

Operator Product Expansion Coefficients of the 3D Ising Criticality via Quantum Fuzzy Spheres

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 (Received 26 March 2023; accepted 21 June 2023; published 21 July 2023)

Conformal field theory (CFT) plays a crucial role in the study of various critical phenomena. While much attention has been paid to the critical exponents of different universalities, which correspond to the conformal dimensions of CFT primary fields, other important and intricate data such as operator product expansion (OPE) coefficients governing the fusion of two primary fields, have remained largely unexplored, especially in dimensions higher than 2D (or equivalently, $1 + 1D$). Motivated by the recently proposed fuzzy sphere regularization, we investigate the operator content of 3D Ising criticality from a microscopic perspective. We first outline the procedure for extracting OPE coefficients on the fuzzy sphere and then compute 13 OPE coefficients of low-lying CFT primary fields. Our results are highly accurate and in agreement with the numerical conformal bootstrap data of 3D Ising CFT. Moreover, we were able to obtain 4 OPE coefficients, including $f_{T_{\mu\nu}T_{\rho\sigma}}$, which were previously unknown, thus demonstrating the superior capabilities of our scheme. Expanding the horizon of the fuzzy sphere regularization from the state perspective to the operator perspective opens up new avenues for exploring a wealth of new physics.

DOI: [10.1103/PhysRevLett.131.031601](https://doi.org/10.1103/PhysRevLett.131.031601)

Critical phenomena attract broad interests in various fields of physics, ranging from condensed matter to high energy physics. An intriguing reason is, enlarged symmetry often emerges at the low energy and large length scale in the vicinity of a critical point [1,2]. One well-developed scenario is that at a continuous phase transition where scale invariance is enhanced to the conformal invariance [3–5], the low-energy physics is well captured by the conformal field theory (CFT) [6,7], a field theory that is invariant under the conformal symmetry. The nature of a CFT, hence its corresponding critical point, is largely determined by a set of conformal data, consisting of a list of conformal dimensions and operator product expansion (OPE) coefficients [3,8,9]. The complete conformal data is able to reproduce many universal properties of the phase transition and determine the stability of fluctuations close to a fixed point [2]. For example, the scaling dimensions are critical exponents and OPE coefficients give rise to a universal (charge or thermal) conductivity [10] which are experimentally measurable at a critical point. It remains an outstanding challenge for various fields of physics to accurately determine the conformal data of interacting CFTs beyond 2D (or, equivalently, $1 + 1D$).

Generally speaking, there are two strategies to study the conformal data of an interacting critical system. The first strategy, adopted by the conformal bootstrap program [11],

utilizes general axioms of CFTs to constrain the conformal data. Although the bootstrap approach can produce a great deal of conformal data accurately for certain critical theories such as the 3D Ising [12–14], it very often loses its power for many interesting universalities since its starting point is too general to make contact with a specific universality. In contrast, an alternative strategy, adopted in condensed matter research, is to study the microscopic model that directly realizes the universality of interest [15]. The weakness of this down-to-earth strategy is that it can access only a very limited number of conformal data. For example, in recent progress on computing OPE coefficients of 3D classical transitions, either directly from three-point correlators [16,17], or indirectly from off-critical two-point correlators [18], due to various complications a very limited number of (only two or three) OPE coefficients were accessible in these computations [19].

