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# Approximate degeneracy of $J = 1$ spatial correlators in high temperature QCD

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We study spatial isovector meson correlators in $N_f = 2$ QCD with dynamical domain-wall fermions on $32^3 \times 8$ lattices at temperatures $T = 220$–380 MeV. We measure the correlators of spin-one ($J = 1$) operators including vector, axial-vector, tensor and axial-tensor. Restoration of chiral $U(1)_A$ and $SU(2)_L \times SU(2)_R$ symmetries of QCD implies degeneracies in vector–axial-vector ($SU(2)_L \times SU(2)_R$) and tensor–axial-tensor ($U(1)_A$) pairs, which are indeed observed at temperatures above $T_c$. Moreover, we observe an approximate degeneracy of all $J = 1$ correlators with increasing temperature. This approximate degeneracy suggests emergent $SU(2)_{CS}$ and $SU(4)$ symmetries at high temperatures, that mix left- and right-handed quarks.

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## I. INTRODUCTION

One of the key questions of QCD, that is of crucial importance both for astrophysics and cosmology, is the nature of the strongly interacting matter at high temperatures. This question attracts enormous experimental and theoretical efforts. It is established in *ab initio* QCD calculations on the lattice that there is a transition to the chirally symmetric regime where the quark condensate, an order parameter of $SU(2)_L \times SU(2)_R$ chiral symmetry, vanishes. At the same time there is strong evidence from calculations with manifestly chirally-invariant lattice fermions that above the critical temperature also the $U(1)_A$ symmetry gets restored and a gap opens in the spectrum of the Dirac operator [1–3].

To get some new insight on the symmetry structure of high temperature QCD we calculate spatial correlators of all possible spin $J = 0$ and $J = 1$ local isovector operators using the chirally invariant domain-wall Dirac operator in two-flavour QCD at different temperatures up to 380 MeV. All correlators that are connected by $SU(2)_L \times SU(2)_R$ or $U(1)_A$ transformations are the same within errors at temperatures above the critical one, which is in agreement with restoration of both symmetries at high temperature. Surprisingly, we also observe an approximate degeneracy of some correlators that are connected neither by $SU(2)_L \times SU(2)_R$ nor by $U(1)_A$ transformations.

The observed approximate degeneracies at high temperature are in agreement with emergent $SU(2)_{CS}$ (chiral-spin) and $SU(4)$ symmetries [4,5], which contain transformations that mix the left- and right-handed components of quarks. These symmetries have been observed before in $T = 0$ dynamical calculations upon artificial truncation of the near-zero modes of the overlap Dirac operator [6–9]. The near-zero modes of the Dirac operator on the lattice are strongly suppressed at high temperatures [1,2], which motivates our present exploration of the correlators and symmetries at high temperatures, this time without truncating the Dirac eigenmodes.

## II. SIMULATION

### A. Lattices

The gauge configurations used in the numerical simulation of QCD are generated using the Symanzik gauge action and two degenerate flavors of Möbius domain wall fermions [10,11]. The gauge links are stout smeared three times before the computation of the Dirac operator. The length in the fifth direction $L_s$ is chosen to achieve precise chiral symmetry. The boundary conditions for quarks are set antiperiodic in $t$-direction, and periodic in spatial directions. The ensembles and parameters including the lattice spacing $a$ are listed in Table I. The degenerate up and down quark masses $m_{ud}$ are set to 2–15 MeV, which is essentially negligible at the temperatures we studied, i.e. $T \simeq 220$–380 MeV. More details on the chiral properties for this set of parameters are given in [2,12]. We study

 

TABLE I. Gauge ensembles for $32^3 \times 8$ lattices used in this work. $L_s$ denotes the length of the fifth dimension in the domain wall fermion formulation. The critical temperature for this set of parameters is $T_c = 175 \pm 5$ MeV.

| $\beta$ | $m_{ud}a$ | $a$ [fm] | # configs | $L_s$ | $T$ [MeV] | $T/T_c$ |
|---|---|---|---|---|---|---|
| 4.10 | 0.001 | 0.113 | 800 | 24 | ∼220 | ∼1.2 |
| 4.18 | 0.001 | 0.096 | 230 | 12 | ∼260 | ∼1.5 |
| 4.30 | 0.001 | 0.075 | 260 | 12 | ∼320 | ∼1.8 |
| 4.37 | 0.005 | 0.065 | 120 | 12 | ∼380 | ∼2.2 |

spatial ($z$-direction) correlators as was first suggested in Ref. [13] (see also [14]).

