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# $SU(2N_F)$ symmetry of confinement in QCD and its observation at high temperature.

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**Abstract.** In this talk we first overview lattice results that have led to the observation of new $SU(2)_{CS}$ and $SU(2N_F)$ symmetries upon artificial truncation of the near-zero modes of the Dirac operator at zero temperatute and at high temperature without any truncation. These symmetries are larger than the chiral symmetry of the QCD Lagrangian and contain chiral symmetries $SU(N_F)_L \times SU(N_F)_R$ and $U(1)_A$ as subgroups. In addition to the standard chiral transformations the $SU(2)_{CS}$ and $SU(2N_F)$ transformations mix the right- and left-handed components of the quark fields. It is a symmetry of the confining chromo-electric interaction while the chromo-magnetic interaction manifestly breaks it. Emergence of these symmetries upon truncation of the near-zero modes of the Dirac operator at T=0 means that all effects of the chromo-magnetic interaction are located exclusively in the near-zero modes, while confining chromo-electric interaction is distributed among all modes. Appearance of these symmetries at high T, where the temperature suppresses the near-zero modes, has radical implications because these symmetries are incompatible with the asymptotically free deconfined quarks at increasing temperature. The elementary objects in the high-temperature phase of QCD should be quarks bound by the pure chromo-electric field that is not accompanied by the chromo-magnetic effects.

## 1 Introduction

The QCD Lagrangian with $N_F$ massles quarks

$$\mathcal{L} = \bar{\Psi}(x)(i\gamma_\mu D_\mu)\Psi(x) - \frac{1}{2}Tr(G^{\mu\nu}G_{\mu\nu}), \tag{1}$$

has the chiral symmetry:

$$U(N_F)_L \times U(N_F)_R = SU(N_F)_L \times SU(N_F)_R \times U(1)_A \times U(1)_V. \tag{2}$$

The $U(1)_V$ symmetry is responsible for the vector current conservation and is irrelevant to our subject. The $U(1)_A$ symmetry is an invariance upon the axial flavor-neutral transformations

$$\Psi(x) \to e^{i\alpha\gamma_5}\Psi(x); \quad \bar{\Psi}(x) \to \bar{\Psi}(x)e^{i\alpha\gamma_5}. \tag{3}$$

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The $SU(N_F)_L \times SU(N_F)_R$ chiral symmetry is an invariance under independent flavor $SU(N_F)$ rotations of the left- and right-handed components of quarks. These transformations contain the flavor (isospin for $N_F = 2$) rotations as well as the axial flavor transformations

$$\Psi(x) \rightarrow e^{i\gamma_5 \frac{\vec{\lambda}\cdot\vec{\alpha}}{2}} \Psi(x); \quad \bar{\Psi}(x) \rightarrow \bar{\Psi}(x) e^{i\gamma_5 \frac{\vec{\lambda}\cdot\vec{\alpha}}{2}}, \tag{4}$$

where $\vec{\lambda}$ are $SU(N_F)$ generators.

The $U(1)_A$ symmetry is broken anomalously, which is due to a noninvariance of the integration measure in the functional integral under a local $U(1)_A$ transformation [1]. The $SU(N_F)_A$ "symmetry" (the transformations (4) do not form a closed subgroup of the chiral group) is broken spontaneously (dynamically). The ground state, the vacuum, is not invariant under the transformation (4). This is encoded in the quark condensate, $< 0|\bar{\Psi}(x)\Psi(x)|0 > \neq 0$.

The quark condensate of the vacuum is connected to the density of the near-zero modes of the Euclidean Dirac operator via the Banks-Casher relation [2]

$$\lim_{m \to 0} < 0|\bar{\Psi}(x)\Psi(x)|0 >= -\pi\rho(0) . \tag{5}$$

Historically it was believed that the mass of hadrons in the $u, d$ quark sector is generated via the spontaneous breaking of chiral symmetry. For example, within the linear sigma-model the nucleon acquires its mass via its coupling with the chiral order parameter. Similar, within the Nambu and Jona-Lasinio model it is the vacuum fermion condensate that is responsible for a generation of a large mass of initially massless fermions. This view was one of the reasons to assume that at high temperatures, where chiral symmetry is restored, there should appear the quark-gluon plasma phase, contrary to the hadron phase at low temperatures.