In the 1980s, Cardy outlined a scheme to compute the conformal data in microscopic models by making use of the state-operator correspondence, i.e., radial quantization, on the cylinder geometry $S^{d-1} \times \mathbb{R}$ [20,21]. The OPE coefficients can be computed directly via the one-point expectation value $\langle \phi_\alpha | \phi_\beta | \phi_\gamma \rangle$. This scheme not only greatly reduces computational complexities, but also, in principle, enables the access of all the conformal data that are infinitely many. For the 2D (i.e., $1 + 1D$) critical theory,

this scheme has been applied extensively [22–27] since one just needs to study a lattice quantum Hamiltonian defined on a circle (S^1 , i.e., periodic boundary chain). Moving to critical theories in higher dimensions, a fundamental obstacle arises because spatial spherical geometry $S^{m \geq 2}$ has a curvature such that any regular lattice cannot fit in. This fundamental obstacle has recently been removed by the two of us using an innovative idea, the *fuzzy (non-commutative) sphere regularization* [28]. The correspondence between scaling dimensions and energy spectra of the 3D Ising CFT has been demonstrated convincingly. In this Letter we are extending the horizon of the fuzzy sphere regularization to the precise determination of OPE coefficients. In particular, we have accurately computed 13 different OPE coefficients of the 3D Ising CFT, including four OPE coefficients that were unavailable in the existing literature. Our successful access of CFT operator contents in the 3D Ising criticality paves the way for studying higher dimensional CFTs through the fuzzy sphere regularization.

3D Ising CFT on the fuzzy sphere.—We consider spinful electrons moving on the sphere S^2 in the presence of $4\pi \cdot s$ monopole placed in the sphere origin. Because of the monopole, the single-electron states form quantized Landau levels, with $2s + 1$ degenerate orbitals in the lowest Landau level (LLL) [29]. In the case that the LLL is partially filled and the gap between the LLL and higher LLs is much larger than other interaction scales in the system, we can effectively project the system into the LLL. After the LLL projection, coordinates of electrons are not commuting any more, $[\tilde{x}_\mu, \tilde{x}_\nu] = i(R/s)\epsilon_{\mu\nu\rho}\tilde{x}_\rho$. Thus, we end up with a system defined on a fuzzy (noncommutative) two-sphere [30]. Interestingly, these Landau orbitals form a spin- s irreducible representation of $SO(3)$, and the sphere radius $R \sim \sqrt{s}$. The thermodynamic limit is reached by $s \rightarrow \infty$.

In practice, the Hamiltonian on the fuzzy sphere can be written in the second quantized form using the basis of Landau orbitals. A 2 + 1D Ising transition can be realized by the following Hamiltonian [28]

$$H = \sum_{m_1, 2, 3, 4 = -s}^s V_{m_1, m_2, m_3, m_4} [(\mathbf{c}_{m_1}^\dagger \mathbf{c}_{m_4})(\mathbf{c}_{m_2}^\dagger \mathbf{c}_{m_3}) - (\mathbf{c}_{m_1}^\dagger \sigma^z \mathbf{c}_{m_4})(\mathbf{c}_{m_2}^\dagger \sigma^z \mathbf{c}_{m_3})] - h \sum_{m=-s}^s \mathbf{c}_m^\dagger \sigma^x \mathbf{c}_m, \quad (1)$$

where $\mathbf{c}_m^\dagger = (c_{m\uparrow}^\dagger, c_{m\downarrow}^\dagger)$ is the fermion creation operator on the m_{th} Landau orbital and $\sigma^{x,y,z}$ is the Pauli matrix. The element V_{m_1, m_2, m_3, m_4} is connected to the Haldane pseudopotential V_l [29] by $V_{m_1, m_2, m_3, m_4} = \sum_l V_l (4s - 2l + 1) \binom{s \ s \ 2s-l}{m_1 \ m_2 \ -m_1 - m_2} \binom{s \ s \ 2s-l}{m_4 \ m_3 \ -m_3 - m_4}$, where $\binom{j_1 \ j_2 \ j_3}{m_1 \ m_2 \ m_3}$ is the Wigner $3j$ symbol. In this Letter we will only consider ultra-local density-density interactions in real space, corresponding to non-zero Haldane pseudopotentials V_0, V_1 . It is important to note that our model is indeed local in real space

(see Ref. [28]), and, like any local model, it becomes nonlocal in orbital (momentum) space as shown in Eq. (1).

