Scaling Variable for Nuclear Shadowing in Deep-Inelastic Scattering

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Abstract

A new scaling variable is introduced in terms of which nuclear shadowing in deep-inelastic scattering is universal, i.e. independent of $A$, $Q^2$ and $x$. This variable can be interpreted as a measure of the number of gluons probed by the hadronic fluctuations of a virtual photon during their lifetime. The shadowing correction grows at small $x$ substantially less steeply than is suggested by the eikonal approximation. This results from the fact that shadowing is dominated by soft rather than hard interactions.
1. Introduction

A large amount of new high precision data on nuclear shadowing in deep-inelastic scattering is now available. The comparison with theory is not easy since the data are taken for a variety of nuclei at different values of $Q^2$ and $x$. One of the goals of this paper is to find a scaling variable, dependent on $Q^2$, $x$ and $A$, which makes shadowing an universal function of this variable. Our guess of the correlations between $Q^2$, $x$ and $A$ entering the scaling variable is inspired by the prejudice about the underlying QCD dynamics of shadowing.

Deep-inelastic scattering is usually interpreted in two alternative frames of reference, in the infinite momentum frame of the proton in terms of its structure function, or, in the proton rest frame in terms of hadronic fluctuations of the photon. In the first case, nuclear shadowing looks like a result of overlap in the longitudinal direction of the parton clouds originated from different bound nucleons. The latter approach is more in line with the familiar vector dominance model \[1\]. It seems to be advantageous treating shadowing at small $x$ as the consequence of the hadronic fluctuations of the photon. On the technical level our approach is similar to that of \[2\].

Prior to the detailed discussion, we introduce the scaling variable motivated by the Glauber model \[3\] as an average number of bound nucleons taking part in interacting with the hadronic fluctuation of the photon,

$$n(x, Q^2, A) = \frac{1}{4} \frac{\langle \sigma_h^2 \rangle}{\langle \sigma_h \rangle} \langle T_A \rangle F_A^2(q_L)$$  \hspace{1cm} (1)

In this expression $\sigma_h$ is the total cross section of interaction of the hadronic fluctuation $h$ of the photon with a nucleon. We choose the basis of states $h$ to be eigenstates of the interaction, rather than of the mass matrix. This eliminates the off diagonal amplitudes in the double dispersion relation \[4\] for deep-inelastic scattering. The concrete choice of the eigenstate basis is done below. The averaging in eq. (1) is weighted with the probability to find the fluctuation $h$ in the photon. $\langle T_A \rangle$ is the mean nuclear thickness function of the nucleus. The nuclear formfactor $F_A(q_L)$ depends on the longitudinal momentum transfer in the diffractive dissociation $\gamma N BhN$.

In what follows, we try to justify our choice of the scaling variable and provide model-dependent estimates of $n(x, Q^2, A)$. 

\[2\]
2. A model for $n(x, Q^2, A)$

At high energies the lifetime of the hadronic fluctuations of the photon may substantially exceed the nuclear radius, $2\nu/Q^2 \gg R_A$. In this case, the nuclear photoabsorption cross section can be represented in the same form as in the hadron-nucleus interaction [5, 6, 2],

$$
\sigma^{\gamma^* A}_{\text{tot}}(x, Q^2) = 2 \int d^2 b \left( 1 - \left[ 1 - \frac{\sigma(\rho, x)T(b)}{2A} \right]^A \right) \tag{2}
$$

Here $T(b) \approx \int_{-\infty}^{\infty} dz \rho_A(b, z)$ is the nuclear thickness function, where $\rho_A(b, z)$ is the nuclear density, which depends on the impact parameter $b$ and the longitudinal coordinate $z$. The dipole cross section $\sigma(\rho, x)$ of the interaction with a nucleon of a $q\bar{q}$ pair depends on its transverse separation $\rho$ and the energy related to $x = Q^2/2m_N\nu$. The averaging over $\rho$ and $\alpha$, the fraction of the photon light-cone momentum carried by the quark, is weighted with the photon wave function squared. One should be cautious with such a definition of the averaging [2] in the case of the photon, because its wave function, being different from the hadronic one, is not normalized to one. At small $\rho$ it has a form $|\Psi_{\gamma^*}(\rho, \alpha)|^2_T = 6\alpha_{\text{em}}/(2\pi)^2 \Sigma f^2 e^2 [1 - 2\alpha(1 - \alpha)] e^2 K_1^2(\epsilon \rho)$, where $\epsilon^2 = \alpha(1 - \alpha)Q^2 + m_q^2$. The mean transverse separation of the $q\bar{q}$ fluctuation is given by $\langle \rho^2 \rangle \propto 1/\epsilon^2$ and is of the order $1/Q^2$. However, at the endpoints of the kinematical region $\alpha$ or $1 - \alpha \sim m_q^2/Q^2$, the $q\bar{q}$ fluctuations acquire a large transverse size, $\rho^2 \sim 1/m_q^2$ [4, 5, 6]. For light quarks such a big size may substantially exceed the confinement radius, and one should put a cut off on the integration over $\alpha$. This is equivalent to a replacement of the quark mass by the cut off $\lambda$.