One can artificially restore the flavor-nonsinglet chiral symmetry by removing the quark condensate by hands, i.e. via subtraction of the near-zero modes of the Dirac operator on the lattice [3]. It turned out that the nucleon, the rho-meson and some other hadrons, except for the pion, survive this "unbreaking" of the chiral symmetry and their mass remains large [4]. This tells that while the chiral symmetry breaking in the vacuum is important for the eventual shape of the hadron spectra, the chiral symmetry breaking is not the main mechanism of the mass generation of hadrons such as the $\rho$-meson or the nucleon. Still it was unclear what happens with the $U(1)_A$ symmetry upon truncation of the near-zero modes. This was a motivation for a systematic study of the hadron spectra upon truncation of the near-zero modes within $N_F = 2$ dynamical calculations with the chirally-invariant overlap Dirac operator[5–8] . This study has led to discovery of new $SU(2)_{CS}$ and $SU(4)$ symmetries that are larger than the chiral symmetry of the QCD Lagrangian and that allow mixing of the left- and right-handed components of the quark field [9, 10].

Below we consider observation of these symmetries and their implication. Then we report recent results on these symmetries at high temperature without any truncation [11] and their consequences for the nature of the strongly interacting matter at high T.

## 2 Truncation of the near-zero modes and naive expectations

The hermitian Euclidean Dirac operator, $i\gamma_\mu D_\mu$, where $D_\mu = \partial_\mu + ig\frac{t^a}{2}A^a_\mu$ has in a finite volume $V$ a discrete spectrum with real eigenvalues $\lambda_n$:

$$i\gamma_\mu D_\mu \Psi_n(x) = \lambda_n \Psi_n(x). \tag{6}$$

We subtract from the quark propagator $S_{Full}$ the $k$ lowest eigenmodes of the Dirac operator

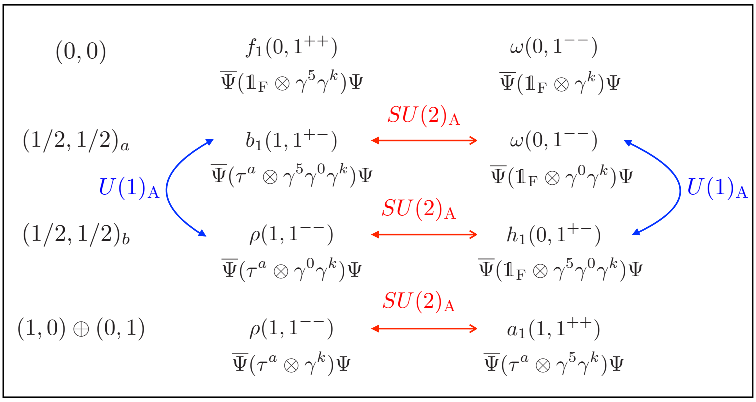


**Figure 1.** $SU(2)_L \times SU(2)_R$ classification of the $J = 1$ meson operators. Operators that are connected by the $SU(2)_L \times SU(2)_R$ and $U(1)_A$ transformations are connected by the red and blue arrows, respectively.

$$S(x, y) = S_{Full}(x, y) - \sum_{n=1}^{k} \frac{1}{\lambda_n + im} \Psi_n(x)\Psi_n^\dagger(y). \tag{7}$$

Then we apply standard procedures to extract hadron spectra using the variational approach. We perform $N_F = 2$ dynamical lattice calculations with the overlap Dirac operator with the gauge configurations generated by the JLQCD collaboration [12], for details see refs. [5, 6].

Upon truncation of the near-zero Dirac modes, that are connected to the quark condensate of the vacuum, we can expect chiral $SU(2)_L \times SU(2)_R$ symmetry in correlators. If the hadron states survive this truncation, i.e. an exponential decay of the correlators is observed, then we can expect a mass degeneracy of chiral partners. The chiral partners of the $J = 1$ mesons on Fig. 1 are linked by the red arrows.

