

Dissecting Jet Modification in the Quark-Gluon Plasma with Multipoint Energy Correlators

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Energy correlators have recently attracted significant attention in the study of heavy ion collisions due to their potential to robustly connect experimental measurements with an underlying quantum field theoretic description. While theoretical studies have so far primarily focused on the simplest two-point correlator, mapping out the dynamics of the quark-gluon plasma (QGP) will require developing a theoretical understanding of multipoint energy correlators. In this Letter, we present a systematic theoretical study of multipoint energy correlators for jets fragmenting in a dense quark-gluon plasma, accounting for both the medium's perturbative modification to the jet and its hydrodynamical backreaction. We consider both the scaling behavior of projected correlators, as well as the shape dependent three-point correlator, highlighting how both provide insight into interactions with the QGP. We discuss the parametric dependence of modifications on the medium scales, opening new opportunities to experimentally separate jet modifications from the medium response. Our results open the door to a systematic exploration of multipoint energy correlators in heavy ion collisions.

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Introduction—A primary goal of high-energy nuclear physics is the study of ordinary matter under extreme conditions to search for emergent phenomena in the standard model. On Earth, high-energy heavy ion collisions (HICs) are the prime experimental avenue to explore the finite temperature and density properties of QCD, allowing for the production of droplets of the quark-gluon plasma (QGP). Deciphering the underlying physics of the QGP from HICs requires the combination of different soft and hard probes to construct a tomographic picture of the emergent QCD matter, see, e.g., Ref. [1] for a recent overview. Although a nearly comprehensive qualitative understanding of these events has been achieved, a significant knowledge gap remains in connecting various experimental findings to a first-principles description from the underlying theory of QCD.

A natural observable for achieving a direct connection between experimental measurements and QFT is matrix

elements [2–5] of the average null energy (ANE) operator [6–9]:

$$\mathcal{E}(\mathbf{n}) = \lim_{r \rightarrow \infty} r^2 \int_0^\infty dt n^i T^{0i}(\mathbf{n}), \quad (1)$$

which experimentally corresponds to an idealized calorimeter cell on a sphere of radius r along the direction set by the normal vector $\mathbf{n}$. Because of their operator definition in terms of the stress tensor of the underlying theory, one can hope to use symmetries [10,11] or the operator product expansion [12] to make robust statements about their behavior in a nuclear environment.

Energy correlators were introduced as jet substructure observables in [2–5,8,13–15] and were first applied as an approach to study the QGP in [16]. They have since been measured in proton-proton [15,17,18], proton-nucleus [19], and nucleus-nucleus [20] collisions, and clear nuclear modification has been observed. This has prompted significant theoretical efforts to understand these observables, see, e.g., Refs. [12,16,21–29].

In the HIC context, most of the discussion has been focused on the simplest two-point energy correlator. Much in analogy with cosmology, while the two-point correlator is able to detect the presence of a scale, higher point

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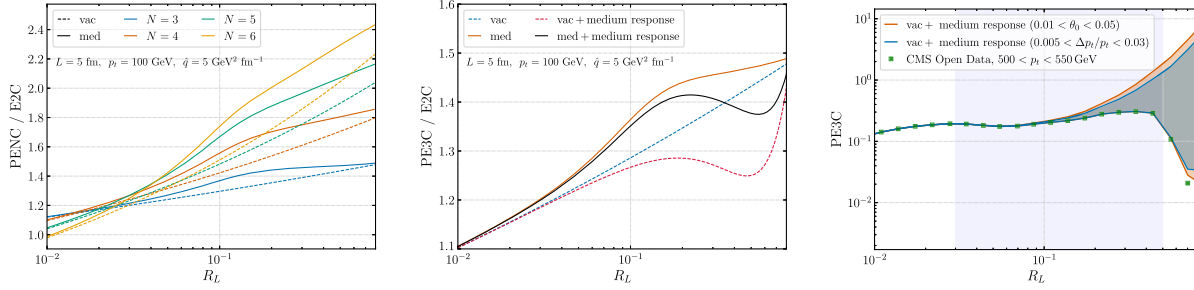


