

# Imprints of octupole collectivity in uranium-238 on relativistic heavy-ion flow observables

Chunjian Zhang<sup>1,2,3,\*</sup> Jiangyong Jia<sup>3,4,†</sup> Jinhui Chen<sup>1,2,‡</sup> Chun Shen<sup>5,6,§</sup> and Lumeng Liu<sup>7,||</sup>

<sup>1</sup>Key Laboratory of Nuclear Physics and Ion-Beam Application (MOE), and Institute of Modern Physics, Fudan University, Shanghai 200433, China

<sup>2</sup>Shanghai Research Center for Theoretical Nuclear Physics, NSFC and Fudan University, Shanghai 200438, China

<sup>3</sup>Department of Chemistry, Stony Brook University, Stony Brook, New York 11794, USA

<sup>4</sup>Physics Department, Brookhaven National Laboratory, Upton, New York 11976, USA

<sup>5</sup>Department of Physics and Astronomy, Wayne State University, Detroit, Michigan 48201, USA

<sup>6</sup>RIKEN BNL Research Center, Brookhaven National Laboratory, Upton, New York 11973, USA

<sup>7</sup>Physics Department and Center for Particle Physics and Field Theory, Fudan University, Shanghai 200438, China



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Some atomic nuclei exhibit enhanced octupole collectivity, reflected in finite reflection-asymmetric multipole correlations rather than necessarily in a rigid static pear-shaped ground state. Low-energy studies indicate finite octupole strength in uranium-238, commonly interpreted as soft or vibrational in nature, in addition to its large prolate quadrupole collectivity [F. K. McGowan and W. T. Milner, *Nucl. Phys. A* **571**, 569 (1994).; T. Kibédi and R. H. Spear, *At. Data Nucl. Data Tables* **80**, 35 (2002)]. Here, we investigate how such octupole correlations can be encoded in the initial geometry of relativistic heavy-ion collisions and mapped to final-state flow observables. Using state-of-the-art hydrodynamic calculations, we demonstrate quantitative sensitivity to octupole-induced features encoded in the initial-state geometry and suggest a modest effective octupole-correlation strength in uranium-238 that is compatible with the latest high-energy experimental measurements [B. E. Aboona *et al.* (STAR), *Rep. Prog. Phys.* **88**, 108601 (2025)]. These findings provide a complementary probe of odd-order nuclear collectivity and help constrain quark-gluon plasma initial conditions.

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## I. INTRODUCTION

The atomic nucleus, as a strongly interacting many-body system, exhibits a variety of shapes from spherical to superdeformed [1,2]. Nuclear deformation has been traditionally inferred from low-energy spectroscopic measurements and models of reduced transition probability between low-lying rotational states, less than a few tens of MeV/nucleon beam energies [3]. More recently, it has been demonstrated that nuclear shapes can also be probed in high-energy nuclear collisions in yoctosecond-scale ( $10^{-24}$  s) resolution of nucleon spatial distributions by taking advantage of the strong responses of the hydrodynamics collective flow to the shape and size of the initial state [4,5]. This paradigm has been validated by measurements at Relativistic Heavy Ion Collider (RHIC) [6,7] and at the Large Hadron Collider (LHC) [8,9], aided by significant theoretical efforts [10–22].

In modeling high-energy collisions, the nuclear density is commonly parametrized by a deformed Woods-Saxon (WS) profile [23,24]

$$\begin{aligned}\rho(r, \theta) &\propto (1 + \exp[(r - R(\theta, \phi))/a])^{-1}, \\ R(\theta, \phi) &= R_0(1 + \beta_2[\cos \gamma Y_{2,0} + \sin \gamma Y_{2,2}] \\ &\quad + \beta_3 Y_{3,0} + \beta_4 Y_{4,0}),\end{aligned}\quad (1)$$

where  $R_0 \approx 1.2A^{1/3}$  denotes the nuclear radius and  $a$  is the surface diffuseness (skin depth). The nuclear surface  $R(\theta, \phi)$  is expanded in spherical harmonics  $Y_{l,m}(\theta, \phi)$  in the intrinsic frame, including only the most relevant axially symmetric quadrupole ( $\beta_2$ ), octupole ( $\beta_3$ ), and hexadecapole ( $\beta_4$ ) deformations, while triaxiality parameter  $\gamma$  ( $0^\circ \leq \gamma \leq 60^\circ$ ) characterizes deviations from axial symmetry [25,26]. In this work, the parameters  $\beta_n$  in Eq. (1) are used as effective amplitudes controlling multipole correlations in event-by-event coordinate-space sampling. In particular,  $\beta_3$ , an independent degree of freedom orthogonal to even order, should not be interpreted as evidence for a rigid static octupole deformation of the  $^{238}\text{U}$  ground state. Rather, it is a model-dependent effective amplitude parametrizing the rms strength of reflection-asymmetric octupole correlations that contributes to the initial geometry sampled by an ultrarelativistic collision and should not be equated with a unique low-energy nuclear-structure parameter.

