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Azimuthal harmonics of color fields in a high energy nucleus

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ABSTRACT

Recent experimental results have revealed a surprisingly rich structure of multiparticle azimuthal correlations in high energy proton–nucleus collisions. Final state collective effects can be responsible for many of the observed effects, but it has recently been argued that a part of these correlations are present already in the wavefunctions of the colliding particles. We evaluate the momentum space 2-particle cumulant azimuthal anisotropy coefficients $v_n\{2\}$, $n = 2, 3, 4$ from fundamental representation Wilson line distributions describing the high energy nucleus. These would correspond to the flow coefficients in very forward proton–nucleus scattering. We find significant differences between Wilson lines from the MV model and from JIMWLK evolution. The magnitude and qualitative transverse momentum dependence of the $v_n\{2\}$ values suggest that the fluctuations present in the initial fields are a significant contribution to the observed anisotropies.

1. Introduction

One of the most surprising results from the LHC proton–nucleus collision experiments has been the kind of azimuthal multiparticle correlation structures [1–7] (see also RHIC results from deuteron-gold collisions [8,9]) that have, in larger collision systems, been attributed to hydrodynamical flow. The particle multiplicities in these collision systems are large enough for some collective effects to take place. Many of these structures have indeed been successfully described by hydrodynamical calculations [10,11]. This agreement requires, however, a very specific model of the geometry of the initial state [12]. It is also not clear whether these small systems are within the regime of validity of a hydrodynamical description with realistic values of the energy density, viscosity and system size [13].

The primary collisions leading to energy deposition in the central rapidity region are, at the high energies reached at the LHC, characterized by very strong nonlinear color fields [14]. These fields are, to leading order in the coupling constant, boost invariant. This immediately leads to the presence of long range azimuthal correlations in particle production [15–22]. In larger collision systems, the structure of these correlations in azimuthal angle and transverse momentum is strongly influenced by collective behavior in the later evolution stages of the system. However, in smaller systems, such as proton–nucleus collisions, these collective effects are presumably less significant than in nucleus–nucleus collisions. This raises the intriguing possibility that in proton–nucleus collisions also the azimuthal structure of the initial stage color fluctuations could be directly visible in the measurable particle spectrum. We will here argue that at least they need to be considered as an initial contribution for further collective effects when analyzing correlations in small systems.

We do not yet have a very solid quantitative understanding of the relative importance of initial color field and later evolution effects for generating anisotropies in particle production. A complete calculation of azimuthal anisotropies in this context requires complicated modeling that includes the color field and nucleonic scale fluctuations in the nucleus [23] and in the proton [24], combined with a calculation of the time evolution of the initial color fields and eventual matching to a hydrodynamical description [25]. We will not attempt to carry out this whole program here, but concentrate in this paper only on a part of it, namely the anisotropies produced when a bunch of valence quark-like particles in the fundamental representation of the gauge group scatter off the color field of a large nucleus. The physical picture of particle production (see [20,21] and more recently [26–29]) in our calculation is that of valence quarks from the probe deflected in a preferred transverse direction by a domain in the target color field. This generates a multiparticle correlation that probes the spatial fluctuations of the target. Our calculation extends the work in [28,29] in two significant ways. Firstly, we perform the Fourier-transform

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from coordinate to momentum space, in order to get an azimuthal harmonic coefficient corresponding to particles with a specific transverse momentum. Secondly, unlike [28,29], we correlate the particles in a given p_T-bin with a reaction plane determined by all the produced particles using the 2-particle cumulant method.

2. Azimuthal correlations in CGC fields

In the "hybrid formalism" for particle production in the dilute-dense limit [30–34] the quark spectrum produced in a collision is proportional to the Fourier-transform of the two point function of Wilson lines in the color field of the target

$$\frac{\mathrm{d}N}{\mathrm{d}^2\mathbf{p}_T} \propto \int\limits_{\mathbf{x}_T,\mathbf{y}_T} e^{-\mathbf{p}_T\cdot(\mathbf{x}_T-\mathbf{y}_T)} \frac{1}{N_\mathrm{c}} \operatorname{Tr} V^\dagger_{\mathbf{x}_T} V_{\mathbf{y}_T}. \tag{1}$$

The Wilson lines $V(\mathbf{x}_T)$ in (1) are, in the Color Glass Condensate (CGC) description, stochastic random SU(3) matrices in the representation of the projectile. To calculate the single inclusive cross section one has to average Eq. (1) by the appropriate probability distribution of Wilson lines.

