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

L.Ya. Glozman

1Institute of Physics, University of Graz, 8010 Graz, Austria
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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 analyse temporal and spatial correlators in this interval and demonstrate that they indicate simultaneously the chiral spin symmetry as well as 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 a $SU(2)_{\text{color}}$ 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.
I. INTRODUCTION

It is understood that QCD at low temperatures and vanishing baryon chemical potential is a hadron (meson) gas. At temperatures between 100 MeV 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 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 macroscopic 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. Another evidence that QCD is a 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 $3T_c$.

At the same time fluctuations of conserved charges [3–8] 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 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 either [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 RHIC and LHC hints that indeed the magnetically induced electric current either vanishes or very small [18]. If confirmed, it would be an experimental verification of emerged chiral spin symmetry above $T_c$.

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 absence of $\rho$-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 $T_c$ as well as because of the energy density and pressure and perfect fluidity properties. In this paper we reanalyse existing correlators [9–11] and demonstrate that while correlators show clear patterns of the chiral spin symmetry, they are simultaneously consistent with absence of the interquark interactions in channels constrained by a conserved current. Hence these correlators are compatible with known results about fluctuations of conserved charges and absence of bound states as seen by dilepton production. This suggests that we have an evidence of confinement of quarks and at the same time their deconfined-like 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 $SU(2)_{color}$ subgroup of $SU(3)_{color}$, it does not happen in the $SU(3)_{color}/SU(2)_{color}$. This way we automatically obtain the chiral spin symmetry of correlators and at same time absence of interactions in channels constrained by a conserved current. This assumption allows to reconcile all empirical facts listed above.

FIG. 1. Left: The ratio of fourth and second order cumulants of net-baryon number fluctuations versus 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. Grey bands show continuum estimate. From ref. [8].
II. CHIRAL SPIN SYMMETRY AS A SYMMETRY OF THE COLOR CHARGE

The $SU(2)_{CS}$ chiral spin (CS) transformations \cite{12 13} are defined by

$$\psi(x) \to \exp \left( \frac{i \vec{\Sigma} \vec{\epsilon}}{2} \right) \psi(x), \quad \bar{\psi}(x) \to \bar{\psi}(x) \gamma_4 \exp \left( -\frac{i \vec{\Sigma} \vec{\epsilon}}{2} \right) \gamma_4,$$

where $\vec{\epsilon}$ are the rotation parameters. For the generators $\vec{\Sigma}$ there are four different choices $\vec{\Sigma} = \vec{\Sigma}_k$ 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 different angular momentum. The generators are given by

$$\vec{\Sigma}_k = \{ \gamma_k, -i\gamma_5\gamma_k, \gamma_5 \},$$

and the $su(2)$ algebra is satisfied for any choice $k = 1, 2, 3, 4$. Here we use the set $\gamma_\mu, \mu = 1, 2, 3, 4$ of hermitian Euclidean $\gamma$-matrices that fulfill the anti-commutation relations

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

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

$$Q^a_c = \int d^3x \bar{\psi}(x) \frac{t^a}{2} \psi(x), \quad a = 1, 2, ..., N^2_c - 1,$$

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 $\bar{\psi} \gamma^i \vec{A}_i \psi$. 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 gluonic field one needs to fix a reference frame since they transform through each other 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:

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

where

$$D^i_\mu \psi = \left( \partial_\mu - ig \frac{t^a}{2} A^a_\mu \right) \psi.$$

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

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

with the chromo-electric component $A^a_0$ of the gluonic field. It is invariant under $SU(2)_{CS}$. We stress that the $SU(2)_{CS}$ transformations defined in Eq. \cite{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

\footnote{This symmetry was reconstructed from a hadron degeneracy observed on the lattice upon artificial truncation of the near-zero modes of the Dirac operator \cite{19 20}.}
one needs to fix a reference frame. At high temperatures Lorentz invariance is broken and a natural frame to discuss physics is the rest frame of the medium.

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

$$\{ (\bar{\tau} \otimes 1_D), (1_F \otimes \bar{\Sigma}_k), (\bar{\tau} \otimes \bar{\Sigma}_k) \}.$$

The corresponding transformations are a straightforward generalization of Eq. (1) obtained by replacing the generators $\bar{\Sigma}$ 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 r.h.s. of Fig. 3 we show temporal correlators

$$C_T(t) = \sum_{x,y,z} \langle O_T(x,y,z,t) O_T(0,0) \rangle,$$

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 $O_T(x,y,z,t)$ 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 apriori 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 anomaly is a pertinent property of QCD its effect is determined by the topological charge density. There are strong evidences from the lattice that the $U(1)_A$ symmetry is also effectively restored above $T_c$ [24]. The $U(1)_A$ restoration is signalled by degeneracy of correlators of the isovector scalar (S) and isovector pseudoscalar (PS) operators connected by the $U(1)_A$ transformation. It is also evidenced by degeneracy of the correlators with $b_1$ and $\rho(1,1/2,1/2)$ operators.

