Multiple parton interactions and forward double pion production in pp and dA scattering

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We estimate the contributions by double-parton interactions to the cross sections for $pp \to \pi^0\pi^0 X$ and $dA \to \pi^0\pi^0 X$ at the Relativistic Heavy Ion Collider (RHIC). We find that such contributions become important at large forward rapidities of the produced pions. This is, in particular, the case for dA scattering, where they strongly enhance the azimuthal-angular independent pedestal component of the cross section, providing a natural explanation of this feature of the RHIC dA data. We argue that the discussed processes open a window to studies of double quark distributions in nucleons. We also briefly address the roles of shadowing and energy loss in dA scattering, which we show to affect the double-inclusive pion cross section much more strongly than the single-inclusive one. We discuss the implications of our results for the interpretation of pion azimuthal correlations.

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I. INTRODUCTION

Cross sections for the production of identified hadrons at large transverse momentum play a crucial role at RHIC, for both the spin and the heavy-ion program. In the latter they serve as important probes of phenomena such as shadowing, gluon saturation, or parton energy-loss [1]. For single-inclusive hadron production in pp scattering, $pp \to hX$, it was found that next-to-leading order (NLO) perturbative QCD [2] provides a very good description of the RHIC data over wide ranges of transverse momentum, rapidity, and beam energy [3]. Striking suppression effects with respect to the pp baseline have been observed on the other hand for scattering involving nuclei, among them in dA scattering at forward rapidities [4–6]. These experimental studies at RHIC were also extended to the production of two forward pions, in both pp and dA scattering [5,7,8]. Of particular interest here are correlations between the pions in the difference of their azimuthal angles, $\Delta\varphi$. As expected, strong peaks in the distributions at $\Delta\varphi = 0, \pi$ were observed in pp scattering. These are also present in peripheral dA collisions. However, in central dA collisions, the "backward" peak at $\Delta\varphi \sim \pi$ is strongly depleted. It has been suggested that this depletion is due to gluon saturation effects in the Color Glass Condensate of the gold nucleus [9,10]. At the same time, all distributions show a very significant $\Delta\varphi$-independent pedestal that is much higher in the dA case than in pp, a feature that has received somewhat less attention.

In this paper, we will demonstrate that double-parton interactions for which two leading partons of the "projectile" proton (or deuteron) interact with the "target" naturally make large contribution in the forward kinematics studied at RHIC, often dominating over the leading-power contribution. They are particularly important in dA scattering. We will find that they could well be responsible for the pedestals in the $\Delta\varphi$ correlations and impact the interpretation of the observed correlations in both the pedestal and the backward peak regions.

Apart from their relevance for forward scattering at RHIC, double-parton interactions are also of wider interest in QCD as they provide a novel window on strong-interaction dynamics, including correlations of leading partons inside nucleons or nuclei [11]. As a result, they have received an ever-growing attention over the past few years [12–15]. In addition, understanding of double- and multiparton interactions is also important for a proper modeling of the structure of the final state for central pp collisions at the LHC [11,14,16] and hence for the search for new particles. Current experimental studies of multiparton interactions involve selection of events with two back-to-back pairs of jets (or, a jet and a photon); see e.g. [17,18]. The fact that RHIC may provide a unique way to learn about multiparton interactions, without having to use the more traditional double-scattering observables, is remarkable.

Our paper is organized as follows. In Sec. II, we discuss double-inclusive-pion production in pp scattering. We first demonstrate that in the leading-twist (LT) approximation, the cross section at forward rapidities involves incoming quarks with very high momentum fraction. As a result, we find that double-parton processes in which two quarks each with relatively moderate $x \sim 0.3$–0.4 scatter independently, become competitive over a fairly wide kinematic range at RHIC. In Sec. III, we study the double-parton mechanism for dA collisions, which we find in the impulse approximation to be significantly enhanced as compared to the LT mechanism. We then discuss the impact of double-scattering contributions on the pedestal and the peak of the $\Delta\varphi$ distribution, along with generic features of gluon shadowing and parton energy loss, and argue that the suggested mechanisms allow to describe the bulk features of the data. Finally, we summarize our results in Sec. IV.

II. TWO-PION PRODUCTION IN pp SCATTERING

In this section, we explore the main features of $pp \to \pi^0\pi^0 X$ through the LT mechanism based on a single hard scattering, and through double-parton interactions. We choose the case of pp collisions, both because of its potential for studying new aspects of high energy QCD and nucleon correlation structure, but also because it sets the baseline for our later discussion of nuclear scattering. In the following, $p_{T,1}$, η_1 are the transverse momentum and pseudorapidity of the "trigger" pion. The corresponding variables of the second "associated" pion are denoted by $p_{T,2}$, η_2. We consider collisions at center-of-mass energy $\sqrt{S} = 200$ GeV and 500 GeV.

A. Leading-twist mechanism

We start with the LT mechanism for which two partons collide in a single hard scattering. The generic expression for the LT $pp \to \pi^0\pi^0 X$ cross section is given in factorized form by

$$\frac{d^4\sigma_{\rm LT}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} = \sum_{abcd}\int dx_a dx_b dz_c dz_d f_a^p(x_a) f_b^p(x_b) \times \frac{d^4\hat{\sigma}_{ab\to cdX}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} D_c^{\pi^0}(z_c) D_d^{\pi^0}(z_d), \quad (1)$$

where the sum runs over all partonic channels, with f_a^p, f_b^p denoting the usual parton distribution functions of the proton, $D_c^{\pi^0}$, $D_d^{\pi^0}$ the pion fragmentation functions for partons c, d, and $\hat{\sigma}_{ab\to cdX}$ the corresponding partonic hard-scattering cross sections. The latter may be computed in QCD perturbation theory, starting at lowest order (LO) from $2 \to 2$ scattering $ab \to cd$. Even though the NLO corrections are available in the literature [19], we will restrict ourselves for this study to LO computations. We shall comment on this point below. We have for simplicity suppressed in Eq. (1) the dependence of the various functions on the factorization/renormalization scale μ. Throughout our studies we choose the CTEQ6L parton distribution functions [20] and the LO de Florian-Sassot-Stratmann set of fragmentation functions [21].

Since $\eta_1 + \eta_2 = \log(x_a/x_b)$, production of two pions at relatively forward rapidities must arise from "imbalanced" collisions where a large-x parton from one proton hits a small-x parton from the other [22]. These will typically be collisions of a valence quark and a gluon. Figure 1 shows the distributions of the integrand in Eq. (1) in x_a for various values of η_1 and bins in η_2. Here, we have chosen $p_{T,1} = 2.5$ GeV and $1.5\ \text{GeV} \le p_{T,2} \le p_{T,1}$ and the scale $\mu^2 = (p_{T,1}^2 + p_{T,2}^2)/2$. The inserts in the figure show the distributions on a linear scale, normalized in such a way that their integral is unity in each case. One observes overall that with increasing η_1 or η_2, the distributions are shifted to higher values of x_a. In particular, we see that the average value of $x_a = x_{\rm quark}$ for typical kinematics of the RHIC forward measurements is very high.

B. Double-scattering mechanism

The results shown in Fig. 1 suggest that "double-scattering" contributions, with two separate hard