In the leading order CGC treatment that we use here, multiparticle correlations can be calculated from the higher order moments of the Wilson line operators in Eq. (1). When the correlation is calculated from the so-called "glasma graphs" [18,19,22], the target color field domain structure is built into the k_T-dependent unintegrated gluon distributions describing the colliding particles. Note that the intrinsic k_T in these distributions results from a Fourier-transform of the coordinate dependence of the fields. The same physics of color field domains appears very clearly in the hybrid formalism. The target nucleus is represented by a sheet of color magnetic and color electric fields, which have a characteristic length scale $1/Q_\mathrm{s}$ in the transverse plane. When a small enough probe (comparable in size to the domain size) hits this target, the resulting particle production has a preferred direction given by the direction of the color field in the domain. Since this direction fluctuates from event to event, there is of course no anisotropy on average, but the existence of a preferred direction in individual events shows up in a global angular correlation among all of the produced particles, similarly to hydrodynamical flow. We are neglecting here "connected" or "BFKL"-like correlations [22], that give rise to a back-to-back peak in the two-particle correlation. These correlations are typical "nonflow" correlations that involve only a few particles, which the experimental analyses of azimuthal anisotropy try to exclude. We will not discuss them further here, see however Refs. [22,35] for more studies on these lines.

It is evident from the above discussion that we expect the correlation to be very sensitive to the transverse size of the probe. In the case of calculating the initial condition for an ion–ion collision the probe is large, with the consequence that the correlation is washed away by the sum over many independent domains in the transverse plane. Thus, in contrast to the correlations generated by collective flow, the effect discussed here is stronger in small collision systems than in large ones.

The purpose of this paper is to analyze the azimuthal correlation structure of particle production using Eq. (1) in more detail. In particular, we want to study its dependence on the harmonic n, transverse momentum, and the transverse size of the probe. The practical procedure used here is the following. We first divide the p_T range accessible on the lattice into bins. We use here 50 bins, but we have checked that the results are independent of the size of the bin. We then define the Fourier coefficient of the single particle spectrum as

$$b_n(p_T) \equiv \int\limits_{|\mathbf{p}_T|\in\mathrm{bin}} \mathrm{d}^2\mathbf{p}_T e^{in\varphi_{\mathbf{p}_T}} \int\limits_{\mathbf{x}_T,\mathbf{y}_T} e^{-i\mathbf{p}_T\cdot(\mathbf{x}_T-\mathbf{y}_T)} \times S_p(\mathbf{x}_T-\mathbf{b}_T) S_p(\mathbf{y}_T-\mathbf{b}_T) \frac{1}{N_\mathrm{c}} \operatorname{Tr} V^\dagger_{\mathbf{x}_T} V_{\mathbf{y}_T}. \tag{2}$$

The transverse coordinate profile of the probe has been taken as a Gaussian

$$S_p(\mathbf{x}_T-\mathbf{b}_T) = \exp\left\{\frac{-(\mathbf{x}_T-\mathbf{b}_T)^2}{2B}\right\} \tag{3}$$

around an impact parameter $\mathbf{b}_T$ chosen randomly in the transverse plane of the target. The product of the two Gaussian profiles in Eq. (2) could be interpreted as the Wigner distribution for a quark localized in an area $\sim B$ in the transverse plane, Fourier-transformed into a function of two coordinates $\mathbf{x}_T$ and $\mathbf{y}_T$. We will present results for different values of the parameter B characterizing the size of the probe. Note that the coefficients (2) need not be normalized, since we will eventually divide by the angular average spectrum b_0 to construct the Fourier harmonic coefficient. We want to calculate the angular correlations with respect to an event plane defined by all the produced particles, which form the "reference" that we correlate individual particles with. This is done following the procedure used in the experimental analysis (see e.g. the 2-particle cumulant method in [4]). For this we need to calculate also the reference coefficients

$$b_n(\mathrm{ref}) \equiv \int \mathrm{d}^2\mathbf{p}_T e^{in\varphi_{\mathbf{p}_T}} \int\limits_{\mathbf{x}_T,\mathbf{y}_T} e^{-i\mathbf{p}_T\cdot(\mathbf{x}_T-\mathbf{y}_T)} \times S_p(\mathbf{x}_T-\mathbf{b}_T) S_p(\mathbf{y}_T-\mathbf{b}_T) \frac{1}{N_\mathrm{c}} \operatorname{Tr} V^\dagger_{\mathbf{x}_T} V_{\mathbf{y}_T} \tag{4}$$

integrated over all momenta.