At the half-filling (i.e., electron number $N = 2s + 1$), the transverse field h triggers a phase transition from a quantum Hall ferromagnet [31] with spontaneous Z_2 symmetry broken to a quantum paramagnet, which falls into the 2 + 1D Ising universality class. Hereafter, we consider the critical point at $V_1/V_0 = 4.75$, $h/V_0 = 3.16$ that is determined by the order parameter scaling [28]. At the critical point, the state-operator correspondence, a one-to-one correspondence between Hamiltonian eigenenergies and the scaling dimensions of the CFT operators, has been convincingly demonstrated [28], which falls into the 3D Ising universality class.

Operators on the fuzzy sphere.—One can appreciate the beauty of the fuzzy sphere regularization by contrasting it with the familiar lattice regularization/model: The former is defined in the continuum with the full space symmetry [i.e., $SO(3)$ rotation of sphere], while the latter is defined on a discrete lattice with only discrete symmetries such as the lattice translation and rotation. More importantly, the continuum nature of the fuzzy sphere model does not lead to any UV divergence, we still have a finite dimensional Hilbert space to study, thanks to the representation in the Landau orbital basis [28,32]. We can easily translate between operators (and states) on the continuous sphere and those on the discrete Landau orbitals, e.g.,

$$\psi_a(\mathbf{\Omega})^\dagger = \sum_{m=-s}^s c_{m,a}^\dagger Y_{s,m}^{(s)}(\mathbf{\Omega}). \quad (2)$$

Here we are using the spin-weighted spherical Harmonics (also called monopole Harmonics) $Y_{s,m}^{(s)}(\mathbf{\Omega})$ [33] because of the monopole in the origin of the sphere. The more interesting operators are spin operators, which are the particle-hole pairs of electrons. For example, the simplest spin operator is

$$n^a(\mathbf{\Omega}) = \psi^\dagger(\mathbf{\Omega}) \sigma^a \psi(\mathbf{\Omega}) = N \sum_{l=0}^{2s} \sum_{m=-l}^l n_{l,m}^a Y_{l,m}(\mathbf{\Omega}), \quad (3)$$

$$n_{l,m}^a = \frac{1}{2s+1} \int d\mathbf{\Omega} \bar{Y}_{l,m}(\mathbf{\Omega}) n^a(\mathbf{\Omega}). \quad (4)$$

Here we are using spherical Harmonics $Y_{l,m}(\mathbf{\Omega})$ instead of spin-weighted spherical harmonics because the spin is not seeing the monopole placed in the origin. We note that the charge excitation [e.g., created by $\psi^\dagger(\mathbf{\Omega})$] gap is always nonzero but spin degree of freedom undergoes a phase transition [28].

Any gapless spin operator $\hat{\mathcal{O}}(\mathbf{\Omega})$ in our microscopic model can be expressed as a linear combination of CFT scaling operators $[\hat{\phi}_\alpha(\mathbf{\Omega})]$ including primaries and descendants,

$$\hat{\mathcal{O}}(\tau = 0, \mathbf{\Omega}) = \sum_\alpha c_\alpha \hat{\phi}_\alpha(\tau = 0, \mathbf{\Omega}). \quad (5)$$

c_α is a nonuniversal and operator dependent numerical factor, which is generically nonzero as long as the spin operator has the same symmetry (i.e., Ising \mathbb{Z}_2 and parity symmetry) as the CFT scaling operator. Now we can use Eq. (5) and cylinder correlators of CFT to compute OPE coefficients. For any primary state $|\phi_\alpha\rangle$, $\langle\phi_\alpha|\hat{\phi}_\beta|0\rangle$ is nonvanishing only if ϕ_β equals to ϕ_α or its descendant, giving rise to [34]

$$\langle\phi_\alpha|\mathcal{O}(\tau=0, \mathbf{\Omega})|0\rangle = \sum_{n=0}^{\infty} \frac{c_n h_n(\mathbf{\Omega})}{R^{\Delta_\alpha+n}}, \quad (6)$$