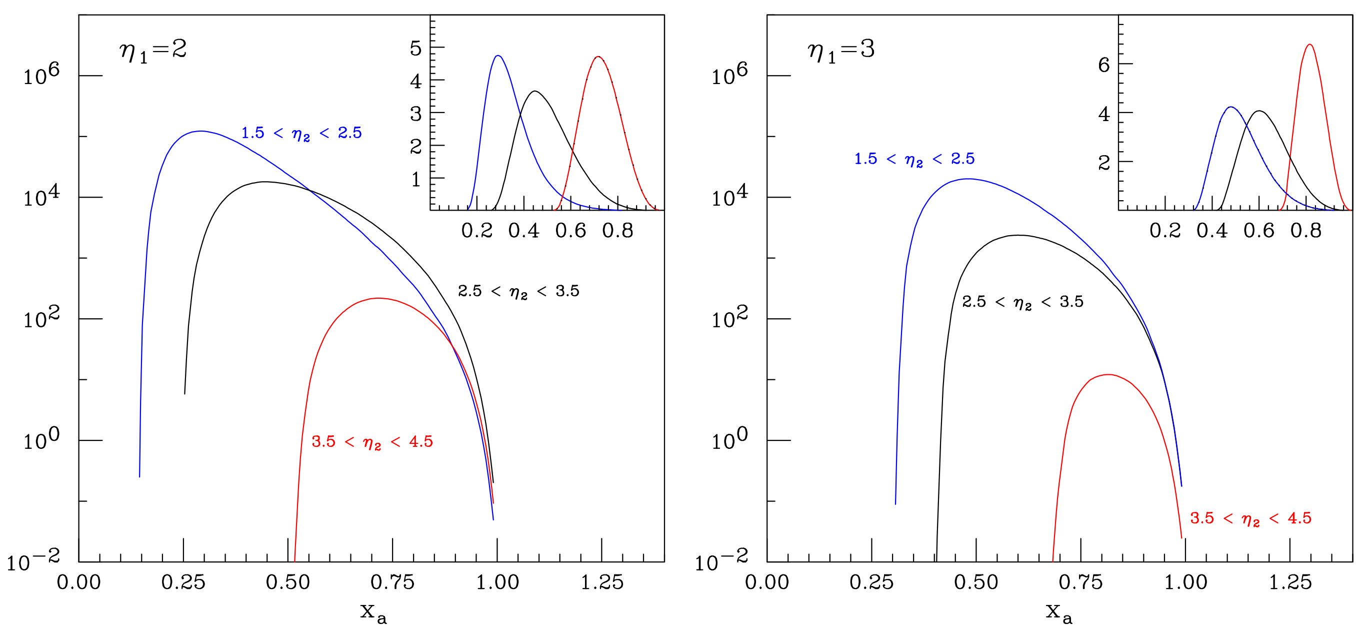

FIG. 1 (color online). Distributions of the leading-power LO cross section for $pp \to \pi^0\pi^0 X$ (see Eq. (1)) at $\sqrt{S} = 200$ GeV in momentum fraction x_a, for $\eta_1 = 2$ (left) and $\eta_1 = 3$ (right). We have chosen $p_{T,1} = 2.5$ GeV and integrated over $1.5\ \text{GeV} \le p_{T,2} \le p_{T,1}$ and various bins in η_2. Units are arbitrary. The inserts show the corresponding normalized distributions on a linear scale.

interactions in a single pp collision, could become relevant at forward rapidities. Here the idea is that each of the two hard interactions produces a pion. The double-scattering contributions are power-suppressed or "higher-twist". Like for the leading-twist contribution in Eq. (1), each of the two interactions will proceed primarily by a high-x (valence) quark scattering off a small-x gluon. However, compared to the leading-twist part, the momentum fractions of the quarks participating in the double-scattering contributions will on average be much smaller, because for the latter the kinematics of the two recoiling "unobserved" partons is unconstrained. This makes the double-scattering contributions potentially dominant at forward rapidities.

A proper treatment of the double-parton mechanism would involve use of "two-parton generalized parton distributions" (2pGPDs) in the proton; see [13] for summary and references. This would allow to calculate the dimensional factor between single- and double-inclusive scattering, which characterizes the transverse spread of the double-parton distributions in the colliding protons. At present, knowledge about 2pGPDs is overall not sufficient to fully apply this formalism to the case of hadron pair production in hadronic collisions. We therefore resort to simple physically motivated estimates of the double-scattering contribution. Here, we are guided by the observation that in the case when the partons in each of the scattering protons are completely uncorrelated, we will have

$$\frac{d^4\sigma_{\text{double,uncorr.}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} = \frac{1}{\pi R_{\text{int}}^2}\frac{d^2\tilde{\sigma}_{\text{LT}}}{dp_{T,1}d\eta_1}\frac{d^2\tilde{\sigma}_{\text{LT}}}{dp_{T,2}d\eta_2} \quad (2)$$

for the double-scattering contribution to $pp \to \pi^0\pi^0 X$, where $\tilde{\sigma}_{\text{LT}}$ denotes the leading-twist *single-inclusive* cross section for $pp \to \pi^0 X$, given by

$$\frac{d^2\tilde{\sigma}_{\text{LT}}}{dp_T d\eta} = \sum_{abc}\int dx_a dx_b dz_c f_a^p(x_a) f_b^p(x_b) \times \frac{d^2\hat{\sigma}_{ab\to cX}}{dp_T d\eta} D_c^{\pi^0}(z_c), \quad (3)$$

with single-inclusive partonic cross sections $\hat{\sigma}_{ab\to cX}$. Furthermore, in Eq. (2) πR_{int}^2 is an "effective" transverse area covered by the two correlated partons (it was denoted as σ_{eff} in a number of experimental papers and some of the theoretical papers, although it has little to do with an interaction cross section). In the approximation of partons uncorrelated in the transverse plane, it can be expressed through a convolution of usual generalized parton distributions in the hadrons [13,23]. Hence, if we assume for simplicity that the partons' transverse spread does not depend on their momentum fractions x, we can write

$$\frac{d^4\sigma_{\text{double}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} = \frac{1}{\pi R_{\text{int}}^2}\sum_{abca'b'c'}\int dx_a dx_b dz_c dx_{a'} dx_{b'} dz_{c'} f_{aa'}^p(x_a, x_{a'}) f_b^p(x_b) f_{b'}^p(x_{b'}) \frac{d^2\hat{\sigma}^{ab\to cX}}{dp_{T,1}d\eta_1}\frac{d^2\hat{\sigma}^{a'b'\to c'X'}}{dp_{T,2}d\eta_2} \times D_c^{\pi^0}(z_c) D_{c'}^{\pi^0}(z_{c'}). \quad (4)$$

Here, $f_{aa'}^p(x_a, x_{a'})$ is a "double-parton" distribution for partons a, a' in the same proton, which we will model in the following. If the partons are not correlated, it is equal to the product of two ordinary parton distributions, $f_{aa'}^p(x_a, x_{a'}) = f_a^p(x_a) f_{a'}^p(x_{a'})$, and Eq. (4) reverts to (2). As shown in Eq. (4), we neglect for simplicity correlations in the target, i.e., in the proton probed at small-x.

The double-parton distribution has to obey the kinematic constraint $x_a + x_{a'} \leq 1$. We implement this condition for all partonic channels. Beyond that, we only consider double-parton correlations for valence quarks, that is, for the case $a, a' \equiv q, q'$, with $q = u, d$. For these, we make the ansatz

$$f_{qq'}^p(x_q, x_{q'}) = \frac{1}{2}\left[f_q^p(x_q) \times \phi\left(\frac{x_{q'}}{1 - x_q}\right) + (q \leftrightarrow q')\right]. \quad (5)$$

The picture we have in mind here is that the first interaction involves a valence quark with its distribution $f_q(x_q)$. The distribution of a second valence quark that participates in the second hard-scattering is then expected to be modified relative to the usual parton distribution. For instance, if the first hard-scattering involves an up quark, then fewer up quarks will be available for the second interaction. We assume this effect to be described by a single function ϕ, given by

$$\phi(\xi) = \frac{c}{\sqrt{\xi}}(1 - \xi)^n, \quad (6)$$

with $c = 3/4$ and $n = 1$. The latter value follows from counting rule arguments; scaling violations would be expected to increase it somewhat. The normalization factor c in Eq. (6) is determined from the baryon number sum rule $\int_0^1 d\xi\phi(\xi) = 1$. Since the expression for the cross section is symmetric in x_q, $x_{q'}$ we perform symmetrization of Eq. (5). For all partonic combinations not involving valence quarks, we also use Eq. (5), but with ϕ replaced by the usual parton distribution function $f_{a'}^p(x_{a'}/(1 - x_a))$. Here, the modified argument guarantees that the kinematic constraint $x_a + x_{a'} \leq 1$ is respected. Our procedure should be compared to the model of [24] where it was assumed that also for valence quarks the function $\phi(\xi)$ is given by the usual distribution $f_q^p(\xi)$, which has $n \sim 3$, and no symmetrization was performed. While our ansatz arguably is physically better motivated, we do not find much numerical difference between the two models, except very close to the phase space boundary at high rapidities and/or transverse momenta. Rather than the precise choice of ϕ, it is the kinematic constraint $x_a + x_{a'} \leq 1$ that matters most in our numerical studies, reducing the cross section.

