Small-$x$ Parton Distributions of Large Hadronic Targets

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Abstract

A simple and intuitive calculation, based on the semiclassical approximation, demonstrates how the large size of a hadronic target introduces a new perturbative scale into the process of small-$x$ deep inelastic scattering. The above calculation, which is performed in the target rest frame, is compared to the McLerran-Venugopalan model for scattering off large nuclei, which has first highlighted this effect in the infinite momentum frame. It is shown that the two approaches, i.e., the rest frame based semiclassical calculation and the infinite momentum frame based McLerran-Venugopalan approach are quantitatively consistent.
1 Introduction

It has been suggested by McLerran and Venugopalan \[1\] that for very large nuclei the parton distribution functions at small $x$ are perturbatively computable. Their argument relies on the high density of partons per unit area per unit rapidity which serves as a new hard scale in the problem. Extending these ideas, an evolution equation has been formulated \[2\] to follow the growth of this potentially large density as $x$ decreases. These calculations are formulated in the infinite momentum frame of the hadron.

In an apparently quite different approach, small-$x$ deep inelastic scattering is described in the target rest frame, where the long-lived partonic fluctuations of the very energetic photon scatter off the target hadron \[3\]. Explicit formulae focusing on the $q\bar{q}$ component of the photon can be found in \[4\]. In the semiclassical framework, the scattering of the partonic fluctuations off the target is modeled by an eikonal interaction with the target colour field \[5–7\].

In the present paper, the semiclassical approach in the target rest frame is applied to the case of very large hadronic targets, such as heavy nuclei. The perturbative calculability of parton densities advertised previously emerges in a very simple and intuitive way. It is the result of the limited transverse size of the relevant $q\bar{q}$ fluctuations of the virtual photon. For an ordinary hadronic target, leading twist contributions to deep inelastic scattering come from both small and large size $q\bar{q}$ configurations. For larger targets, the $q\bar{q}$ pair has to penetrate an extended colour field and the geometric limit of the cross section is already reached for small size pairs. Accordingly, non-perturbative large size $q\bar{q}$ configurations do not contribute significantly to the total photon hadron cross section.

We then go on to compare the semiclassical approach with the McLerran-Venugopalan approach in greater detail. It is shown that they can be mapped onto each other once the non-Abelian eikonal factors characteristic of the target colour field are recognised as the basic building blocks of both methods. We confront the different perspectives inherent to the rest frame and infinite momentum frame calculation.

Performing an average over different colour field configurations, explicit formulae are obtained for the dipole cross section $\sigma(\rho)$, which describes the interaction of a $q\bar{q}$ pair of transverse size $\rho$ with the target.

Although the above results are only valid in the limit of very large targets, we believe they are useful in a much wider sense. They provide a theoretical laboratory in which, starting from a non-perturbative yet technically manageable description of the target, different small-$x$ processes can be investigated. In particular, they establish a non-trivial
limit in which the semiclassical approach can be justified in the framework of QCD.

The paper is organized as follows. After summarising the target rest frame view on deep inelastic scattering, a simple derivation of the perturbative scale associated with large targets is given in Sect. 2. The equivalence to the McLerran-Venugopalan approach is demonstrated in Sect. 3 and explicit models for the averaging over the gluon field are described in Sect. 4. The conclusions are given in Sect. 5.

2 Calculation in the target rest frame

The total cross section for the scattering of a virtual photon off a hadronic target can be obtained from the Compton scattering amplitude via the optical theorem. In the target rest frame, the leading order partonic process is the fluctuation of the photon into a $q\bar{q}$ pair (see Fig. 1). To keep calculations as simple as possible, we consider a longitudinally polarized photon coupled to scalar quarks with one unit of electric charge. As far as the $Q^2$-behaviour is concerned ($Q^2 = -q^2$), this is analogous to the standard partonic process where a transverse photon couples to spinor quarks [8].