The $U(1)_A$ transformations connect other operators that are linked by the blue arrows. If the whole chiral symmetry of QCD $SU(2)_L \times SU(2)_R \times U(1)_A$ is restored, then we should expect a degeneracy of four mesons from the $(1/2, 1/2)_a$ and $(1/2, 1/2)_b$ chiral representations, and in addition a degeneracy of the $\rho$ and $a_1$ mesons from the $(1, 0) + (0, 1)$ chiral representation. [1]

Consequently, given only the $SU(2)_L \times SU(2)_R \times U(1)_A$ chiral symmetry we should expect a degeneracy of all mesons connected by the arrows on Fig. 1.

## 3 Results

Upon truncation of some amount of the low-lying Dirac modes a very clean exponential decay of all $J = 1$ correlators is observed, which means that there are physical states. For the pion it is the other way round. After truncation of a few modes there is no exponential decay of the pion correlator, which implies that the pion does not survive truncation. The quark condensate is crucially important for the existence of the pseudo Goldstone bosons. After truncation the exponential decay of all $J = 1$ correlators is much cleaner than in the untruncated (real) world. The reason for this is intuitively clear:

[1] In the chirally symmetric world there are two independent and orthogonal $\rho$-mesons that belong to two different chiral representations. In a world with the chiral symmetry breaking two different $\rho$ operators couple to one and the same $\rho$-meson, because its wave function is a mixture of two chiral representations.

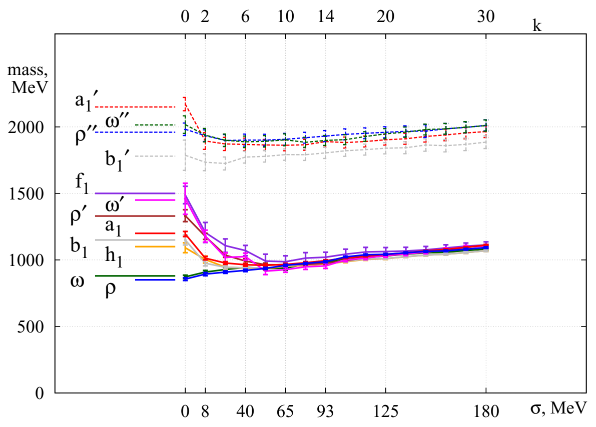


**Figure 2.** $J = 1$ meson mass evolution as a function of the truncation number $k$. $\sigma$ shows energy gap in the Dirac spectrum.

After truncation there are no pion fluctuations in the system. We can conclude that mesons (which are bound states now) survive the truncation.

The evolution of meson masses upon truncation of $k$ lowest eigenmodes is shown in Fig. 2. Chiral (and $U(1)_A$) symmetry restoration happens at k=10-20. At the same time the hadron mass is large, of the order of 1 GeV. In this regime the mass generation is obviously not connected to the chiral symmetry breaking and to the quark condensate. We conclude that while the spontaneous chiral symmetry breaking is important for hadron mass in the real world and for the shape of observed hadron spectra, the hadron mass arises mostly not from the chiral symmetry breaking.

Starting with k=10-20 we clearly see a larger degeneracy than the chiral $SU(2)_L \times SU(2)_R \times U(1)_A$ symmetry of the QCD Lagrangian. This implies that there is a symmetry in the system that is higher than $SU(2)_L \times SU(2)_R \times U(1)_A$. What does it mean?! The same results persist for the $J = 2$ mesons [7] and baryons [8].

## 4 $SU(2)_{CS}$ and $SU(4)$ symmetries

First we need to understand what symmetry group corresponds to the observed degeneracy [9].

Given the standard quantum numbers we can construct basis vectors for all irreducible representations of the chiral group in Fig. 1.