For our calculations, we choose $\pi R_{\rm int}^2 = 15$ mb [17,18] in Eq. (2). This experimental value is smaller than $\pi R_{\rm int}^2 = 34$ mb obtained in the mean field approximation for the 2pGPDs when partons are not correlated in the transverse plane [23]. The value ~34 mb is an upper limit on $\pi R_{\rm int}^2$ when only momentum fractions larger than $\sim 10^{-2}$ are relevant, provided there is no repulsion between the partons. A smaller experimental value of $\pi R_{\rm int}^2$ indicates the presence of transverse correlations among partons. In principle, these correlations can depend on the flavors and momentum fractions of the partons involved in the double-scattering interaction. Qualitatively, we expect that picking two leading quarks would select configurations with reduced transverse separation between the quarks, making it more natural to use the experimental value for $\pi R_{\rm int}^2$ than the uncorrelated estimate. It is of

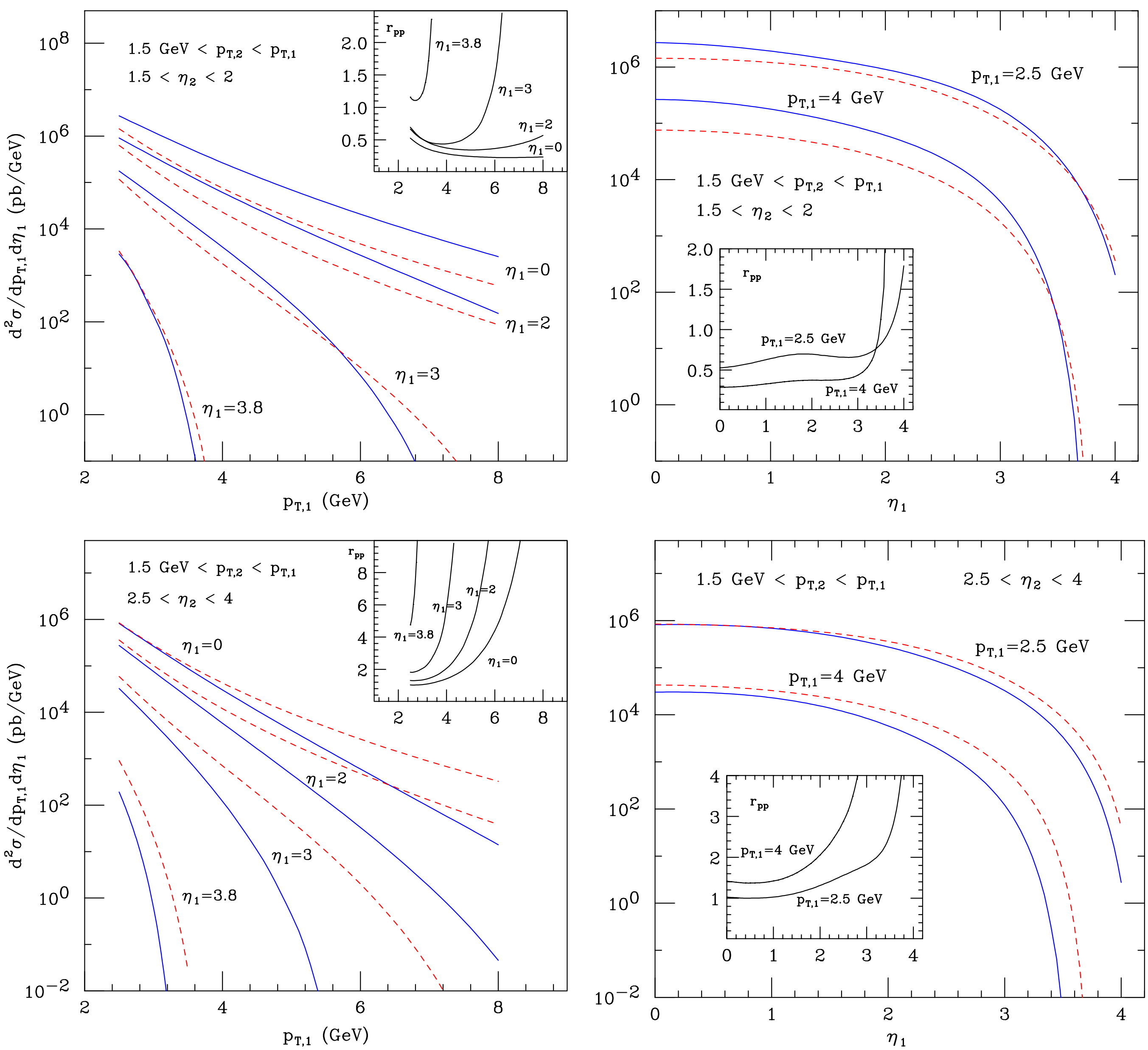

FIG. 2 (color online). Comparison of the leading-twist cross section for $pp \to \pi^0\pi^0 X$ (solid, see Eq. (1)) and the double-interaction contribution estimated from Eq. (4) (dashed), as functions of $p_{T,1}$ (left) and η_1 (right). The plots in the upper row are for $1.5 < \eta_2 < 2$, the ones in the lower row for $2.5 < \eta_2 < 4$. For all plots, we have chosen 1.5 GeV $< p_{T,2} < p_{T,1}$. The inserts in each plot show the ratio r_{pp} of the double-interaction contribution to the leading-twist one, see Eq. (8).

interest here that in the limit when the typical transverse separation between the quarks is much smaller than for the small-x gluons, one can derive based on [13]

$$\pi R_{\rm int}^2 = \left[\int \frac{d^2\Delta}{(2\pi)^2} F_{2g}^2(\Delta)\right]^{-1} = \frac{12\pi}{m_g^2} \approx 14 \text{ mb}. \quad (7)$$

Here, $F_{2g}(\Delta) \approx 1/(\Delta^2/m_g^2 + 1)^2$ with $m_g^2(x \sim 0.01) = 1.1\ \text{GeV}^2$ is the two-gluon form factor of the nucleon. A larger value of $\pi R_{\rm int}^2$ would evidently reduce the size of our estimates for the double-interaction contribution.

Figure 2 shows our results for the leading-twist cross section $d\sigma_{\rm LT}$ for $pp \to \pi^0\pi^0 X$ in Eq. (1) and for the double-interaction contribution $d\sigma_{\rm double}$ according to Eq. (4), as functions of the trigger pion's transverse momentum and rapidity. For the associated pion, we have integrated the cross sections over $1.5\ \text{GeV} \le p_{T,2} \le p_{T,1}$ and $1.5 \le \eta_2 \le 2$ (upper row) or $2.5 \le \eta_2 \le 4$ (lower row). $d\sigma_{\rm LT}$ has been calculated as before; for $d\sigma_{\rm double}$ in Eq. (4), we have chosen the same parton distributions and fragmentation functions, and the scales $\mu = p_{T,i}$. All calculations are done at LO. As one can see from Fig. 2, the estimated double-scattering contribution shows the typical features of a higher-twist (power-suppressed) contribution. It tends to increase relative to the leading-twist cross section towards lower transverse momenta. Near midrapidity and for moderately high $p_{T,1}$, double-scattering is essentially negligible. On the other hand, it also increases towards the kinematic boundaries at high rapidities and transverse momenta. Therefore, it is likely to play a significant role for much of the kinematic regime relevant in the studies of two-pion correlations at forward rapidities at RHIC. Here it would affect also the distributions in the difference $\Delta\varphi$ of the azimuthal angles of the two pions, where it should enhance both the backward peak at $\Delta\varphi \sim \pi$ and the pedestal at $\Delta\varphi < \pi$. It is worth emphasizing already at this point that the lowest-order LT part only contributes at $\Delta\varphi = \pi$, whereas the double-scattering piece will uniformly contribute at all $\Delta\varphi$. Hence, as the LT cross section receives contributions at $\Delta\varphi < \pi$ only at higher orders in perturbation theory, the double-scattering is expected to dominate even more strongly away from the backward peak. For future reference, we also define the ratio of the double-scattering contribution to the leading-twist one:

$$r_{pp} \equiv \frac{d^4\sigma_{\rm double}}{d^4\sigma_{\rm LT}}. \quad (8)$$

We show the results for r_{pp} in the inserts in Fig. 2. This number should be compared with the ratio of the areas under the pedestal and under the backward peak at $\Delta\varphi \sim \pi$, which is of the order of 2.