The target Wilson lines are drawn from a completely homogeneous and isotropic distribution that fills the whole transverse lattice with periodic boundary conditions, and the probe is azimuthally symmetric. Thus there is no geometrical (i.e. originating in the shape of the probe or the target) origin for azimuthal anisotropy present in the calculation. Since the probability distribution of Wilson lines is azimuthally symmetric (although the individual configurations are not), the correlations among the coefficients b_n are diagonal:

$$\left\langle b_n^*(p_T) b_m(q_T)\right\rangle \propto \delta_{m,n}, \tag{5}$$

where $\langle\rangle$ denotes averaging over the configurations of Wilson lines in the target. Note that the single particle spectrum Eq. (1) is explicitly real, configuration by configuration, leading to $b_n = b^*_{-n}$. This can be shown by taking the complex conjugate of Eq. (1) and exchanging the integration variables $\mathbf{x}_T$ and $\mathbf{y}_T$. The two particle pair correlation function is now

$$\frac{\mathrm{d}N_\mathrm{pair}}{\mathrm{d}\Delta\varphi} \propto \sum_{n=-\infty}^{\infty} \left\langle b_n^*(p_T) b_n(q_T)\right\rangle \cos(n\Delta\varphi). \tag{6}$$

From this we can identify the correlation function Fourier coefficients (using the notation of [4])

$$V_{n\Delta}(p_T,q_T) = \frac{\left\langle b_n^*(p_T) b_n(q_T)\right\rangle}{\left\langle b_0^*(p_T) b_0(q_T)\right\rangle}, \tag{7}$$

and define the 2-particle cumulant azimuthal harmonic as in [4] as

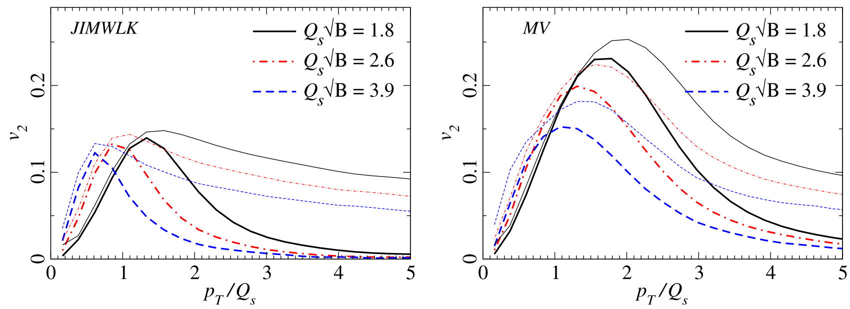

Fig. 1. Second harmonic coefficient $v_2\{2\}$ calculated with JIMWLK-evolved (left) and MV-model (right) Wilson line configurations. The thin lines represent the coefficients $v_2\{\text{bp}\}$ (see Eq. (11)) calculated with respect to the event plane in the p_T bin only.

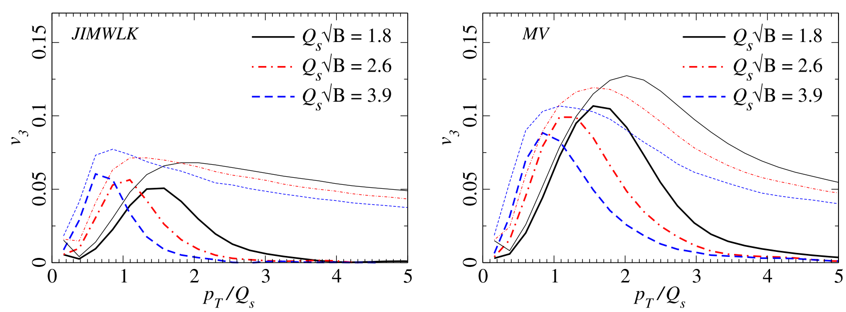

Fig. 2. Third harmonic coefficient $v_3\{2\}$ calculated with JIMWLK-evolved (left) and MV-model (right) Wilson line configurations. The thin lines represent the coefficients $v_3\{\text{bp}\}$ (see Eq. (11)) calculated with respect to the event plane in the p_T bin only.

$$v_n\{2\} = \frac{V_{n\Delta}(p_T, \text{ref})}{\sqrt{V_{n\Delta}(\text{ref}, \text{ref})}} \tag{8}$$

$$= \frac{\frac{\langle b_n^*(p_T) b_n(\text{ref})\rangle}{\langle b_0^*(p_T) b_0(\text{ref})\rangle}}{\sqrt{\frac{\langle b_n^*(\text{ref}) b_n(\text{ref})\rangle}{\langle b_0^*(\text{ref}) b_0(\text{ref})\rangle}}}. \tag{9}$$