### B. Operators

The observables of interest are correlators of non-singlet local operators $\mathcal{O}_\Gamma(x) = \bar{q}(x)\Gamma\frac{\vec{\tau}}{2}q(x)$, where $\Gamma$ might be any combination of $\gamma$-matrices, i.e., the Clifford algebra, containing 16 elements; $\tau_a$ are the isospin Pauli matrices.

A zero-momentum projection is done by summation over all lattice points in slices orthogonal to the measurement direction. When measuring in $z$-direction this means

$$C_\Gamma(n_z) = \sum_{n_x,n_y,n_t} \langle \mathcal{O}_\Gamma(n_x,n_y,n_z,n_t)\mathcal{O}_\Gamma(\mathbf{0},0)^\dagger\rangle. \quad (1)$$

For the vector and axial-vector operators $\Gamma$ has the following components:

$$\mathbf{V} = \begin{pmatrix} \gamma_1 = V_x \\ \gamma_2 = V_y \\ \gamma_4 = V_t \end{pmatrix}, \qquad \mathbf{A} = \begin{pmatrix} \gamma_1\gamma_5 = A_x \\ \gamma_2\gamma_5 = A_y \\ \gamma_4\gamma_5 = A_t \end{pmatrix}. \quad (2)$$

Conservation of the vector current requires that $V_z$ does not propagate in $z$-direction. As the axial vector current $j_5^\mu$ is not conserved at zero temperature, the relevant component $\gamma_3\gamma_5$ of the Axial-vector does propagate at zero temperature and eventually couples to the Pseudoscalar. Above the critical temperature—after $U(1)_A$ and $SU(2)_L \times SU(2)_R$ restoration—$A_z$ behaves as its parity partner $V_z$ and does not propagate in $z$-direction. For propagation in $z$-direction the tensor elements $\sigma_{\mu\nu}$ of the Clifford algebra are organized in the following way in components of tensor- and axial-tensor vectors:

$$\mathbf{T} = \begin{pmatrix} \gamma_1\gamma_3 = T_x \\ \gamma_2\gamma_3 = T_y \\ \gamma_4\gamma_3 = T_t \end{pmatrix}, \qquad \mathbf{X} = \begin{pmatrix} \gamma_1\gamma_3\gamma_5 = X_x \\ \gamma_2\gamma_3\gamma_5 = X_y \\ \gamma_4\gamma_3\gamma_5 = X_t \end{pmatrix}. \quad (3)$$

Table II summarizes our operators and gives the $U(1)_A$ and $SU(2)_L \times SU(2)_R$ relations of these operators. Given restoration of the $U(1)_A$ and $SU(2)_L \times SU(2)_R$ symmetries at high $T$ we expect degeneracies of correlators calculated with the corresponding operators.

TABLE II. Bilinear operators considered in this work and their transformation properties (last column). This classification assumes propagation in $z$-direction. The open vector index $k$ denotes the components 1,2,4, i.e., $x$, $y$, $t$.

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

For measurements at zero temperature the three components of the vectors give the same expectation value due to the $SO(3)$ symmetry in continuum. In our finite temperature setup this rotational symmetry is broken and the residual $SO(2)$ symmetry in the $(x,y)$-plane connects $V_x \leftrightarrow V_y$, $A_x \leftrightarrow A_y$ etc. operators.

On the lattice at finite temperature the symmetry is $D_{4h}$ where the vector components belong to one two-dimensional ($V_x$,$V_y$) and one one-dimensional ($V_t$) irreducible representations, and similar for $\mathbf{A}$, $\mathbf{T}$, $\mathbf{X}$. This is discussed in more detail in Appendix (see also [15]). Operators from different irreducible representations are not connected by the $D_{4h}$ transformations and consequently the $D_{4h}$ symmetry of the lattice does not predict the $E_1$, $E_2$, $E_3$ multiplet structures discussed in Sec. III. The $x$ and $y$ components of $\mathbf{V}$ have degenerate energy levels, and correspondingly those for the other Dirac structures. We therefore show only $x$- and $t$-components in the plots.

## III. RESULTS

Figure 1 shows the spatial correlation functions normalized to 1 at $n_z = 1$ for the operators given in Table II. As argument we show $n_z$ which is proportional to the dimensionless product $zT$ for fixed $N_t$, the temporal extent of the lattice.