As there will be high luminosity runs at RHIC with polarized protons at $\sqrt{S} = 500$ GeV, we have also performed calculations at this energy. Figure 3 shows the results as functions of $p_{T,1}$ and η_1, integrated over $2\ \text{GeV} < p_{T,2} < p_{T,1}$ and $1.5 < \eta_2 < 4$. As expected, the effects are overall smaller than at $\sqrt{S} = 200$ GeV, but remain significant at large rapidities.

As we have hinted at earlier, there are still considerable uncertainties in the computation of the double-scattering contribution. We remind the reader that one would in principle need to set up a framework based on 2pGPDs.

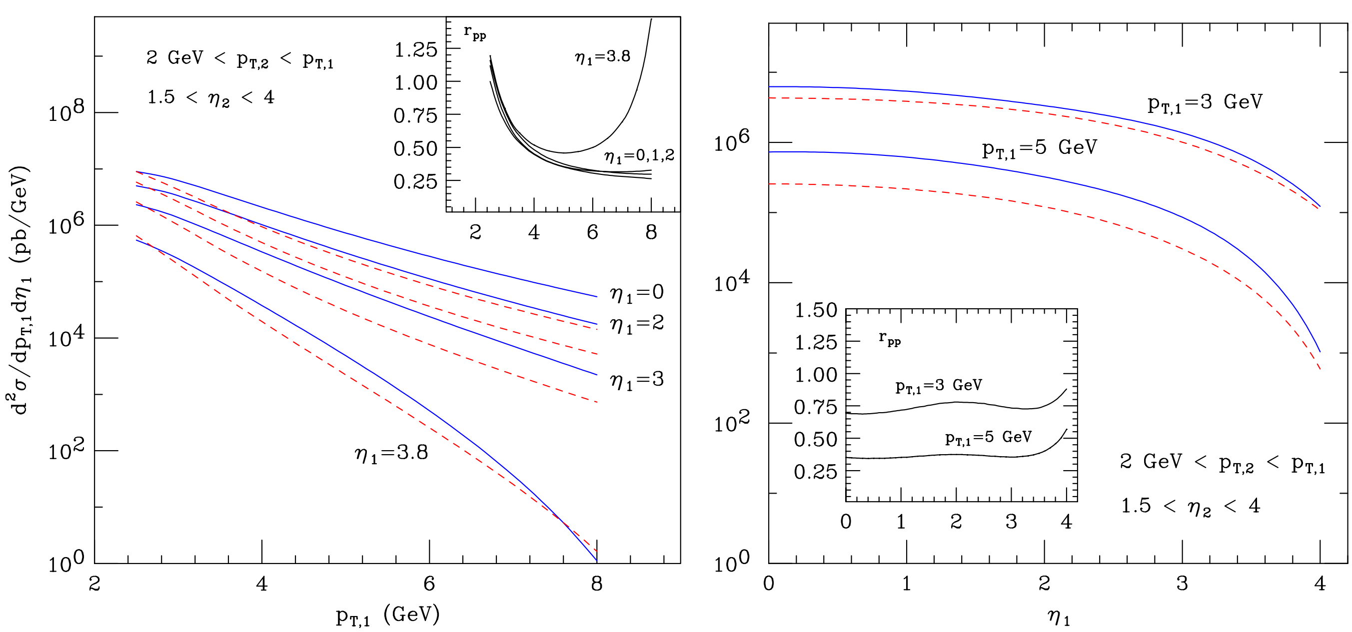

FIG. 3 (color online). As Fig. 2, but at $\sqrt{S} = 500$ GeV and for $2\ \text{GeV} < p_{T,2} < p_{T,1}$ and $1.5 < \eta_2 < 4$.

Even within our ansatz in Eq. (2), there is some uncertainty regarding the value for $\pi R_{\rm int}^2$ and the model used for the double-parton correlation functions. Furthermore, at forward rapidities the fragmentation functions are probed at rather high momentum fractions z, where they are not known accurately. On top of this, one needs to address the role of higher-order QCD corrections. The double-inclusive and single-inclusive leading-twist cross sections carry significant dependence on the renormalization/factorization scales. While the NLO corrections are available for both the double-inclusive leading-twist cross section [19] and for the single-inclusive one in Eq. (3) [2], it is not guaranteed that the form of $d\sigma_{\rm double}$ in Eq. (4) carries over to higher orders of perturbation theory, since particle radiation will tend to correlate the two separate hard interactions (this effect should be small, however, for configurations which dominate in the mean field uncorrelated approximation, since in this case the bulk of the parton cross sections originates from quark transverse separations much larger than $1/p_T$). That said, each of the cross sections $d\sigma_{\rm LT}$, $d\tilde{\sigma}_{\rm LT}$ in Eqs. (1) and (2) is known to receive positive NLO radiative corrections of $\gtrsim 50\%$ or so for RHIC kinematics, so that it appears likely that *QCD corrections will overall enhance the relevance of the double-scattering contribution*.

The uncertainties inherent in the present calculations somewhat limit the possibilities to achieve a better determination of $\pi R_{\rm int}^2$ from RHIC measurements. Nonetheless, if our phenomenological predictions are correct, it might be possible to identify the double-scattering contributions from detailed studies of the dependence of the two-pion cross section on transverse momenta and rapidities. A further possible test of this picture would be to measure a third pion at "recoil kinematics" $\eta_3 \sim 0$. This could serve to further enhance the double-scattering contribution over the leading-twist one, since the latter can give rise to a third pion only at higher orders in perturbation theory, whereas the double interactions naturally give rise to a third (and even a fourth) recoiling "jet." Obviously, the study of polarization effects in two-pion correlations would be of interest as well in the context of the double-scattering mechanism.

III. TWO-PION PRODUCTION IN *dA* SCATTERING

A. Introductory remarks

An important finding at RHIC [4–6] is that the rate of forward pion production at relatively large transverse momenta, where perturbative QCD describes the corresponding pp data, is suppressed in dA scattering by a large factor as compared to the impulse approximation result. This suppression is expressed by the "nuclear modification factor" R_{dA}, which effectively compares the observed production rates for a given centrality trigger to the prediction based on the approximation that the parton density in nuclei at an impact parameter b is equal to the additive sum of the parton densities of individual nucleons at this impact parameter. A more formal way to formulate the latter assumption is to define the impact-parameter dependent parton distribution of the nucleus, $f_a^A(x, Q^2, b)$, which coincides with the corresponding diagonal generalized parton distribution (GPD) in impact parameter representation [25]. In the discussed approximation, $f_a^A(x, Q^2, b)$ is given by

$$f_a^A(x, Q^2, b) = f_a^N(x, Q^2) T_A(b), \tag{9}$$

where $T_A(b)$ is the standard nuclear profile function which is given by an integral of the density function over the longitudinal direction

$$T(b) = \int \rho_A(\sqrt{b^2 + z^2}) dz. \tag{10}$$

T_A is normalized to $\int d^2 b T(b) = A$. The experimental data show that the suppression becomes stronger with increase of rapidity η. It is found that R_{dA} is typically of the order $1/3$ for forward kinematics. Furthermore, the suppression becomes stronger with decrease of b and is strongest for $b \sim 0$.

The analysis [22] has demonstrated that the dominant mechanism for single-inclusive-pion production in the forward kinematics explored at RHIC is scattering of a leading quark of one (projectile) nucleon off a gluon in the other (target) nucleon. The median value of momentum fraction x_g of the gluon was found to be in the range $x_g \sim 0.01$–0.03, depending on the rapidity of the pion. The nuclear gluon density for such values of x_g is known to be close to the incoherent sum of the gluon fields of the individual nucleons since the coherence length in the interaction is rather modest for the distances involved. As a result, the leading-twist nuclear shadowing effects cannot explain the observed suppression [22], and one needs a novel dynamical mechanism to explain the suppression of pion production in such collisions.

An important additional piece of information comes from the study of correlations of the leading forward pion with an additional pion produced at central rapidities [5,26]. In this case, the dominant contribution comes from the scattering off gluons with $x_g \sim 0.01$–0.02. An extensive analysis performed in [27] has demonstrated that the strengths of such forward-central correlations are similar in dA and in pp scattering once one corrects for the contribution of soft interactions to the pion yield at $\eta \sim 0$, and that in dA the dominant source of leading pions is scattering at large impact parameters. This conclusion is supported by the observation of the STAR experiment [28] that the associated multiplicity of soft hadrons in events with a forward pion is a factor of 2 smaller than in minimum-bias dA events. This reduction factor is consistent with the estimate of [27]. Overall, the patterns observed in forward inclusive-pion production and forward-central correlations are consistent with the picture of effective energy losses which we further discuss in Sec. III C. Hence we will use it below for the numerical estimates of the deviation form the impulse approximation. We note in passing that the above-mentioned features of the forward pion production data represent a challenge for the $2 \to 1$ scattering mechanism [29,30] that dominates in the color glass condensate model. In this mechanism, forward pions are predominantly produced at central impact parameters without producing recoil pions at central rapidities.