Using the “imaging-by-smashing” method, the STAR experiment at RHIC has extracted effective shape parameters

\*Contact author: chunjianzhang@fudan.edu.cn

†Contact author: jiangyong.jia@stonybrook.edu

‡Contact author: chenjinhui@fudan.edu.cn

§Contact author: chunshen@wayne.edu

||Contact author: liulumeng@fudan.edu.cn

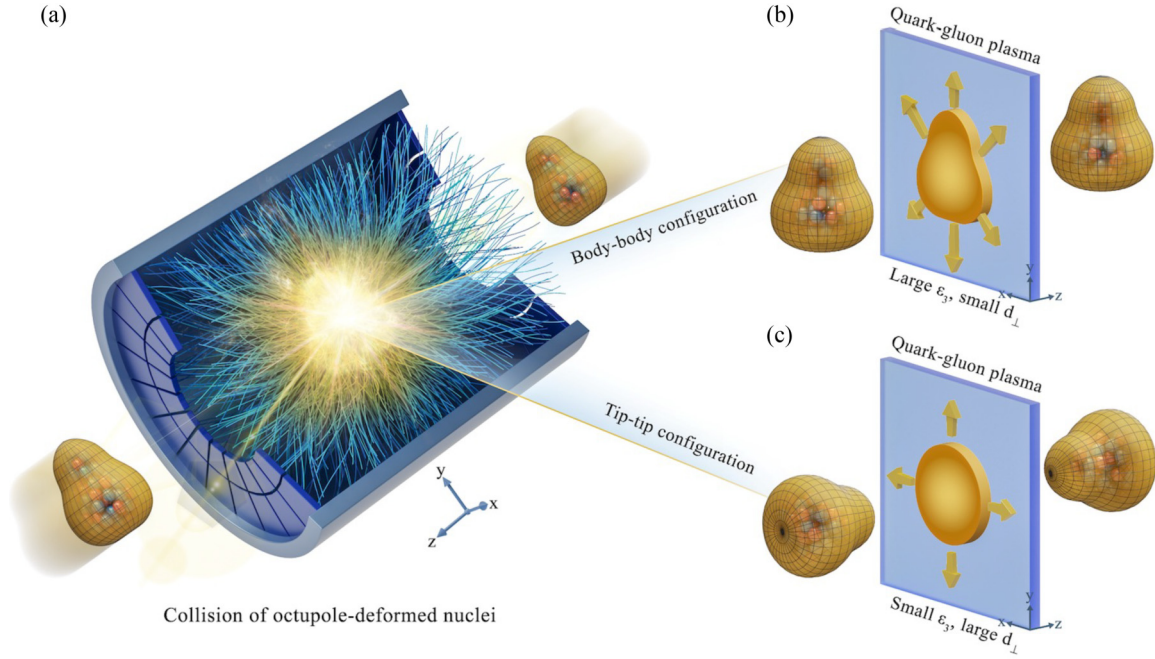


FIG. 1. Illustration of relativistic heavy-ion collisions with sampled octupole-correlated nuclear configurations. Cartoon of collisions involving nuclei whose event-by-event coordinate-space configurations contain reflection-asymmetric octupole correlations. The produced charged particle tracks are measured in the detector [panel (a)]. Head-on body-body [panel (b)] and tip-tip [panel (c)] configurations are shown together with the corresponding 3D profiles of the initial quark-gluon plasma, in which the arrows indicate the pressure gradients.

associated with the ground-state  $^{238}\text{U}$ , indicating a strongly prolate quadrupole deformation ( $\beta_2$ ) and a small but finite triaxiality ( $\gamma$ ) [7], consistent with low-energy constraints. Model studies further suggest sensitivity to hexadecapole deformation  $\beta_4$  in  $^{238}\text{U} + ^{238}\text{U}$  collisions [17,27,28]. Here,  $\beta_4$  can be correlated with the dominant  $\beta_2$  in nuclear structure calculations through multipole couplings in the deformation energy surface [3,29–37].

The octupole collectivity in the heavy  $^{238}\text{U}$  nucleus, away from the best known region of octupole deformed nuclei centered at  $Z \sim 90$ ,  $N \sim 136$ , remains less explored. Low-energy spectroscopic studies of  $^{238}\text{U}$  have long established finite octupole strength through  $E3$  transition observables. However, these data do not imply a rigid static octupole-deformed ground state at zero or low spin [38–46]. The octupole collectivity in  $^{238}\text{U}$  is more commonly interpreted as soft or vibrational in nature [2,43,44,47]. This octupole collectivity plays an important role in nuclear structure models [48], fission dynamics [49–51], with potential implications in searches for atomic electric dipole moments [43,44,47,52,53]. The question addressed here is different from establishing a static pear-shaped ground state. We ask whether the equal-time coordinate-space octupole correlations associated with this collectivity can leave measurable imprints on the initial geometry of relativistic heavy-ion collisions and subsequently on final-state flow observables.