where $h_n(\mathbf{\Omega})$ is a universal function fixed by the conformal symmetry, and $h_n(\mathbf{\Omega}) = 1$ if ϕ_α and ϕ_β are both Lorentz scalar primaries [34]. Similarly, we have

$$\langle\phi_\alpha|\mathcal{O}(\tau=0, \mathbf{\Omega})|\phi_\gamma\rangle = \sum_{\beta} f_{\alpha\beta\gamma} \frac{c_\beta \tilde{h}_{\alpha\beta\gamma}(\mathbf{\Omega})}{R^{\Delta_\beta}}. \quad (7)$$

Here $f_{\alpha\beta\gamma}$ is the universal OPE coefficients [6,8], as it determines the fusion channel of two primary operators $\lim_{y \rightarrow x} \phi_\alpha(\mathbf{x}) \times \phi_\beta(\mathbf{y}) \sim \sum f_{\alpha\beta\gamma} \phi_\gamma(\mathbf{x})$. $\tilde{h}_{\alpha\beta\gamma}(\mathbf{\Omega})$ is again a universal function fixed by the conformal symmetry, and is identity if $\phi_{\alpha,\beta,\gamma}$ are Lorentz scalar primaries [34]. Here the summation is over all the primaries and descendants in the operator expansion Eq. (5). Next we will demonstrate the method of extracting OPE coefficients $f_{\alpha\beta\gamma}$ using Eqs. (6) and (7).

Numerical method.—We calculate the low-energy spectrum and corresponding eigenstates of the Hamiltonian Eq. (1) using exact diagonalization (ED) on finite system sizes with up to $N = 18$ Landau orbitals [37]. To corroborate the finite-size extrapolations, we perform large-scale density-matrix renormalization group (DMRG) simulations [38,39]. With a bond dimension $D = 4000$, we are able to access system sizes up to $N = 48$ Landau orbitals. The DMRG data points beyond the ED data are important, as they allow us to confirm and improve the precision of the extrapolated OPE coefficients.

Extracting OPE coefficients.—To implement the idea of accessing OPE coefficients through spin operators, we need to first check if local spin operators produce the correct scaling in Eq. (6). Below we will show the case for the \mathbb{Z}_2 odd operator $\hat{n}^z(\mathbf{\Omega})$ and two \mathbb{Z}_2 even operators $\hat{n}^x(\mathbf{\Omega})$ and $\hat{O}_\epsilon(\mathbf{\Omega}) \equiv \hat{H}(\mathbf{\Omega}) + 2h\hat{n}^x(\mathbf{\Omega})$, where $\hat{H}(\mathbf{\Omega})$ is the Hamiltonian density, i.e., $\int d\mathbf{\Omega} \hat{H}(\mathbf{\Omega})$ is our Hamiltonian in Eq. (1). Figure 1 shows the numerical data of candidate operators, in good agreement with the CFT expectation, indicating the lowest primary fields will dominate the observables of spin operators when the system size is large enough [40]. In particular, $\hat{n}^z(\mathbf{\Omega})$ is dominated by $\hat{\sigma}$ (with conformal dimension ≈ 0.518) with very small high order correction from its descendants, and $\hat{n}^x(\mathbf{\Omega})$, $\hat{O}_\epsilon(\mathbf{\Omega})$ are dominant by $\hat{\epsilon}$ (conformal dimension ≈ 1.412) [41].