In Ref. [22], we suggested that to study the effects of small-x gluon fields in the initial state one needs to study production of *two leading forward* pions in nucleon-nucleus collisions. Recently, such data were taken in dAu collisions [5,7,8]. In the next subsection, we will analyze the role of the double-parton interactions in the kinematics explored at RHIC. The effects associated with suppression of the single-inclusive spectrum mentioned above will be discussed further below in Sec. III C.

B. Double-parton versus single-parton interactions—treatment in the impulse approximation

As we saw in the previous section, large values of the rapidities of the two pions select high x in the projectile hadron, and double-scattering contributions may become very significant. As measurements in this kinematic domain have been carried out in dAu scattering at RHIC [5,7,8], it is of much interest to see in how far the double-interaction contributions are further enhanced in the reaction $dA \to \pi^0\pi^0X$. Compared to the pp case, it is clear that the presence of many nucleons in the scattering process will offer more possibilities for multiple-parton interactions.

We may distinguish three contributions to the double-parton mechanism in dA scattering, as shown in Fig. 4:

(a) Two (valence) quarks from one of the nucleons in the deuteron participate in the hard-scattering, striking the same nucleon in the heavy nucleus [Fig. 4(a)].

(b) Independent scattering of the deuteron's proton and neutron off separate nucleons in the heavy nucleus. Each of the two collisions produces one of the observed pions [Fig. 4(b)].

(c) Same as (a), but with the double interaction occurring off two different nucleons in the heavy nucleus. Again each of the two collisions produces one of the observed pions [Fig. 4(c)].

We now proceed to make estimates for these contributions. For our more illustrative purposes, we neglect effects of nuclear (anti-)shadowing for the heavy nucleus. Also, we treat the heavy nucleus as roughly isoscalar. For our estimates, we need to take into account the distribution of nucleons in a heavy nucleus. Since the experiments are performed with a centrality trigger, it is useful to first write the double-inclusive cross section in a form where the integral over impact-parameter b is kept explicitly [12]. We write all expressions for N-nucleus scattering, where $N = (p+n)/2$ denotes an isoscalar combination of proton and neutron. Since they are bound in a deuteron they propagate at similar impact parameters. We further assume that the impulse approximation is valid for the interaction with the nucleus. For any contribution that involves scattering off only one of the target nucleons, we then have the generic formula

$$\frac{d^4\sigma^{NA}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} = \int d^2b T(b) \frac{d^4\sigma^{NN}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} \quad (11)$$

for the two-pion cross section. Here, $T(b)$ is the nuclear thickness factor defined above in Eq. (10). Equation (11) holds for contribution (a), but evidently also for the leading-twist piece. Hence, if we consider a fixed impact parameter and take their ratio, the factor $T(b)$ will cancel:

$$r_{\rm a}(b) \equiv \frac{d^4\sigma^{NA}_{\rm double,(a)}}{d^4\sigma^{NA}_{\rm LT}} = \frac{d^4\sigma^{NN}_{\rm double}}{d^4\sigma^{NN}_{\rm LT}} = r_{NN}. \quad (12)$$

In the last step, we have used our definition in Eq. (8) for the ratio of the double-interaction contribution to the leading-twist one, now adapted to the case of NN collisions. Apart from trivial (and small) isospin modifications related to the fact that there are contributions here from pp, pn, np, and nn scattering, r_{NN} is identical to r_{pp} considered in the previous section.

The situation is different, however, for contributions (b) and (c), for which two target nucleons are involved in the scattering, so that the square of $T(b)$ will appear in the expressions [12]. For contribution (b), we have

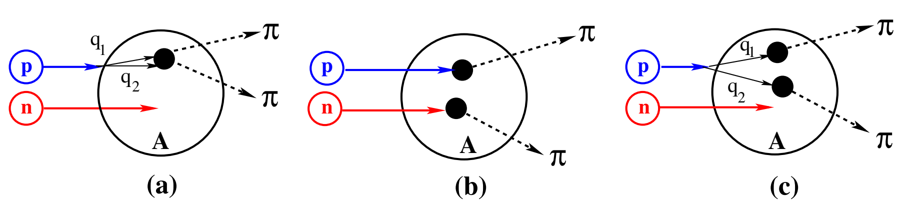

FIG. 4 (color online). Contributions to two-pion production in dA collisions through the double-interaction mechanism.

$$\frac{d^4\sigma^{NA}_{\text{double,(b)}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} = \int d^2b \frac{A-1}{A} T^2(b) \times \frac{1}{2}\left[\frac{d^2\tilde{\sigma}^{pN}_{\text{LT}}}{dp_{T,1}d\eta_1}\frac{d^2\tilde{\sigma}^{nN}_{\text{LT}}}{dp_{T,2}d\eta_2} + (p_{T,1},\eta_1 \leftrightarrow p_{T,2},\eta_2)\right], \tag{13}$$

where $\tilde{\sigma}_{\text{LT}}$ again denotes a LT single-inclusive cross section as introduced in Eq. (3). Here, we have neglected for simplicity any correlations between quarks in the two projectile nucleons. As indicated, we need to properly symmetrize (13), since either the pN or the nN interaction can produce a given pion. Taking again the ratio to the leading-twist term at fixed impact-parameter, one factor of $T(b)$ cancels, and we have for large A

$$r_{\text{b}}(b) = \frac{d^4\sigma^{NA}_{\text{double,(b)}}}{d^4\sigma^{NA}_{\text{LT}}} = \frac{T(b)[d^2\tilde{\sigma}^{pN}_{\text{LT}} d^2\tilde{\sigma}^{nN}_{\text{LT}} + (p_{T,1},\eta_1 \leftrightarrow p_{T,2},\eta_2)]}{2d^4\sigma^{NN}_{\text{LT}}}. \tag{14}$$

Finally, for contribution (c) we define analogously

$$r_{\text{c}}(b) \equiv \frac{d^4\sigma^{NA}_{\text{double,(c)}}}{d^4\sigma^{NA}_{\text{LT}}}. \tag{15}$$

Here, the numerator is again proportional to $T^2(b)$, while the denominator is linear in $T(b)$. With the help of Eq. (4), we find at fixed impact parameter

$$r_{\text{c}}(b) = T(b)\pi R^2_{\text{int}} \frac{d^4\sigma^{NN}_{\text{double}}}{d^4\sigma^{NA}_{\text{LT}}} = T(b)\pi R^2_{\text{int}} r_{NN}, \tag{16}$$

again up to small isospin corrections.

Before presenting more detailed numerical results for the ratios in Eqs. (12), (14), and (16), we discuss their relative size. From (12) and (16) we see immediately that to very good approximation

$$\frac{r_a}{r_c} = \frac{1}{T(b)\pi R^2_{\text{int}}}. \tag{17}$$

As before, we use $\pi R^2_{\text{int}} = 15$ mb. For heavy nuclei with $b \sim 0$, we have $T(0) \approx 2.2\ \text{fm}^{-2} = 1/(4.54\ \text{pb})$. Therefore, we have $r_a/r_c \approx 0.3$. The ratio of r_b and r_c will be close to 1 at midrapidity, where correlations and valence-gluon scattering are not very important. Toward large rapidities, however, r_b must become much larger than r_c, since it is not subject to the constraint $x_a + x_{a'} \leq 1$ because of the fact that for (b) the proton and the neutron scatter independently.

Figure 5 shows the sum

$$r_{dA} \equiv r_a + r_b + r_c \tag{18}$$

at impact-parameter $b = 0$. It gives the ratio of the full double-scattering contribution to the leading-twist one.[1] One can see that double-parton interactions in dA scattering appear to lead to very significant enhancements of the cross section over the leading-twist one, much stronger than in pp scattering. The inserts in the figure show the corresponding ratios r_a/r_c and r_c/r_b, which show the trend discussed above.