FIG. 1. Left: PENC/E2C ratios for the vacuum (dashed) and in-medium (solid). Center: PE3C/E2C including the medium response. Right: PE3C extracted from CMS Open Data presented in [15], with a theory band for the medium response, obtained by convoluting with the E2/3C data. The θ_0 band has $\Delta/p_t = 0.02$ while the Δ/p_t band has $\theta_0 = 0.025$. The purple shaded area indicates the region where the perturbative calculation of the EEC is valid; the points for $R_L > 0.5$ lie outside the cone of the jet where the PE3C is extracted.

correlators are able to probe the underlying dynamics. Higher point correlators will, therefore, be crucial in accessing different underlying properties of the QGP, including the geometric characteristics of the medium response and the transport scales responsible for perturbative modifications to the jet cascade. Ultimately, we envision a program, much in analogy with cosmology, of mapping out the shapes of higher point correlators in different phases of QCD matter produced in HICs. In vacuum, the three- [30] and four-point [31] correlator have been computed, analyzed [10], and measured [32]. An initial exploration of the three-point correlator in the QGP was presented in [28], highlighting its utility for mapping out the shape of the wake of jets. A related discussion of energy loss effects on higher-point correlators was presented in [33]. Further expanding our understanding of jet substructure in the QGP will require developing a theoretical understanding of multipoint correlators and their modification by the QCD matter produced in HICs.

While it may seem that a robust understanding of multipoint energy correlators in HICs is currently beyond reach, two recent developments provide optimism. First, [12] initiated a discussion of the light-ray operator product expansion (OPE) to the study of energy correlators in nuclear environments, showing on general symmetry grounds that medium modification could imprint itself with well defined scaling laws into the energy correlators. Second, [34] showed how the energy correlator observables could be understood in hydrodynamical states, also leading to robust scaling laws. Combined, these two approaches provide the means to analyze both perturbative and hydrodynamic medium response of multipoint energy correlators.

In this Letter, we present the first theoretical calculation of higher-point energy correlators in the QGP. We consider both the full shape-dependent three-point correlator (E3C), as well as the projected multi-point correlators (PENCs) [14] (To facilitate comparison with the notation commonly used in the phenomenological literature and adopted in [28], our full E3C and projected PE3C correspond to their EEEEC and E3C, respectively; more generally, our PENC corresponds to their ENC. We use the prefix (P) to

distinguish projected correlators explicitly). Accounting for both the medium induced perturbative modifications to the jet cascade and the hydrodynamically driven back-reaction, we show how the QGP imprints itself onto the (P) ENCs' structure set by the symmetries of the $\mathcal{E}$ operators in the OPE limit. We argue that this mapping depends very weakly on the particular assumptions made in the present calculation and is underlined by universal features. Although such conclusions are natural from the perturbative side [12], where we derive a closed form for the OPE coefficients, we show a nontrivial correspondence between the OPE limit and the medium response's structure, dictated by purely geometrical correlations [34]. This promotes the geometric medium response contribution from a two-point effect to a broader organizing principle, whose small-angle expansion parallels the perturbative twist structure and encodes genuinely higher point information. Furthermore, the medium response and the perturbative contributions have distinct parametric dependencies on the medium's scales, allowing for a future first theory driven approach to disentangle them.

Projected energy correlators—We begin by considering the simpler case of the projected energy correlators [14], which are obtained by integrating out the shape information, keeping only the information on the largest angle. The PENCs exhibit a scaling, which is sensitive to the twist-2 anomalous dimensions in the theory.

In the collinear limit and at leading accuracy, the PENC quark distribution, $d\Sigma_P^{(N)}$, reads at fixed coupling [14],

$$\frac{d\Sigma_P^{(N)}}{dR_L} = \frac{\alpha_s R_L^{\frac{\alpha_s}{2} \gamma(N+1)}}{\pi R_L} \int_0^1 dz [1 - z^N - (1 - z)^N] P_{gq}, \quad (2)$$

with the anomalous dimension $\gamma(j) \equiv -\int_0^1 dz z^{j-1} \hat{P}_{gq}(z)$ of spin $j = N + 1$, and $\hat{P}_{gq}$ the regularized splitting function. The PENC/E2C ratio gives direct access to the perturbative anomalous scalings in QCD, which we illustrate in the top of Fig. 1 in the dashed lines for $N = 3-6$.

