Abstract

We analyze $x_F$ dependence of the Drell-Yan transverse momentum broadening in hadron-nucleus collisions. In terms of generalized factorization theorem, we show that the $x_F$ dependence of the transverse momentum broadening, $\Delta\langle q_T^2 \rangle(x_F)$, can be calculated in perturbative QCD. We demonstrate that $\Delta\langle q_T^2 \rangle(x_F)$ is a good observable for studying the effects of initial-state multiple scattering and extracting quark-gluon correlation functions.

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Parton multiple scattering is responsible for many interesting and important phenomena in high energy collisions involving nucleus, such as transverse momentum broadening, energy loss, as well as the nuclear suppression of quarkonium states. Understanding the mechanism of parton multiple scattering and its effects is crucial for making precise predictions of the nuclear suppression of quarkonium productions, which maybe a potential signal for the quark-gluon plasma in relativistic heavy ion collisions [1]. Parton multiple scattering can happen both at the initial state or at the final state. The Drell-Yan pair production in hadron-nucleus collisions provide an excellent place to study effects of initial state parton multiple scattering. The nuclear dependence in the Drell-Yan transverse momentum spectrum for the large $q_T$ region has been studied in QCD perturbation theory [2,3]. In terms of generalized factorization theorem in QCD [4], the effects of multiple scattering can be expressed in terms of multiparton correlation functions [5], which are as fundamental as the parton distributions. However, these parton correlation functions are not well determined yet.

In this letter, we derive the $x_F$ dependence of the Drell-Yan transverse momentum broadening. We show that $\Delta \langle q_T^2 \rangle (x_F)$ can be used as a good observable to study the effects of initial-state multiple scattering and parton energy loss. In addition, it can be used as an excellent observable for extracting information on multi-parton correlation functions.

Consider the Drell-Yan process in hadron-nucleus collisions, $h(p') + A(p) \rightarrow \ell^+ \ell^- (q) + X$, where $q$ is the four-momentum for the virtual photon $\gamma^* \rightarrow \ell^+ \ell^-$ which decays into the lepton pair. $p'$ is the momentum for the incoming beam hadron and $p$ is the momentum per nucleon for the nucleus with the atomic number $A$. Let $q_T$ be the transverse momentum of the Drell-Yan pair, we define the averaged transverse momentum square as

$$\langle q_T^2 \rangle^h_A = \int dq_T^2 \cdot q_T^2 \cdot \frac{d\sigma^h_A}{dQ^2 dq_T^2} / \frac{d\sigma^h_A}{dQ^2}.$$

In Eq. (1), $Q$ is the total invariant mass of the lepton pair with $Q^2 = q^2$. Since single hard scattering is localized in space, only multiple scattering (at least, double scattering) are sensitive to the nuclear size (or $A^{1/3}$ type dependence). Therefore, in order to extract the effect due to multiple scattering, we introduce the nuclear enhancement of the Drell-Yan $\langle q_T^2 \rangle$ as

$$\Delta \langle q_T^2 \rangle \equiv \langle q_T^2 \rangle^h_A - \langle q_T^2 \rangle^h_N,$$

which is often called the transverse momentum broadening. The broadening of the transverse momentum square defined in Eq. (2) should be sensitive to parton multiple scattering between nucleons inside a large nucleus.

In a perturbatively calculable hard scattering process, having an extra scattering between physical partons is suppressed by a power of the hard scale [5]. Therefore, multiple scattering in momentum space between physical partons correspond to an expansion in power series of $1/Q^2$. In this letter, we limit ourselves to double scattering in momentum space. With only single and double scattering, the broadening of the Drell-Yan transverse momentum square, $\Delta \langle q_T^2 \rangle$, can be parameterized as

$$\Delta \langle q_T^2 \rangle = a + b A^{1/3},$$

(3)
which is consistent to existing data [3, 4]. In Eq. (3), \( bA^{1/3} \) term represents the contribution directly from the double scattering which is explicitly proportional to the nuclear size \( (\propto A^{1/3}) \), with \( A \) the atomic weight of the nucleus target.