![Figure 1: The Compton scattering amplitude within the semiclassical approach.](image)

In analogy to [4], the longitudinal cross section can be written as

$$\sigma_L = \int d^2 \varrho_\perp \sigma(\varrho) W_L(\varrho),$$

with the square of the wave function of the virtual photon given by

$$W_L(\varrho) = \frac{3\alpha_{em}}{4\pi^2} \int d\alpha N^2 K_0^2(N\varrho).$$

Here $\varrho = |\varrho_\perp|$ is the transverse size of the $q\bar{q}$ pair, $\alpha$ is the longitudinal momentum fraction of the photon carried by the antiquark, $N^2 = \alpha(1 - \alpha)Q^2$, and $K_0$ is a modified Bessel function. Note that, in contrast to [4], we have defined $W_L$ to include the integration over $\alpha$.

The cross section for the realistic case of a transverse photon and spinor quarks is obtained by substituting $K_0^2(\varrho N)$ with $2[\alpha^2 + (1 - \alpha)^2]K_1^2(\varrho N)$ in Eq. (2). None of the qualitative results derived below is affected by this substitution.
Within the semiclassical approach, the dipole cross section \( \sigma(\varrho) \) is given by

\[
\sigma(\varrho) = \frac{2}{3} \int d^2 x_\perp \text{tr} \left[ 1 - U(x_\perp)U^\dagger(x_\perp + \varrho_\perp) \right],
\]

where the SU(3) matrices \( U \) and \( U^\dagger \) represent the non-Abelian eikonal factors associated with the quark and antiquark propagating through the target colour field and averaging over all field configurations is implicit.

The functional form of \( \sigma(\varrho) \) is shown qualitatively in Fig. 2. For conventional hadrons of size \( \sim 1/\Lambda \) (where \( \Lambda \sim \Lambda_{\text{QCD}} \)) its typical features are the quadratic rise at small \( \varrho \) \( (\sigma(\varrho) \sim \varrho^2 \) with a proportionality constant \( O(1) \)) and the saturation at \( \sigma(\varrho) \sim 1/\Lambda^2 \), which occurs at \( \varrho \sim 1/\Lambda \). Consider now the idealized case of a very large target of size \( \eta/\Lambda \) with \( \eta \gg 1 \) (\( \eta \sim A^{1/3} \) for a nucleus). It is easy to see that at small \( \varrho \) the functional behaviour is given by \( \sigma(\varrho) \sim \eta^2 \varrho^2 \) while saturation has to occur at \( \sigma(\varrho) \sim \eta^2/\Lambda^2 \) for geometrical reasons. It follows that the change from quadratic rise to constant behaviour takes place at \( \varrho \sim 1/\sqrt{\eta} \Lambda \), i.e., at smaller \( \varrho \) than for conventional targets.

Figure 2: Qualitative behaviour of the function \( \sigma(\varrho) \).

From the above behaviour of \( \sigma(\varrho) \) we will now derive the dominance of small transverse distances in the convolution integral of Eq. (I). To do this, a better understanding of the function \( W_L(\varrho) \) is necessary. Recalling that \( K_0(x) \sim \ln(1/x) \) for \( x \ll 1 \) while being exponentially suppressed for \( x \gg 1 \), it is easy to see that \( W_L(\varrho) \sim Q^2 \ln^2(1/\varrho^2 Q^2) \) for \( \varrho \ll 1/Q \) and \( W_L(\varrho) \sim 1/\varrho^4 Q^2 \) for \( \varrho \gg 1/Q \). Here numerical constants and non-leading terms have been suppressed.

Under the assumption \( \Lambda^2 \ll \eta \Lambda^2 \ll Q^2 \), the integral in Eq. (I) can now be estimated by decomposing it into three regions with qualitatively different behaviour of the functions \( W_L(\varrho) \) and \( \sigma(\varrho) \),

\[
\sigma_L = \sigma_L^I + \sigma_L^{II} + \sigma_L^{III} = \left( \int_0^{1/Q^2} + \int_{1/Q^2}^{1/\eta \Lambda^2} + \int_{1/\eta \Lambda^2}^{\infty} \right) \pi d\varrho^2 \sigma(\varrho) W_L(\varrho).
\]
Of the three contributions

\[
\sigma^I_L \sim \int_0^{1/Q^2} d\bar{\rho}^2 \eta^3 \bar{\rho}^2 Q^2 \ln^2(1/\bar{\rho}^2 Q^2) \sim \frac{\eta^3}{Q^2} \\
\sigma^II_L \sim \int_1/Q^2 \eta^2 \bar{\rho}^2 \frac{1}{\bar{\rho}^4 Q^2} \sim \frac{\eta^3}{Q^2} \ln(Q^2/\eta^2) \\
\sigma^III_L \sim \int_{1/\eta^2}^{\infty} d\bar{\rho}^2 \eta^2 \frac{1}{\bar{\rho}^4 Q^2} \sim \frac{\eta^3}{Q^2}
\]

the second one dominates, giving the total cross section

\[
\sigma_L \sim \frac{\eta^3}{Q^2} \ln(Q^2/\eta^2). \tag{6}
\]