A nice interpretation of Eq. (9) can be obtained by writing it as a product of three terms,

$$v_n\{2\} = v_n\{\text{bp}\} \frac{R_n(p_T, \text{ref})}{R_0(p_T, \text{ref})}. \tag{10}$$

Here we denote by

$$v_n\{\text{bp}\}^2 = \frac{\langle b_n^*(p_T) b_n(p_T)\rangle}{\langle b_0^*(p_T) b_0(p_T)\rangle} \tag{11}$$

the flow coefficient for particles in the p_T bin with respect to the event plane of that p_T bin ("bp" stands for "bin plane"). This is the equivalent (although here in momentum, not position space) of the quantity calculated in [28]. Note also that the ALICE analysis [1] correlates particles within a p_T bin when determining the flow coefficient, similarly to the $v_n\{\text{bp}\}$ defined here. The "bin plane" flow coefficient is then corrected by two "correlation coefficients". The first one is the correlation coefficient between the reference reaction plane and the p_T-bin reaction plane:

$$R_n(p_T, \text{ref}) \equiv \frac{\langle b_n^*(p_T) b_n(\text{ref})\rangle}{\sqrt{\langle b_n^*(p_T) b_n(p_T)\rangle \langle b_n^*(\text{ref}) b_n(\text{ref})\rangle}} \leq 1, \tag{12}$$

where the inequality follows from the Schwartz inequality. The interpretation of this correction is clear: for a fixed anisotropy with respect to the p_T-bin reaction plane, a decorrelation of the p_T-bin reaction plane from the reference reaction plane decreases the flow coefficient $v_n\{2\}$. The other correlation coefficient factor in (10)

$$R_0(p_T, \text{ref}) \equiv \frac{\langle b_0^*(p_T) b_0(\text{ref})\rangle}{\sqrt{\langle b_0^*(p_T) b_0(p_T)\rangle \langle b_0^*(\text{ref}) b_0(\text{ref})\rangle}} \leq 1 \tag{13}$$

is related to the multiplicity and appears in the denominator, increasing $v_n\{2\}$. This can be understood as follows: with larger fluctuations in the p_T-bin multiplicity that are independent of the reference multiplicity, a fixed correlation between $b_n(p_T)$ and $b_n(\text{ref})$ implies a larger correlation between p_T-bin and reference reaction planes. In other words, since $b_n \sim v_n b_0$, for a given correlation between (ref and p_T) b_n's, the smaller the correlation between b_0's, the larger must the correlation between v_n's be.

We take the Wilson lines $V(\mathbf{x}_T)$ appearing in Eq. (1) either from the McLerran–Venugopalan [36–38] (MV) model or resulting from JIMWLK evolution of the distribution of Wilson lines. Both are discretized on a 1024^2 transverse lattice. For the MV model we use a (fundamental representation) saturation scale of $Q_s a = 0.119$, where a is the lattice spacing. The JIMWLK calculation starts with an MV model at $Q_s a = 0.0220$ and, after $y = 10$ units of evolution in rapidity (with running coupling) ends up with $Q_s a = 0.117$. The MV model Wilson lines are constructed following the procedure described in more detail in [39] and the running coupling JIMWLK evolution performed using the algorithm of [40]. The parameter values used here are exactly the same as for the 1024^2-lattice in [41]. Note that we are only averaging two-point functions of the coefficients b_n, not ratios $\sim b_n/b_0$. This makes the averaging procedure numerically quite stable and is physically the correct thing to do, since the pair correlation function Eq. (6) is the correct inclusive observable to be obtained via the target average [42,17].

3. Results and discussion

The numerical evaluations of the first four anisotropy coefficients are shown in Figs. 1, 2 and 3. The results are presented in scaling units, as v_n plotted against p_T/Q_s for different size probes, i.e. different $\sqrt{B}Q_s$. Also shown are the "bin plane" coefficients $v_n\{\text{bp}\}$, defined by Eq. (11). To set the scale of the parameters in physical units we note that the fundamental representation saturation scale around midrapidity at the LHC should be [43] around $Q_s \sim 1$ GeV and the typical size of a proton in hard particle pro-