In the conformal limit of QCD, the structure of the correlators is set by the symmetries of $\mathcal{E}$, allowing for an

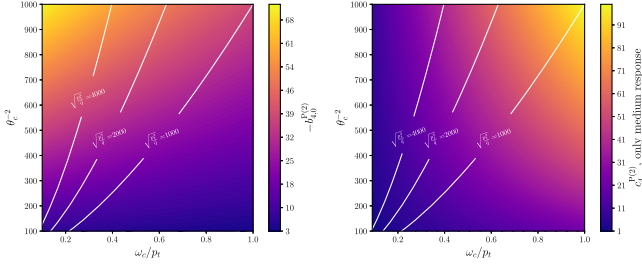


FIG. 2. OPE limit coefficients for the medium perturbative corrections, $b_{4,0}^{P(2)}$, and the response, $c_4^{P(2)}$, neglecting hard correlations. White lines have $p_t^3/\hat{q}$ constant.

operator product expansion [9]. In particular, this imposes that the PENC can be expanded about the collinear limit as

$$\frac{d\Sigma_P^{(N)}}{dR_L} = \sum_{k=1} a_{\tau,0}^{P(N)} R_L^{\tau-3} \Big|_{\tau=2k}. \quad (3)$$

Note that Eq. (2) satisfies this form at leading power.

We first discuss how the QGP imprints onto Eqs. (3) and (9), focusing on perturbative modifications to fragmentation and the medium backreaction. We assume a static and homogeneous QGP, indicating when these assumptions become relevant.

At leading order, the perturbative corrections can be absorbed into a modification of the splitting function entering Eq. (2). One can write $P_{gq} = P_{gq}^{\text{vac}} + P_{gq}^{\text{med}}$, where the medium induced splitting functions P^{med} have been widely discussed in the literature [35–39] and are detailed in Supplemental Material [40]. Importantly, one can show that, for a static and homogeneous medium, $P^{\text{med}} = P^{\text{med}}(\omega_c, \theta_c)$, where $\omega_c \propto \hat{q}L^2$ is the critical gluon frequency below which the QCD Landau-Pomeranchuk-Migdal (LPM) effect is prevalent and $\theta_c^{-2} \propto \hat{q}L^3$ is the respective characteristic angle, with L being the medium longitudinal size and $\hat{q}$ being the jet quenching transport coefficient. Using the explicit form of P^{med} , the PENC distribution in the OPE limit, with the (vacuum) leading-twist contribution subtracted, reads

$$\frac{d\Sigma_{P,\text{med(no vac)}}^{(N)}}{dR_L} = \sum_{k=2} b_{\tau,0}^{P(N)}(\omega_c/p_t, \theta_c) R_L^{\tau-3} \Big|_{\tau=2k}, \quad (4)$$

where the coefficients $b_{\tau,0}^{P(N)}(\omega_c/p_t, \theta_c)$ are given in Supplemental Material [40], depending only on two parameters, with p_t being the jet energy. As an illustration, the leading OPE coefficient, $b_{4,0}$, is shown in Fig. 2, as a function of (θ, ω_c) ; note that $a_{\tau,0}$ would be constant in a similar plot. The parametric dependence in Eq. (4) becomes more complex in an evolving medium but can still be derived [41,42]. Importantly, although the explicit parametric dependence of the coefficients in Eq. (4) was

realized for a specific model (BDMPS-Z [43] in a static, structureless medium), the resulting collinear twist expansion is general since it follows from the OPE structure of energy correlators. Thus, higher order perturbative corrections are expected to modify the coefficients but not the symmetry-based twist organization. Critically, following typical jet quenching theoretical considerations, all odd twist coefficients vanish since CP-odd interactions are neglected and the medium is structureless. Measurements yielding nonvanishing odd coefficients would then give access to nonperturbative features of the jet-medium interactions and the spatial structure of the QGP, not yet addressed by any other jet observable. Further discussion on these aspects is left to a future publication. In the left of Fig. 1, we show the perturbative medium contributions to the leading PENC ratios in solid lines. Compared to the vacuum, there is an enhancement increasing with N at large R_L , reflecting the region where $b_{\tau,0} > a_{\tau,0}$.

