be exactly the same) is identical with its free quark counterpart. In this correlator, an integration over time is performed. From this equivalence, it follows that the spatial equal time correlator in full QCD will also coincide with its free quark counterpart. Since the spatial equal time correlators contain complete information about fluctuations of conserved charges, it follows that soon above T_c the fluctuations of conserved charges should demonstrate an absence of the interquark interactions, in agreement with Fig. 1.

Now we summarize both this and the previous sections. While the temporal and spatial correlators above T_c but below $3T_c$ demonstrate clear patterns of the $SU(2)_{CS}$ and $SU(4)$ symmetries, which are symmetries of confinement and are incompatible with free quarks, at the same time these correlators show very soon above T_c an absence of the interquark interactions in channels that are constrained by a current conservation. The latter channels are responsible for both dilepton production and for fluctuations of conserved charges. How do we reconcile these seemingly contradicting facts? This will be discussed in the subsequent section.

V. DECONFINEMENT IN $SU(2)_c$ AND CONFINEMENT IN $SU(3)_c/SU(2)_c$ IN QCD

We certainly do not know a microscopic mechanism that leads to the present puzzling situation. However, from the symmetry grounds we can obtain some solid insights.

A confining interquark interaction in QCD persists until the color charge of quarks is not screened. Once it is screened, then there is no color force between quarks and they behave as free particles. We can judge from the chiral-spin [and $SU(2N_F)$] symmetry of observables the existence of confinement in QCD at high temperatures. They persist in the $T_c - 3T_c$ interval [9–11]. The chiral-spin $SU(2)_{CS}$

⁴The author is thankful to T. Cohen for pointing this out and to F. Karsch for correspondence on this issue.

symmetry is a symmetry in any $a = 1, \dots, 8$ component of the color charge (4). Consequently, observation of this symmetry in the $T_c - 3T_c$ interval implies that the color charge is not screened at least in some of the eight components. The absence of the strong interquark interactions in channels constrained by a current conservation requires that in these channels the color charge should be screened. This could happen if the color charge (4) was screened in an $SU(2)_c$ subgroup of $SU(3)_c$.

A strong color force is a consequence of a local gauge invariance. If the color charge is screened within an $SU(2)_c$, then the $SU(2)_c$ symmetry becomes global. A global $SU(2)_c$ symmetry implies via Noether theorem a conserved colored vector current,

$$\partial_\mu \left(\bar{\psi} \frac{\tau^a}{2} \gamma^\mu \psi \right) = 0, \quad a = 1, 2, 3. \quad (18)$$

The conserved baryon charge and isospin current also satisfy a continuity equation:

$$\partial_\mu (\bar{\psi} \gamma^\mu \psi) = 0, \quad (19)$$

$$\partial_\mu \left(\bar{\psi} \frac{\tau^a}{2} \gamma^\mu \psi \right) = 0, \quad a = 1, 2, 3. \quad (20)$$

Then, one needs a condition that would connect the color current conservation within $SU(2)_c$ with the baryon (19) and isospin current (20) conservations. Such a condition would be provided by the color-isospin locking, i.e., by the locking of the color index of quarks with their isospin index within an $SU(2)_c$ subgroup, and such a locking should be absent in $SU(3)_c/SU(2)_c$. This locking would automatically provide an equivalence of Eq. (18) to Eq. (20). I.e., the absence of a strong confining interquark interaction happens only in channels with conserved current. What a microscopical reason for such locking could be is an open question.

This kind of locking supplemented with some dynamics was discussed in a different context for the color-flavor locking at extremely high baryon density [22,23]. We stress that we use a locking only as a mathematical device that allows us to combine all facts discussed in this paper. Its precise physical meaning should be clarified in subsequent studies.

It is not unreasonable to assume a locking soon after T_c . The above scenario of the color-isospin locking guarantees that a deconfinement within an $SU(2)_c$ subgroup of $SU(3)_c$ can be seen only in channels with conserved current. Equation (18) is not true for all eight generators of $SU(3)_c$. Consequently, within the $SU(3)_c$, confinement is

still there. In this way, we obtain simultaneously a chiral-spin symmetry of correlators (i.e., confinement) and correlators of operators with a conserved current behave as if quarks were free. This means that the color-isospin locking filters out an $SU(2)_c$ subgroup, and the $SU(3)_c/SU(2)_c$ part of the color dynamics is not seen via vector and axial vector operators.

Needless to add that this scenario implies that a dual superconductor picture of confinement in QCD [24–26] that is an Abelian $U(1)$ confinement is quite far from real life in QCD because it implies that above a critical temperature quarks are free. In QCD, there is still confinement in $SU(3)_c$.

A general possibility that there could be at high temperatures a partial deconfinement, i.e., that an $SU(M)$ subgroup of the $SU(N_c)$, $M < N_c$ is deconfined above T_c , was also discussed in Refs. [27–30].

VI. CONCLUSIONS

We have reanalyzed the results on spatial and temporal correlators above T_c [9–11] and demonstrated that while these correlators exhibit chiral-spin $SU(2)_{CS}$ and $SU(2N_F)$ symmetries at temperatures $T_c - 3T_c$, they also show the absence of the interquark interactions in channels constrained by a current conservation. The chiral spin and $SU(2N_F)$ are not symmetries of the Dirac action and hence are incompatible with free deconfined quarks. They are symmetries of the color charge in QCD and indicate that the theory is still in the confining regime.

The absence of the interquark interactions in channels with a conserved current explains why fluctuations of conserved charges soon after T_c behave as if quarks were free. It also explains why no bound states are seen via dileptons above T_c .

In order to combine all these facts, we have conjectured that in the temperature region $T_c - 3T_c$, an $SU(2)_{\text{color}}$ subgroup of $SU(3)_{\text{color}}$ is deconfined; i.e., a screening happens in three components of the color charge, and five generators of $SU(3)_{\text{color}}$ are still confined, i.e., not screened. Assuming an $SU(2)_{\text{color}}-SU(2)_{\text{isospin}}$ locking in one of the $SU(2)_{\text{color}}$ subgroups of $SU(3)_{\text{color}}$, we have explained why the chiral-spin symmetry persists, and at the same time, there is no $SU(2)_{\text{color}}$ color-mediated strong force in channels with conserved current. These channels are still in a confining mode since there is confinement in $SU(3)_{\text{color}}/SU(2)_{\text{color}}$.

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