## II. MAPPING OCTUPOLE CORRELATIONS TO TRIANGULAR FLOW

The “imaging-by-smashing” approach focuses on head-on (near-zero impact parameter) collisions of octupole-deformed

nuclei [Figs. 1(a)–1(c)]. Each incident nucleus is a Lorentz-contracted disk, generating a quark-gluon plasma (QGP)—a hot, dense state of deconfined quarks and gluons [54,55]. The shape and size of the QGP overlap region are determined by the transverse ( $xy$ ) plane projection of the colliding systems, which directly mirror the shape of those colliding nuclei encoded in the initial-state modeling. Body-body collisions generate a larger, triangular QGP profile, whose pressure-gradient-driven asymmetric expansion is illustrated by the arrows in Fig. 1(b). In contrast, tip-tip collisions yield a compact, circular QGP, leading to a faster but azimuthally symmetric expansion [Fig. 1(c)].

The sensitivity to nuclear deformation arises from the approximately linear mapping between the initial-state geometry and final-state anisotropic flow [7]. When deformed nuclei collide in random orientations, they imprint distinct spatial anisotropies on the initial energy density profile, characterized by eccentricity coefficients  $\varepsilon_{n>1} \equiv |\int r^n e^{in\phi} e(r, \phi) d^2r / \int r^n e(r, \phi) d^2r|$ , where the local energy density  $e(r, \phi)$  in the  $xy$  plane is derived from the overlap region, although its precise magnitude depends on the hydrodynamic model used. Driven by pressure gradients, the QGP undergoes hydrodynamic collective expansion, transferring the  $\varepsilon_n$  into azimuthal anisotropy of final-state hadrons, described by a Fourier expansion of the particle azimuthal angle distribution  $dN/d\phi \propto 1 + 2 \sum v_n \cos n\phi$  [56–59]. The elliptic ( $v_2$ ) and triangular ( $v_3$ ) flows linearly respond to the initial eccentricities  $\varepsilon_n$  [60],  $v_n \propto \varepsilon_n$ , making them sensitive to quadrupole and octupole collectivity.

Eventwise average transverse momentum ( $\langle p_T \rangle$ ) quantifies QGP compactness via the inverse area of the overlap  $d_{\perp} \propto 1/\sqrt{\langle x^2 \rangle \langle y^2 \rangle}$  [61] and drives radial flow, following a linear

relation,  $\delta p_T \propto \delta d_\perp$ , where  $\delta p_T = \langle p_T \rangle - \langle\langle p_T \rangle\rangle$  and  $\delta d_\perp = d_\perp - \langle d_\perp \rangle$  denote eventwise deviations from mean values [7,62]. Here,  $\langle\langle p_T \rangle\rangle$  represents the event-averaged transverse momentum in event ensemble.

In collisions of spherical nuclei,  $\varepsilon_2$  primarily reflects the impact-parameter-driven overlap ellipticity, while nonzero  $\varepsilon_3$  arises solely from random fluctuations of nucleons [63–65], leading to an anticorrelation  $\langle v_3^2 \rangle \sim 1/A$  in the absence of octupole deformation ( $\beta_3 = 0$ ). The presence of  $\beta_3$  and  $\beta_4$  enhances the variance of  $v_n$ , following a quadratic form [66] for  $n = 2, 3$ ,

$$\begin{aligned}\langle v_2^2 \rangle &\approx a_2 + b_{2,2}\beta_2^2 + b_{2,3}\beta_3^2 + b_{2,4}\beta_2\beta_4, \\ \langle v_3^2 \rangle &\approx a_3 + b_{3,3}\beta_3^2 + b_{3,4}\beta_4^2.\end{aligned}\quad (2)$$

$\beta_4$  can affect  $v_2$  by its nonlinear coupling with  $\beta_2$  [17], while  $v_3$  is not expected to be affected by  $\beta_2$  [66]. Here,  $a_n$  encodes the system size and centrality dependence for spherical nuclei, while the coefficients  $b_{n,m}$  are expected to be  $A$  independent. These relations enable a data-driven method to constrain the  $\beta_n$  parameters by comparing collisions of two species of similar sizes (see Refs. [10,66] for details).

In this article, we identify experimental signatures sensitive to octupole collectivity in  $^{238}\text{U}$  using state-of-the-art hydrodynamic calculations. We provide quantitative constraints on the magnitude of  $\beta_{3,\text{U}}$  through two novel observables, triangular flow and its correlation with transverse momentum, offering new insights into pear-shaped collectivity that can be contrasted with recent high-precision experimental measurements [67,68].

We emphasize that for a rotational invariant Hamiltonian, the one-body density of a  $J^\pi = 0^+$  ground state is spherical in the laboratory frame [69]. While intrinsic deformations are not directly observable, collective multipole correlations in even-even nuclei can be inferred from electric transition rates ( $B(E\ell)$ ), within theoretical models [17]. In our approach, the deformation parameters  $\beta_n$  serve as effective intrinsic-shape descriptors in event-by-event hydrodynamic modeling, characterizing the strength of multipole correlations rather than fixed laboratory-frame orientation or static symmetry breaking [23,24]. The observables studied here are multiparticle correlation measures, sensitive to event-by-event fluctuations and deformation collectivity encoded in the collision's initial state, for which orientation-interference effects decay exponentially with increasing mass number [70].

