DECORRELATION OF FORWARD DIJETS IN DIS OFF NUCLEI

V.R. Zoller
ITEP, Moscow, Russia

Abstract

Based on the color dipole QCD approach we discuss the multiple scattering mechanism (MSM) of suppression of back-to-back azimuthal correlations of high $p_T$ forward dijets in DIS off nuclei. We quantify the effect in terms of acoplanarity/decorrelation momentum of jets. For hard dijets the decorrelation momentum is found to be of the order of the nuclear saturation momentum $Q_A$. Minijets with the transverse momentum below the saturation scale are proved to be completely decorrelated. The analysis of DIS indicates that in heavy ion collisions the contribution of MSM to the reactions dynamics can be of the order of magnitude of the decorrelation effect observed at RHIC.

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1 Introduction

Recent observation of gradual disappearance of azimuthal back-to-back correlations of high $p_T$ particles with centrality of Au-Au collisions at RHIC is presently viewed as a consequence of the jet energy loss in hot quark-gluon plasma produced in central collisions. Another possible explanation is that initial and final state interactions associated with multiple parton scatterings on hard stage of the process produce uncorrelated monojets.

To get an idea of the size of the multiple scatterings (MS) effect at RHIC we start with the QCD description of the breakup of photons into forward dijets in small-$x$ deep inelastic scattering (DIS) off nuclei in the saturation regime. In we reported a derivation of the general formula for the two-body transverse momentum distribution. Our formalism, based on the technique, includes consistently the diffractive attenuation of color dipoles and effects of transitions between different color states of the $q\bar{q}$-pair propagating through the nucleus. Opacity of nuclei brings in a new scale $Q_A$ which separates the regimes of opaque nuclei and weak attenuation. For hard dijets the decorrelation momentum is found to be of the order of the nuclear saturation momentum $Q_A$. For parton momenta below the saturation scale $Q_A$ the evolution of sea from gluons was shown to be dominated by the anti-collinear, anti-DGLAP splitting. As a result, minijets with the transverse momentum below the saturation scale are proved to be completely decorrelated.

Turning back to the RHIC observation we argue that MS mechanism may contribute substantially to the effect of vanishing of back-to-back azimuthal correlations of high $p_T$ hadrons in central nucleus-nucleus collisions.

2 Breakup of photons into hard dijets on nuclear targets

We consider DIS at $x \lesssim x_A = 1/R_A m_N \ll 1$ which is dominated by interactions of $q\bar{q}$ Fock states of the photon and make use of the conventional approximation of two t-channel gluons in DIS off free nucleons. The two-gluon exchange approximation amounts to neglecting diffractive DIS off free nucleons which is justified by a small fraction of diffractive DIS, $\eta_D \ll 1$. We sum unitarity cuts of the forward Compton scattering amplitude which describe the transition from the color-neutral $q\bar{q}$ dipole to the color-octet $q\bar{q}$ pair. It should be emphasized that we are interested in the unitarity cuts which correspond to the genuine inelastic DIS with color excitation of the target nucleus.

For $x \lesssim x_A$ the propagation of the $q\bar{q}$ pair inside nucleus can be treated in the straight-path approximation. Let $b_+$ and $b_-$ be the impact parameters of the quark and antiquark, respectively, and $S_A(b_+, b_-)$ be the S-matrix for
interaction of the $qq$ pair with the nucleus. Regarding the color states $|c\rangle = |1\rangle, |8\rangle$ of the $qq$ pair, we sum over all octet and singlet states. Then the 2-body density matrix enter four straight-path trajectories $\delta$ fraction $z$ of excitation of the nucleon $gN$ inclusive spectrum is calculated in terms of the 2-body density matrix as

\[ \frac{d\sigma_{in}}{dzd^2p_+d^2p_-} = \frac{1}{(2\pi)^4} \int d^2b'_+d^2b'_-d^2b_+d^2b_- \times \exp[-ip_+(b_+-b'_+)-ip_-(b_-b'_-)]\Psi^*\Psi \times \{ \sum_{A} \sum_c \langle 1; A|S^*_A|A; c\rangle \langle c; A|^S_A|A; 1 \rangle \}