In turn, the imprint of the jet-induced medium response can be computed for any PENC. This follows from the fact that the response's ANE is purely classical, at least in all the present theoretical descriptions of the response, and, thus, the respective PENCs are dominated by geometrical correlations, as discussed in detail for the E2C in [34]. Restricting, for simplicity, the discussion to the PE3C/E2C ratio, we split the ANE operator $\mathcal{E}(\mathbf{n}) = \hat{\mathcal{E}}_h(\mathbf{n}) + \mathcal{E}_c(\mathbf{n})$. The first term represents the fluctuating part of the flow operator associated to the hard (h) jet fragmentation, just discussed above, while the second term accounts for the classical (c) flux, including the medium response. The latter contribution is intended to capture the mean (classical) contribution [44] of the medium response to the jet, which is usually modeled with linearized hydrodynamics or kinetic theory. The corresponding contributions to the PE3C (ccc , chh , and hhh) can be expressed through the leading perturbative correlators of the h component combined with the c piece. Note that, since, by definition, $\langle \hat{\mathcal{E}}_h \rangle = 0$, terms like cch do not contribute. Here, we do not attempt to constrain particular details of the classical energy flux in an event with a jet, which would go beyond the current theoretical control but rather rely on a two-parameter model, $\mathcal{E}_c(\mathbf{n}) = (\Delta/\pi\theta_0^2)e^{-\theta^2/\theta_0^2}$, where θ_0 controls the angular size of the classical energy flux and Δ accounts for the amount of the energy that is propagating in the medium. It should be stressed that the exact functional form of $\mathcal{E}_c$ does not qualitatively affect the leading collinear behavior of energy correlators. This is rather determined by the magnitude of the energy propagating in the matter, controlled by Δ , and the respective characteristic angular region θ_0 . Using $(d\Sigma_{P,\text{vac}}^{(2,3)}/dR_L) \simeq a_{2,0}^{P(2,3)} R_L^{\gamma(3,4)-1}$, where $3a_{2,0}^{P(2)} = 2a_{2,0}^{P(3)}$ at the leading order, we find that the leading twist vacuum result convoluted with the classical flux reads

$$\left. \frac{\text{PE3C}}{\text{E2C}} \right|_{\text{vac+response}} = \left[\frac{a_{2,0}^{P(3)}}{R_L^{1-\gamma(4)}} + \frac{a_{2,0}^{P(3)} \Delta}{6\pi p_t} C R_L^{1+\gamma(3)} + \frac{\Delta^3}{\theta_0^2 p_t^3} \left(\frac{2}{9} - \frac{\sqrt{3}}{6\pi} \right) R_L^3 \right] \times \left(\frac{a_{2,0}^{P(2)}}{R_L^{1-\gamma(3)}} + \frac{\Delta^2 R_L}{p_t^2 \theta_0^2} \right)^{-1}, \quad (5)$$

where $\mathcal{C}$ corresponds to the chh terms and contains a collinear divergence numerically regulated at small angles, dominated by hadronization effects in two-point correlators, and, for simplicity, we assume that only the direction of the classical flux varies from event to event, while other fluctuations can be straightforwardly accounted for; for further details on this computation, see Supplemental Material [40]. In the center of Fig. 1, we evaluate Eq. (5) for the vacuum result convoluted with the classical response (red) and compare it to the purely vacuum result (blue), the vacuum result including the perturbative medium correction from Eq. (4) (orange), and the combined result with both the perturbative correction and the classical flux (black). We use $\theta_0 = 0.025$, $p_t = 100$ GeV, $\theta_{\text{cut}} = \Lambda_{\text{QCD}}/p_t = 0.002$, $\Delta = 0.01 p_t$, $a_{2,0}^{P(3)} = 0.286$, and $\alpha_s = 0.3$. Including the response leads to a characteristic imprint in the ratio's shape, with a depletion at large angles, not observed when considering only the perturbative result. This is in qualitative agreement with recent Hybrid model results [28,46–48]. To better gauge the phenomenological impact of these contributions, in the right of Fig. 1, we combine the medium response with the CMS data analysis in [15], resulting in a large angle enhancement of the PE3C, comparable to the perturbative counterparts. We reiterate that the shape of PE2/3C depends weakly on the functional form of the classical flux, while being sensitive to the parameters taken. Similarly to the perturbative contributions, we provide the leading response coefficient c_4 in Fig. 2 as a function of (θ, ω_c) , displaying the different scaling of the response with respect to the perturbative terms. Following the estimate of [34], we assume that the leading response coefficient scales as $c_4 \propto \hat{q} L^4 T^2/p_t$, or equivalently $c_4 \propto (\omega_c/p_t) \theta_c^{-4/3}$ for $T \propto \hat{q}^{1/3}$.