C. Impact on interpretation of pion azimuthal correlations

We expect our findings in Fig. 5 to be also relevant for the interpretation of the azimuthal distributions of the pions mentioned earlier. Such distributions have recently been investigated by the STAR and PHENIX experiments at RHIC [5,7,8]. What is measured is the distribution in the difference $\Delta\varphi$ of the azimuthal angles of the two pions. The distributions are normalized relative to the total number of trigger events, that is, given a high-$p_{T,1}$ pion with rapidity η_1 that passes the selection cuts, the $\Delta\varphi$ distribution gives the probability for finding a second pion in a given azimuthal bin. The (still preliminary) data show peaks corresponding to nearside ($\Delta\varphi \sim 0$) and awayside ($\Delta\varphi \sim \pi$) correlations, on top of a broad pedestal that extends over all $\Delta\varphi$. The pedestal is significantly higher in dA than in pp scattering. Also, it is found that in central dA the awayside peak is strongly depleted when both pions are produced at forward rapidities, $\eta_i \sim 3$ [5,7,8].

In view of the relatively early stage the data are in, our discussion will be overall more qualitative here. Also, the theoretical framework is not sufficiently developed for a full quantitative study. The leading-twist calculations we have done in the previous section were entirely in the framework of collinear factorization. Here, the ingredients for a full calculation are essentially available, including next-to-leading order corrections (even though for simplicity we did not use these). In the case of the correlation function in $\Delta\varphi$, however, the calculation is much more involved. Away from $\Delta\varphi = \pi$, the leading-twist part will be dominated by $2 \to 3$ processes, which are available. However, near $\Delta\varphi = \pi$—the most interesting region—any finite order of perturbation theory will fail because of the presence of large Sudakov double-logarithms. A resummation of these logarithms to all orders in perturbation theory is required here, which unfortunately so far has not been worked out. To perform this resummation is, of course, well outside the scope of this paper. In addition, also nonperturbative contributions will be present very close to $\Delta\varphi = \pi$. It seems to us that none of the theoretical studies of the correlation function addresses these contributions at an appropriate level. We could follow a standard procedure and attempt to model perturbative and

[1]Note that our r_{dA} is not to be confused with the usual nuclear modification factor R_{dA} mentioned in Sec. III A.

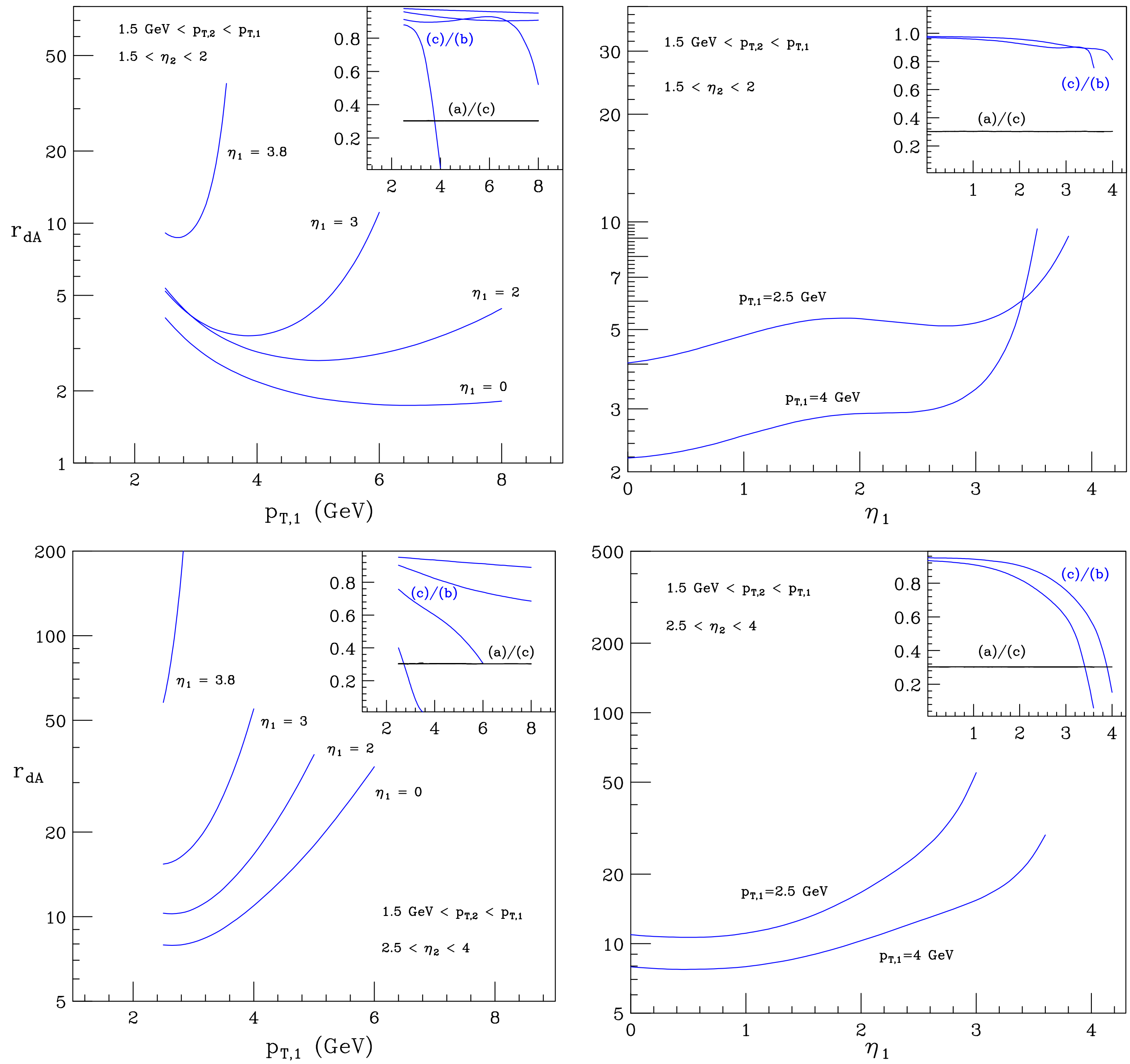

FIG. 5 (color online). Ratio r_{dA} (defined in Eq. (18)) of double-parton and leading-twist contributions in $dA \to \pi^0\pi^0X$. The plots in the upper row are for $1.5 < \eta_2 < 2$, the ones in the lower row for $2.5 < \eta_2 < 4$. For all plots we have chosen $1.5\ \text{GeV} < p_{T,2} < p_{T,1}$. The inserts show the ratios r_a/r_c and r_c/r_b.

nonperturbative contributions to the $\Delta\varphi$ correlation function using Gaussian smearing in parton transverse momenta; however, we refrain from such a rather *ad hoc* approach and stick to a more qualitative discussion that captures the main physics.

Our first observation is that, as discussed in Sec. II A, for the leading-twist mechanism the two pions will predominantly be produced back-to-back in azimuthal angle, that is, around $\Delta\varphi = \pi$. Pure $2 \to 2$ scattering produces the pions at $\Delta\varphi = \pi$; the region away from the backward peak, around say $\Delta\varphi \sim \pi/2$, can only be filled by $2 \to 3, 4, \ldots$ scattering, which are of higher order in the strong coupling α_s. These features are in contrast to the double-scattering mechanism, for which the two pions are produced essentially uncorrelated in $\Delta\varphi$ and which hence is expected to uniformly fill the $\Delta\varphi$ distribution. Since we found in the previous subsections that double-scattering is prevalent at forward angles in pp and in particular dA scattering, we conclude that the *numerator* of the pedestal around $\Delta\varphi \sim \pi/2$ should be almost entirely due to the double-scattering mechanism. The *denominator*, on the other hand, is the single-inclusive trigger cross section