- \langle (1; A|S^*_A|A; 1) (1; A|S_A|A; 1) \rangle \} . \tag{1} \]

In (1) $\Psi = \Psi(Q^2, z, b_+ - b_-)$ stands for the wave function of the $qq$ Fock state of the photon with virtuality $Q^2$ and the photon light-cone momentum fraction $z$ carried by the quark. Notice, that the calculation of the 2-body density matrix enter four straight-path trajectories $\delta^{(2)}(b'_+), \delta^{(2)}(b'_-)$ and $S_A$ and $S'^*_A$ describe the propagation of two quark-antiquarks pairs, $qq$ and $q'q'$, inside a nucleus. In the integrand of (1) we subtracted the coherent diffractive component of the final state.

Upon the application of closure to sum over nuclear final states $A'$ the integrand of (1) can be considered as an intranuclear evolution operator for the 2-body density matrix (for the related discussion see refs. [12], [13])

\[ \sum_{A'} \sum_c \langle A| \{ \langle 1|S^*_A|c\rangle \langle A'|c\rangle \{ \langle c|S_A|1\rangle \} |A \rangle = \langle A| \{ \sum_c \langle 1|S^*_A|c\rangle \langle c|S_A|1\rangle \} |A \rangle . \tag{2} \]

Let the eikonal for the quark-nucleon and antiquark-nucleon QCD gluon exchange interaction be $T^+_\chi(b)$ and $T^-\chi(b)$, where $T^+_\chi$ and $T^-\chi$ are the $SU(N_c)$ generators for the quark and antiquarks states, respectively. The vertex $V_a$ for excitation of the nucleon $g^aN \rightarrow N^*_a$ into color octet state is so normalized that after application of closure the vertex $g^a\bar{g}^bNN$ is $\delta_{ab}$. Then, to the two-gluon exchange approximation, the $S$-matrix of the $(qq)$-nucleon interaction equals

\[ S_N(b_+, b_-) = 1 + i[T^+_\chi(b_+) + T^-\chi(b_-)]V_a - \frac{1}{2}[T^+_\chi(b_+) + T^-\chi(b_-)]^2 . \tag{3} \]

The profile function for interaction of the $qq$ dipole with a nucleon is $\Gamma(b_+, b_-) = 1 - S_N(b_+, b_-)$. For a color-singlet dipole $(T^+_\chi + T^-\chi)^2 = 0$ and the dipole cross section for interaction of the color-singlet $qq$ dipole with the nucleon equals

\[ \sigma(b_+ - b_-) = 2 \int d^2b_+ (N|\Gamma(b_+, b_-)|N) \]
\[
N^2 - 1 = \frac{1}{2Nc} \int d^2b_+ [\chi(b_+) - \chi(b_-)]^2.
\]

The nuclear S-matrix of the straight-path approximation is

\[
S_A(b_+, b_-) = \prod_{j=1}^A S_N(b_+ - b_j, b_- - b_j),
\]

where the ordering along the longitudinal path is understood. We evaluate the nuclear expectation value in (2) in the standard dilute gas approximation. To the two-gluon exchange approximation, per each and every nucleon \( N_j \) only the terms quadratic in \( \chi(b_j) \) must be kept in the evaluation of the single-nucleon matrix element \( \langle N_j|S_N S_N|N_j \rangle \) which enters the calculation of \( S_A S_A \).