In general, the E2C associated with the medium response, in the collinear limit, takes the following form:

$$\frac{d\sigma_{\text{P, response}}^{(2)}}{dR_L} = \sum_{k=2} c_\tau^{P(2)} \left[\mathcal{E}_c \right] R_L^{\tau-3} \Big|_{\tau=2k}, \quad (6)$$

where the coefficients are functionals of the classical energy flux, given by angular averages of its powers and derivatives. Remarkably, the expansion in Eq. (6), independent of the OPE for the ANE, has the same structure as Eqs. (3) and (4). More, the leading coefficients in the expansion for the

response's PE3C also follow the OPE limit structure [see numerator of Eq. (5)]; coefficients again only depend on $\mathcal{E}_c$. These relations, i.e., Eqs. (5) and (6), hold for any N for the leading τ terms, and we conjecture they hold for all τ , assuming the classical energy flux to be cylindrically symmetric, ignoring azimuthal fluctuations. However, we notice that such asymmetries may appear due to the medium evolution and anisotropy both in the in-medium splitting function and medium response, which can be especially important in small HIC systems [25,49–53]. These lead to finite even contributions to Eq. (6)—see Supplemental Material [40]—suggesting prospective avenues to disentangle the in-medium physics from vacuum QCD and providing a tool to access the azimuthal jet substructure.

Finally, all these observations explicitly show that the perturbative terms and the medium response have direct competing contributions to the PENCs, requiring a systematic analysis to be disentangled. Equations (4) and (6) provide the full parametric dependence of the coefficients in the OPE limit, indicating the path for a theoretically driven decomposition of these effects in experiment.

Shape dependent three-point correlator—We now consider the full shape dependent three-point correlator. In the collinear limit, it can be obtained by integrating the $1 \rightarrow 3$ splitting functions [10,30],

$$J_3 = \int d\Phi_3 \frac{g^4 P_{3z_1 z_2 z_3}}{2s_{123}^2} \Big|_{\sum_i z_i=1}^{s_{123}=\sum_{i \neq j} s_{ij}}, \quad (7)$$

where $(4\pi)^5 d\Phi_3 = ds_{12} ds_{13} ds_{23} \prod_i dz_i 4(-D)^{-1/2}$ denotes the three-body phase space [54], $z_i \equiv E_i/p_t$ is the energy fraction carried by the final state particle $i = 1, 2, 3$, and flavor dependence is implicit. The momentum variables read $s_{ij} = z_i z_j p_t^2 \theta_{ij}^2$, while $D = (z_3 s_{12} - z_1 s_{23} - z_2 s_{13})^2 - 4z_1 z_2 s_{13} s_{23} < 0$. We shall restrict the discussion to the quark channel, where the vacuum splitting kernels are $P_{2,\text{vac}} = P_{gq}$, $P_{3,\text{vac}} = C_F^2 P_{3,\text{vac}}^{\text{Abel}} + C_F C_A P_{3,\text{vac}}^{\text{non-Abel}}$ [55], with the Abelian and non-Abelian terms given in Supplemental Material [40].

Since our theoretical description of $P_{3,\text{med}}$ in HICs is still limited, [56] in this work we also consider a factorized *cascade* approximation [57], where the $1 \rightarrow 3$ process is given by a succession of $1 \rightarrow 2$ branchings:

$$\frac{P_{3,\text{vac}}^{\text{cascade}}}{s_{123}} = \frac{P_{gq}^{2,1+3} P_{gq}^{1,3}}{s_{13}} + (1 \leftrightarrow 2) + \frac{P_{gq}^{1+2,3} P_{gq}^{1,2}}{s_{12}}, \quad (8)$$

where we used the spin averaged $1 \rightarrow 2$ splitting functions and employed the notation, e.g., $P_{gq}^{1,2}(x) \equiv C_F [1 + (1-x)^2]/x$, and $P_{gg}^{1,2} = 2C_A \{ [x/(1-x)] + [(1-x)/x] + x(1-x) \}$, with $x = [E_1/(E_1 + E_2)]$. The mismatch between the exact result and the cascade approximation, at the E3C level, is mild. We verified the cascade form does not interfere with the medium effects discussed below,