It is crucial that the third integral is dominated by contributions from its lower limit. Therefore, the overall result is not sensitive to values of \( \bar{\rho} \) that are larger than \( 1/\sqrt{\eta \Lambda} \). Phrased differently, for sufficiently large targets the transverse size of the \( q\bar{q} \) component of the photon wave function stays perturbative.

### 3 Calculation in the infinite momentum frame

The above considerations can be related to the infinite momentum frame calculation of parton distributions within the McLerran-Venugopalan model \[9\]. Here we use the term infinite momentum frame for a reference frame where the ‘+’-component of the hadron momentum and the ‘−’-component of the photon momentum are large (\( k_\pm = (k_0 \pm k_3)/\sqrt{2} \)). We start with the operator definition of scalar parton distributions

\[
q(\xi) = \frac{\xi P^2}{\pi} \int dy_+ dx_- d^2x_\perp e^{-ixP_+y} < P | T \phi^\dagger(0, x_- + y, x_\perp) U_{x_+ + y_-, x_-} \phi(0, x-, x_\perp) | P >. \tag{7}
\]

Note, that we use a time ordered product in this definition (see e.g. \[10\]) and choose to normalize the hadron state by demanding \( <P|P>=1 \). The matrix \( U \) is the standard link operator required for gauge invariance of the above definition. The corresponding space time picture for a hadron localized at \( z_- = 0 \) is illustrated in Fig. \[3\].

It is understood that only the connected diagrams contribute to the definition of \( q(\xi) \) in Eq. (7), which is equivalent to subtracting the vacuum expectation value of \( \phi^\dagger \phi \). It follows that we just need to calculate the propagator of the scalar field in the presence of the hadron and to subtract the propagator in the vacuum,

\[
q(\xi) = \frac{\xi P^2}{\pi} \int dy_+ dx_- d^2x_\perp e^{-ixP_+y} \text{tr} [U(x_\perp)G(x, x+y) - G_0(x, x+y)] . \tag{8}
\]
Figure 3: Space-time picture in the infinite momentum frame.

Here $x = (0, x_-, x_{\perp})$, $y = (0, y_-, 0)$ and $G_0$ is the conventional Feynman propagator.

At very small $x$, the contribution of the constituent quarks to the quark distribution is negligible. It is therefore sufficient to calculate the propagator in the background of a certain colour field and to average over all field configurations of the hadronic state. In a gauge in which the potential vanishes outside the hadron (say, outside the shaded region of Fig. 3) the propagating particle is simply rotated in colour space at the moment when it penetrates the field. The high energy limit makes the region with non-zero potential arbitrarily narrow in $z_{\perp}$-direction so that, for $y_- > 0$ and $x_- < 0$, the propagator can be written as

$$\begin{align*}
G(x, y) &= \int d^4z \, G_0(x, z) \left[ 2\delta(z_-)U^\dagger(z_{\perp}) \frac{\partial}{i\partial z_+} \right] G_0(z, y). 
\end{align*}$$

(9)

Note that this corresponds to the propagation of a positive energy antiquark from $x$ to $y$. The non-Abelian eikonal factor at transverse position $z_{\perp}$ is explicitly given by

$$\begin{align*}
U(z_{\perp}) &= \text{P} \exp -ig \int_{-\infty}^{+\infty} dz_- A_+(z_-, z_{\perp}),
\end{align*}$$

(10)

where $A_\mu$ is the gauge potential of the hadron. An expression analogous to Eq. (9) holds in the region $y_- < 0$ and $x_- > 0$, while for the remaining regions where $x_-$ and $y_-$ are either both positive or both negative the propagator is simply the free one. This intuitive physical picture is slightly less apparent if one uses an $A_+ = 0$ gauge as in [9]. Then the propagators contain an additional eikonal factor corresponding to the relative gauge transformation while the link operator in Eq. (7) is absent.

