Revealing the Harmonic Structure of Nuclear Two-Body Correlations in High-Energy Heavy-Ion Collisions

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Smashing nuclei at ultrarelativistic speeds and analyzing the momentum distribution of outgoing debris provides a powerful method to probe the many-body properties of the incoming nuclear ground states. Within a perturbative description of initial-state fluctuations in the quark-gluon plasma, we express the measurement of anisotropic flow in ultracentral heavy-ion collisions as the quantum-mechanical average of a specific set of operators measuring the harmonic structure of the two-body azimuthal correlations among nucleons in the colliding states. These observables shed a new light on spatial correlations in atomic nuclei while enabling us to test the complementary pictures of nuclear structure delivered by low- and high-energy experiments on the basis of state-of-the-art theoretical approaches rooted in quantum chromodynamics.

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Atomic nuclei are among the most complex quantum systems observed in nature. They exhibit a rich spectrum of many-body phenomena, leaving measurable fingerprints on processes spanning subnuclear [1,2] to astrophysical [3] scales, including exotic interactions relevant to searches for physics beyond the standard model [4–10]. Despite nearly a century of investigations, a clear understanding of many nuclear properties, even in their ground states, remains elusive. Recent years have, however, witnessed important steps forward on both the theoretical and the experimental forefronts.

Breakthroughs in many-body methods and access to unprecedented computing power have led to tremendous progress in the development of first-principles, i.e., *ab initio*, approaches to nuclear structure [11–13]. Such first-principles description rely on the application of scalable many-body frameworks [14–19] on the basis of internucleon interactions explicitly rooted into low-energy effective field theories of the strong force described via quantum chromodynamics (QCD) [20–24]. Although systematic computations across the nuclear chart are still in their infancy [25–35], *ab initio* calculations have recently proven able to address both closed- and open-shell systems in unprecedented regimes of nuclear masses [36–47], which are expected to extend further in the near future. In this context, an understanding of many-body correlations in atomic nuclei rooted in the elementary description of the strong force is within reach.

On the experimental front, traditional low-energy nuclear structure experiments are being complemented by high-energy experiments conducted at the Relativistic Heavy Ion Collider (RHIC) or the Large Hadron Collider (LHC). Because of the completely different timescales at play, this novel platform to study nuclear structure [48] reveals more vividly the internucleon correlations characterizing the ground state of the collided species [49] through the analysis of the momentum distributions of hadrons emitted from the quark-gluon plasma (QGP) [50,51]. Analysis of these experiments via a rigid-rotor model of the colliding ions have already revealed key information regarding collective correlations in these nuclei in terms of intrinsic deformations, including quadrupole [52–56], triaxial [2,57,58], octupole [59], and hexadecapole [60,61] deformations. In parallel, a wide array of more or less sophisticated nuclear models, used as input to heavy-ion collision simulations, have proven successful in capturing the experimental observations [62–83].

In spite of the progress made, our understanding of the observed flow phenomena remains largely based on a

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simple classical picture where the orientation average of the intrinsic shape acts as the sole source of correlations among nucleons. Although effective in capturing the experimental results, this picture ultimately gives little insight into the nuclear force or the detailed origin of nuclear collectivity. Given the great potential of high-energy collisions in shedding a complementary view of nucleon dynamics in nuclear ground states, this calls for a paradigm shift. As originally argued in Ref. [49], one would rather formulate multiparticle correlation observables accessible in heavy-ion collisions within the more rigorous language of many-body operators and quantum correlation functions of nucleons, that encompass the notion of the intrinsic shape (including shape coexistence and fluctuations) while linking more directly to the underlying theory. To achieve this, there are two possible complementary approaches: (i) understanding on a more rigorous ground the connection between the measured observables and nuclear ground-state properties and (ii) performing a systematic analysis of such observables based on state-of-the-art *ab initio* nuclear-structure calculations.