Next we turn to extract the OPE coefficients. Let us take the OPE coefficient $f_{\sigma\sigma\epsilon}$ as an example. As we explained

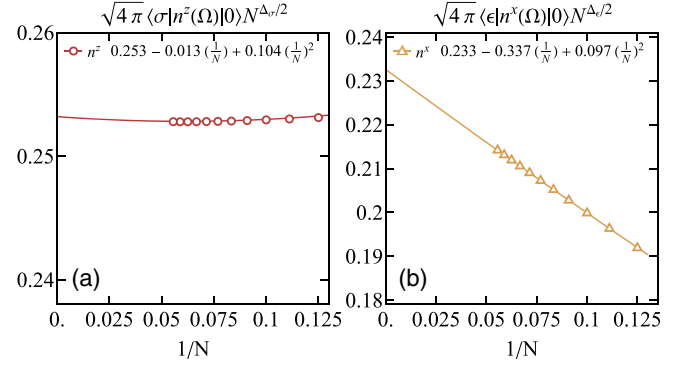


FIG. 1. Local operator content. Extrapolation of (a) $\langle\sigma|\hat{n}^z(\mathbf{\Omega})|0\rangle \times N^{\Delta_\sigma/2}$, (b) $\langle\epsilon|\hat{n}^x(\mathbf{\Omega})|0\rangle \times N^{\Delta_\epsilon/2}$. The finite-size correction of N^{-1} , N^{-2} , respectively, comes from the descendant fields as shown in Eq. (6).

above, to approximate $\hat{\sigma}$ a natural candidate is $n^z(\mathbf{\Omega})$, and its operator expansion should naturally contain all the \mathbb{Z}_2 odd primaries $\sigma, \sigma_{\mu\nu}, \sigma', \dots$ as well as their descendants $\partial_\mu\sigma, \square\sigma, \partial_\mu\sigma_{\mu\nu}, \dots$. So we shall have,

$$\langle\sigma|n^z(\mathbf{\Omega})|0\rangle = \frac{1}{R^{\Delta_\sigma}} \left(c_\sigma + \sum_{n=1}^{\infty} \frac{a_n}{R^{2n}} \right), \quad (8)$$

$$\langle\sigma|n^z(\mathbf{\Omega})|\epsilon\rangle = \frac{f_{\sigma\sigma\epsilon}}{R^{\Delta_\sigma}} \left(c_\sigma + \sum_{n=1}^{\infty} \frac{\tilde{a}_n}{R^{2n}} \right) + \dots, \quad (9)$$

where \dots stands for the contribution from other primaries and associated descendants. We note that only scalar scaling operators can contribute to these two observables. Therefore, we can compute $f_{\sigma\sigma\epsilon}$ using

$$\frac{\langle\sigma|n^z(\mathbf{\Omega})|\epsilon\rangle}{\langle\sigma|n^z(\mathbf{\Omega})|0\rangle} = f_{\sigma\sigma\epsilon} + \frac{a_1 - \tilde{a}_1}{c_\sigma R^2} + O(R^{-4}). \quad (10)$$

In principle, no finite size extrapolation is needed once the system size R is large enough. In the small size simulation like our case, we will perform a simple linear extrapolation with respect to R^{-2} (i.e., N^{-1}) that includes the contribution from the descendant field $\square\sigma$. Similarly, other OPE coefficient $f_{\sigma\phi_1\phi_2}$ can be computed by the \mathbb{Z}_2 odd spin operator $n^z(\mathbf{\Omega})$ using $f_{\sigma\phi_1\phi_2} \approx \langle\phi_1|n^z(\mathbf{\Omega})|\phi_2\rangle / \langle\sigma|n^z(\mathbf{\Omega})|0\rangle$ and OPE coefficients $f_{\epsilon\phi_1\phi_2}$ can be computed by the \mathbb{Z}_2 even spin operators $\hat{n}^x(\mathbf{\Omega})$, $\hat{O}_\epsilon(\mathbf{\Omega})$. One should be cautious that certain OPE coefficient may have a different subleading term (see [34] for details).