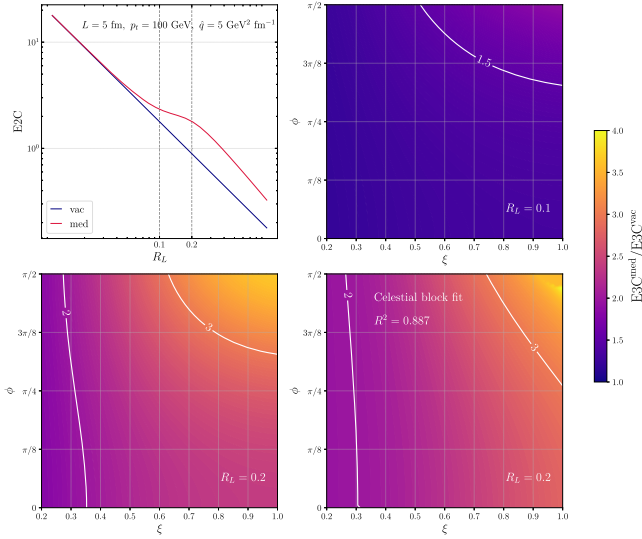


FIG. 3. E2C distribution (top left) indicating the R_L values used for the E3C/E2C ratios in medium (top right, bottom left). Fit to the E3C/E2C ratio with $R_L = 0.2$ (bottom right), following the celestial block structure.

which are more prominent and emerge in different sectors in the (ξ, ϕ) plane, where $\xi = (\theta_{13}/\theta_{23})$, $\sin^2\phi = 1 - [(R_L - \theta_{23})^2/\theta_{13}^2]$, with $0 < 2\phi < \pi$ and $0 < \xi < 1$ in the region $R_L = \theta_{12} > \theta_{23} > \theta_{13}$. The deviations between the cascade and exact solutions are mainly constrained to regions where the three detectors are either in flattened or squeezed configurations, with a $\mathcal{O}(30\%-50\%)$ magnitude; further details on this comparison can be found in Supplemental Material [40]. It will be interesting to extend this calculation using the full medium modified $1 \rightarrow 3$ splitting functions [57].

The leading perturbative medium induced correction to J_3 is obtained straightforwardly, generalizing the cascade form in Eq. (8) to the medium. The ratio of the in-medium and the vacuum E3C distributions is shown in Fig. 3, for two values of R_L chosen by examining the respective E2C, as illustrated. The medium effects lead to a characteristic enhancement in the region where $R_L \approx \theta_{13} \approx \theta_{23}$. At larger R_L , the enhancement extends to lower values of ϕ at large ξ . The cascade approximation only mildly affects this sector.

The three-point energy correlator can be expanded in celestial blocks [10,11,30]:

$$\frac{R_L \theta_{23} \theta_{13} d\Sigma^{(3)}}{\sqrt{-D} dR_L d\theta_{13} d\theta_{23}} = \sum_{k=1} a_{\tau,0}^{(3)} G_{\tau,0}^q(z, \bar{z}) \Big|_{\tau=2k}, \quad (9)$$

where the complex variables $z, \bar{z}$ are defined as $R_L^2 z \bar{z} = \theta_{13}^2$, $R_L^2 (1-z)(1-\bar{z}) = \theta_{23}^2$, while the single-valued $G_{\tau,0}^q$ can be written in terms of hypergeometric functions, whose explicit form can be found in [10]. We fitted the data for $R_L = 0.2$ to the functional form $[\sum_{i=1}^3 n_i G_{2i,0}(\phi, \xi) / \sum_{i=1}^3 d_i G_{2i,0}(\phi, \xi)]$, following the structure in Eq. (9).