### III. HYDRODYNAMIC MODEL AND OBSERVABLES

To identify experimental signatures of octupole deformation in  $^{238}\text{U}$ , we simulate event-by-event (2+1)D boost-invariant  $^{238}\text{U} + ^{238}\text{U}$  collisions at  $\sqrt{s_{\text{NN}}} = 193$  GeV and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions at  $\sqrt{s_{\text{NN}}} = 200$  GeV employing the IP-Glasma + MUSIC + UrQMD framework. This model consists of an initial state IP-Glasma model [71,72], coupled with the hybrid framework of relativistic viscous hydrodynamics, MUSIC [73–75], followed by hadronic transport, UrQMD [76,77]. The energy density  $e(r, \phi)$  from the Yang-Mills evolution is used to initialize the hydrodynamical evolution. This model has been tuned to successfully repro-

duce collective flow in both small and large systems at RHIC and the LHC [78].

For  $^{197}\text{Au}$ , an odd- $A$  nucleus, we adopt  $\beta_{2,\text{Au}} = 0.14$  and triaxiality  $\gamma_{\text{Au}} = 45^\circ$ , guided by conservative low-energy estimates and model systematic comparisons with neighboring nuclei [79,80]. While its odd- $A$  character introduces polarization effects associated with the unpaired nucleon that complicate spectroscopic extraction of deformation parameters [81], the bulk geometry of heavy nuclei is dominated by collective mean-field properties [3,79,82].

To isolate the response to octupole correlations in  $^{238}\text{U}$ , we vary an effective octupole amplitude, denoted by  $\beta_{3,\text{U}}$  from 0 to 0.2 while fixing  $\beta_{2,\text{U}} = 0.28$  and  $\gamma_{\text{U}} = 0^\circ$  in Eq. (1) (Table I). Hexadecapole effects are systematically tested by scanning  $\beta_{4,\text{U}}$  between 0 and 0.09 [3,83,84]. The minimum nucleon distance in nuclei  $d_{\text{min}}$  is 0.9 fm in these two collision systems. For each configuration, 100k–400k high-statistics hydrodynamic events are generated, allowing the precise determination of shape differences between these two species. Each event is oversampled at least 100 times during the UrQMD stage to ensure enough resolution for anisotropic flow vectors and minimize statistical fluctuations in the hadronic transport.

$\beta_n$  parametrizes the intrinsic-shape component of the event-by-event initial geometry in deformed WS sampling. Mapping it to laboratory-frame spectroscopic observables requires additional nuclear-structure input, such as symmetry restoration or rotational model relations [17], which is beyond the scope of this article.

We study two-particle triangular flow,  $\langle v_3^2 \rangle$ , and its correlation with transverse momentum,  $\langle v_3^2 \delta p_T \rangle$ , defined as

$$\begin{aligned}\langle v_3^2 \rangle &= \langle \cos[3(\phi_i - \phi_j)] \rangle_{i,j}, \\ \langle v_3^2 \delta p_T \rangle &= \langle \cos[3(\phi_i - \phi_j)] (p_{T,k} - \langle\langle p_T \rangle\rangle) \rangle_{i,j,k}.\end{aligned}\quad (3)$$

Here,  $\phi_i$  and  $p_{T,i}$  denote the azimuthal angle and transverse momentum of particle  $i$ . The particle indices  $i, j$ , and  $k$  are chosen such that only unique combinations are used. The averages are first over unique particle multiplets in an event and then over all events in a fixed centrality class.  $\langle v_3^2 \rangle$  is obtained using the two-subevent method with pseudorapidity gap  $|\Delta\eta| > 1$  and the covariance  $\langle v_3^2 \delta p_T \rangle$  is calculated via the standard method. Results are verified across four different  $p_T$  intervals. The observables in Eq. (3) are calculated in isobarlike  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions with near  $A$ , from which we construct the ratios

$$R_{v_3^2} = \frac{\langle v_3^2 \rangle_{\text{U+U}}}{\langle v_3^2 \rangle_{\text{Au+Au}}} \approx \frac{a_{3\text{U}}}{a_{3\text{Au}}} + \frac{b_{3,3}}{a_{3\text{Au}}} \beta_3^2 + \frac{b_{3,4}}{a_{3\text{Au}}} \beta_4^2, \quad (4)$$

$$R_{v_3^2 \delta p_T} = \frac{\langle v_3^2 \delta p_T \rangle_{\text{U+U}}}{\langle v_3^2 \delta p_T \rangle_{\text{Au+Au}}} \approx a - b \beta_2 \beta_3^2. \quad (5)$$

These ratios suppress most system-dependent effects, including final-state and odd- $A$  specific effects, sensitive to the rms strength of octupole correlations rather than to the sign or fixed orientation of a static octupole moment. Their right-hand sides, linked to two- and three-body nucleon distributions in the intrinsic frame from Eq. (1), have a simple parametric dependence on shape parameters [66,85], with  $a$  and  $b$  in Eq. (5)

TABLE I. Woods-Saxon parameters used in IP-Glasma + MUSIC + UrQMD simulations of  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions.  $d_{\min}$  represents the nucleon minimum distance.