**(i) (0,0):**

$$|(0,0);\pm;J\rangle = \frac{1}{\sqrt{2}}|\bar{R}R \pm \bar{L}L\rangle_J. \tag{8}$$

**(ii)** $(1/2,1/2)_a$ **and** $(1/2,1/2)_b$**:**

$$|(1/2,1/2)_a;+;I=0;J\rangle = \frac{1}{\sqrt{2}}|\bar{R}L + \bar{L}R\rangle_J, \tag{9}$$

$$|(1/2,1/2)_a;-;I=1;J\rangle = \frac{1}{\sqrt{2}}|\bar{R}\vec{\tau}L - \bar{L}\vec{\tau}R\rangle_J, \tag{10}$$

$$|(1/2,1/2)_b;-;I=0;J\rangle = \frac{1}{\sqrt{2}}|\bar{R}L - \bar{L}R\rangle_J, \tag{11}$$

$$|(1/2, 1/2)_b; +; I = 1; J\rangle = \frac{1}{\sqrt{2}}|\bar{R}\vec{\tau}L + \bar{L}\vec{\tau}R\rangle_J. \tag{12}$$

**(iii) (0,1)⊕(1,0):**

$$|(0, 1) + (1, 0); \pm; J\rangle = \frac{1}{\sqrt{2}}|\bar{R}\vec{\tau}R \pm \bar{L}\vec{\tau}L\rangle_J, \tag{13}$$

The new symmetry transformations must connect all these basis vectors. The latter can be achieved if one allows a mixing of the left- and right-handed quarks. Consequently, the symmetry group that we are looking for must contain as a subgroup the $SU(2)_{CS}$ *chiralspin* rotations that act on the following fundamental doublets:

$$\mathrm{U} = \begin{pmatrix} u_L \\ u_R \end{pmatrix}, \qquad \mathrm{D} = \begin{pmatrix} d_L \\ d_R \end{pmatrix}. \tag{14}$$

An imaginary three-dimensional space where these rotations are performed is refered to as the *chiralspin* space. The rotations in the chiralspin space mix the right- and left-handed components of the fermion fields independently from the quark flavor. It is similar to the well familiar isospin space: Rotations in the isospin space mix particles with different electric charges.

We can construct explicit representations of the $SU(2)_{CS}$ transformations that act on Dirac bispinors [10]. Then the $SU(2)_{CS}$ chiralspin rotations are generated through

$$\mathbf{\Sigma} = \{\gamma^k, -i\gamma^5\gamma^k, \gamma^5\} , \tag{15}$$

where $\gamma^k,\ k = 1, ..., 4$ is any of the Dirac matrices in Euclidean space,

$$\gamma^i\gamma^j + \gamma^j\gamma^i = 2\delta_{ij}; \quad \gamma^5 = \gamma^1\gamma^2\gamma^3\gamma^4. \tag{16}$$

The $SU(2)$ algebra

$$[\Sigma^\alpha, \Sigma^\beta] = 2i\epsilon^{\alpha\beta\gamma}\,\Sigma^\gamma , \tag{17}$$

is satisfied with any $k = 1, ..., 4$ in (15). Different $k$ define different $SU(2)_{CS}$ irreducible representations of dim=2. The Dirac spinor transforms under a given $SU(2)_{CS}$ representation as

$$\Psi \rightarrow \Psi' = e^{i\varepsilon\cdot\mathbf{\Sigma}/2}\Psi . \tag{18}$$

Upon the $SU(2)_{CS}$ rotations of the Dirac spinors with $k = 4$ the operators from the Fig. 1 transform as triplets or singlets, see Fig. 3. The $SU(2)_{CS}$ symmetry requires that mesons within triplets should be degenerate.

If one combines the $SU(2)_{CS}$ and the flavor $SU(2)_F$ group into one larger group then one arrives at the $SU(4)$ group with the fundamental vector

$$\Psi = \begin{pmatrix} u_\mathrm{L} \\ u_\mathrm{R} \\ d_\mathrm{L} \\ d_\mathrm{R} \end{pmatrix}. \tag{19}$$

The $SU(4)$ group contains at the same time $SU(2)_L \times SU(2)_R$ and $SU(2)_{CS} \supset U(1)_A$ as subgroups and has the following set of generators:

$$\{(\tau^a \otimes 1_D), (1_F \otimes \Sigma^i), (\tau^a \otimes \Sigma^i)\} .$$

The global $SU(4)$ transformations of the Dirac spinor are defined through

$$\Psi \rightarrow \Psi' = e^{i\epsilon\cdot\mathbf{T/2}}\Psi . \tag{20}$$

The $SU(4)$ transformations connect all operators from the 15-plet, see Fig. 3.

For related studies see also [13, 14].