Following the technique developed in [4, 5] we can reduce the calculation of the evolution operator for the 2-body density matrix (2) to the evaluation of the S-matrix \( S_{4A}(b_+, b_-, b'_+, b'_-') \) for the scattering of a fictitious 4-parton state composed of the two quark-antiquark pairs in the overall color-singlet state. Because \( (T^{\sigma}_{2g})^* = -T^{\sigma}_2 \), within the two-gluon exchange approximation the quarks entering the complex-conjugate \( S_A^* \) in (2) can be viewed as antiquarks, so that

\[
\sum_c \langle 1|S_A^*|c \rangle \langle c|S_A|1 \rangle = \sum_{c'c} \delta_{c'c} \langle c'|c|S_{4A}|1 \rangle,
\]

where \( S_{4A}(b'_+, b'_-, b_+, b_-) \) is an S-matrix for the propagation of the two quark-antiquark pairs in the overall singlet state. While the first \( q\bar{q} \) pair is formed by the initial quark \( q \) and antiquark \( \bar{q} \) at impact parameters \( b_+ \) and \( b_- \), respectively, in the second \( q'\bar{q}' \) pair the quark \( q' \) propagates at an impact parameter \( b'_- \) and the antiquark \( \bar{q}' \) at an impact parameter \( b'_+ \).

If \( \sigma_4(b'_+, b'_-, b_+, b_-) \) is the color-dipole cross section operator for the 4-body state, then the evaluation of the nuclear expectation value for a dilute gas nucleus in the standard approximation of neglecting the size of color dipoles compared to a radius of heavy nucleus gives [11]

\[
S_{4A}(b'_+, b'_-, b_+, b_-) = \exp\{ -\frac{1}{2} \sigma_4(b'_+, b'_-, b_+, b_-) T(b) \},
\]

where \( T(b) = \int d\rho_2 n_A(b_2, b) \) is the optical thickness of a nucleus at an impact parameter \( b = \frac{1}{4}(b_+ + b'_+ + b_- + b'_-) \), and \( n_A(b_2, b) \) is nuclear matter density with the normalization \( \int d^2b T(b) = A \). The single-nucleon S-matrix \( S_A(b_2, b) \) contains transitions from the color-singlet to the both color-singlet and color-octet \( q\bar{q} \) pairs. However, only color-singlet operators contribute to \( \langle N_j|S_N S_N|N_j \rangle \), and hence the matrix \( \sigma_4 \) only includes transitions between the \( |11 \rangle \) and \( |88 \rangle \) color-singlet 4-parton states, the \( |18 \rangle \) states are not allowed.
After some color algebra, we find matrix elements $\sigma_{11} = \langle 11|\sigma_4|11 \rangle$, $\sigma_{18} = \langle 11|\sigma_4|88 \rangle$ and $\sigma_{88} = \langle 88|\sigma_4|88 \rangle$ (details of the derivation are presented in 2). For forward hard jets with the momenta $p_+^2 \gg Q_A^2$, which are produced from interactions with the target nucleus of small color dipoles in the incident photon the two eigenvalues of the operator $\sigma_4$ are

$$\Sigma_2 \approx \sigma_{11}$$

and

$$\Sigma_1 \approx \sigma_{88}.$$  

For small color-singlet dipoles $\Sigma_2 \approx \sigma_{11} \approx 0$ and the nuclear distortion factor takes on a simple form

$$\sum_{A^*} \sum_{c} \langle A|\{S_A^*|c\} \{c| S_A|1\}|A \rangle - \langle 1; A|S_A^*|A; 1 \rangle \langle 1; A|S_A|A; 1 \rangle$$

$$\approx \sqrt{N_c^2 - 1} \sigma_{18} \sigma_{88} \left\{ 1 - \exp \left\{ -\frac{1}{2} \sigma_{88} T(b) \right\} \right\}.$$  (7)