| Species           | $R$ (fm) | $a$ (fm) | $d_{\min}$ (fm) | $\beta_2$ | $\beta_3$                        | $\beta_4$ | $\gamma$ (°) |
|-------------------|----------|----------|-----------------|-----------|----------------------------------|-----------|--------------|
| $^{197}\text{Au}$ | 6.62     | 0.52     | 0.9             | 0.14      | 0                                | 0         | 45           |
| $^{238}\text{U}$  | 6.81     | 0.55     | 0.9             | 0.28      | 0.00, 0.05<br>0.10, 0.15<br>0.20 | 0.09, 0   | 0            |

being empirical positive coefficients varying with centrality [85]. The sensitivity of  $R_{v_3^2 \delta p_T}$  to octupole correlations requires the presence of large quadrupole component, as in  $^{238}\text{U}$ , making it a suitable system to probe  $\beta_3$ . This particular sensitivity would not appear in  $^{96}\text{Zr} + ^{96}\text{Zr}$  versus  $^{96}\text{Ru} + ^{96}\text{Ru}$  comparison to constrain the  $\beta_{3,\text{Zr}}$ , given the small values of  $\beta_{2,\text{Zr}}$ .

In Eqs. (4) and (5), higher-order deformations in  $^{197}\text{Au}$  are assumed negligible based on low-energy studies, together with the odd- $A$  polarization and higher-multipole assumptions noted above. The constraints quoted below should therefore be understood as constraints within the adopted  $^{238}\text{U} / ^{197}\text{Au}$  reference scheme. Otherwise, the extracted shape information corresponds to the difference between the two species, e.g., replacing  $\beta_3^2$  by  $\beta_{3,\text{U}}^2 - \beta_{3,\text{Au}}^2$ .

In the absence of deformation, the smaller  $^{197}\text{Au} + ^{197}\text{Au}$  system is expected to exhibit larger fluctuations than  $^{238}\text{U} + ^{238}\text{U}$ , yielding  $R_{v_3^2} \approx a_{3,\text{U}}/a_{3,\text{Au}} < 1$ , as well as  $R_{v_3^2 \delta p_T} \approx a < 1$ . The presence of  $\beta_3$  then increases the values of  $R_{v_3^2}$  in central collisions to exceed unity—an inversion of the hierarchy that provides a unique signature for  $\beta_3$ —and reduces  $R_{v_3^2 \delta p_T}$  further below unity, eventually inducing anticorrelation behavior at large  $\beta_3$  magnitudes. Meanwhile,  $\beta_4$  is expected to increase  $R_{v_3^2}$  in noncentral collisions [66], while its role in  $R_{v_3^2 \delta p_T}$  is unclear. Our findings are discussed in light of these mechanisms.

#### IV. RESULTS AND DISCUSSIONS

Three distinct features are identified: (1) For  $\beta_{3,\text{U}} \gtrsim 0.05$ , the ordering of  $\langle v_3^2 \rangle$  in ultracentral collisions (UCC, 0%–2% centrality) is reversed between  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$ , i.e.,  $R_{v_3^2} > 1$ . (2) In the 2%–20% centrality range,  $\langle v_3^2 \rangle$  is still strongly enhanced by  $\beta_3$ . (3)  $\beta_4$  increases the values of  $\langle v_3^2 \rangle$  in 2%–15% centrality, but has no impact in 0%–2% centrality, consistent with findings from the initial Glauber model [66]. It should be noted that a similar behavior has also been observed in a multiphase transport model calculations as cross-checks [86]. These trends follow the expectation of Eq. (4) well, making them robust features that can be verified experimentally. Such a hierarchical reversal feature in 0%–2% centrality can be used to set a reliable limit on  $\beta_3$ .

The result of Fig. 2(a) demonstrates that finite octupole-correlation strength in  $^{238}\text{U}$  enhances the event-by-event initial-state triangularity ( $\varepsilon_3$ ) beyond the fluctuation-driven baseline. The statistical uncertainties of current simulations, represented by the shaded bands, are small enough to be

sensitive to a modest  $\beta_{3,\text{U}}$ , when compared to experimental measurements [7].

Next, we focus on the UCC 0%–2% centrality range, where  $R_{v_3^2}$  is mainly sensitive to  $\beta_3$ . As shown in Fig. 2(b), the predicted ratios exhibit linear dependence on  $\beta_{3,\text{U}}^2$ . Assuming the magnitude of octupole strength  $\beta_{3,\text{U}} = 0.08$ –0.10 [38,39,41], deduced from low-energy nuclear structure, we predict  $R_{v_3^2} \approx 1.05$ –1.12 indicated by the arrows. This range is consistent with the current STAR measurement [67]. Note that experimental results falling within this range would provide a high-energy indication compatible with octupole collectivity in  $^{238}\text{U}$ . Future measurements with sub-percent-level precision, combined with additional sensitivity analyses, could further tighten constraints on  $\beta_3$ .