In this Letter, we take a major step regarding point (i) in view of addressing point (ii) on a solid footing in the future. We show that measurements of anisotropic flow in the final states of heavy-ion collisions can be formulated in terms of the quantum-mechanical average of a new class of two-body self-adjoint operators in the colliding ground states. These observables give access to the harmonic decomposition of the associated two-body densities, establishing ultimately a new paradigm for understanding collider observations in the language of ground-state many-body correlations of nucleons.

To reach this goal, the route proceeds via a simple theoretical picture of initial-state fluctuations in heavy-ion collisions. As in most phenomenological approaches, a nucleus at high energy is associated with a density of matter in the plane transverse to the beam direction, or *thickness* function:

$$t(\mathbf{x}) \equiv \sum_{i=1}^{A} g(\mathbf{x} - \mathbf{r}_{\perp i}), \tag{1}$$

where A is the nuclear mass number and $\mathbf{x} = (x, y)$ denotes a position vector in the transverse plane with its origin located at the center of mass of the nucleus ($\mathbf{x} = 0$), while $\mathbf{r}_{\perp i}$ is the transverse position of the ith nucleon. The function $g(\mathbf{x})$ parametrizes the two-dimensional shape of a nucleon as seen by a high-energy probe.

Equation (1) describes an individual nucleus in an individual collision. The elements of randomness are the positions of the nucleons in the transverse plane, $\mathbf{r}_{\perp i}$. In our understanding of high-energy collisions, the process resolves the instantaneous structure of the colliding objects down to scales dictated by the momentum transfer, which is on the order of 1 GeV or higher. Therefore, the nucleon positions are *frozen* while the two nuclei cross each other [84,85]. If the ground state is characterized by the A-body wave function (omitting spin and isospin)

$$\Psi(\mathbf{r}_1, \ldots, \mathbf{r}_A) \equiv \langle \mathbf{r}_1, \ldots, \mathbf{r}_A | \Psi \rangle, \tag{2}$$

the probability to find nucleons at coordinates $\mathbf{r}_i = (\mathbf{r}_{\perp i}, r_{zi})$ is given by $|\langle \mathbf{r}_1, \ldots, \mathbf{r}_A | \Psi \rangle|^2$. Thus, the nucleon positions sampled according to this probability distribution reflect up to A-body correlations among them [86].

Given the thickness functions $t(\mathbf{x})$ and $t'(\mathbf{x})$ of two colliding nuclei, generic data-driven arguments are employed to move to the initial entropy density of the subsequently formed QGP. First, the interactions of partons from the colliding nuclei are semihard processes, which in the transverse plane are localized over small fractions of the size of a nucleon. Therefore, the entropy deposition at a given location $\mathbf{x}$ can depend on only the values of the two thickness functions in the immediate vicinity of such a point, i.e., $t(\mathbf{x})$ and $t'(\mathbf{x})$. Second, the QGP is initially so hot that it can be roughly viewed as an ideal gas of quarks and gluons, where the total entropy determines the particle number [91]. As the evolution of the QGP is nearly isentropic, this implies that the final-state multiplicity is dictated by the initial-state entropy. In turn, in the limit of ultracentral collisions (small impact parameters), it is observed experimentally that the final-state multiplicity scales with the mass number of the colliding nuclei [92]. Combining these points, the entropy density at midrapidity is of the form

$$s(\mathbf{x}) = \mathcal{F}[t(\mathbf{x}), t'(\mathbf{x})], \qquad \int_{\mathbf{x}} s(\mathbf{x}) \propto A, \tag{3}$$

where $\mathcal{F}$ is some model-dependent function. The above reasoning underlies the popular TRENTo ansatz for the entropy density of heavy-ion collisions [93] and is supported by the results of a Bayesian analysis of LHC data with more generic initial conditions [94].