First, we discuss the OPE coefficients involving three scalar primary operators. For example, the OPE coefficient $f_{\sigma\sigma\epsilon}$ can be extracted by Eq. (10), where the lowest correction comes from the descendant field $\square\sigma$. Figure 2(a) shows the finite-size extrapolation indeed agrees with Eq. (10), giving $f_{\sigma\sigma\epsilon} \approx 1.0539$ by extrapolation. Alternatively, $f_{\sigma\sigma\epsilon}$ can be obtained by choosing spin operators $\hat{O}(\mathbf{\Omega}) = n^x(\mathbf{\Omega})$ and

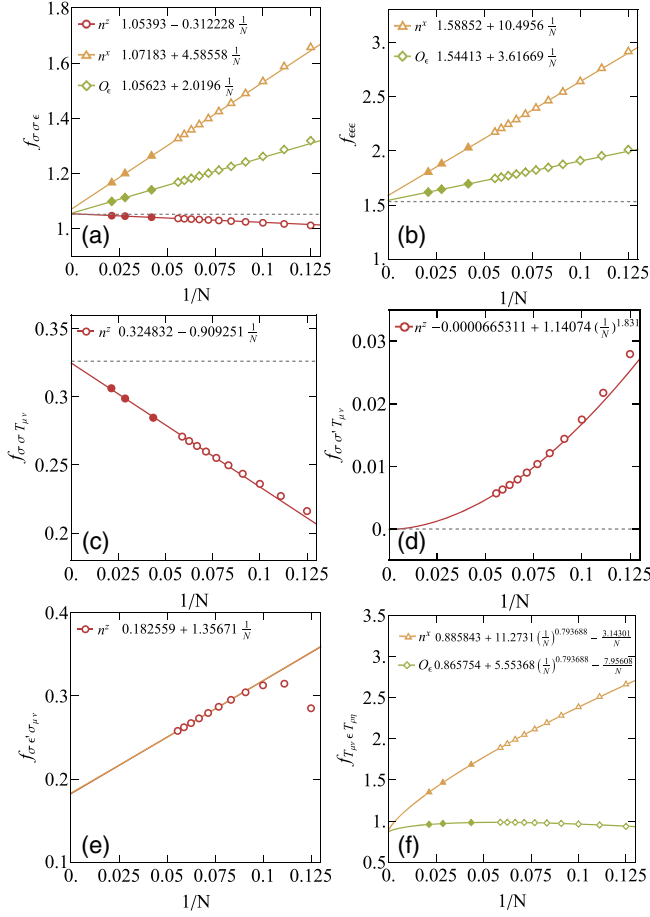


FIG. 2. OPE coefficients of primary operators. (a)–(b) Two representative OPE coefficients involving three scalar primaries $f_{\sigma\sigma\epsilon}$ and $f_{\epsilon\epsilon\epsilon}$, obtained from the finite-size extrapolation via spin operator $n^z(\mathbf{\Omega})$ (red), $n^x(\mathbf{\Omega})$ (yellow), and $O_\epsilon(\mathbf{\Omega})$ (green) (see main text). (c)–(d) Two representative OPE coefficients involving energy-momentum tensor $f_{\sigma\sigma T_{\mu\nu}}$, $f_{\sigma'\sigma T_{\mu\nu}}$ by using $n^z(\mathbf{\Omega})$ operator. The dashed line in (a)–(d) is the value from conformal bootstrap. (e)–(f) Two representative OPE coefficients $f_{\sigma\epsilon'\sigma_{\mu\nu}}$, $f_{T_{\mu\nu}\epsilon T_{\rho\eta}}$ that are not known in conformal bootstrap calculation. In all figures only the data points on largest six sizes are used in the fitting. In (a)–(c) and (f), larger sizes up to $N = 48$ are available by the DMRG method which are labeled by solid symbols.

$\hat{\mathcal{O}}(\mathbf{\Omega}) = H(\mathbf{\Omega}) + 2hn^x(\mathbf{\Omega})$ to approximate the $\hat{\epsilon}$ primary (see Supplemental Material, Sec. III [34]). Figure 2(a) confirms that different local operators give consistent estimates of $f_{\sigma\sigma\epsilon}$. Within the similar procedure, we also obtain a few more OPE coefficients involving scalar primary operators as listed in Tab. I.