$d^2\tilde{\sigma}/dp_{T,1}d\eta_1$, for which we can safely assume that double-scattering contributions play a less important role. In any case, its value is known from past STAR measurements in pp and dA scattering [5]. Therefore, the height of the pedestal for dA is generically given as

$$\text{Ped}_{dA} \approx \frac{d^4\sigma^{dA}_{\text{double}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} \Big/ \frac{d^2\tilde{\sigma}^{dA}}{dp_{T,1}d\eta_1}, \tag{19}$$

and similarly for pp scattering. If our considerations are correct, we can estimate the relative heights of the pedestals in pp and dA scattering:

$$\frac{\text{Ped}_{dA}}{\text{Ped}_{pp}} \approx \frac{d^4\sigma^{dA}_{\text{double}}}{d^4\sigma^{pp}_{\text{double}}} \times \frac{d^2\tilde{\sigma}^{pp}}{d^2\tilde{\sigma}^{dA}} = \frac{d^4\sigma^{dA}_{\text{LT}}}{d^4\sigma^{pp}_{\text{LT}}} \times \frac{r_{dA}}{r_{pp}} \times \frac{d^2\tilde{\sigma}^{pp}}{d^2\tilde{\sigma}^{dA}}, \tag{20}$$

where in the second step we have used Eqs. (8) and (18) to introduce the ratio of leading-twist and double-scattering contributions. The last factor, $d^2\tilde{\sigma}^{pp}/d^2\tilde{\sigma}^{dA}$, corresponds to the inverse of the nuclear modification factor R_{dA} that we mentioned earlier, at trigger transverse momentum $p_{T,1}$ and rapidity η_1. We assume that the first factor, $d^4\sigma^{dA}_{\text{LT}}/d^4\sigma^{pp}_{\text{LT}}$, roughly shows the square of this suppression and hence is of order R^2_{dA}. We shall give a better argument for this below. Then, one factor of R_{dA} cancels in the ratio of the pedestals, and we obtain

$$\frac{\text{Ped}_{dA}}{\text{Ped}_{pp}} \approx R_{dA} \times \frac{r_{dA}}{r_{pp}} \approx 3. \tag{21}$$

Here we have used a typical value of $r_{dA}/r_{pp} \sim 10$ from our previous Figs. 2 and 4, and [4–6] $R_{dA} \sim 1/3$. Obviously, the value we find in (21) can only be a rough estimate; however, we are encouraged by the fact that it is well consistent even quantitatively with the experimental observation of a significant enhancement of the pedestal in central dA scattering [7]. Thus, we conclude that the RHIC experiments may well have discovered the first example of multiparton interactions in many-nucleon systems, with all previous observations having been restricted to pp or $p\bar{p}$ collisions. Data with a finer binning in η_1, η_2 would allow a more detailed check of our expectations.

We can now go one step further and consider the away-side peak around $\Delta\varphi \sim \pi$. In the peak region, the structure of the two-pion correlation in dA scattering is

$$\text{Peak}_{dA} \approx \left(\frac{d^4\sigma^{dA}_{\text{LT}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} + \frac{d^4\sigma^{dA}_{\text{double}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2}\right) \Big/ \frac{d^2\tilde{\sigma}^{dA}}{dp_{T,1}d\eta_1}. \tag{22}$$

As indicated, we here expect to have a contribution also from the leading-twist term. We now subtract the pedestal term given in Eq. (19) and find

$$\text{Peak}_{dA} - \text{Ped}_{dA} \approx \frac{d^4\sigma^{dA}_{\text{LT}}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} \Big/ \frac{d^2\tilde{\sigma}^{dA}}{dp_{T,1}d\eta_1}. \tag{23}$$

Taking again the ratio to the corresponding quantity in pp scattering we obtain

$$\frac{\text{Peak}_{dA} - \text{Ped}_{dA}}{\text{Peak}_{pp} - \text{Ped}_{pp}} \approx \frac{d^4\sigma^{dA}_{\text{LT}}}{d^4\sigma^{pp}_{\text{LT}}} \times \frac{d^2\tilde{\sigma}^{pp}}{d^2\tilde{\sigma}^{dA}}. \tag{24}$$

Compared to the pedestal ratio given in (20), this value does not contain the factor $r_{dA}/r_{pp} \sim 10$. We hence conclude that the *height of the peak above the pedestal* is about $R_{dA} \sim 1/3$ times smaller in dA scattering than in pp. Compared to the relative heights of the pedestals in dA and pp, this is a reduction of even a factor 10. Again, both these findings are consistent with the observations in the data. Evidently, within our more qualitative discussion one cannot rule out that there could also be a contribution due to the $2 \to 1$ broadening mechanism discussed in [9,10] to the pion azimuthal correlation, which in these models constitutes only a small fraction ($\lesssim 1/6$) of the pedestal events. In view of the much larger double-parton mechanism contribution that we find for the same $\Delta\varphi$, it is not clear at the moment how to check experimentally to what degree such a $2 \to 1$ broadening contribution is present.

The step that remains is to investigate the ratios $d^2\tilde{\sigma}^{pp}/d^2\tilde{\sigma}^{dA}$ and $d^4\sigma^{dA}_{\text{LT}}/d^4\sigma^{pp}_{\text{LT}}$ in (20) and (24). The former is given by the nuclear modification factor R_{dA} which, as we mentioned, has been found experimentally at RHIC to show a significant suppression of the dA cross section relative to the pp one at forward rapidities. The mechanism behind this suppression is not conclusively understood so far. As discussed in [22], it could arise, at leading-twist level, from a combination of two effects. The first is the leading-twist shadowing phenomenon [31] whose impact on R_{dA} we computed in [22]. Using the nuclear parton distribution functions (nPDFs) of [32], we found relatively small shadowing effects, because the relevant gluon momentum fractions in the target are on average not very small for single-inclusive hadron production, $\langle x_g \rangle \sim 0.02$, even at forward rapidities. For such x_g, gluon shadowing is predicted in [32] to be relatively moderate. That said, little is known experimentally about gluon shadowing, and the recent set [33] of nPDFs proposes a stronger shadowing effect. For our present study, we will stick to the use of the nPDFs of [32].

The second effect is energy loss of partons. It was shown in [34] that partons propagating through the target nucleus in kinematics close to the black disk regime suffer "fractional" energy losses. The interactions near the black disk regime select configurations in which the parton has split into two or more partons. In [22], we pointed out that even a relatively small energy loss of order 5 to 10%, which is consistent with the estimated magnitude of this effect [27], can explain the observed patterns of suppression in forward

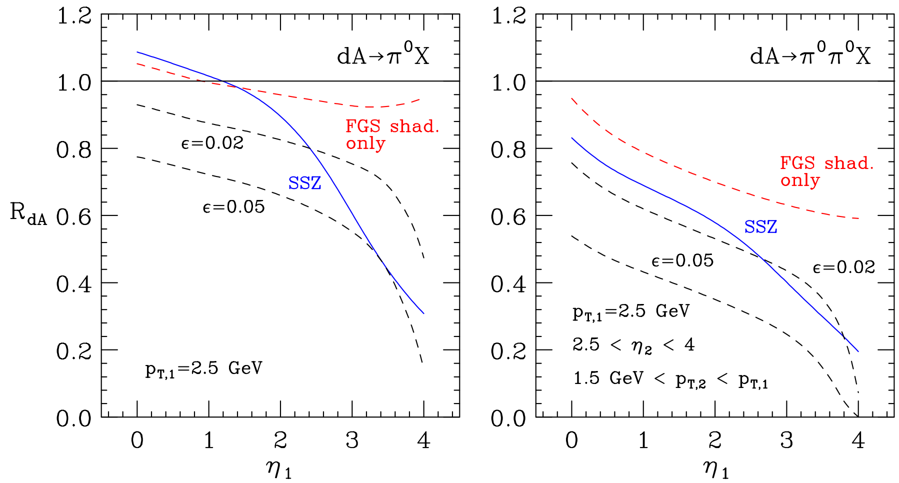

FIG. 6 (color online). Left: Nuclear modification factor R_{dA} for single-inclusive leading-twist pion production as a function of rapidity η_1 at $p_{T,1} = 2.5$ GeV. The upper dashed line shows the effect of leading-twist shadowing for the Frankfurt-Guzey-Strikman (FGS) nuclear parton distributions [32]. The solid line includes shadowing and the "medium-modified" fragmentation functions of Sassot-Stratmann-Zurita (SSZ) [36]. The lower dashed lines show the results for two simple energy-loss models, see text. Right: Same for double-inclusive pion production.

dA-scattering at RHIC. Energy loss effects are also typically embodied in an effective way in "nuclear-modified" fragmentation functions; see for example [35–39], which are fitted or compared to RHIC dA and AA data. These may hence also serve as useful tools for investigating suppression effects in a leading-twist calculation of single-inclusive or double-inclusive particle production at RHIC.