In Ref. [4], Qiu and Sterman argued that the factorization theorem for hadron-hadron scattering [8] should also be valid at the first non-leading power in momentum transfer, which is essential for systematically calculating the double scattering in QCD perturbation theory. According to this generalized factorization theorem, we can expand the numerator in Eq. (1) as

\[
\int dq_T^2 dq_T^2 \frac{d\sigma_{hA}}{dQ^2 dq_T^2} = \sum_{a,b} \phi_{a/A}(x) \otimes C^{(0)}_{ab\rightarrow ll}(x, Q^2) \otimes \phi_{b/h}(x') + \frac{1}{Q^2} \sum_{a,b} \left[ T_{a/A}(x) \otimes C^{(2)}_a(x, Q^2) \otimes \phi_{b/h}(x') + \phi_{a/A}(x) \otimes \bar{C}^{(2)}(x, Q^2) \otimes T_{b/h}(x') \right] + \ldots
\]

\[
\equiv H^{(0)}_A + H^{(2)}_A + \bar{H}^{(2)}_A + \ldots .
\]

where \( \otimes \) represents convolutions over partonic momenta, and “...” represents the terms that are suppressed by higher powers of \( 1/Q^2 \). In Eq. (4), \( C^{(0)}, C^{(2)}, \) and \( \bar{C}^{(2)} \) are perturbatively calculable hard parts. \( \phi_{b/h}(x') \) is the parton distribution of the beam hadron, and \( \phi_{a/A}(x) \) is the parton distribution in the nucleus normalized by the atomic number \( A \). \( T_{a/A}(x) \) and \( T_{b/h}(x') \) are the four-parton correlation functions [5] in the nucleus and the beam hadron, respectively.

Because of the well-known EMC effect, as well as effects of nuclear shadowing and Fermi motion, the nuclear dependence of the \( \phi_{a/A} \) is nontrivial. However, if we parameterize the effective nuclear parton distribution \( \phi_{a/A} \) into \( A^\alpha \) times corresponding nucleon parton distributions, we found a very small power of \( \alpha \) for a wide range of \( x \) [9]. Taking the EKS98 parameterization of nuclear parton distributions [10] as an example, we found that \( \alpha \sim \pm (0.02 - 0.03) \) for the \( x \)-range relevant to this study, which is much smaller than \( 1/3 \) for the \( A^{1/3} \)-type enhancements. Therefore, \( H^{(0)}_A \) in Eq. (4), which is proportional to \( \phi_{a/A} \), should have a very weak nuclear dependence. Similarly, \( \bar{H}^{(2)}_A \) in Eq. (4) also has a weak nuclear dependence. On the other hand, the nuclear parton correlation function \( T_{a/A} \) has an explicit dependence on the nuclear size \( (\propto A^{1/3}) \) [3, 11], so as the \( H^{(2)}_A \) in Eq. (4).

According to the factorization theorem [3], the denominator in Eq. (1) can also be expand in terms of power series:

\[
\frac{d\sigma_{hA}}{dQ^2} = \sum_{a,b} \phi_{a/A}(x, \mu^2) \otimes \frac{d\sigma_{ab\rightarrow ll}^{(0)}}{dQ^2}(x, x', \mu^2/Q^2, \alpha_s(\mu^2)) \otimes \phi_{b/h}(x', \mu^2) \left[ 1 + O\left( \frac{1}{Q^2} \right) \right]
\]

\[
\equiv \sigma_A^{(0)} \left[ 1 + O\left( \frac{1}{Q^2} \right) \right],
\]

where \( \mu \) represents both renormalization and factorization scale. In Eqs. (4) and (5), all quantities are normalized by the atomic number \( A \). Substituting Eqs. (4) and (5) into Eq. (1), we obtain
In deriving Eq. (3), we kept only terms up to \( O(A^{1/3}/Q^2) \), and dropped all power correction terms without the \( A^{1/3} \) nuclear enhancement. Similarly, for a nucleon target, we have

\[
\langle q_T^2 \rangle_{hN} \approx \frac{H_N^{(0)}}{\sigma_N^{(0)}} \left[ 1 + O \left( \frac{A^0}{Q^2} \right) \right] .
\]

Substituting above Eqs. (6) and (7) into our definition of the nuclear broadening of the transverse momentum square in Eq. (2), we derive

\[
\Delta \langle q_T^2 \rangle \approx \left[ \frac{H_A^{(0)}}{\sigma_A^{(0)}} - \frac{H_N^{(0)}}{\sigma_N^{(0)}} \right] + \frac{H_A^{(2)}}{\sigma_A^{(0)}},
\]

where we neglected terms of \( O(A^0/Q^2) \). The first term in Eq. (8) should have a very weak nuclear dependence due to the fact that effective nuclear parton distributions have a very week \( A^\alpha \) dependence. Therefore, the first term contributes to the \( a \)-term in Eq. (3). In addition, the first term should be numerically very small due to the cancellation between the two terms inside the bracket. On the other hand, the second term in Eq. (8) represents the double scattering contribution and it gives the \( A^{1/3} \)-type enhancement, and therefore, it contributes to the \( bA^{1/3} \)-term in Eq. (3).