As we describe in more detail later, we find that all correlators connected by $SU(2)_L \times SU(2)_R$ and $U(1)_A$ transformations coincide within small deviations at $T > 220$ MeV, which means that at these temperatures both chiral symmetries get restored. More interestingly, there are additional degeneracies of correlators. In total we observe three different multiplets:

$$E_1\colon\ PS \leftrightarrow S \quad (4)$$

$$E_2\colon\ V_x \leftrightarrow T_t \leftrightarrow X_t \leftrightarrow A_x \quad (5)$$

$$E_3\colon\ V_t \leftrightarrow T_x \leftrightarrow X_x \leftrightarrow A_t. \quad (6)$$

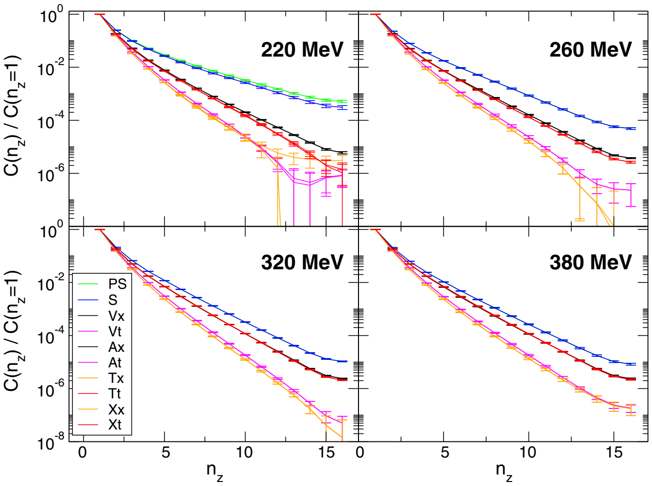


FIG. 1. Normalized spatial correlators. The temperatures correspond to the ensembles listed in Table I.

$E_1$ describes the pseudoscalar-scalar multiplet connected by the $U(1)_A$ symmetry, that is realized in the absence of chiral zero-modes [16]. Note that we only consider the isospin triplet channels so $S$ corresponds to the $a_0$- rather than the $\sigma$-particle. The $E_2$ and $E_3$ multiplets on the other hand contain some operators that are not connected by either $SU(2)_L \times SU(2)_R$ or $U(1)_A$ transformations.

Figure 2 shows the correlators of the $E_1$ and $E_2$ multiplets in detail at the highest available temperature $T = 380$ MeV. Here we also show correlators calculated with noninteracting quarks. The noninteracting (*free*) data have been generated on the same lattice sizes using a unit gauge configuration [17]. Due to the small quark mass the difference between chiral partners is negligible for the free case, therefore they are omitted.

We observe a precise degeneracy between $S$ and $PS$ correlators, which is consistent with the effective $U(1)_A$ restoration on these lattice ensembles [1]. The logarithmic slope of the interacting (*dressed*) $S$ and $PS$ correlators is substantially smaller than that for free quarks. In the latter case the slope is given by $2\pi/N_t$. A system of two free quarks cannot have "energy" smaller than twice the lowest Matsubara frequency [13]. For the $E_2$ multiplet we observe asymptotic slopes that are quite close to $2\pi/N_t$ in agreement with previous studies [18].

Figure 3 shows normalized ratios of $X_t$ and $T_t$ correlators on the left, as well as of $V_x$ and $T_t$ correlators on the right side. The $U(1)_A$ symmetry is restored, as is evident from the left side of this Figure, where a ratio of correlators in the $X_t$ and $T_t$ channels is plotted. We also find a similar degeneracy between $V_x$ and $A_x$ due to $SU(2)_A$ (See also, e.g. [19,20]).

Figures 2 and 3 suggest a possible higher symmetry ($SU(2)_{CS}$ symmetry, see next section) that connects $V_x$ and $T_t$ channels. The right panel of Fig. 3 shows the corresponding ratio, which demonstrates an approximate degeneracy at the level of 5% above $T \simeq 320$ MeV. We notice that this degeneracy is not expected in the free quark limit which is plotted by a dashed curve. This unexpected symmetry requires that the cross-correlator calculated with the $V_x$ and $T_t$ operators (both create the $1^{--}$ states) should vanish. We have carefully checked that it indeed vanishes to high accuracy.