The expression for the propagator provided by Eq. (9) is exact for a $\delta$-function-like source. Within the present context, the highly Lorentz contracted target has a $\delta$-function-like appearance on scales larger than $1/P_+$. The use of Eq. (9) for the quark density calculation according to Eq. (8) is justified since typical values of $y_- \sim 1/\xi P_+$ are much larger than $1/P_+$. 

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Note furthermore, that the gauge potential in Eq. (11) is extremely large in the infinite momentum frame. However, this is compensated by the smallness of the relevant integration region in $z_\perp$. Clearly, the explicit formula for the eikonal factor Eq. (11) is invariant with respect to arbitrary boosts along the $z_3$-axis.

Inserting our expression for the propagator into Eq. (8) and performing the necessary integrations [9] the following result is obtained

$$\xi q(\xi) = \frac{3}{(2\pi)^4} \int_0^{\Lambda_{UV}^2} N^2 dN^2 \int d^2 q_\perp K^2_0(Nq)\sigma(q). \quad (11)$$

The dipole cross section $\sigma(q)$ (see Eq. (3)) emerges naturally due to the presence of two eikonal factors from the combination of Eqs. (7) and (9). The ultraviolet cutoff $\Lambda_{UV} \ll Q$ is introduced to suppress the contribution from partons with very high transverse momenta (compare the discussion in [11]).

Next, we want to explain in which sense Eq. (11) is equivalent to the cross section formula Eq. (1). The cross section $\sigma_L$ receives leading twist contributions from two regions, the aligned jet region, where $\alpha \ll 1$ or $(1 - \alpha) \ll 1$ and $q$ is large, and the high-$p_\perp$ region, where $\alpha \sim 1/2$ and $q$ is small. The small-$\alpha$ contribution of Eq. (1),

$$\sigma_{L,q} = \frac{3\alpha_{em}}{4\pi^2 Q^2} \int_0^{\Lambda_{UV}^2} dN^2 \int d^2 q_\perp N^2 K^2_0(Nq)\sigma(q), \quad (12)$$

corresponds, in parton model language, to the contribution from the quark distribution of the target (hence the index $q$). The latter can be identified by comparing with the parton model formula (for scalar partons)

$$\sigma_{L,q} = \frac{4\pi^2 \alpha_{em}}{Q^2} xq(x). \quad (13)$$

As expected, the obtained quark distribution is in agreement with Eq. (11).

The perturbative nature of the scattering process off large nuclei, as derived in Sect. 3 from generic features of $\sigma(q)$, is implemented within the McLerran-Venugopalan model via the colour field average implicit in Eq. (11). This is subject of the next section.

4 Explicit models for $\sigma(q)$

Consider first a particularly simple model that can be analysed with minimal calculational effort.

\footnote{Note that a sign error in Eq. (47) of [9] leads to a spurious cancellation which affects their result.}
According to Eq. (3), the expectation value of $UU^\dagger$ in the background of a hadron is required. The averaging over the field configurations of an ordinary hadron with radius $\sim 1/\Lambda$ is symbolized by $\langle \cdots \rangle_{1/\Lambda}$. For values of $x_\perp$ that correspond to a central collision and in the limit of small $\varrho$, we have

$$\left\langle U_{ij}(x_\perp)U^\dagger_{kl}(x_\perp + \varrho_\perp) \right\rangle_{1/\Lambda} \simeq \delta_{ij}\delta_{kl} - B\varrho^2\Lambda^2 T^a_{ij}T^a_{kl}. \quad (14)$$

Here $i, j, k, l$ are colour indices, the matrices $T^a$ form the conventional basis of the SU(3) Lie algebra and $B$ is a constant. The simple colour structure follows from the fact that, at small $\varrho$, the path ordered exponentials implicit in $UU^\dagger$ can be expanded in powers of the gauge field $A^a$. The leading contribution comes from $\langle A^a A^b \rangle \sim \delta_{ab}$.