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

On the l.h.s 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 l.h.s. and the pattern on the r.h.s. of Fig. 3 is obvious. While for free (i.e., deconfined quarks) on the l.h.s. of Fig. 3 we observe multiplets of $U(1)_A$ and $SU(2)_L \times SU(2)_R$ groups, on the r.h.s. 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 one additional and 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 r.h.s and on the l.h.s. are very close. The vector $\rho^{(1,0)+(0,1)}$ operator is constrained by the current conservation:

$$\partial_\mu (\bar{\psi} \gamma^\mu \gamma^a \psi) = 0.$$  \hspace{1cm} (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 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 $O_\Gamma(0,0)$

$$C_\Gamma(z) = \sum_{x,y,t} \langle O_\Gamma(x, y, z, t) O_\Gamma(0, 0)^\dagger \rangle$$  \hspace{1cm} (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 [1].
| Name                  | Dirac structure | Abbreviation |
|-----------------------|-----------------|--------------|
| **Pseudoscalar**      | $\gamma_5$      | $PS$         | $U(1)_A$   |
| Scalar                | $1$             | $S$          | $SU(2)_A$  |
| **Axial-vector**      | $\gamma_k\gamma_5$ | $A$         | $SU(2)_A$  |
| Vector                | $\gamma_k$      | $V$          | $U(1)_A$   |
| **Tensor-vector**     | $\gamma_k\gamma_3\gamma_5$ | $T$         | $U(1)_A$   |
| **Axial-tensor-vector** | $\gamma_k\gamma_3\gamma_5$ | $X$         | $U(1)_A$   |

**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 $z$-direction. The open vector index $k$ here runs over the components $1, 2, 4$, i.e., $x, y$ and $t$.

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

\[
(V_x, V_y) ; (A_x, A_y, T_t, X_t) , \quad (V_t) ; (A_t, T_x, T_y, X_x, X_y) .
\]

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

\[
(V_x, V_y, A_x, A_y, T_t, X_t) , \quad (V_t, A_t, T_x, T_y, X_x, X_y) .
\]

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).

On 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. The dashed lines represent 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 $U(1)_A$. Equility of the $V_x$ and $A_x$ correlators is because of the restored $SU(2)_L \times SU(2)_R$ 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 $U(1)$ and $SU(2)_L \times SU(2)_R$ symmetries. All these features have been discussed in detail in Refs. [9, 10].

On 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 the factor 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 normalization of all correlators to 1 at $n_z = 1$ on Fig. 5.

What distinguishes the $V_t$ (and $A_t$) operator from the other ones on Fig. 4 and Fig. 5? This operator is constrained by the current conservation $[10]$. Given results depicted on Fig. 4 and Fig. 5 as well as 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 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 that 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 fluctuations of conserved charges that have been mentioned in 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 $Q^a_F$ are defined as spatial integrals of the corresponding charge densities:

$$ Q = \int dxdydz \bar{\psi}(x,y,z,t)\gamma^0\psi(x,y,z,t), $$

$$ Q^a_F = \int dxdydz \bar{\psi}(x,y,z,t)\gamma^0\gamma^a_F\psi(x,y,z,t), \quad a = 1, 2, ..., N_F^2 - 1, $$

Note 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.

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

The author is thankful to T. Cohen for pointing this out and to F. Karsch for correspondence on this issue.
where $\tau^a$ 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 $Q$ and the flavor charge $Q_P$ 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 fluctuations of conserved charges.

The spatial correlator $V_t$ of conserved charge shown in Fig. 3 (the flavor singlet correlator of the quark charge density, that is not shown on Fig. 3 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 absence of the interquark interactions, in agreement with Fig. 1.

Now we summarize both this and 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 $T_c$ 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 to 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.

Confinement in QCD persists until the color charge is not screened. Once it is screened, then there is no color force between quarks and they behave as free particles. About existence of confinement in QCD at high temperatures we can judge from the chiral spin (and $SU(2N_F)$) symmetry of observables. They persist in the $T_c - 3T_c$ interval [9–11]. The chiral spin $SU(2)_{CS}$ symmetry is a symmetry in any $a = 1, \ldots, 8$ component of the color charge (1). 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. 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 (1) was screened in a $SU(2)_c$ subgroup of $SU(3)_c$.

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

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

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

$$\partial_\mu (\bar{\psi} \gamma^\mu \psi) = 0,$$  \hspace{1cm} (19)

$$\partial_\mu (\bar{\psi} \frac{\tau^a}{2} \gamma^\mu \psi) = 0, \quad a = 1, 2, 3.$$  \hspace{1cm} (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 a $SU(2)_c$ 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) with eqs. (19)-(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 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 a $SU(2)_c$ subgroup of $SU(3)_c$ can be seen only in channels with conserved current. In this
way we obtain simultaneously a chiral spin symmetry of correlators (i.e. confinement) and absence of the interquark interaction in channels with a current conservation.

The eq. \( \text{[18]} \) is not true for all eight generators of \( SU(3)_c \). Consequently within the \( SU(3)_c \) confinement is still there and only the \( SU(3)_c \)-singlets survive a gauge averaging and can freely propagate.

Needless to add that this scenario implies that a dual superconductor picture of confinement in QCD \([21, 22]\), 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 \( SU(3)_c \). It is a well known fact that on the lattice only the color-singlet states survive a gauge averaging above \( T_c \).

A general possibility that there could be at high temperatures a partial deconfinement, i.e. that a \( SU(M) \) 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 reanalysed results on spatial and temporal correlators above \( T_c \) \([3, 11]\) and demonstrated that while these correlators exhibit a chiral spin \( SU(2)_{CS} \) and \( SU(2N_F) \) symmetries at temperatures \( T_c - 3T_c \), they also show 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. Only the color-singlet objects can freely propagate.

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 \( T_c - 3T_c \), a \( SU(2)_{color} \) subgroup of \( SU(3)_{color} \) is deconfined, i.e. a screening happens in three components of the color charge, and five generators of \( SU(3)_{color} \) are still confined, i.e. not screened. Assuming a \( SU(2)_{color} - SU(2)_{isospin} \) locking in one of \( SU(2)_{color} \) subgroups of \( SU(3)_{color} \) we explain why the chiral spin symmetry persists and at the same time there is no color-mediated strong force in channels with conserved current. Since there is still confinement in \( SU(3)_{color}/SU(2)_{color} \) only the \( SU(3)_{color}\)-singlet states can propagate.

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