At the leading order, the inclusive cross section is

\[
\sigma_A^{(0)} \equiv \frac{d\sigma_{hA \to \ell^+ \ell^-}}{dQ^2} = \sigma_0 \sum_q c_q^2 \int dx' \phi_{q/A}(x') \int dx \phi_{q/h}(x) \delta(Q^2 - xx's),
\]

with \( s = (p + p')^2 \) and the Born cross section

\[
\sigma_0 = \frac{4\pi \alpha_s^2}{9Q^2} .
\]

Also at the leading order, the Drell-Yan differential cross section is given by

\[
\frac{d\sigma_{hA \to \ell^+ \ell^-}}{dQ^2 dq_T^2} = \frac{d\sigma_{hA \to \ell^+ \ell^-}}{dQ^2} \delta(q_T^2). 
\]

Therefore, \( H_A^{(0)} \) and \( H_N^{(0)} \) in Eq. (8) vanish at the leading order, because \( \int dq_T^2 q_T^2 \delta(q_T^2) = 0 \), and consequently, the first term in Eq. (8), or the \( a \)-term in Eq. (3), vanishes at the leading order in \( \alpha_s \).

The double scattering contribution to the Drell-Yan transverse momentum broadening, \( H_A^{(2)} \) in Eq. (8), was derived in Ref. [13]. At the leading order in \( \alpha_s \), it is given by

\[
H_A^{(2)} = \sigma_0 \left( \frac{4\pi^2 \alpha_s}{3} \right) \sum_q \int dx' \phi_{q/h}(x') \int dx T_{q/A}^{(1)}(x) \delta(Q^2 - xx's),
\]

where the four-parton correlation function is defined as
\[ T_{q/A}(x) = \int \frac{dy^-}{2\pi} e^{ixp^+y^-} \int \frac{dy_1^- dy_2^-}{2\pi} \theta(y^- - y_1^-) \theta(-y_2^-) \times \frac{1}{2} \langle p_A | F_+^+(y_2^-) \bar{\psi}_q(0) \gamma^+ \psi_q(y^-) F_+^-(y_1^-) | p_A \rangle. \]  

These parton correlation functions are not well measured yet. By comparing the operator definitions of the correlation functions and the definitions of the normal twist-2 parton distributions, Luo, Qiu, and Sterman (LQS) proposed the following model [3,13]:

\[ T_{f/A}(x) = \lambda^2 A^{1/3} \phi_{f/A}(x), \]  

where \( \lambda \) is a free parameter to be fixed by experimental data, and was estimated in Ref. [13] as \( \lambda^2 = 0.01 \text{ GeV}^2 \) by using the Drell-Yan data from NA10 and E772 experiment [7,14].

In order to obtain the leading order \( x_F \)-dependence of \( \Delta\langle q_T^2 \rangle \), we multiply \( \delta(x_F - (x' - x)) \, dx_F \) to the right-hand side of Eq. (12), and obtain

\[ \frac{dH_A^{(2)}(x_F)}{dx_F} = \sigma_0 \left( \frac{4\pi^2 \alpha_s}{3} \right) \sum_q \phi_{q/h}(x_1) T_{q/A}^{(I)}(x_2) \frac{1}{(x_1 + x_2)s}, \]  

with

\[ x_1 = (\sqrt{x_F^2 + 4\tau + x_F})/2, \quad \text{and} \quad x_2 = (\sqrt{x_F^2 + 4\tau - x_F})/2, \]  

where \( \tau = Q^2/s \). Combining Eqs. (3), (4), and (14), we obtain the \( x_F \) dependence of the Drell-Yan transverse momentum broadening \( \Delta\langle q_T^2 \rangle \) at the leading order in \( \alpha_s \):

\[ \frac{d\Delta\langle q_T^2 \rangle(x_F)}{dx_F} = \left( \frac{4\pi^2 \alpha_s}{3} \right) \sum_q e_q^2 \phi_{q/h}(x_1) T_{q/A}^{(I)}(x_2)/(x_1 + x_2) \frac{1}{\sum_q e_q^2} \int dx' \phi_{q/h}(x') \phi_{q/A}(\tau/x)/x'. \]  

From Eq. (17), we see that \( d\Delta\langle q_T^2 \rangle/dx_F \) directly depends on the quark-gluon correlation functions. Therefore, the broadening is a very good observable for measuring quark-gluon correlation functions.