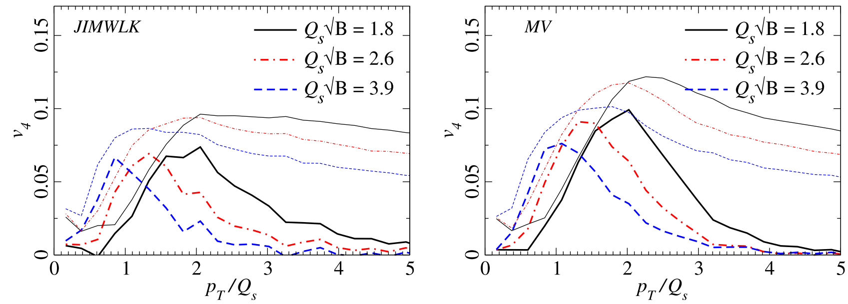

Fig. 3. Fourth harmonic coefficient $v_4\{2\}$ calculated with JIMWLK-evolved (left) and MV-model (right) Wilson line configurations. The thin lines represent the coefficients $v_4\{\text{bp}\}$ (see Eq. (11)) calculated with respect to the event plane in the p_T bin only.

duction at small x around $B \approx 4\ \text{GeV}^{-2}$ [44]. Thus a realistic probe size for LHC pA collisions would very roughly be $\sqrt{B}Q_s \approx 2$.

The first immediate observation from the numerical results is that the color field fluctuations indeed generate anisotropies that are large, of the order of the experimentally measured anisotropy coefficients. It seems therefore plausible that the color field fluctuations do play a sizeable role in the observed anisotropy in small systems, and must be taken into account together with the flow contribution. Also the momentum distribution has the same structure as the observed transverse momentum dependence of the flow, first rising until $\sim Q_s$ and then decreasing. The "bin plane" coefficients $v_n\{\text{bp}\}$ do not decrease nearly as fast at high momentum, from which one can deduce that the decrease of the anisotropy coefficients at large p_T follows from the decorrelation of the event plane in the p_T bin from the reference. This explains why this decrease was not seen in [28], where this decorrelation was not taken into account. The MV model has a gluon spectrum that is more sharply peaked around Q_s, i.e. a narrower distribution of different size color field domains. This shows up in significantly larger values for the v_n coefficients. The main effect of JIWMLK evolution is to add more small color field domains (larger p_T gluons), which decrease the anisotropy of the particle spectrum.

There is, however, an important caveat concerning any direct comparison of these results to experimental values. Namely, we were considering, in Eq. (1), only incoming quarks. For antiquarks one must replace the Wilson line by its Hermitian conjugate, which changes the sign of b_n for odd n. Away from the very forward valence region in the probe, there are an approximately equal amount of quarks and antiquarks present, with contributions to v_3 that therefore cancel. Gluons do not have nonzero odd harmonics in this mechanism, because the adjoint representation is real and thus odd b_n's vanish. Any odd harmonic surviving in the final state around midrapidity must therefore have an origin that is different from the one discussed here.

The other word of caution in interpreting these results is related to the dependence on the size of the probe, parametrized here by the width of the Gaussian $\sqrt{B}$. As anticipated, the magnitude of the correlation, and its dependence on the transverse momentum, depends strongly on the size of the interaction region. Although one can quite well estimate this, it depends on nonperturbative physics in the proton and cannot ultimately be controlled in a weak coupling calculation.

Results for azimuthal correlations in a full Classical Yang–Mills (CYM) simulation have also been presented recently by Schenke, Schlichting and Venugopalan [45,46]. Their calculation includes effects of both color field and nucleonic fluctuations in the probe proton and the target nucleus. The probe and target geometries also have a significant effect through the CYM pre-equilibrium version of the usual hydrodynamical mechanism that converts spatial anisotropy to momentum space, leading to also odd harmonics. These geometrical effects have not been included in our work, which should therefore not be compared directly with experimental data. Our focus here has been, instead, on quantifying the generic observation that fluctuating color fields result in azimuthal anisotropies in multiparticle correlations, even in the absence of anisotropies in the impact parameter dependence.

As a conclusion, we have here studied the momentum space azimuthal anisotropy structure of the "color glass" gluon fields in a high energy nucleus, as they are seen by a small probe consisting of valence-like quarks. We also quantified here the effect of correlating the particles with the event plane determined by all the produced particles, using the two-particle cumulant method at the parton level. In our calculation also high p_T-particles exhibit strong azimuthal correlations, but with respect to an event plane that becomes decorrelated from the lower p_T bulk. Clear experimental indications of this decorrelation have not been reported in the p_T range studied so far (see e.g. [4]). The quantitative results strongly depend on the details of the p_T-distribution of gluons in the CGC wavefunction and on the transverse size of the probe. However, all the results show large contributions to the harmonics from these purely initial state effects. For odd harmonics they largely cancel between quarks and antiquarks, but for even harmonics these are sizeable effects that need to be considered when interpreting the experimental results from proton–nucleus collisions.

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