The introduction of this nuclear distortion factor into (1) gives the hard dijet inclusive cross section:

$$\frac{d\sigma}{d^2b dz d^2p_+ d^2\Delta} = T(b) \sum_{j=0}^{\infty} w_A(b, j) \int d^2\kappa f^{(j)}(\Delta - \kappa) \frac{d\sigma_N}{dz d^2p_+ d^2\kappa},$$  (8)

where the acoplanarity/ decorrelation momentum $\Delta$ is

$$\Delta = p_+ + p_-,$$

the probability of finding $j$ spatially overlapping nucleus in a Lorentz-contracted nucleus is

$$w_A(b, j) = \frac{1}{j!} \frac{\gamma(j+1, 2\lambda_c \nu_A(b))}{2\lambda_c \nu_A(b)},$$

$$\nu_A(b) = \frac{1}{2} \alpha_S \sigma_0 T(b),$$

$$\lambda_c = \frac{N_c^2}{N_c^2 - 1}$$

and $\gamma(j, x) = \int_0^x dy y^{j-1} \exp(-y)$ is an incomplete gamma-function. The function

$$f^{(j)}(\kappa) = \int \prod_{i=1}^j d^2\kappa_i f(\kappa_i) \delta(\kappa - \sum_{i=1}^j \kappa_i), \quad f^{(0)}(\kappa) = \delta(\kappa)$$  (9)
in eq. (8) is a collective gluon field of \( j \) overlapping nucleons introduced in [15].

The probabilistic form of a convolution of the differential cross section on a free nucleon target with the manifestly positive defined distribution \( f^{(1)}(\kappa) \) in (8) can be understood as follows. Hard jets originate from small color dipoles. It is the quantum mechanical interference that suppresses interaction with soft gluons of the small-sized color-singlet \( q\bar{q} \) state. However, the first inelastic interaction inside a nucleus converts the \( q\bar{q} \) pair into the color-octet state, in which color charges of the quark and antiquark do not neutralize each other, rescatterings of the quark and antiquark in the collective color field of intranuclear nucleons become uncorrelated, and the broadening of the momentum distribution with nuclear thickness follows a probabilistic picture.

3 Decorrelation of dijets in DIS off nuclei: numerical estimates

To quantify the azimuthal decorrelation of two forward jets we find it convenient to introduce the mean transverse acoplanarity momentum squared \( \langle \Delta_T^2(b) \rangle \), where \( \Delta_T \) is transverse to an axis of the jet with higher momentum. It is assumed that jets are hard, \( |p_+| \gg Q_A \). The convolution property of the hard dijet cross section (8) suggests

\[
\langle \Delta_T^2(b) \rangle_A \approx \langle \kappa_T^2(b) \rangle_A + \langle \Delta_T^2 \rangle_N ,
\]

where \( \langle \Delta_T^2 \rangle_N \) refers to DIS on a free nucleon, and \( \langle \kappa_T^2(b) \rangle_A \) is the nuclear broadening term. The sign \( \approx \) in (10) reflects the kinematical limitations on \( p_- \) and \( \kappa \) in the practical evaluation of the acoplanarity distribution. In a typical final state it is the harder jet with larger transverse momentum which defines the jet axis and the acoplanarity momentum \( \Delta \) will be defined in terms of components of the momentum of softer jet with respect to that axis, for instance, see [16]). For the sake of definiteness, we present numerical estimates for the Gedanken experiment in which we classify the event as a dijet if the quark and antiquark are produced in different hemispheres, i.e., if the azimuthal angle \( \phi \) between two jets exceeds \( \pi/2 \), the quark jet has fixed \( |p_+| \) and the antiquark jet has higher transverse momentum \( |p_+| \lesssim |p_-| \lesssim 10|p_+| \) (in the discussion of the experimental data one often refers to the higher momentum jet as the trigger jet and the softer jet as the away jet [16]). The free-nucleon quantity \( \langle \Delta_T^2 \rangle_N \) can be estimated starting with the small-\( \Delta \) expansion for excitation of hard, \( p_+^2 \gg \varepsilon^2 = z(1-z)Q^2 \), light flavor dijets from transverse photons

\[
\frac{d\sigma_N}{dzd^2p_+d^2\Delta} \approx \frac{1}{\pi} e_f^2 \alpha_m \alpha_S(p_+^2) \left[ z^2 + (1-z)^2 \right] \times \frac{1}{\Delta^4} \left( \frac{\Delta^2}{\partial \log \Delta^2} \right) \left( \varepsilon^2 + p_+^2 \right) \left( \varepsilon^2 + p_+^2 + \Delta^2 \right),
\]
Figure 1: The mean acoplanarity momentum squared $\langle Q^2_T \rangle_N$ for DIS on a free nucleon target with production of trigger jets with the transverse momentum higher than $p_+^+$ for several values of $Q^2$. The numerical results are for $x = 0.01$ and the input unintegrated gluon SF of the proton is taken from ref. 17.