If we use LQS model for \( T_{q/A} \) given in Eq. (14), from Eq. (17), we can derive a much simpler expression for the broadening,

\[ \frac{d\Delta\langle q_T^2 \rangle}{dx_F} = \left( \frac{4\pi^2 \alpha_s}{3} \right) \lambda^2 A^{1/3} \frac{d\sigma}{dQ^2 dx_F} \int \frac{d\sigma}{dQ^2}. \]  

From Eq. (18), we can see that the Drell-Yan transverse momentum broadening should have similar \( x_F \) dependence as the differential cross section \( d\sigma/dQ^2 dx_F \) if the LQS model for \( T_{q/A}^{(I)}(x) \) is valid. Therefore, by comparing the \( x_F \) dependence of \( d\Delta\langle q_T^2 \rangle/dx_F \) and \( d\sigma/dQ^2 dx_F \), we can provide an immediate test of LQS model for the correlation functions. We emphasize that even if LQS model is not a good approximation for the quark-gluon correlation functions, measuring the \( x_F \)-dependence of the nuclear broadening provides excellent information for extracting the quark-gluon correlation functions \( T_{q/A}(x) \) directly, as shown in Eq. (17).

In the following, we use Eq. (18) to obtain the numerical estimates of the \( x_F \) dependence of the Drell-Yan transverse momentum broadening. Although the value of \( \lambda^2 \) for the
correlation function is not well determined, a different value of $\lambda^2$ corresponds to a simple adjustment to the overall normalization. Therefore, the uncertainty in the value of $\lambda^2$ should not affect our following discussions.

In obtaining our following numerical results, we use the CTEQ4L distribution as the quark distributions in the nucleus. For effective quark distributions in the nucleus, we define $q_{i/A}(x) = q_{i/p}(x)R_i(x, A)$, and use EKS98 for the parameterizations of $R_i(x, A)$ \cite{10}, which fit the data well. We choose the renormalization and factorization scale to be $\mu = Q$, and choose the incoming beam energy $p = 800$ GeV which is the energy used by the Fermilab experiments \cite{15,17}.

In Fig. 1, we plotted $d\Delta(q_{T}^2)/dx_F$ in Eq. (18) as a function of $x_F$ with $A = 184$ and $Q = 5$ GeV and 11 GeV, respectively. The dotted lines correspond to EKS98 parameterizations of effective nuclear parton distributions. In order to separate the nuclear dependence caused by multiple scattering and that caused by the effective nuclear parton distributions, we also plotted $d\Delta(q_{T}^2)/dx_F$ in solid curves with $R_i(x, A) = 1$. The difference between the solid and dotted lines is a direct consequence of the difference between nucleon and nuclear parton distributions. As shown in Eq. (17), $x_2$ represents the momentum fraction of a quark (or antiquark) from the nuclear target. At $x_F = 0$, we have $x_2 = Q/\sqrt{\hat{s}} \approx 0.13$ for $Q = 5$ GeV, and $\approx 0.28$ for $Q = 11$ GeV. It is clear that $Q = 5$ GeV and $Q = 11$ GeV cover very different range of $x_2$. For $Q = 5$ GeV, $x_F > 0$ corresponds to $x_2 < 0.1$ or corresponds to the shadowing region, while $x_F < 0$ covers the region of EMC effect. On the other hand, the Drell-Yan pairs of $Q = 11$ GeV are not sensitive to the nuclear shadowing at all. The entire range of $x_F$ values matches the range of EMC effect which include the antishadowing region ($0.1 \leq x_2 \leq 0.2$), and the EMC suppression region ($0.2 \leq x_2 \leq 0.7$), as well as the Fermi motion region for larger $x_2$. Such dependence in effective nuclear quark distributions are clearly shown in Fig. 1. In Fig. 1a, the dotted line is above the solid line in the central region of $x_F$ due to the fact that $x_2$ is in the antishadowing region. When $x_F$ is positive, the dotted line is below the solid line because $q_{i/A}(x_2)$ is now in the shadowing region. When $x_F$ is less than zero, the dotted line is again below the solid line due to the fact that $x_2$ is now in the region of EMC suppression. On the other hand, Fig. 1b shows slightly different relation between the dotted and solid line due to the fact that at $Q = 11$ GeV, $x_2$ covers a different range as shown in Fig. 1. Although the nuclear dependence in effective nuclear parton distributions change the $x_F$ dependence of the Drell-Yan transverse momentum broadening, the change (e.g., the difference between the dotted and solid lines in Fig. 1) is extremely small. Therefore, the $x_F$ dependence of the nuclear broadening, $\Delta(q_{T}^2)(x_F)$ is not very sensitive to nuclear shadowing, and can be an excellent probe of parton multiple scattering.