As a robustness test of the hydrodynamic response, we also perform constrained Skyrme Hartree-Fock-Bogoliubov (SHFB) calculations using the Skyrme-Hartree-Fock-Bogoliubov equations using the transformed harmonic oscillator basis (hfbtho) [87]. These calculations are not used to establish a static octupole-deformed ground state of  $^{238}\text{U}$ , nor do they describe octupole vibrations by themselves. Mean-field predictions of octupole minima in actinides, including  $^{238}\text{U}$ , are known to be functional and model dependent [88]. Instead, for prescribed multipole moments, the constrained SHFB densities provide self-consistent radial density profiles that can be used as alternative inputs to the IP-Glasma + MUSIC + UrQMD framework. The agreement between the deformed WS and constrained SHFB inputs demonstrates that the predicted flow response is not an artifact of the WS radial parametrization. The intrinsic axial multipole moments are extracted from the SHFB density  $\rho_{\text{WS}}(\vec{r})$  as  $Q_\lambda = \int \rho_{\text{WS}}(\vec{r}) r^\lambda Y_{\lambda 0}(\theta) d^3r$ , from which the deformation parameters  $\beta_\lambda^{\text{DFT}} = 4\pi Q_\lambda / (3N_\tau r^\lambda)$  are defined. The SLy4 interaction [89] is employed in the multiconstraint SHFB calculations. We find that WS and realistic SHFB densities lead to consistent results, indicating that WS parametrization is sufficient to capture the deformation features relevant for the  $^{238}\text{U}$  nucleus, as shown in Fig. 2(b).

To gain further insight into nuclear shape, we examine the three-particle correlator  $\langle v_3^2 \delta p_T \rangle$ . Body-body collisions create a small  $d_\perp$  with large  $\varepsilon_3$  QGP shape, while tip-tip collisions form a large  $d_\perp$  with minimal  $\varepsilon_3$  profile, as illustrated in Figs. 1(b) and 1(c), leading to a pronounced anticorrelation between  $d_\perp$  and  $\varepsilon_3$  and resulting in different pressure-gradient-driven expansion patterns. Figure 3(a) shows  $R_{v_3^2 \delta p_T}$  versus centrality for different  $\beta_{3,\text{U}}$ . In central collisions (0%–5% centrality),  $R_{v_3^2 \delta p_T}$  exhibits a near-linear dependence on  $\beta_{2,\text{U}} \beta_{3,\text{U}}^2$ , aligning with the parametric form

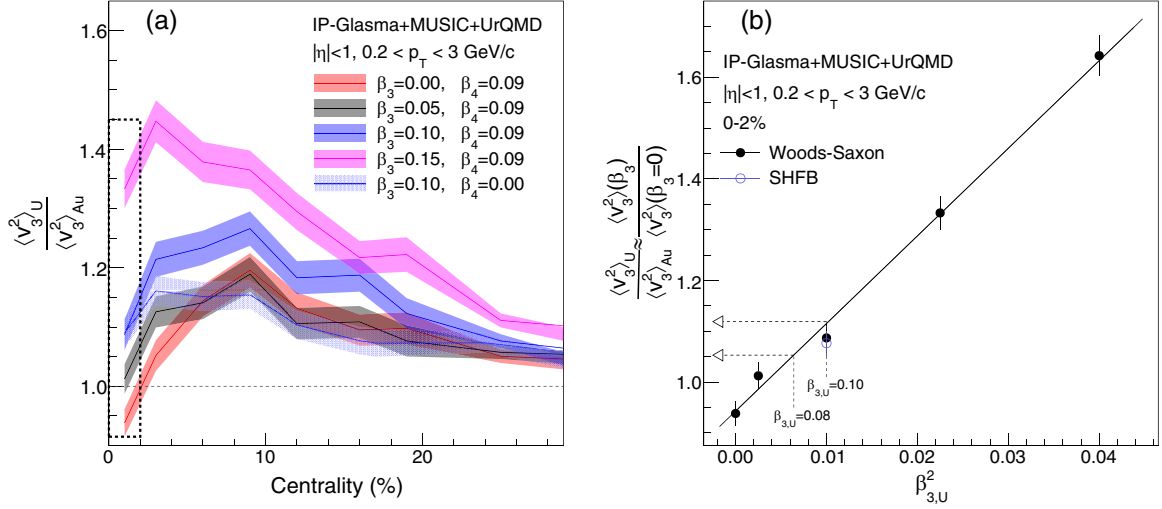


FIG. 2. (a) Ratios of  $\langle v_3^2 \rangle$  between  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions as a function of centrality for different  $\beta_3$  from IP-Glasma + MUSIC + UrQMD using WS densities. Shaded bands denote statistical uncertainties. The dashed box marked the 0%–2% centrality region that is most sensitive to  $\beta_3$ . (b) The ratio  $\langle v_3^2 \rangle_{\text{U}} / \langle v_3^2 \rangle_{\text{Au}}$  vs  $\beta_{3,\text{U}}^2$  in 0%–2% centrality, assuming  $\beta_{3,\text{Au}}^2 = 0$ . Arrows mark  $\beta_{3,\text{U}} = 0.08$  and  $0.10$  from low-energy studies [38,41], comparable with recent STAR measurements [67]. Results with SHFB density yield consistent trends.

predicted by Eq. (5), and provides complementary insights into  $\beta_{3,\text{U}}$  constraints. The observed suppression is qualitatively consistent with the expected negative contribution of  $\beta_{3,\text{U}}$  to  $\langle v_3^2 \delta p_T \rangle$ .