Second, for the OPE coefficients involving spinning operators, the calculation is largely parallel except that we need to consider the angle $(\mathbf{\Omega})$ dependence [34]. The angle dependence can be eliminated by integrating over the sphere against the spherical Harmonics, for example, one representative OPE coefficient $f_{\sigma\sigma T_{\mu\nu}}$ should follow (see [34])

TABLE I. List of OPE coefficients of primary operators obtained on the fuzzy sphere. The primary operators in consideration have scaling dimensions $\Delta_\sigma \approx 0.51815$, $\Delta_\epsilon \approx 1.4126$, $\Delta_{\epsilon'} \approx 3.8297$, $\Delta_{\sigma'} \approx 5.2906$, $\Delta_{T_{\mu\nu}} = 3$, and $\Delta_{\sigma_{\mu\nu}} \approx 4.1803$. The conformal bootstrap (CB) data is from Ref. [14], where some unavailable data are labeled by “NA.” We note that our convention for $f_{T_{\mu\nu}\epsilon T_{\rho\eta}}$ is $f_{T_{\mu\nu}\epsilon T_{\rho\eta}} = (1/\sqrt{4\pi}) \int d\mathbf{\Omega} \langle T_{\mu\nu}, m = 0 | \hat{\epsilon}(\mathbf{\Omega}) | T_{\rho\eta}, m = 0 \rangle$ [34].

Operators	Spin	Z_2	$f_{\alpha\beta\gamma}$ (Fuzzy sphere)	$f_{\alpha\beta\gamma}$ (CB)
σ	0	–	$f_{\sigma\sigma\epsilon} \approx 1.0539(18)$	$f_{\sigma\sigma\epsilon} \approx 1.0519$
ϵ	0	+	$f_{\epsilon\epsilon\epsilon} \approx 1.5441(23)$	$f_{\epsilon\epsilon\epsilon} \approx 1.5324$
ϵ'	0	+	$f_{\sigma\sigma\epsilon'} \approx 0.0529(16)$	$f_{\sigma\sigma\epsilon'} \approx 0.0530$
σ'	0	–	$f_{\epsilon\epsilon\epsilon'} \approx 1.566(68)$	$f_{\epsilon\epsilon\epsilon'} \approx 1.5360$
			$f_{\sigma'\sigma\epsilon} \approx 0.0515(42)$	$f_{\sigma'\sigma\epsilon} \approx 0.0572$
			$f_{\sigma'\sigma\epsilon'} \approx 1.294(51)$	NA
			$f_{\sigma'\epsilon\epsilon'} \approx 2.98(13)$	NA
$T_{\mu\nu}$	2	+	$f_{\sigma\sigma T} \approx 0.3248(35)$	$f_{\sigma\sigma T} \approx 0.3261$
			$f_{\sigma'\sigma T} \approx -0.00007(96)$	$f_{\sigma'\sigma T} = 0$
			$f_{\epsilon\epsilon T} \approx 0.8951(35)$	$f_{\epsilon\epsilon T} \approx 0.8892$
			$f_{T\epsilon T} \approx 0.8658(69)$	NA
$\sigma_{\mu\nu}$	2	–	$f_{\sigma\epsilon\sigma_{\mu\nu}} \approx 0.400(33)$	$f_{\sigma\epsilon\sigma_{\mu\nu}} \approx 0.3892$
			$f_{\sigma\epsilon'\sigma_{\mu\nu}} \approx 0.18256(69)$	NA

$$\sqrt{\frac{15}{8}} \frac{\langle \sigma | \int d\mathbf{\Omega} \bar{Y}_{2,0}(\mathbf{\Omega}) \hat{n}^z(\mathbf{\Omega}) | T_{\mu\nu}, m = 0 \rangle}{\langle \sigma | \int d\mathbf{\Omega} \bar{Y}_{0,0}(\mathbf{\Omega}) \hat{n}^z(\mathbf{\Omega}) | 0 \rangle} \approx f_{\sigma\sigma T_{\mu\nu}} + \frac{a}{N}. \quad (11)$$

The corresponding finite-size extrapolations are shown in Fig. 2(c) and the estimates of OPE are in Table I.