The left part of Fig. 6 makes our observations more quantitative. We show results for R_{dA} (for $A = Au$), computed from the leading-twist single-inclusive cross section $d\tilde{\sigma}^{dA}_{\rm LT}/p_{T,1}d\eta_1$ in Eq. (3) and normalized to the corresponding pp cross section. The upper dashed line shows the effect of including leading-twist shadowing of [32] which, as discussed above, is quite small. The solid line shows the result when using the same shadowing and in addition the set of nuclear-modified fragmentation functions of Ref. [36].[2] As one can see, R_{dA} is suppressed, except for the midrapidity region, where antishadowing effects are relevant. The suppression grows with η_1 and is of order $1/3 - 1/2$ at forward rapidities of the pion, in line with the experimental observations. This is expected since the fragmentation functions of [36] have been fitted to the RHIC forward single-pion production data.

Interestingly, a simple model of energy loss generically yields results of the same size, as shown by the two lower dashes lines. Here we have, in the spirit of the earlier discussion of fractional energy loss, simply rescaled the momentum fraction x_b of the parton in the gold nucleus by $x_b \longrightarrow x_b(1 + \epsilon)$, and similarly for the fragmenting parton. At forward rapidities, the results for $\epsilon = 0.02$ and $\epsilon = 0.05$ roughly span the one obtained for the nuclear-modified fragmentation functions. At midrapidity, they are lower and fail to reproduce the antishadowing effects seen in the data. This is not a surprise, however, since our simple energy loss estimate is only expected to work at larger rapidities where the produced parton has to traverse the largest amount of strongly-interacting matter and where one is closer to the black disk regime. What is surprising is that even rather small values of ϵ generate relatively large suppression effects. This suggests that, regardless of its precise mechanism, energy loss will always be expected to play a significant role in dA scattering. Note that in our picture, energy losses are fractional only in the proximity of the black disk regime. Consequently, for fixed transverse momentum and increasing rapidity we expect ϵ to increase. In our rough estimates, we have neglected this effect which obviously will work to amplify further the suppression effect. At the same time, the energy losses are expected to be energy independent far away from the black disk regime [40], which explains the absence of suppression of the forward-central correlations. We point out that the resummation of nuclear-enhanced power corrections to the leading-twist cross section also results effectively in a shift of the momentum fraction of the initial projectile quark [41], whose size depends on kinematics. This approach was shown to be quantitatively consistent with the forward suppression of R_{dA}.

Strikingly, the effects we find for single pions are amplified for double-inclusive scattering. The corresponding results are shown in the right part of Fig. 6. One reason for

[2] We note that only *NLO* sets of nuclear-modified fragmentation functions have been presented in the paper [36]. In order to avoid any mismatch with the de Florian-Sassot-Stratmann set [21] that we use for the ordinary fragmentation functions, the solid curves in Fig. 6 have been computed by also using the NLO set of [21] for the calculation of the pp cross section in the denominator of R_{dA}.

the additional suppression in this case is that significantly smaller momentum fractions are probed in the target, down to $x_g \sim 10^{-3}$ [22], where gluon shadowing is stronger. This effect is seen from the upper dashed line in Fig. 6. Furthermore, since two fragmentation functions are present for double-inclusive pion production, the energy loss effect is much more prominent, as shown by the solid and lower dashed lines. Indeed, as we anticipated, the overall suppression of the dA double-inclusive leading-twist cross section is roughly given by the square of that for the single-inclusive one, that is, by R_{dA}^2. This feature is likely generic for any kind of shadowing and energy loss mechanisms present in dA scattering. In this sense, a strong depletion of the backward peak in the pion azimuthal correlation in forward dA scattering (which is of the same magnitude as the experimentally observed depletion) is very natural, given the previously found milder suppression of single-inclusive pion production. It is of interest that the observed suppression of the double-inclusive cross section at central impact parameters as compared to the impulse approximation is close to its lower bound, corresponding to the probability that a quark passing through the nucleus encounters only one nucleon at its impact parameter, which for the case of scattering off gold at $b \sim 0$ is about $1/20 - 1/10$ [42].

We note that for pA, scattering the dominant mechanism (b) of Subsection III B would be absent, so that the double-scattering contributions would remain a bit closer in size to what we found in the pp case. The pion azimuthal correlation should then have a less pronounced pedestal, but a similarly suppressed backward peak.

We finally briefly address the "nearside" correlation of two hadrons produced at $\Delta\varphi \sim 0$, both with large rapidity. Experimentally, this correlation shows a strong peak that (unlike the backward peak at $\Delta\varphi \sim \pi$) is not suppressed in dA scattering. One may wonder what can be said in the context of our present calculation about the region $\Delta\varphi \sim 0$. It is clear that the nearside correlation receives contributions from a rather different physics mechanism: the two hadrons produced at similar azimuthal angle and rapidity may originate from just a *single* fragmentation process, through leading-twist double fragmentation of a high-p_T final-state parton, described by di-hadron fragmentation functions [43]. Unfortunately, little is known at present about the latter. Kinematically, such a contribution is rather similar to a single-inclusive cross section. It is therefore not expected to show the suppression $\sim R_{dA}^2$ that we found above for the leading-twist double-inclusive piece, but rather a suppression of order R_{dA}. As a result, one would naturally expect the nearside peak in dA to show little suppression. To be more specific, we denote the contribution to two-hadron production arising from double-fragmentation by $d^4\sigma_{\mathrm{df,LT}}^{dA}$. It adds to the (roughly $\Delta\varphi$-independent) double-scattering pedestal contribution. Subtracting the pedestal as before, we find the following structure of the two-pion correlation in dA scattering in the nearside peak region:

$$(\mathrm{Peak}_{dA} - \mathrm{Ped}_{dA})_{\text{near-side}} \approx \frac{d^4\sigma_{\mathrm{df,LT}}^{dA}}{dp_{T,1}d\eta_1 dp_{T,2}d\eta_2} \Big/ \frac{d^2\tilde{\sigma}^{dA}}{dp_{T,1}d\eta_1}, \tag{25}$$

where as before $d^2\tilde{\sigma}^{dA}$ is the ordinary single-inclusive trigger piece. Taking again the ratio to the corresponding quantity in pp scattering, we obtain

$$\left(\frac{\mathrm{Peak}_{dA} - \mathrm{Ped}_{dA}}{\mathrm{Peak}_{pp} - \mathrm{Ped}_{pp}}\right)_{\text{near-side}} \approx \frac{d^4\sigma_{\mathrm{df,LT}}^{dA}}{d^4\sigma_{\mathrm{df,LT}}^{pp}} \times \frac{d^2\tilde{\sigma}^{pp}}{d^2\tilde{\sigma}^{dA}}, \tag{26}$$

to be compared to Eq. (24) for the awayside correlation at $\Delta\varphi \sim \pi$. *If* the double-fragmentation contribution indeed behaves similar to a single-inclusive cross section—which we consider to be quite natural given its origin from fragmentation of a single parton—the two factors on the right-hand side of Eq. (26) would be R_{dA} and $1/R_{dA}$, respectively, and hence cancel. The whole ratio would then be unity, which is indeed what the data [7] show. Even though this discussion is again rather qualitative, it offers a natural and straightforward explanation of the nonsuppression of the nearside peak in dA scattering.

IV. CONCLUSIONS

We have investigated the role of double-scattering contributions to double-inclusive pion production in pp and dA scattering at RHIC. We have found that these become important at large rapidities of the produced pions. This is, in particular, the case for central dA scattering, where the double-scattering contribution can exceed the leading-twist one by large factors. Further detailed studies of double-inclusive pion production at RHIC may provide a unique way for studying parton correlations in the nucleon. This is remarkable because traditionally only four-particle final states were considered as possible probes of double-parton scattering. These would typically not be viable at forward rapidities.

The double-scattering contributions appears to play a critical role in the interpretation of the pion azimuthal correlations observed experimentally at RHIC. They primarily produce pions that are uncorrelated in azimuthal angle and hence are expected to strongly dominate the pedestals seen in the distributions. We have shown that the relative heights of the pedestals in pp and dA scattering can be qualitatively understood in this way. We have furthermore shown that once the pedestal is subtracted, the remaining backward correlation peak in dA scattering is strongly affected by shadowing and energy loss effects. These are found to be much stronger for double-inclusive scattering compared to single-inclusive, giving rise to a depletion of the backward peak in dA, consistent with the observations at RHIC. Overall, in the light of our results, the patterns observed in the pion azimuthal correlations at RHIC find a natural qualitative explanation.

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