It is observed that  $R_{v_3^2 \delta p_T}$  also exhibits sensitivity to  $\beta_{4,\text{U}}$ . A pure  $\beta_{4,\text{U}} = 0.09$  has a negative contribution comparable to that from  $\beta_{3,\text{U}} = 0.10$  alone. However, when both  $\beta_{4,\text{U}}$  and  $\beta_{3,\text{U}}$  are present, their combined impact is larger. This suggests that terms like  $-\beta_2 \beta_3 \beta_4$  and  $-\beta_2 \beta_4^2$  should be added to Eq. (5) in noncentral collisions. These couplings imply that the signature of  $\beta_{3,\text{U}}$  and  $\beta_{4,\text{U}}$  requires larger  $\beta_{2,\text{U}}$ , making  $^{238}\text{U}$  an ideal nucleus. For moderate values of  $\beta_{3,\text{U}}$  and  $\beta_{4,\text{U}}$  ( $\beta_{3,\text{U}} \lesssim 0.1$ ,  $\beta_{4,\text{U}} \approx 0.05$ – $0.1$ ),  $R_{v_3^2 \delta p_T}$  falls in the range of 0.5–1 in central collisions, underscoring the need for combined

observables, such as  $R_{v_3^2}$  and  $R_{v_3^2 \delta p_T}$ , to constrain the correlated region of the  $\beta_3$  and  $\beta_4$  parameter space.

We validate the linear scaling of  $R_{v_3^2 \delta p_T}$  in 0%–5% centrality, shown in Fig. 3(b) as a function of  $\beta_{2,\text{U}} \beta_{3,\text{U}}^2$ . This linear scaling serves as another significant constraint on octupole deformation in  $^{238}\text{U}$  in high-energy collisions, compared to the recent STAR measurements [67]. Importantly,  $R_{v_3^2 \delta p_T}$  also exhibits consistent agreement between WS and SHFB density input.

Figure 4 shows the centrality dependence of  $R_{v_3^2}$  [panel (a)] and  $R_{v_3^2 \delta p_T}$  [panel (b)] calculated in four  $p_T$  intervals with fixed  $\beta_{3,\text{U}} = 0.10$ . No variations with the choice of  $p_T$  interval are observed. This implies that the final-state effects are almost completely canceled, leaving the model uncertainties

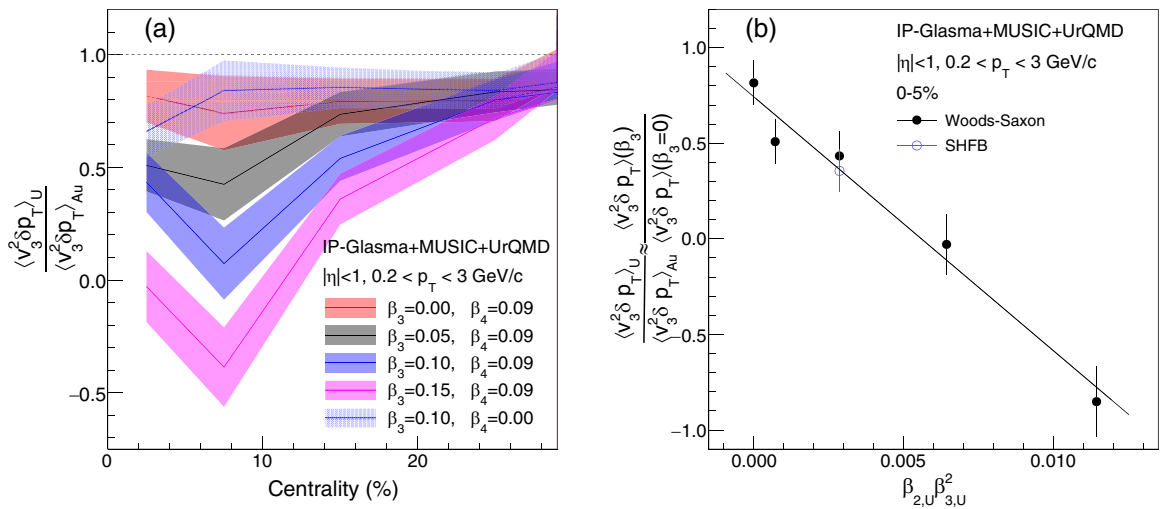


FIG. 3. (a) Ratios of  $\langle v_3^2 \delta p_T \rangle$  between  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions vs centrality for different  $\beta_{3,\text{U}}$  from IP-Glasma + MUSIC + UrQMD calculations. Shaded bands denote the statistical uncertainties. (b) Ratio  $\langle v_3^2 \delta p_T \rangle_{\text{U}} / \langle v_3^2 \delta p_T \rangle_{\text{Au}}$  vs  $\beta_{2,\text{U}} \beta_{3,\text{U}}^2$  for 0%–5% centrality. Results with SHFB density yield consistent trends.