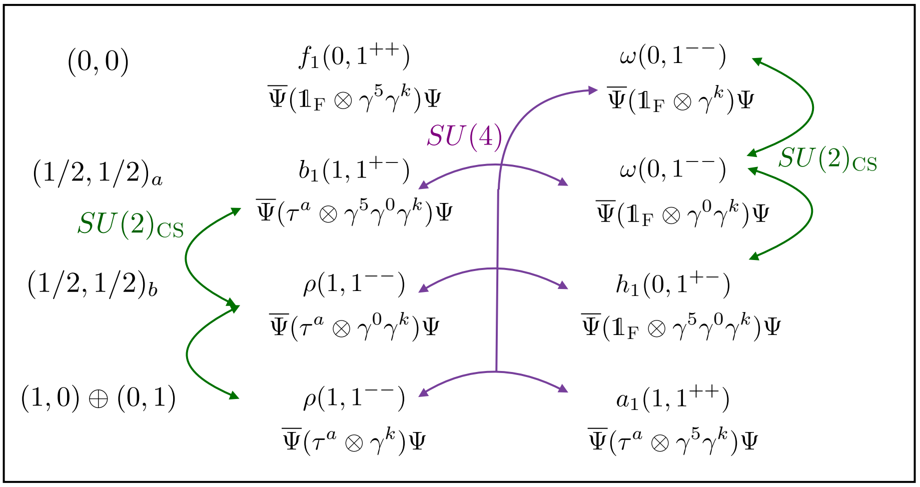


**Figure 3.** The green arrows connect operators that belong to the $SU(2)_{CS}$ triplets. The $f_1$ and $a_1$ operators are the $SU(2)_{CS}$ singlets. The purple arrows show the $SU(4)$ 15-plet. The $f_1$ operator is a singlet of $SU(4)$.

## 5 Chromo-electric versus chromo-magnetic interactions in QCD

While the QCD Lagrangian does not have the $SU(2)_{CS}$ and $SU(4)$ symmetries one clearly observes emergence of these symmetries upon elimination of the near-zero modes. Their emergence tells that there is some dynamics in QCD that is $SU(2)_{CS}-$ and $SU(4)-$symmetric. At the same time there is another dynamics that breaks both symmetries and the latter dynamics is intrinsically connected only to the low-lying modes of the Dirac operator.

Consider the interaction part of the QCD Lagrangian:

$$\overline{\Psi} i\gamma^\mu D_\mu \Psi = \overline{\Psi} i\gamma^0 D_0 \Psi + \overline{\Psi} i\gamma^i D_i \Psi \ . \tag{21}$$

The first (temporal) term is an interaction of the quark charge density $\rho(x) = \bar{\Psi}(x)\gamma^0\Psi(x) = \Psi(x)^\dagger\Psi(x)$ with the chromo-electric part of the gluonic field. The second (spatial) term contains a quark kinetic term and an interaction of the spatial current density with the chromo-magnetic field. The temporal part is invariant under any unitary transformation that can be defined in the Dirac spinor space. In particular it is invariant under the $SU(2)_{CS}$ and $SU(4)$ transformations. At the same time the kinetic quark term and the magnetic part of the interaction Lagrangian do not admit these higher symmetries (because the magnetic interaction explicitly distinguishes the left and the right and is not invariant upon the $L \leftrightarrow R$ operation) and are invariant only with respect to global $SU(N_F)_L \times SU(N_F)_R \times U(1)_A$ chiral transformations [10]. Consequently QCD has, given the $U(1)_A$ anomaly, only the $SU(N_F)_L \times SU(N_F)_R$ chiral symmetry.

This symmetry classification of different parts of the interaction Lagrangian together with emergence of the $SU(2)_{CS}$ and $SU(4)$ symmetries upon truncation of the near-zero modes implies that effect of the chromo-magnetic interaction in QCD is located exclusively in the near-zero modes, while confining chromo-electric interaction is distributed among all modes of the Dirac operator. Obviously some not yet known microscopic dynamics should responsible for this phenomenon.

To summarize this section, the chromo-electric interaction, that is responsible for confinement in QCD, is $SU(2)_{CS}$ and $SU(2N_F)$ symmetric. One can speculate that the confining dynamics is realized via a dynamical chromo-electric string [9, 10].