Figure 4 shows the $E_3$ multiplet. Here again we observe a precise degeneracy in all $SU(2)_L \times SU(2)_R$ and $U(1)_A$ connected correlators, as well as the approximate degeneracy in all four correlators. We also see qualitatively different data between free and dressed correlators at $n_z \geq 11$, as also seen in [21]. This implies that we do not observe free non-interacting quarks but instead systems with some interquark correlation, which is in accordance with the known results for energy density and pressure at high temperatures [22].

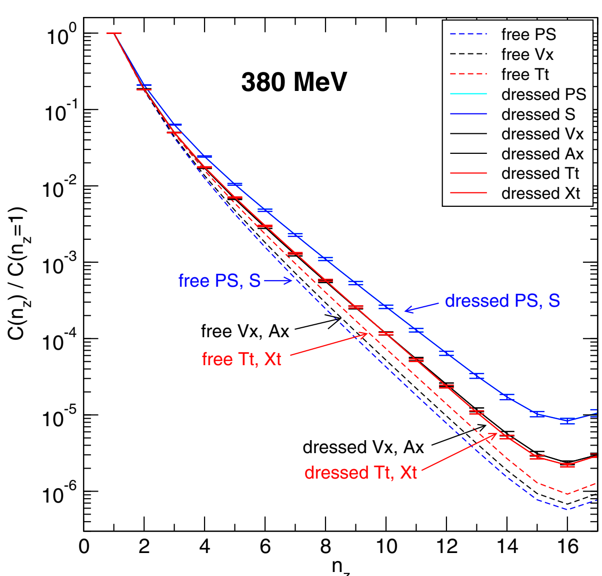


FIG. 2. $E_1$ and $E_2$ multiplets (4)–(5) for interacting (*dressed*) and noninteracting (*free*) calculations at $T = 380$ MeV.

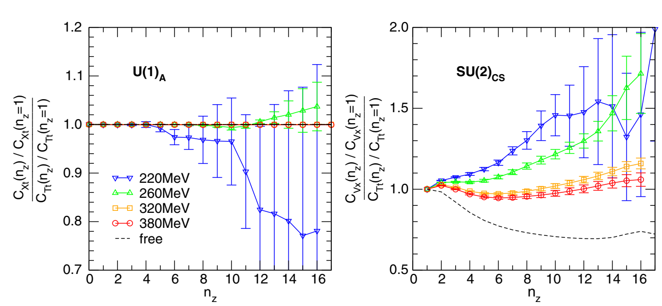


FIG. 3. Ratios of normalized correlators from Fig. 1, that are related by $U(1)_A$ and $SU(2)_{CS}$ symmetry.

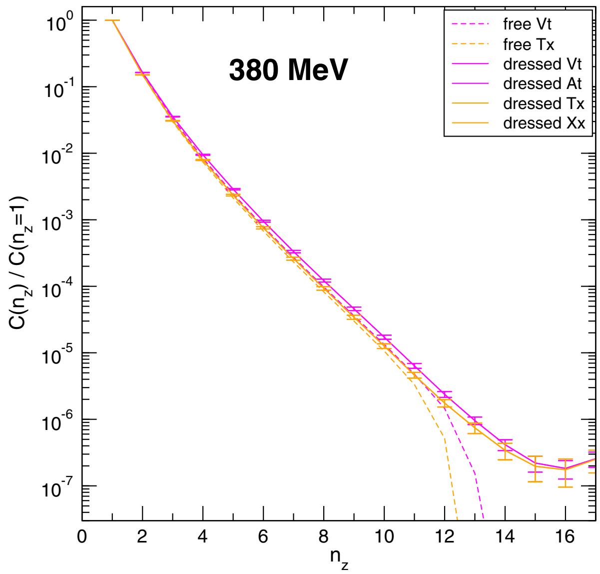


FIG. 4. $E_3$ multiplet (6) for interacting (*dressed*) and noninteracting (*free*) calculations.

## IV. $SU(2)_{CS}$ AND $SU(4)$ SYMMETRIES

In this section we introduce the $SU(2)_{CS}$ and $SU(4)$ transformations, which connect operators from multiplet $E_2$ (5) as well as from multiplet $E_3$ (6) and contain chiral transformations as a subgroup. The basic ideas of $SU(2)_{CS}$ and $SU(4)$ symmetries at zero temperature are given in [5]. Here we adapt the group structure to our setup.