To extend Eq. (14) to the case of a large hadron, an additional assumption has to be made. It is assumed that the volume of the hadron can be split into mutually uncorrelated regions of size $\sim 1/\Lambda$, as the naive geometrical picture of a large nucleus built from nucleons suggests. If the linear extension of the hadron is $\sim \eta/\Lambda$ and $n \simeq \eta$ is an integer, the $q\bar{q}$ pair passes approximately $n$ such regions on its way through the hadron (see Fig. 1). The expectation value of $UU^\dagger$ is given by

$$\left\langle U(x_\perp)U^\dagger(x_\perp + \varrho_\perp) \right\rangle_{\eta/\Lambda} \simeq \left( \left\langle U(x_\perp)U^\dagger(x_\perp + \varrho_\perp) \right\rangle_{1/\Lambda} \right)^n. \quad (15)$$

Here the appropriate contraction of the colour indices on the r.h. side is implied. Using Eq. (15) the trace of Eq. (14) can be calculated for the large hadron case. For sufficiently large $\eta$ and small $\varrho$, the product of $U$’s exponentiates giving

$$\left\langle \text{tr} \left[ U(x_\perp)U^\dagger(x_\perp + \varrho_\perp) \right] \right\rangle_{\eta/\Lambda} \simeq N_c \left( 1 - C_F B\varrho^2\Lambda^2 \right)^n \simeq N_c \exp\left[ -\eta C_F B\varrho^2\Lambda^2 \right] \quad (16)$$

with $C_F = (N_c^2 - 1)/2N_c$. Again, $x_\perp$ is chosen such as to describe a central collision. Therefore, the above is valid for an $x_\perp$ area of size $\sim \eta^2/\Lambda^2$. Inserting this into Eq. (3) the following result is obtained

$$\sigma(\varrho) \sim \frac{\eta^2}{\Lambda^2} N_c \left( 1 - \exp\left[ -\eta C_F B\varrho^2\Lambda^2 \right] \right). \quad (17)$$

Clearly, this simple model is consistent with the generic features of $\sigma(\varrho)$ discussed in Sect. 2. The above expansion in powers of the gauge potential is only valid if $\varrho \ll 1/\Lambda$. The exponentiation in Eq. (16) is valid in the limit of large $\eta$ for $\varrho \simeq 1/\sqrt{\eta}\Lambda$. Both conditions are fulfilled in deep inelastic scattering in the limit of very large hadronic targets. The simple formula of Eq. (17), which includes the numerical constant $B$ and an unknown overall normalization, provides a model for $\sigma(\varrho)$ in this limit.

Let us now consider the same quantity within the framework of [1,9]. Here, one of the central issues was to provide an explicit prescription for the averaging over the target’s
colour fields. The initial suggestion, the McLerran-Venugopalan model, used a Gaussian statistical weight for the colour charge in the target that generates these fields (compare also the discussion in [12]). Later, an evolution equation was derived to determine such a statistical weight and its \(x\) dependence [2]. These results suggest that, although a Gaussian weight is not a solution of the evolution equation, it may be used at not too small \(x\), as long as the parton densities are not too high. In the present calculation we satisfy ourselves with the simple Gaussian weight introduced in [1].

As in [1, 2], we calculate the classical field for a static, i.e. \(x_+\) independent, colour source of the form

\[
J_\mu^a(x) = \rho^a(x_-, x_\perp) g_{\mu_-},
\]  

(18)

where \(\rho\) parametrises the adjoint charge density of the hadron (not to be confused with the size \(\rho\) of the \(q\bar{q}\) pair used above). In the \(A_- = 0\) gauge the corresponding gauge field is given explicitly by

\[
\nabla_\perp^2 A_+(x_-, x_\perp) = g \rho(x_-, x_\perp); \quad A_+(x_-, x_\perp) = 0,
\]  

(19)

where \(g\) is the gauge coupling.

Note, that here we are concerned with the space time region where the potential in Eq. (11) is nonzero. From the point of view of Eq. (9), where much larger distances in \(x_-\) dominate, this corresponds to resolving the structure of the \(\delta\)-function in \(x_-\).