## 6 Zero modes and $SU(2)_{CS}$

The QCD Lagrangian is not invariant under $SU(2)_{CS}$ and $SU(4)$ transformations. This is because the Dirac operator does not commute with all $SU(2)_{CS}$ generators. The $SU(2)_{CS}$ and $SU(4)$ symmetries are obtained in lattice simulations upon subtraction of the near-zero modes of the Dirac operator. This hints that dynamically the $SU(2)_{CS}$ breaking in QCD might be related to topology and zero modes.

Consider the zero modes of the Dirac equation,

$$\gamma_\mu D_\mu \Psi_0(x) = 0. \tag{22}$$

With the standard antiperiodic boundary conditions for the quark field along the time direction, the zero modes are solutions of the Dirac equation with a gauge configuration of a nonzero global topological charge. The difference of numbers of the left-handed and right-handed zero modes is fixed, according to the Atiyah-Singer theorem, by the global topological charge $Q$ of the gauge configuration:

$$n_L - n_R = Q. \tag{23}$$

Some $SU(2)_{CS}$ transformations rotate the right-handed spinor into the left-handed one and vice versa. The $SU(2)_{CS}$ symmetry is possible only if the amount of the left-handed and right-handed zero modes is the same. Consequently, within a fixed $Q \neq 0$ sector the $SU(2)_{CS}$ should be broken.

While the exact zero modes of the Dirac operator are irrelevant for observables in the thermodynamical limit, it might happen that via fluctuations of the gluonic field this asymmetry between the left and the right related to topology becomes a property of the near-zero modes.

## 7 Observation of approximate $SU(2)_{CS}$ and $SU(4)$ symmetries at high temperatures [11]

So far we have discussed $SU(2)_{CS}$ and $SU(4)$ symmetries that emerge upon artificial truncation of the near-zero Dirac modes in $T = 0$ calculations. The near-zero modes of the Dirac operator are naturally suppressed at high temperature. It is well established on the lattice that above the critical temperature the chiral $SU(2)_L \times SU(2)_R$ symmetry gets restored and the quark condensate of the vacuum vanishes, i.e. a density of the quasi-zero modes is zero. Even more, there are clear indications from the simulations with the chirally invariant Dirac operator that above $T_c$ also $U(1)_A$ is restored and a gap opens in the Dirac spectrum [15, 16]. Then, one can expect that at high temperatures not only the $SU(2)_L \times SU(2)_R \times U(1)_A$ symmetry of the QCD Lagrangian is manifest but also the $SU(2)_{CS}$ and $SU(4)$ symmetries emerge with far reaching consequences [17].

Given this expectation spatial ($z$-direction) correlators

$$C_\Gamma(n_z) = \sum_{n_x, n_y, n_t} \langle O_\Gamma(n_x, n_y, n_z, n_t) O_\Gamma(\mathbf{0}, 0)^\dagger \rangle \tag{24}$$

of all possible $J = 0, 1$ local isovector operators $O_\Gamma(x) = \bar{q}(x)\Gamma \frac{\vec{\tau}}{2} q(x)$ have been calculated at temperatures above $T_c$ up to 380 MeV with the domain wall fermions [11, 18]. The operators as well as their $SU(2)_L \times SU(2)_R$ and $U(1)_A$ transformation properties are presented in Table 1.

Figure 4 shows the spatial correlators normalized to 1 at $n_z = 1$ for all operators in Table 1. The argument $n_z$ is proportional to the dimensionless product $zT$. We observe three distinct multiplets:

$$E_1: \quad PS \leftrightarrow S \tag{25}$$

$$E_2: \quad V_x \leftrightarrow T_t \leftrightarrow X_t \leftrightarrow A_x \tag{26}$$

$$E_3: \quad V_t \leftrightarrow T_x \leftrightarrow X_x \leftrightarrow A_t. \tag{27}$$

$E_1$ is the Pseudoscalar-Scalar multiplet connected by the $U(1)_A$ symmetry. The $E_2$ and $E_3$ multiplets contain however some operators that are connected by neither $SU(2)_L \times SU(2)_R$ nor $U(1)_A$ transformations. Below we demonstrate that the symmetries responsible for emergence of the $E_2$ and $E_3$ multiplets are $SU(2)_{CS}$ and $SU(4)$.