We use the $\gamma$-matrices given by

$$\gamma_i\gamma_j + \gamma_j\gamma_i = 2\delta^{ij}; \qquad \gamma_5 = \gamma_1\gamma_2\gamma_3\gamma_4. \tag{7}$$

The transformations and generators of the $SU(2)_{CS}$ chiral-spin group, defined in the Dirac spinor space and diagonal in flavor space, are given by

$$q \to \exp\left(\frac{i}{2}\vec{\epsilon}\cdot\vec{\Sigma}\right)q, \qquad \vec{\Sigma} = \{\gamma_k, -i\gamma_5\gamma_k, \gamma_5\}. \tag{8}$$

$SU(2)_{CS}$ contains $U(1)_A$ as a subgroup. The $\mathfrak{su}(2)$ algebra $[\Sigma_\alpha, \Sigma_\beta] = 2i\epsilon^{\alpha\beta\gamma}\Sigma_\gamma$ is satisfied with any $k = 1, 2, 3, 4$. The $SU(2)_{CS}$ transformations mix the left- and right-handed components of the quark field. It is not a symmetry of the free massless quark Lagrangian. For $z$-direction correlators the following representations of $SU(2)_{CS}$ are relevant:

$$R_1\colon \{\gamma_1, -i\gamma_5\gamma_1, \gamma_5\} = \{\sigma^{23}i\gamma_5\gamma_4, \sigma^{23}\gamma_4, \gamma_5\}, \tag{9}$$

$$R_2\colon \{\gamma_2, -i\gamma_5\gamma_2, \gamma_5\} = \{\sigma^{31}i\gamma_5\gamma_4, \sigma^{31}\gamma_4, \gamma_5\}. \tag{10}$$

Those differ from the representation $\{\gamma_4, -i\gamma_5\gamma_4, \gamma_5\}$ relevant for $t$-direction correlators [5] by the rotations $\sigma^{23} = \frac{i}{2}[\gamma_2, \gamma_3]$ and $\sigma^{31} = \frac{i}{2}[\gamma_3, \gamma_1]$.

These $R_1$ and $R_2$ $SU(2)_{CS}$ transformations connect the following operators from the $E_2$ multiplet:

$$R_1\colon V_y \leftrightarrow T_t \leftrightarrow X_t, \tag{11}$$

$$R_2\colon V_x \leftrightarrow T_t \leftrightarrow X_t, \tag{12}$$

as well as the operators from the $E_3$ multiplet:

$$R_1\colon V_t \leftrightarrow T_y \leftrightarrow X_y, \tag{13}$$

$$R_2\colon V_t \leftrightarrow T_x \leftrightarrow X_x. \tag{14}$$

Our lattice symmetry group includes both the permutation operator $\hat{P}_{xy}$ and $\mathbb{1}$ transformations, which form a group $S_2$. $\hat{P}_{xy}$ permutes $\gamma_1$ and $\gamma_2$, and transforms $\gamma_5$ to $-\gamma_5$. Then $P_{xy}R_1$ is isomorphic to $R_2$. This means that $S_2 \times SU(2)_{CS}$ contains multiplets

$$(V_x, V_y, T_t, X_t); \qquad (V_t, T_x, T_y, X_x, X_y). \tag{15}$$

The degeneracy between $\mathbf{V}$ and $\mathbf{A}$ means $SU(2)_L \times SU(2)_R$ symmetry. A minimal group that includes $SU(2)_L \times SU(2)_R$ and $SU(2)_{CS}$ is $SU(4)$. The 15 generators of $SU(4)$ are the following matrices:

$$\{(\tau_a \otimes \mathbb{1}_D), (\mathbb{1}_F \otimes \Sigma_i), (\tau_a \otimes \Sigma_i)\} \tag{16}$$

with flavor index $a = 1, 2, 3$ and $SU(2)_{CS}$ index $i = 1, 2, 3$. Predictions of $S_2 \times SU(4)$ symmetry for isovector operators are the following multiplets:

$$(V_x, V_y, T_t, X_t, A_x, A_y); \qquad (V_t, T_x, T_y, X_x, X_y, A_t). \tag{17}$$

$S_2 \times SU(4)$ multiplets include in addition the isoscalar partners of $V_x$, $V_y$, $T_t$ and $X_t$ operators for the first multiplet in (17) as well as of $V_t$, $T_x$, $T_y$, $X_x$, $X_y$ for the second multiplet in (17).