In addition to the transverse momentum broadening, parton multiple scattering also causes the energy loss. In recent experiments, the ratio of the cross section per nucleon from the Drell-Yan pair production in $p - A$ collisions has been used as a direct measurement of the parton energy loss in nuclear medium \cite{15}. However, as pointed out in Ref. 15, a substantial fraction of the variation in the cross section ratios versus $x_1$ comes from the shadowing of $\phi_{j/A}(x_2)$ at small $x_2$, and therefore it is difficult to extract precise information on parton energy loss from the cross section ratios. On the other hand, BDMPS showed that for a jet produced in nuclear matter, the energy loss per unit length $-\frac{dE}{ds}$ and the transverse momentum broadening $\Delta p_T^2$ has the following relation \cite{18}:
\[- \frac{dE}{dz} = \frac{\alpha_s N_c}{4} \Delta p_{\perp}^2. \] 

(19)

Similar to Eq. (19), one can use the transverse momentum broadening of the Drell-Yan pair \( \Delta \langle q_{T}^2 \rangle (x_F) \) to estimate the parton energy loss due to the initial state multiple scattering. As we demonstrated in Fig. 1, the variation in the shape of \( \Delta \langle q_{T}^2 \rangle (x_F) \) due to the shadowing is small. Therefore, \( \Delta \langle q_{T}^2 \rangle (x_F) \) is also a good probe for the parton energy loss.

Cross sections on nuclear targets are often parameterized as \( A^\alpha \) times corresponding cross sections on a nucleon target. In principle, the power \( \alpha \) can be a function of \( A \), \( q_T \), \( x_F \), as well as other physical observables. The nuclear dependence of \( \alpha \) is often used to study the effects of multiple scattering. In recent experimental studies of the \( J/\Psi \) and \( \Psi' \) suppression in \( p - A \) collisions [16], data was presented in terms of \( \alpha \) as a function of the transverse momentum \( p_T \) in different \( x_F \) regions. It was found [16] that the shapes of \( \alpha (p_T) \) are very similar in different regions of \( x_F \), and it was concluded [16] that the parton energy loss is independent of the \( c \bar{c} \) energy.

Using our result of \( \Delta \langle q_{T}^2 \rangle (x_F) \), we can also estimate the \( \alpha \) as a function of \( q_T \) in different \( x_F \) region for the Drell-Yan production. We define the parameter \( \alpha (q_T) \) for the Drell-Yan production as

\[
\frac{d\sigma_{hA}}{dQ^2 dq_{T}^2} = A^{\alpha(q_T)} \times \frac{d\sigma_{hN}}{dQ^2 dq_{T}^2}. \tag{20}
\]

If \( q_T \) is not too large, the \( q_T \) spectrum of the Drell-Yan pairs can be approximately parameterized by a Gaussian form [19],

\[
\frac{d\sigma_{hN}}{dQ^2 dq_{T}^2} \propto \frac{1}{\langle q_{T}^2 \rangle_{hN}} \exp \left[ -\frac{q_{T}^2}{\langle q_{T}^2 \rangle_{hN}} \right]; \tag{21}
\]

and

\[
\frac{1}{A} \frac{d\sigma_{hA}}{dQ^2 dq_{T}^2} \propto \frac{1}{\langle q_{T}^2 \rangle_{hA}} \exp \left[ -\frac{q_{T}^2}{\langle q_{T}^2 \rangle_{hA}} \right]. \tag{22}
\]