Next, the notion of *anisotropic flow* of hadrons [95] is introduced to connect properties of the initial entropy density $s(\mathbf{x})$ to measurable quantities in the final states of the collisions. In a collision, hadrons are emitted at midrapidity with some azimuthal distribution

$$\frac{dN}{d\phi} \propto 1 + 2v_2 \cos[2(\phi - \phi_2)] + 2v_3 \cos[3(\phi - \phi_3)] + \cdots. \tag{4}$$

The set of complex Fourier coefficients $V_n = v_n e^{in\phi_n}$, $n \geq 2$, defines the anisotropic flow of the system. Because of the hydrodynamic nature of the QGP, the anisotropy coefficients in momentum space, V_n, are determined by the anisotropies in coordinate space, $\mathcal{E}_n$, characterizing the large-scale structure of the initial entropy density field [96]. These quantities are defined by [with $\mathbf{x} = (|\mathbf{x}|, \phi_x)$]

$$\mathcal{E}_n \equiv -\frac{\int_{\mathbf{x}} |\mathbf{x}|^n e^{in\phi_x} s(\mathbf{x})}{\int_{\mathbf{x}} |\mathbf{x}|^n s(\mathbf{x})}, \qquad \varepsilon_n = |\mathcal{E}_n|. \tag{5}$$

Given the probabilistic nature of the nucleon distributions in the colliding nuclei, the field $s(\mathbf{x})$ and, consequently, the value of $\mathcal{E}_n$ fluctuate event by event. Hydrodynamic studies show, then, that v_n in the final state is linearly correlated with the value of ε_n in the initial state: If we observe a fluctuation in the value of ε_n, a similar variation will be also observed in the value of v_n, i.e., $v_n \propto \varepsilon_n$ [97]. Following Ref. [98], a background-fluctuation splitting is, thus, introduced according to

$$s(\mathbf{x}) = \langle s(\mathbf{x})\rangle + \delta s(\mathbf{x}), \qquad \langle \delta s(\mathbf{x})\rangle = 0, \tag{6}$$

where the brackets denote a statistical average over ultracentral collision events. Inserting Eq. (6) into the expression of $\mathcal{E}_n$ in Eq. (5), expanding in powers of δs, and keeping terms up to second order, one obtains

$$\langle v_n^2\rangle \propto \langle \varepsilon_n^2\rangle = \frac{\int_{\mathbf{x},\mathbf{y}} |\mathbf{x}|^n |\mathbf{y}|^n e^{in(\phi_x-\phi_y)} C_2(\mathbf{x},\mathbf{y})}{(\int_{\mathbf{x}} |\mathbf{x}|^n C_1(\mathbf{x}))^2}, \tag{7}$$

where the one-point and connected two-point functions of the entropy density field have been introduced:

$$\begin{aligned} C_1(\mathbf{x}) &\equiv \langle s(\mathbf{x})\rangle, \\ C_2(\mathbf{x},\mathbf{y}) &\equiv \langle s(\mathbf{x}) s(\mathbf{y})\rangle - C_1(\mathbf{x}) C_1(\mathbf{y}). \end{aligned} \tag{8}$$

As shown in Ref. [49], Eq. (7) is accurate at the percent level for collisions of nuclei with $A \sim 100$. This completes our leading-order picture of initial-state and spatial-anisotropy fluctuations in heavy-ion collisions. The variations of a quantity that is experimentally accessible, the flow coefficient v_n, are determined by the variations of an initial-state quantity, ε_n, whose fluctuations are, in turn, related to the correlation functions of the entropy density field, $\langle s(\mathbf{x})\rangle$ and $\langle s(\mathbf{x}) s(\mathbf{y})\rangle$, via a perturbative expansion.

One important link is now added to the above scheme by connecting the fluctuations of the entropy density to the many-body wave functions of the colliding nuclei. In practice, a model of $s(\mathbf{x})$ satisfying the generic requirements outlined above is needed. The simplest choice, amenable to an analytical treatment, corresponds to an arithmetic average [99]:

$$s(\mathbf{x}) = s_0 \frac{t(\mathbf{x}) + t'(\mathbf{x})}{2}, \tag{9}$$

where s_0 is a normalization typically inferred from experimental data. When the two colliding nuclei are nearly identical, one has $s(\mathbf{x}) \propto t(\mathbf{x})$. The entropy density, thus, becomes akin to that discussed in so-called independent source model calculations, which have been studied at large in heavy-ion collisions [98,106–112]. In our case, A sources determine the scaling of the total entropy with the size of the colliding nuclei, but their *independent* character is abandoned given that their density in coordinate space is identified with the ground-state probability $|\langle \mathbf{r}_1, \ldots, \mathbf{r}_A|\Psi\rangle|^2$. This has a profound conceptual implication given that it allows one to commute the statistical average of the entropy density field into a quantum-mechanical expectation in the nuclear ground state. As shown in Supplemental Material [100], Eq. (9) leads to

$$\begin{aligned} C_1(\mathbf{x}) &= s_0 A \int_{\mathbf{r}_1} \rho^{(1)}(\mathbf{r}_1) g(\mathbf{x}-\mathbf{r}_{1\perp}), \\ C_2(\mathbf{x},\mathbf{y}) &= \frac{s_0^2}{2}\Big[A \int_{\mathbf{r}_1} \rho^{(1)}(\mathbf{r}_1) g(\mathbf{x}-\mathbf{r}_{1\perp}) g(\mathbf{y}-\mathbf{r}_{1\perp}) \\ &\quad + (A^2-A)\int_{\mathbf{r}_1,\mathbf{r}_2} \rho^{(2)}(\mathbf{r}_1,\mathbf{r}_2) g(\mathbf{x}-\mathbf{r}_{1\perp}) g(\mathbf{y}-\mathbf{r}_{2\perp}) \\ &\quad - A^2 \int_{\mathbf{r}_1} \rho^{(1)}(\mathbf{r}_1) g(\mathbf{x}-\mathbf{r}_{1\perp}) \int_{\mathbf{r}_2} \rho^{(1)}(\mathbf{r}_2) g(\mathbf{y}-\mathbf{r}_{2\perp})\Big], \end{aligned} \tag{10}$$

where $\rho^{(1)}$ and $\rho^{(2)}$ denote the local one- and two-body ground-state densities, respectively, defined through

$$\rho^{(n)}(\mathbf{r}_1, \ldots, \mathbf{r}_n) \equiv \int_{\mathbf{r}_{n+1},\ldots,\mathbf{r}_A} |\Psi(\mathbf{r}_1, \ldots, \mathbf{r}_A)|^2. \tag{11}$$

As shown in Supplemental Material [100], inserting Eqs. (10) into Eq. (7) and considering that the nucleon form factor $g(\mathbf{x})$ is a highly localized function that can be approximated by a Dirac delta, we arrive at the following key expression:

$$\begin{aligned} \langle \varepsilon_n^2\rangle &= \frac{1}{2A^2} \frac{1}{\left(\int_{\mathbf{r}} \rho^{(1)}(\mathbf{r}) \hat{R}_n(\mathbf{r})\right)^2} \Big[A \int_{\mathbf{r}} \rho^{(1)}(\mathbf{r}) \hat{R}_{2n}(\mathbf{r}) \\ &\quad + (A^2-A)\int_{\mathbf{r}_1,\mathbf{r}_2} \rho^{(2)}(\mathbf{r}_1,\mathbf{r}_2) \hat{\mathcal{E}}_n(\mathbf{r}_1,\mathbf{r}_2)\Big] \\ &= \frac{1}{2A} \frac{1}{\langle \hat{R}_n\rangle^2} \Big[\langle \hat{R}_{2n}\rangle + (A-1)\langle \hat{\mathcal{E}}_n\rangle\Big], \end{aligned} \tag{12}$$

involving the ground-state expectation value of the one- and two-body operators

$$\hat{R}_n(\mathbf{r}) \equiv r_{1\perp}^n = (r_{1x}^2 + r_{1y}^2)^{n/2}, \tag{13a}$$

$$\hat{\mathcal{E}}_n(\mathbf{r}_1,\mathbf{r}_2) \equiv (r_{1x} + i r_{1y})^n (r_{2x} - i r_{2y})^n \tag{13b}$$

$$= r_{1\perp}^n \, r_{2\perp}^n \, e^{in(\phi_1-\phi_2)} [\text{polar coords}]. \tag{13c}$$