Third, we can calculate a few more OPE coefficients which are unknown before. Figures 2(e) and 2(f) show the results of $f_{\sigma\epsilon'\sigma_{\mu\nu}}$ and $f_{T_{\mu\nu}\epsilon T_{\rho\eta}}$, the latter is defined as

$$\frac{\langle T_{\mu\nu}, m = 0 | \int d\mathbf{\Omega} \hat{\mathcal{O}}(\mathbf{\Omega}) | T_{\mu\nu}, m = 0 \rangle - \langle 0 | \int d\mathbf{\Omega} \hat{\mathcal{O}}(\mathbf{\Omega}) | 0 \rangle}{\langle \epsilon | \int d\mathbf{\Omega} \hat{\mathcal{O}}(\mathbf{\Omega}) | 0 \rangle} \approx f_{T_{\mu\nu}\epsilon T_{\rho\eta}} + \frac{a}{N^{0.7937}} + \frac{b}{N}, \quad (12)$$

where $\hat{\mathcal{O}} = \hat{n}^x, \hat{O}_\epsilon$. The estimated values are shown in Table I [42].

To quantify the data we get, we compare these results with the numerical bootstrap data in Table I. Most of OPE coefficients match known results from the conformal bootstrap [11,14] with discrepancies smaller than 2%. Importantly, we also compute four OPE coefficients which have not been computed before. The ability to access OPEs which are not known before shows the superiority of fuzzy sphere method. In addition, we present a systematical error analysis in Supplemental Material, [34], where the mean values are estimated by \hat{n}^z or \hat{O}_ϵ , which have smaller finite-size effect, and the error bars are estimated by different fitting processes.

Summary and discussion.—We have explained the procedure of extracting the conformal data of emergent 3D conformal field theory using the recently proposed fuzzy sphere regularization. By exploiting the state-operator correspondence, we accurately compute operator product expansion (OPE) coefficients of CFT primary fields with low-lying conformal dimensions in a $(2 + 1)$ -D quantum Ising model. As far as we know, most of the OPE coefficients reported here have never been studied in a microscopic model before. Some of OPE coefficients are not even computed by the conformal bootstrap calculations. More importantly, our work exposes the operator perspective of 3D CFTs on the fuzzy sphere, which enables a lot more important physics to explore. For example, we can directly compute 3D CFT four-point correlators—the core object in the classic and modern story of conformal bootstrap. Another interesting direction is to study the operator growth of 3D CFTs under a quantum quench, which will lead to new insights on the quantum chaos of CFTs. We foresee that the fuzzy sphere scheme will be a powerful tool to study 3D CFTs in a broad universality class [e.g., Wilson-Fisher $O(N)$ critical point and critical gauge theories.].

We acknowledge the useful discussion with Yijian Zou. We thank C. Han, E. Huffman, and J. Hofmann for collaboration on a related project. L. D. H. and W. Z. were supported by National Natural Science Foundation of China (No. 92165102, No. 11974288), National key R&D program (No. 2022YFA1402204), and the Key R&D Program of Zhejiang Province (2021C01002). Research at Perimeter Institute is supported in part by the Government of Canada through the Department of Innovation, Science and Industry Canada and by the Province of Ontario through the Ministry of Colleges and Universities.

Note added.—Recently, we noticed that OPE coefficient f_{TeT} was recently computed by bootstrapping a five-point correlator [43]. Their numerical estimation is close to our result.

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