Substituting Eqs. (21) and (22) into Eq. (20), we derive

\[
\alpha (q_T) = 1 + \frac{1}{\ln(A)} \left[ \ln \left( \frac{1}{1 + \chi} \right) + \chi \frac{q_{T}^2}{\langle q_{T}^2 \rangle_{hN}} \right], \tag{23}
\]

where \( \chi \equiv \Delta \langle q_{T}^2 \rangle / \langle q_{T}^2 \rangle_{hN} \). In deriving Eq. (23), we used Eq. (2). In order to estimate the value of \( \alpha (q_T) \), we use the value \( \langle q_{T}^2 \rangle_{hN} = 1.2 \text{ GeV}^2 \) for the cross section per nucleon [17]. And we use Eq. (18) to integrate over different ranges of \( x_F \) value to obtain the \( \Delta \langle q_{T}^2 \rangle \) in different ranges of \( x_F \). For \( A = 184 \), we choose three different \( x_F \) ranges, which are the same as the regions used in Ref. [14], small \( x_F \) (SXF: \(-0.1 \leq x_F \leq 0.3\)), intermediate \( x_F \) (IXF: \(0.2 \leq x_F \leq 0.6\)), and large \( x_F \) (LXF: \(0.3 \leq x_F \leq 0.93\)). For these three \( x_F \) regions, we obtain the corresponding values of \( \Delta \langle q_{T}^2 \rangle \) as \([0.10, 0.043, 0.026] \text{ GeV}^2 \) for \( Q = 5 \text{ GeV} \), and \([0.078, 0.043, 0.027] \text{ GeV}^2 \) for \( Q = 11 \text{ GeV} \), respectively. It is clear that the value of \( \chi \) should be very small, and the power \( \alpha (q_T) \) in Eq. (23) can be approximated as
\[ \alpha(q_T) \approx 1 + \frac{\chi}{\ln(A)} \left[ -1 + \frac{q_T^2}{\langle q_T^2 \rangle^{hN}} \right]. \]  

From Eq. (24), we conclude that if \( q_T \) is not too large, \( \alpha(q_T) \) should show linear dependence on \( q_T^2 \) (or quadratic dependence on \( q_T \)). For the Drell-Yan production, we expect an extremely small coefficient due to a weak transverse momentum broadening (or a small \( \chi \) value in Eq. (24)).

In the above discussion, we used the same averaged value of \( \langle q_T^2 \rangle^{hN} \) for different \( x_F \) ranges. In principle \( \langle q_T^2 \rangle^{hN} \) should also have \( x_F \) dependence due to kinematics. The larger \( |x_F| \), the smaller \( \langle q_T^2 \rangle^{hN} \). Hence \( \langle q_T^2 \rangle^{hN} \) should have similar \( x_F \) dependence as \( \Delta \langle q_T^2 \rangle(x_F) \) shown in Fig. 1. Therefore, \( \chi = \Delta \langle q_T^2 \rangle / \langle q_T^2 \rangle^{hN} \) should have even smaller dependence on \( x_F \).

In Fig. 2, we plotted our predictions for the value of \( \alpha(q_T) \) defined in Eq. (23) as a function of \( q_T \) in three different \( x_F \) regions. From Fig. 2, it is clear that the shape of \( \alpha(q_T) \) in different \( x_F \) ranges are similar to what was observed in recent data on J/Ψ and Ψ' suppression [16]. Such similarity is natural because both \( q_T \) spectrum of the Drell-Yan and Charmonium production can be approximated by a Gaussian form when \( q_T \) is not too large.

In summary, we derived the \( x_F \) dependence of the Drell-Yan transverse momentum broadening in terms of quark-gluon correlation functions. We predicted that \( \Delta \langle q_T^2 \rangle(x_F) \) has a strong \( x_F \) dependence. In particular, if the LQS model for the correlation functions is valid, \( \Delta \langle q_T^2 \rangle(x_F) \) should have similar \( x_F \) dependence as the differential cross section \( d\sigma/dQ^2dx_F \). We also demonstrated that \( \Delta \langle q_T^2 \rangle(x_F) \) is a very sensitive observable to study the effects of initial state multiple scattering and the parton energy loss. In fact, \( \Delta \langle q_T^2 \rangle(x_F) \) itself is an excellent observable for extracting direct information on the nuclear quark-gluon correlation functions.

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FIG. 1. The transverse momentum broadening of the Drell-Yan pair, $d\Delta q_T^2/dx_F$, as a function of $x_F$ for 800 GeV proton beam on nuclear target $A=184$, at $Q = 5$ GeV (a) and $Q = 11$ GeV (b).

FIG. 2. The $\alpha(q_T)$ as a function of $q_T$ for small, intermediate, and large $x_F$ region, respectively. The curves are for $Q = 5$ GeV and 800 GeV beam energy.