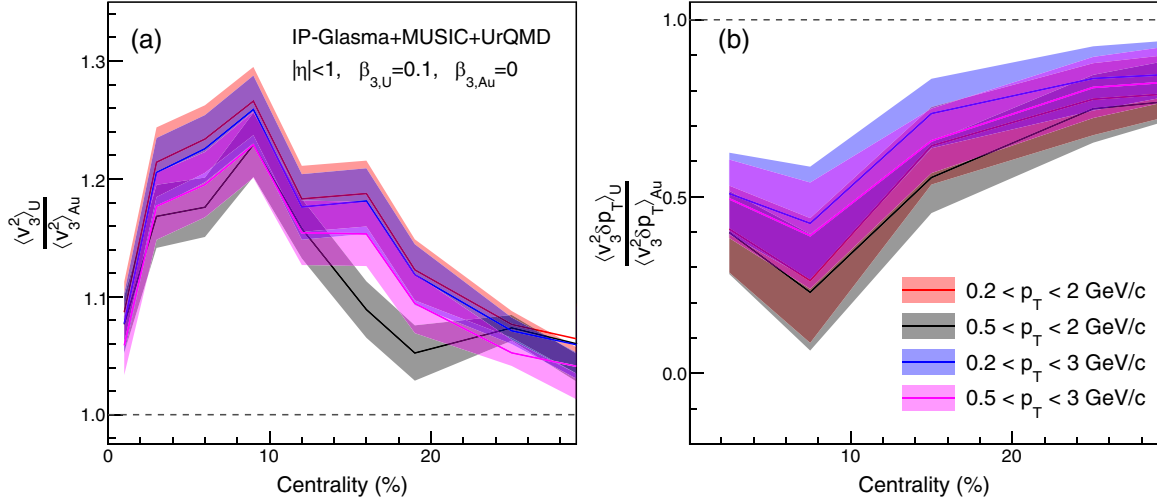


FIG. 4. Centrality dependence of ratios of  $\langle v_3^2 \rangle$  [panel (a)] and  $\langle v_3^2 \delta p_T \rangle$  [panel (b)] between  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  collisions from the IP-Glasma + MUSIC + UrQMD model for four  $p_T$  intervals. The shaded bands denote the statistical uncertainties.

mainly due to the initial conditions. The 20% difference in the mass numbers between  $^{197}\text{Au}$  and  $^{238}\text{U}$  does not lead to any appreciable difference in the hydrodynamic response. The stability of the ratio against varying  $p_T$  intervals is crucial to demonstrate that the impact of nuclear deformation can be measured without the need to impose stringent  $p_T$  selections.

## V. SUMMARY

We have investigated how effective octupole correlations in  $^{238}\text{U}$  can be encoded in the initial geometry of relativistic heavy-ion collisions and manifested in final-state flow observables. The parameter  $\beta_3$  used here represents an effective rms octupole-correlation amplitude in event-by-event initial-state modeling, rather than a rigid static octupole deformation of the  $^{238}\text{U}$  ground state. By constructing ratios of the same observable between collisions of  $^{238}\text{U} + ^{238}\text{U}$  and  $^{197}\text{Au} + ^{197}\text{Au}$  within IP-Glasma + MUSIC + UrQMD framework, we identify quantitative signatures associated with octupole collectivity characterized by  $\beta_{3,U}$ . In particular, a modest  $\beta_{3,U}$  leads to a characteristic reverse ordering of  $\langle v_3^2 \rangle$  ratio in UCC 0%–2% centrality, while hexadecapole collectivity parametrized by  $\beta_{4,U}$  primarily affects noncentral collisions. The recent STAR measurements are compatible with a finite effective octupole-correlation strength,  $\beta_{3,U} \sim 0.08$ –0.1 [67], providing a valuable high-energy interpretation. We further show that the ratio of  $\langle v_3^2 \delta p_T \rangle$  exhibits complementary sensitivity to both  $\beta_{3,U}$  and  $\beta_{4,U}$  in a centrality-dependent manner.  $\langle v_3^2 \rangle$  increases linearly with  $\beta_3^2$ , while  $\langle v_3^2 \delta p_T \rangle$  exhibits a previously unexplored suppression, scaling linearly with  $\beta_2 \beta_3^2$ . Importantly, we confirm that WS and SHFB densities yield consistent results in heavy nuclei.

Taken together, the combined analysis of these ratios enables systematic constraints on higher-order collectivity, including an upper limit on  $\beta_{4,U}$ , once  $\beta_{3,U}$  is determined. Their ratios are insensitive to  $p_T$  cuts, demonstrating robustness against kinematic selections. Overall, these observables can

be combined with Bayesian analysis to further refine quantitative constraints.

By providing new constraints on initial-state modeling, this “imaging-by-smashing” approach contributes to reducing uncertainties in the description of QGP initial conditions, and their subsequent dynamics and transport properties [4,90,91]. Moreover, these findings provide a complementary probe of nonzero octupole collectivity in heavy nuclei [44,45], complementing traditional nuclear-structure studies. While low-energy measurements indicate finite octupole collectivity in  $^{238}\text{U}$  [38,39], the interpretation in terms of static rigid or soft vibrational modes remains model dependent and under debate [92–94]. Future high-energy measurements of multi-particle flow fluctuations, such as  $\langle v_3\{4\}/v_3\{2\} \rangle^4$ , as explored in Refs. [95,96], provide a complementary and experimentally accessible probe of the sensitivity to these scenarios.

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## DATA AVAILABILITY

The data that support the findings of this article are available from the corresponding author on request.

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