where $\partial G(x, \Delta^2)/\partial \log \Delta^2 = F(x, \Delta^2)$ is the unintegrated gluon structure function (SF) of the free nucleon 17. Then a quick estimate

$$\langle Q^2_T \rangle_N \approx p_+^2 \frac{F(x, p_+^2)}{2G(x, p_+^2)},$$

(12)
correctly describes the numerical results shown in fig. 1. As far as the dijets are hard, $p_+^2 \gtrsim z(1-z)Q^2 \sim \frac{1}{4}Q^2$, the acoplanarity momentum distribution would not depend on $Q^2$, which holds still better if one considers $\sigma_T + \sigma_L$. This point is illustrated in fig. 1, where we show $\langle Q^2_T \rangle_N$ at $z = 1/2$ for several values of $Q^2$. Because of this weak dependence on $Q^2$ here-below we make no distinction between DIS and real photoproduction, $Q^2 = 0$. In the practical evaluations of the nuclear contribution $\langle \kappa_2^2(b) \rangle_A$ one can use the eq. 8 which gives the result

$$\langle \kappa_2^2(b) \rangle_A \approx \frac{1}{2} \lambda_c Q_A^2(b) \left[ \log \frac{2p_+}{\sqrt{\lambda_c Q_A(b)}} - 1 \right],$$

(13)

where

$$Q_A^2(b) \approx \frac{4\pi^2}{N_c} \alpha_S(Q^2)G(x, Q^2)T(b).$$

(14)
For average DIS on heavy nuclei the reference value is $\langle Q_{Au}^2(b) \rangle = 0.9$ (GeV/c)$^2$, see [2]. The atomic mass number dependence of nuclear broadening $\langle \kappa_T^2(b) \rangle_A$ for jets with $p_+ = 4$ GeV/c in average DIS off nucleus is shown in fig. 2. The principal reason why $\langle \kappa_T^2 \rangle_A$ is numerically small compared to $\langle Q_{Au}^2(b) \rangle$ is that even for such a heavy nucleus as $^{197}$Au the no-broadening probability in average DIS is large, $\langle w_{Au}(b,0) \rangle \approx 0.5$. A comparison of the free nucleon broadening $\langle \Delta_T^2 \rangle_N$ from fig. 1 with the nuclear contribution $\langle \kappa_T^2(b) \rangle_A$ from fig. 2 shows that the nuclear mass number dependence of azimuthal decorrelation of dijets in average DIS off nuclei will be relatively weak.

However, nuclear broadening will be substantially stronger for a subsample of central DIS events at $b \sim 0$. In fig. 3 we show for the gold, $^{197}$Au, target a dependence of the averaged nuclear broadening $\langle \kappa_T^2(b) \rangle_A$ on the impact parameter at several values of $p_+$.  

One can enhance $Q_{Au}^2$ and nuclear contribution $\langle \kappa_T^2(b) \rangle_A$ still further selecting DIS events when the photon breaks up into the $q\bar{q}$ pair on the front face of a nucleus. Experimentally, precisely such events are isolated by selecting very large multiplicity or very high transverse energy of produced secondary particles ([16] and references therein).