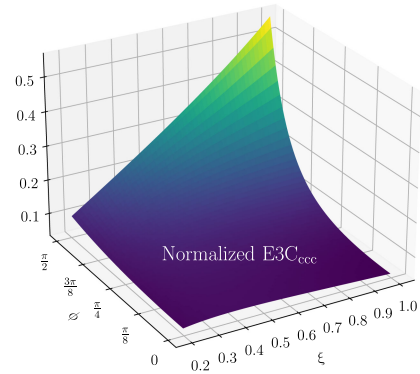


FIG. 4. Normalized E3C for the ccc term from hydrodynamic response.

This is shown in the bottom right plot of Fig. 3. We obtained a good fit to data, with $R^2 = 0.887$, while we find that the numerator's coefficients, for $i > 1$, are approximately an order of magnitude larger than the denominator's ones. This points to the dominance of higher-twist terms in medium, as expected. It will be interesting to explore this in more detail in different models.

The E3C's extraction including the full response is far from trivial; nonetheless, insight into the distribution's shape can be gained from the ccc term:

$$\frac{d\Sigma_{ccc}^{(3)}}{dR_L d\xi d\phi} = \frac{4\pi^2 p_t^{-2} R_L^3 \xi \int d\theta_1 \mathcal{E}^3(\cos\theta_1) \sin\theta_1}{(1 + \xi \cos\phi)^3 \sqrt{4 + 4\xi \cos\phi - \xi^2 \sin^2\phi}},$$

shown in Fig. 4. The distributions' maximizer is nearly identical to the one for the perturbative terms in Fig. 3, further evidencing the overlap of these two effects, as for the PENCs. Furthermore, this result closely matches the numerical result in [28]; here, it is solely determined by geometrical correlations. Finally, writing the ccc term in $(z, \bar{z})$ coordinates, $(d\Sigma_{ccc}^{(3)}/dR_L d\text{Re}z d\text{Im}z) = (2\pi^2 R_L^3/p_t^2) [\int d\theta_1 \mathcal{E}^3(\cos\theta_1) \sin\theta_1]$, and decomposing it in the celestial blocks, we find in the (collinear) limit $z \rightarrow 0$ at fixed $\bar{z}$ that the only contribution comes from the leading spin and $\tau = 4$ block, demonstrating the overlap with higher terms in the twist expansion.

Conclusion—In this Letter, we have presented the first theory study of the collinear limit of higher-point energy correlators in HICs, including both the perturbative modifications to the jet and the medium response. The main result is the development of a multipoint energy correlator framework in the QGP, based on the universal power structure derived in the OPE limit, and a procedure to systematically separate different medium-induced effects in jets. For the PENCs, we have demonstrated the universal structure,

$$\frac{d\Sigma_P^{(N)}}{dR_L} = \sum_{\tau} \left\{ a_{\tau,0}^{P(N)} R_L^{-\delta_{\tau,2}\gamma(N+1)} + b_{\tau,0}^{P(N)} + c_{\tau}^{P(N)} \right\} R_L^{\tau-3},$$

for even τ , with $b_{2,0}^{P(N)} = c_2^{P(N)} = 0$, assuming an azimuthal trivial background and jet substructure; odd τ terms can be associated with nonperturbative effects or nontrivial QGP profiles. This extends the classical energy-flow interpretation of the hydrodynamically driven medium response identified in [34] with PENCs. Further, we have separated the parametric dependence of the different coefficients, permitting a first theory driven identification of the PENC structure. In Fig. 2, we schematically show how to distinguish these elements, which exhibit distinct evolutions along (white) lines of constant $\hat{q}/p_t^3$. Experimental proxies for such curves would allow us to map the PENCs in the QGP; see also, e.g., Ref. [58], a necessary and critical step for a clear interpretation of the experimental data.

At the level of the E3C, we have analytically shown that the perturbative and the medium response contributions overlap, as for the E2C case, and cannot be directly distinguished, requiring a treatment analogous to the PENCs. More, we have explained the nature and shape of the medium response imprint on the E3C, previously only observed in numerical simulations [28], clarifying their physical origin. We find strong evidences that the E3C follows the pattern set by the celestial block decomposition, with the QGP perturbatively imprinting itself on the higher-twist terms, with the response following the same structure. This further exposes the universal character of energy correlator observables in the QGP, complementing the recent discussions on the properties of these objects, e.g., [59–63].

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Data availability—There are no publicly available research data or software supporting this manuscript.

Requests for further information or data should be sent to the authors.

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