Thus, the mean squared anisotropy $\langle \varepsilon_n^2 \rangle$ is expressed in terms of the ground-state expectation of a one-body transverse radius operator $\hat{R}_n(\mathbf{r})$ and of an intriguing two-body *eccentricity operator* $\hat{\mathcal{E}}_n(\mathbf{r}_1, \mathbf{r}_2)$, which measures the azimuthal anisotropy of the local two-body density in the transverse plane. It should be noted that there is a strong similarity between Eq. (13c) and the definition of the mean squared elliptic flow, $\langle v_n^2 \rangle = \langle e^{in(\phi_{1p} - \phi_{2p})} \rangle$. For the latter observable, one measures a two-body azimuthal (ϕ_p) correlation of hadrons in momentum space. In Eq. (13c), we are instead interested in two-body azimuthal correlations of nucleons in coordinate space. The two are, thus, connected to each other via the hydrodynamic evolution of the system.

It is insightful to rewrite the eccentricity operator using spherical coordinates typically employed in the context of nuclear structure. Introducing nth-order (complex) spherical harmonics in their maximal projections,

$$Y_n^{\pm n}(x, y, z) \equiv c_{\pm n} \frac{(x \pm iy)^n}{(x^2 + y^2 + z^2)^{n/2}}, \tag{14}$$

where c_n is a real coefficient, and moving to spherical coordinates with $\mathbf{r} = (r, \Omega)$, one arrives at

$$\begin{aligned} \hat{\mathcal{E}}_n(\mathbf{r}_1, \mathbf{r}_2) &= c_n^{-1} r_1^n Y_n^n(\Omega_1) c_{-n}^{-1} r_2^n Y_n^{-n}(\Omega_2) \\ &\equiv \hat{F}_n(\mathbf{r}_1) \hat{F}_{-n}(\mathbf{r}_2), \end{aligned} \tag{15}$$

whose derivation is also reported in Supplemental Material [100]. We stress that the appearance of the maximal projections of the spherical harmonics is due to the Lorentz contraction of the three-dimensional nuclear state. Indeed, in a classical picture, this observable measures the squared magnitude of the nth-order eccentricity of a two-dimensional image obtained by projecting a deformed three-dimensional shape onto a plane (in this case, the x-y plane).

We note that the operator $\hat{\mathcal{E}}_2(\mathbf{r}_1, \mathbf{r}_2)$ bears some resemblance with the two-body Kumar quadrupole operator [113]. However, while no low-energy experiment has been identified to directly measure Kumar's observables in nuclear ground states, the novel set of eccentricity operators can be measured through the elliptic flow in symmetric [114] nucleus-nucleus collisions at high energy. Note that, much as higher-order Kumar parameters enable one to extract additional information about nuclear collectivity, it will also be possible to derive an expression akin to Eq. (12) for other types of correlators that will probe the local three-body density [49].

Let us briefly discuss how the discussion is modified in the scenario where the collisions occur at nonzero impact parameter. The coordinates $\mathbf{x} = (x, y)$ can be parametrized in the so-called *reaction plane* (RP), where x is the direction of the impact parameter. The mean squared elliptic anisotropy receives a contribution from the eccentricity of the average density profile, $\varepsilon_{\mathrm{RP}}$, as follows (see, e.g., Appendix A in [116]):

$$\langle v_2^2 \rangle \propto \langle \varepsilon_2^2 \rangle = \sigma_2^2 + \frac{\int_{\mathbf{x}} (y^2 - x^2) C_1(\mathbf{x})}{\int_{\mathbf{x}} |\mathbf{x}|^2 C_1(\mathbf{x})} \left(\equiv \varepsilon_{\mathrm{RP}}^2 \right), \tag{16}$$