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

**Table 1.** Bilinear operators and their transformation properties. This classification assumes propagation in $z$-direction. The index $k$ denotes the components $1, 2, 4$, *i.e.* $x, y, t$.

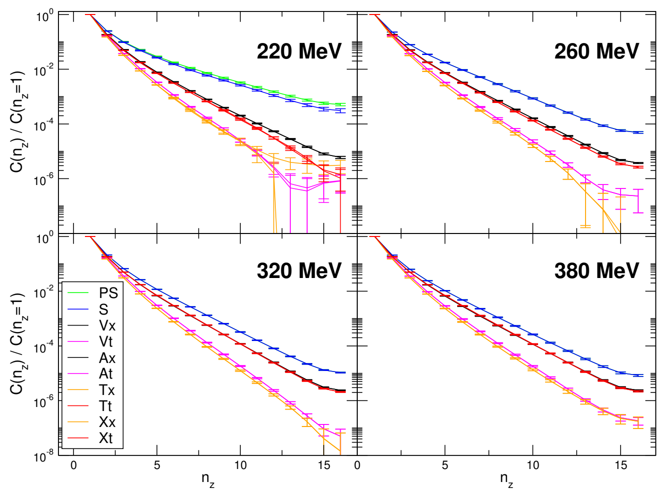


**Figure 4.** Normalized spatial correlators.

Consider the following dim=2 representations of $SU(2)_{CS}$ that are defined in the Dirac bispinor space:

$$R_1 : \{\gamma_1, -i\gamma_5\gamma_1, \gamma_5\}, \tag{28}$$

$$R_2 : \{\gamma_2, -i\gamma_5\gamma_2, \gamma_5\}. \tag{29}$$

Those differ from the representation $\{\gamma_4, -i\gamma_5\gamma_4, \gamma_5\}$ relevant for $t$-direction correlators by rotations. The $R_1$ and $R_2$ $SU(2)_{CS}$ transformations of the quark fields combine the following operators from the $E_2$ multiplet into triplets:

$$R_1 : \qquad V_y \leftrightarrow T_t \leftrightarrow X_t, \tag{30}$$

$$R_2 : \qquad V_x \leftrightarrow T_t \leftrightarrow X_t, \tag{31}$$

as well as from the $E_3$ multiplet:

$$R_1 : \qquad V_t \leftrightarrow T_y \leftrightarrow X_y, \tag{32}$$

$$R_2 : \qquad V_t \leftrightarrow T_x \leftrightarrow X_x. \tag{33}$$

Given the $S_2$ symmetry ($x \leftrightarrow y$ permutations) one obtains the following $S_2 \times SU(2)_{CS}$ multiplets

$$(V_x, V_y, T_t, X_t); \quad (V_t, T_x, T_y, X_x, X_y). \tag{34}$$

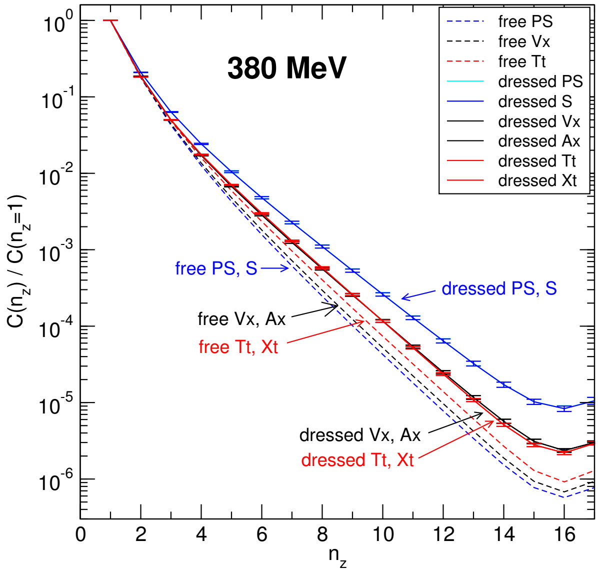


**Figure 5.** $E_1$ and $E_2$ multiplets (25-26) for interacting (*dressed*) and non-interacting (*free*) calculations at $T=$ 380 MeV.