According to [1], the Gaussian average over \(\rho\) is implemented independently at each position in \(x_-\). This means that the average of an arbitrary functional \(O[\rho]\) is given by

\[
\langle O[\rho] \rangle = \int D[\rho] \exp \left(-\int_{\text{target}} dx_- d^2 x_\perp \frac{\rho^2}{2\mu^2} \right) O[\rho]
\]  

\[
= \int D[A_+] \exp \left(-\int_{\text{target}} dx_- d^2 x_\perp \frac{(\nabla_\perp^2 A_+)^2}{2g^2\mu^2} \right) O[\nabla_\perp^2 A_+],
\]  

(20)

where the parameter \(\mu^2 \sim \Lambda^3\) determines the typical value of the colour charge density. The above average implies a correlation function

\[
\langle A_+^a(x_-, x_\perp) A_+^b(y_-, y_\perp) \rangle = g^2 \mu^2 \delta^{ab} \delta(x_- - y_-) \gamma(x_\perp - y_\perp)
\]  

(21)

within the target. Here \(\gamma\) is the inverse of the differential operator in Eq. (20) and is given by

\[
\gamma(x_\perp) = \frac{1}{\nabla_\perp^4(x_\perp)} = \int_{\Lambda} \frac{d^2 k_\perp}{(2\pi)^2} \frac{e^{-i k_\perp x_\perp}}{(k^2_\perp)^2},
\]  

(22)
where an infrared cutoff \( \Lambda \) has been introduced to define the \( k_{\perp} \) integral. It follows that, at sufficiently small \( x_{\perp} \),

\[
\gamma(x_{\perp}) - \gamma(0) \simeq \frac{x_{\perp}^2}{8\pi} \ln(x_{\perp}^2 \Lambda^2). \tag{23}
\]

Non-logarithmic terms depend on the specific way of introducing the cutoff and have been neglected in the above equation.

Using Eqs. (20) and (21), the required average of a product of eikonal factors can be evaluated along the same lines as the gluon distribution in [13],

\[
\langle \text{tr} \left[ U(x_{\perp}) U^\dagger(x_{\perp} + \rho_{\perp}) \right] \rangle_{\eta/\Lambda} = N_c \exp \left( g^4 C_F \bar{\mu}^2 \left[ \gamma(\rho_{\perp}) - \gamma(0) \right] \right) \tag{24}
\]

where \( \bar{\mu}^2 = \int dx_{\perp} \mu^2 \sim \eta \Lambda^2 \) reflects the large linear extension \( \eta \) of the target. According to Eq. (3), an \( x_{\perp} \) integration has to be performed over a transverse area of a size \( \sim \eta^2 / \Lambda^2 \), giving rise to the formula

\[
\sigma(\rho) \sim \frac{\eta^2}{\Lambda^2} N_c \left( 1 - \exp \left[ -g^4 C_F \frac{\Lambda^2}{8\pi} \rho^2 \ln(1/\rho^2 \Lambda^2) \right] \right). \tag{25}
\]

Due to the simple model for \( \gamma \) used above, this formula is only valid for \( \rho \ll 1/\Lambda \). However, as has been discussed in Sect. 4, this is sufficient to determine the virtual photon cross section for large \( \eta \). The result shows the same Glauber type \((1 - \exp)\) structure as Eq. (17) and the gluon distributions calculated in [13, 14].

The present analysis is concerned with an energetic colour dipole created far way from the target and scattering off the target colour fields. It is the combination of the propagation of this dipole, as encoded in the Bessel functions in Eq. (4) or the Feynman propagators of Eq. (1), and its interaction with the target that determines the quark distribution. A similar analysis could be performed replacing the pair of quarks with a pair of gluons. This would result in a contribution to the gluon density beyond the instantaneous interactions considered in [13].

5 Conclusions

The scattering of a virtual photon off a large hadronic target, such as a very large nucleus, has been considered in the high energy limit. The process has been discussed using both the semiclassical target rest frame calculation and the McLerran-Venugopalan approach in the infinite momentum frame. Both approaches show that, for sufficiently large targets, the cross section is dominated by perturbative contributions.

More specifically, the cross section can be described in terms of the \( q\bar{q} \) component of the photon scattering off the target. For large targets, the process receives no lead-
ing order contributions from quark pairs with large transverse size. This means that higher Fock states of the photon are expected to be suppressed by powers of $\alpha_s$ and the semiclassical approach can be justified in QCD.

One of the main problems of the presented approach is certainly the procedure of determining the target colour fields which form the non-perturbative input of the calculation. Two simple models for the averaging over these fields have been outlined but no satisfactory QCD based procedure is presently known.

We would like to thank W. Buchmüller, A. Kovner, L. McLerran and A.H. Mueller for helpful discussions and comments.

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