where σ_2^2 corresponds to the right-hand side of Eq. (7) in the limit of ultracentral collisions. Although we do not derive this here, $\varepsilon_{\mathrm{RP}}$ adds a correction to Eq. (12) which is proportional to $\int_{x,y,z} (y^2 - x^2) \rho^{(1)}(x - b, y, z)$, where b is the impact parameter. This contribution is, thus, maximal for coordinates at a distance b from the center of the nucleus. As is well known, $\varepsilon_{\mathrm{RP}}$ quickly overshadows the fluctuation-driven term σ^2 that involves $C_2(\mathbf{x}, \mathbf{y})$ as soon as b increases. Therefore, while ultracentral collisions are suited to isolate the impact of two- and many-body correlations, off-central collisions probe the radial dependence of the one-body density. This has recently allowed, in particular, extracting the neutron skin of ^{208}Pb from LHC data [72] or isolating imprints of the difference in skin thickness between ^{96}Ru and ^{96}Zr in RHIC data [117].

Back to the ultracentral collisions, in the uncorrelated limit where $\rho^{(2)}(\mathbf{r}_1, \mathbf{r}_2) = \rho^{(1)}(\mathbf{r}_1)\rho^{(1)}(\mathbf{r}_2)$, the second term in the rhs of Eq. (12) vanishes due to the spherical symmetry of the local one-body density of $J = 0$ even-even ground states. One is left with the first term associated with the transverse radius operator corresponding to the independent-source result proportional to $1/A$ originally derived in Ref. [106].

In reality, any ground state displays nonzero two-body correlations that are, thus, captured by the second term in the rhs of Eq. (12) via the two-body operator $\hat{\mathcal{E}}_n(\mathbf{r}_1, \mathbf{r}_2)$. For example, quadrupole correlations among nucleons *in the ground state* are directly accessed in the transverse plane via $\langle v_2^2 \rangle$; i.e., the larger $\langle v_2^2 \rangle$, the larger these correlations. This is, in essence, our main result.

While these novel observables can quantify specific two-body spatial correlations in atomic nuclei, they cannot, *per se*, characterize the inner mechanisms that generate them. This can be achieved only *within* a given theoretical scheme, in a way that depends on such a scheme. Our ambition is eventually to combine the measurement of these collider observables for a portfolio of nuclei and to analyze whether and how a given *ab initio* theoretical scheme reproduces them. For example, expansion many-body methods proceeding through the breaking and restoration of rotational symmetry along with the resummation of missing dynamical correlations [27–29,32] can quantify the extent to which, e.g., quadrupole correlations measured via the fluctuations of the mean squared anisotropy are due to static quadrupole deformation, quadrupole shape fluctuations, and residual noncollective two-body correlations. This is a long-term research program that the present Letter seeks to motivate.

For now, let us proceed through such an analysis based on an elementary nuclear model that has proven useful in heavy-ion collision simulations, i.e., the textbook rigid-rotor model where the incoming nuclei are modeled as (potentially) deformed intrinsic shapes rotating in all possible ways [118,119]. More specifically, a colliding nucleus is treated as a batch of nucleons independently sampled from a deformed intrinsic density that has, in each collision, a random orientation in space. The model assumes, by construction, that spatial correlations in the nucleus can be entirely mocked up via the rotation of nucleons collectively aligned along a preferred direction in the intrinsic frame; i.e., the potentially deformed shape is rigid, and there is no concept of shape fluctuations or noncollective correlations. The deformations of the associated intrinsic density distribution are quantified via standard Bohr multipole parameters β_n. Using such a nuclear model, one obtains [120]

$$\langle \varepsilon_n^2 \rangle = a_n + b_n \beta_n^2, \tag{17}$$

where a_n and b_n are positive coefficients. The first and second terms in the right-hand side of Eq. (17) are naturally identified with the first and second terms in Eq. (12), respectively. Within this model, one can conclude that the operator $\hat{\mathcal{E}}_n$ quantifies the (square of the) effective deformation of the intrinsic rotor [121].