The degeneracy between **V** and **A** requires the $SU(2)_L \times SU(2)_R$ symmetry. Extending $SU(2)_{CS}$ to $SU(4)$ one arrives at the following multiplets of the isovector operators:

$$(V_x, V_y, T_t, X_t, A_x, A_y);\ (V_t, T_x, T_y, X_x, X_y, A_t). \tag{35}$$

$S_2 \times SU(4)$ multiplets include in addition the isoscalar partners of operators.

While the $U(1)_A$ and $SU(2)_L \times SU(2)_R$ symmetries are "exactly" restored at temperatures above 220 MeV, the $SU(2)_{CS}$ and $SU(4)$ multiplets are only approximate. A degree of degeneracy at the highest available temperature 380 MeV can be deduced from the Fig. 5 where we show correlators of the $E_1$ and $E_2$ multiplets in detail. The remaining $SU(2)_{CS}$ and $SU(4)$ breaking is at the level of 5%. We also show there correlators calculated with the noninteracting quarks (abbreviated as "free").

The slopes of the PS and S correlators are substantially smaller than for the free quark-antiquark pair. This can happen only if the system represents a bound meson-like state [19]. In the free quark case the slope is determined by twice of the lowest Matsubara frequency because of the antiperiodic boundary conditions for quarks in time direction. If the quark-antiquark sytem is bound, and of the bosonic nature the periodic boundary conditions (for bosons) do allow the slope to be smaller. For the $J = 1$ correlators the difference of slopes of dressed and free correlators is smaller than for the $J = 0$ correlators, but is still visible.

In order to see a tendency of the $U(1)_A$ restoration and of $SU(2)_{CS}$ and $SU(4)$ emergence we show on Fig. 6 a ratio of the correlators from the multiplet $E_2$ that are connected by the corresponding transformations at different temperatures. We also show a ratio calculated with the free noninteracting quarks. We observe that upon increasing the temperature the correlators approach the $SU(2)_{CS}$-symmetric limit and not the free quark limit as is prescribed by the asymptotic freedom at high T (i.e. by the renormalization group flow equation calculated with perturbative $\beta$-function). This is a striking observation. We plan to verify this behaviour at higher temperatures.

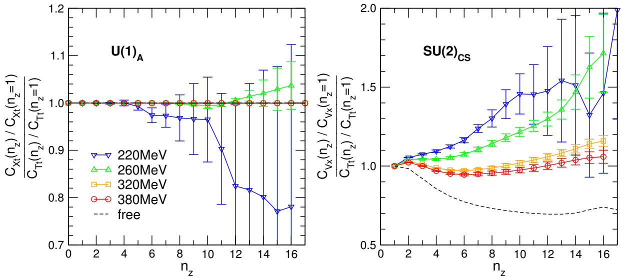


**Figure 6.** Ratios of normalized correlators , that are related by $U(1)_A$ and $SU(2)_{CS}$ symmetries.

## 8 Implications

We conclude that our lattice results are consistent with emergence of the $SU(2)_{CS}$ and $SU(4)$ symmetries by increasing temperature. The correlation functions do not seem to approach the free quark limit.

These results have a direct implication on the nature of the degrees of freedom in the high T phase. Emergence of the $SU(2)_{CS}$ and $SU(4)$ symmetries rules out the possibility that elementary objects are deconfined asymptotically free quarks. Instead such elementary objects should be chiral quarks bound by the pure chromo-electric field that is not accompanied by the magnetic effects (!), a kind of a string. Such a schematic construction automatically incorporates the $SU(2)_{CS}$ and $SU(4)$ symmetries, as it follows from Sec. 4. Absence of the magnetic effects implies that e.g. the spin-orbit force is absent and the string with the left or with the right quark at the end has the same energy [20].

How should such a state of matter be called? It is not a plasma, because according to the standard definition plasma is a system of free color charges where the chromo-electric field is Debye screened. According to our results it follows that there are no free quarks and in addition it is the chromo-magnetic, but not the chromo-electric field, is screend. So conditionally one could call this matter a stringy matter.

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