## V. CONCLUSIONS AND DISCUSSION

Our lattice results are consistent with emergence of approximate $SU(2)_{CS}$ and $SU(4)$ symmetries in spin $J = 1$ correlators by increasing temperature. The considered correlation functions do not seem to approach the free quark limit.

How could these approximate $SU(2)_{CS}$ and $SU(4)$ symmetries arise at high temperatures? They are not symmetries of the QCD Lagrangian. They are both symmetries of the confining chromo-electric interaction in Minkowski space since any unitary transformation leaves the temporal part of the fermion Lagrangian $\bar{q}\gamma^\mu D_\mu q$ invariant. The chromo-magnetic interaction described by the rest of the Lagrangian breaks both symmetries [5]. Consequently the emergence of these symmetries suggests that the chromo-magnetic interaction is suppressed at high temperature while the chromo-electric interaction is still active. This could have implications on the nature of the effective degrees of freedom in the high temperature phase of QCD since these symmetries are incompatible with asymptotically free deconfined quarks.

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## APPENDIX: REPRESENTATIONS OF $D_{4H}$

The symmetry of the $(x, y, t)$–volume (the fixed $n_z$ subvolume, where the discussed operators are defined) is $D_{4h}$ [23]. Consider the transformations of the Euclidean interpolators $O(n) = \bar{q}(n)\Gamma q(n)$ with $n = (n_x, n_y, n_z, n_t)$:

$$\text{rot}:\ \bar{q}\Gamma q \to \bar{q}\exp\left(\frac{i}{2}\omega_{\mu\nu}\sigma^{\mu\nu}\right)\Gamma\exp\left(-\frac{i}{2}\omega_{\mu\nu}\sigma^{\mu\nu}\right)q$$

$$\hat{P}^z:\ \bar{q}\Gamma q \to \bar{q}\gamma_3\Gamma\gamma_3 q \qquad \text{(A1)}$$

under discrete rotations and $\hat{P}^z$ that performs inversion $n \to n_{P^z} = (-n_x, -n_y, n_z, -n_t)$.

The relevant symmetry group $D_{4h}$ has ten classes of group elements and ten irreducible representations identified by characters in Table III: $C_4$ and $C_2$ are rotations around $t$ for $\pi/2$ and $\pi$, respectively, $C_2'$ is a rotation for $\pi$ around $x$, while $C_2''$ is a rotation for $\pi$ around $x + y$. Further five classes are obtained by multiplication of the elements with $\hat{P}^z$ and the characters with $(-1)^{P^z}$.

According to these transformations, the interpolators for **V**, **A**, **T**, **X** operators of Table II transform under the irreducible representations given in Table IV. Note that group elements of $D_{4h}$ transform interpolators only within one box of that table and indeed the observed energy levels for the $x$ and $y$ components of an operator agree. However, no group element of $D_{4h}$ transforms components of **V** to **T** (or **A** to **X**).

TABLE III. The first five columns of the character table of $D_{4h}$ [23]; the superscript in the irreps denotes $P_z$. The first column gives the dimension of the irrep. Another five classes are obtained by multiplication of elements with $P^z$, while the characters get a factor of $(-1)^{P^z}$.

| | Id | $C_4$ | $C_2$ | $C_2'$ | $C_2''$ |
|---|---|---|---|---|---|
| $A_1^\pm$ | 1 | 1 | 1 | 1 | 1 |
| $A_2^\pm$ | 1 | 1 | 1 | −1 | −1 |
| $B_1^\pm$ | 1 | −1 | 1 | 1 | −1 |
| $B_2^\pm$ | 1 | −1 | 1 | −1 | 1 |
| $E^\pm$ | 2 | 0 | −2 | 0 | 0 |

TABLE IV. The interpolators $O(n)$, which are denoted by **V**, **A**, **T**, **X** in Table II, transform according to the irreps in this Table.

| | | |
|---|---|---|
| $V_x, V_y: E^-$ | $V_z: A_1^+$ | $V_t: A_2^-$ |
| $A_x, A_y: E^+$ | $A_z: A_1^-$ | $A_t: A_2^+$ |
| $T_x, T_y: E^-$ | $T_z: A_1^+$ | $T_t: A_2^-$ |
| $X_x, X_y: E^+$ | $X_z: A_1^-$ | $X_t: A_2^+$ |

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