To illustrate more explicitly how the novel observables are interpreted within the rigid-rotor model, the intrinsic density is taken as a simple Gaussian function:

$$\rho(r, \Omega) \equiv \frac{1}{(2\pi)^{3/2} R^3} \exp\left(-\frac{r^2}{2R(\Omega)^2}\right), \tag{18}$$

with the nuclear surface expanded in quadrupole and octupole modes:

$$R(\Omega) \equiv R[1 + \beta_{20} Y_2^0(\Omega) + \beta_{30} Y_3^0(\Omega)]. \tag{19}$$

From a Taylor expansion truncated to first order in the deformation parameters β_{20} and β_{30}, the evaluation of Eq. (12) for $n = 2, 3$ leads to

$$\langle \varepsilon_2^2 \rangle = \frac{1}{A} + \frac{3}{4\pi} \beta_{20}^2, \qquad \langle \varepsilon_3^2 \rangle = \frac{16}{3\pi A} + \frac{2048}{245\pi^3} \beta_{30}^2. \tag{20}$$

In each case, the first term proportional to $1/A$ corresponds to that derived in independent-source calculations with a Gaussian density of sources (see, e.g., Ref. [110]). The second term delivers the correction induced by two-body correlations through intrinsic axial deformation parameters with prefactors matching exactly those obtained in Ref. [118] based on a liquid-drop-like model rotating in space. More sophisticated calculations are reported in Supplemental Material [100], showing that the correction to $\langle \varepsilon_2^2 \rangle$ induced by the quadrupole deformation of the nucleus is robust and independent of the precise modeling of the collisions.

Because correlations impact the two-body density in any nucleus, the rigid-rotor model is *bound* to send back a nonzero effective deformation for all nuclei through the second term of Eq. (17). As a matter of fact, this model has proven very successful in the analysis of heavy-ion collisions even for nuclei that do not show any rotational character, such as ^{197}Au, or the isobars ^{96}Ru and ^{96}Zr. The case of ^{96}Zr is especially interesting: While calculations based on energy density functional theory currently struggle in giving reliable predictions for the ground-state deformation of this nucleus [122,123], data on elliptic and triangular flow in isobar collisions at RHIC appear to be consistent with the picture of an intrinsic rotor with a large octupole deformation parameter β_3 [68–71] extracted from Eq. (20), even though this is at odds with the transitional nature of this isotope. The effectiveness of such a type of modeling suggests that associating a nucleus with a rigid intrinsically deformed density rotating in space enables one to obtain azimuthal correlations consistent with the measured $\langle \hat{\mathcal{E}}_n \rangle$ values. Interestingly, the *picture* extracted from other pertinent [124] theoretical schemes is likely to be different even though the observables may be equally well reproduced. Still, it may be possible, i.e., in *ab initio* projected generator coordinate method calculations explicitly including shape fluctuations [74,77], to successfully map the results of a different theoretical framework back onto a rigid-rotor picture and extract an effective rigid deformation consistent with anisotropic flow measurements.

The above example illustrates how measurements of anisotropic flow in heavy-ion collisions, aimed at quantifying two-body correlations in nuclei, can help elucidate the inner mechanisms that generate such correlations in a given theoretical scheme. This is complementary to low-energy observables such as, e.g., quadrupole transition strengths from the ground state to the first 2^+ excited state, $B(E2; 0^+ \to 2^+)$, that are traditionally used to characterize collectivity in nuclei. To motivate future experimental endeavors, our ambition is to analyze the novel observables within the frame of *ab initio* expansion many-body methods for doubly magic, singly magic, and doubly open-shell nuclei where the mechanisms building up two-body correlations are expected to vary considerably. Particularly interesting candidates for such an endeavor are ^{16}O and ^{20}Ne, which are also relevant for current and upcoming ion runs at colliders [125–127]. Furthermore, the observables introduced in this Letter should be evaluated while performing a global sensitivity analysis of the parameters of the nuclear interaction [128–131], to rigorously quantify how anisotropic flow measurements at colliders constrain the low-energy constants of nuclear forces based on chiral effective field theory.

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