Now we comment on the recent finding by the STAR collaboration of a disappearance of back-to-back high $p_T$ hadron correlation when going from peripheral to central gold-gold collisions at RHIC [17]. The application of
color dipole formalism to hard hadron-nucleus interactions \footnote{1} suggests that our analysis of acoplanarity of forward hard jets can be readily generalized to mid-rapidity jets. One only has to choose an appropriate system of dipoles, for instance, the open heavy flavor production can be treated in terms of the intranuclear propagation of the gluon-quark-antiquark system in the overall color-singlet state. At RHIC energies jets with moderately large \( p_T \) are for the most part due to gluon-gluon collisions. In our language that can be treated as a breakup of gluons into dijets and azimuthal decorrelation of hard jets must be discussed in terms of intranuclear propagation of color-octet gluon-gluon dipoles. For such gluon-gluon dipoles the relevant saturation scale \( Q_{S_A}^2 \) is larger \footnote{18} than that for the quark-antiquark dipoles by the factor \( 2\lambda_c = C_A/C_F = 9/4 \). Arguably, in central nucleus-nucleus collisions distortions in the target and projectile nuclei add up and the effective thickness of nuclear matter is about twice of that in DIS. Then, the results shown in fig. 2 suggest that for central gold-gold collisions the nuclear broadening of gluon-gluon dijets could be quite substantial, \( \langle \kappa^2_T (b = 0) \rangle_{AuAu} \sim (3-4) (GeV/c)^2 \) for average central \( AuAu \) collisions and even twice larger if collisions take place at front surface of colliding nuclei.

The principal effect of nuclear broadening is a reduction of the probability
of observing the back-to-back jets

$$P(b) \approx \frac{\langle \Delta_T^2 \rangle_N}{\langle \kappa_T^2(b) \rangle_A + \langle \Delta_T^2 \rangle_N}$$

(15)

and one needs to compare $\langle \Delta_T^2 \rangle_N$ to $\langle \kappa_T^2(b) \rangle_A$. Our eq. (12) for the free nucleon case holds as well for the gluon-gluon collisions. Then the results shown in fig. 1 entail $\langle \Delta_T^2 \rangle_N \approx \langle \kappa_T^2(0) \rangle_{AuAu} \sim (3-4) (GeV/c)^2$ at the jet momentum $p_+ = p_J = (6-8)$ GeV/c and our nuclear broadening would become substantial for all jets with $p_+$ below the decorrelation threshold momentum $p_J$. In practice, the STAR collaboration studied the azimuthal correlation of two high-$p_T$ hadrons and for the quantitative correspondence between the STAR observable and azimuthal decorrelation in the parent dijet one needs to model fragmentation of jets into hadrons (for the modern fragmentation schemes see 19), here we notice that the cutoff $p_+$ in our Gedanken experiment is related to the momentum cutoff $p_{T,\text{min}}$ of associated tracks from the away jet, whereas our jet of momentum $p_-$ can be regarded as a counterpart of the trigger jet of STAR. The STAR cutoff $p_T = 2$GeV/c corresponds to the parents jets with the transverse momentum $p_+ \sim (2 - 3)p_T = (4 - 6)$ GeV which is comparable to, or even smaller than, the decorrelation threshold momentum $p_J = (6 - 8)$ GeV/c. Then eq. (15) suggests that in the kinematics of STAR the probability to observe the back-to-back away and trigger jets decreases approximately twofold from peripheral to central $AuAu$ collisions,

$$P(0) \approx 0.5,$$

and perhaps even stronger, so that our azimuthal decorrelation may contribute substantially to the STAR effect.

4 Remark on breakup of photons into semihard dijets

In our previous analysis 19 of the single particle spectrum it has been discovered that the transverse momentum distribution of sea quarks is dominated by the anticollinear, anti-DGLAP splitting of gluons into sea, when the transverse momentum of the parent gluons is larger than the momentum of the sea quarks. That suggests strongly a complete azimuthal decorrelation of forward minijets with the transverse momenta below the saturation scale, $p_{\pm} \lesssim Q_A$.

In 21 this limiting case has been considered in detail. The principal point is that the minijet-minijet inclusive cross section depends on neither the minijet nor decorrelation momentum. This observation proves a disappearance of the azimuthal correlation of minijets with the transverse momentum below the saturation scale.
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