Expanding eq. (2), we represent the nuclear photoabsorption cross section in the form

$$
\sigma^{\gamma^* A}_{\text{tot}}(x, Q^2) = A \sigma^{\gamma^* N}_{\text{tot}}(x, Q^2) \left[ 1 - n(x, Q^2, A) + \ldots \right], \tag{3}
$$

where $n(x, Q^2, A) = \langle T(b) \rangle \langle \sigma^2 \rangle / 4 \langle \sigma \rangle$ and $\langle T(b) \rangle = (A - 1)/A^2 \int d^2 b T^2(b)$.

Note that expansion (3) looks similar to that given by the Glauber approximation [3]. In fact we use the eikonal Glauber formalism for projectile states with definite transverse dimension $\rho$, since they are the eigenstates of interaction. However, as we conclude below, after averaging over the photon wave function the first shadowing correction is of the order of $1/\lambda^2$, rather than $1/Q^2$ as expected in the Glauber model. This comparison demonstrates that, in terms of multiple scattering theory, DIS on nuclei is dominated by Gribov’s inelastic
shadowing [3], while the Glauber eikonal contribution [3] vanishes at high $Q^2$. In such a case $n(x, Q^2, A)$ should be interpreted as a measure of a number of gluons probed by the $q\bar{q}$ fluctuation of the photon, rather than a number of nucleons. This is justified at small $\rho$ since the photoabsorption cross section is proportional to $\rho^2$ and the gluon distribution function [10, 11]. The latter is not proportional to the nucleon density because of gluon fusion [12] - [15]. This effect is related to the inelastic corrections corresponding to the excitation of heavy mass intermediate states. On the other hand, if $\rho$ is not small, the $q\bar{q}$ fluctuation experiences additional shadowing interacting with gluons. This results in a rather small unitarity correction to the photoabsorption cross section on a nucleon, but in a substantial correction on nuclei. Taking into account the shadowing corrections coming from large size fluctuations, we may say that such $q\bar{q}$ fluctuations also probe the number of gluons.

In order to evaluate the lowest order nuclear shadowing correction $n(x, Q^2, A)$ in eq. (3), note that the denominator in eq. (1) is directly related to the proton structure function, $\langle \sigma(x, \rho) \rangle = 4\pi^2 \alpha_{em} F^p_2(x, Q^2)/Q^2$. We performed a fit to available data on $F^p_2(x, Q^2)$ from NMC [16], H1 [17, 18] and ZEUS [19, 20] experiments with $x \leq 0.05$ and $Q^2 \geq 0.5/GeV^2$. We used a simple parameterization motivated by the double–leading–log approximation (DLLA) for QCD evolution equations [13] (see also review [21]), $F^p_2(x, Q^2) = f(Q^2)[a \exp(2\sqrt{L})/L + b\sqrt{x}]$, where the first term corresponds to the sea- quark (Pomeron), while the second term originates from the valence - quark (Reggeons) contributions. $L = (4\pi/\beta_0) \ln(c/\alpha_s) \ln(d/x)$, where $\alpha_s = 4\pi/\beta_0 \ln(Q^2/\Lambda^2_{QCD})$ and $\Lambda_{QCD} = 0.2 GeV$, $\beta_0 = 9$ for three active flavors. The factor $f(Q^2) = Q^2/(e + Q^2)$ guarantees that the structure function vanishes in the limit of real photoabsorption. This parameterization fits the data very well ($\chi^2/d.f. = 0.8$) with parameters $a = 0.036 \pm 0.005 \ b = 0.4 \pm 0.08 \ c = 0.59 \pm 0.02 \ d = 0.31 \pm 0.07 \ e = 0.12 \pm 0.08 \ GeV^2$.

Although DLLA does not provide a Regge form of the structure functions, an effective Regge parameterization may be a good approximation, $F^p_2(x, Q^2) \propto \exp[\Delta_{eff}(Q^2)\xi]$. Here $\xi = \ln(1/x)$ and $\Delta_{eff}(Q^2) = d \ln[F^p_2(x, Q^2)]/d\xi$ correspond to the effective Pomeron intercept $\alpha_{eff} = 1 + \Delta_{eff}$. Of course, $\Delta_{eff}$ can be treated as $x$-independent only in a restricted interval of $x$. The values of $\Delta_{eff}(Q^2, x)$ corresponding to the results of our fit are depicted in fig.1 as function of $Q^2$ versus $x$. We see that $\Delta_{eff}(Q^2, x)$ is almost $x$-independent which
justifies the Regge parameterization as a good approximation in the range of $x$ and $Q^2$ under consideration.

![Figure 1: The effective Pomeron intercept as function of $Q^2$ versus $x = 10^{-2,3,4,5}$. The proton structure function $F_2^p(x, Q^2)$ is fitted to available data as is explained in the text.](image)

Remarkably, the values of $\Delta_{\text{eff}}$ in fig. 1 substantially exceed what is known from the energy dependence of the total cross sections of proton-proton interaction \cite{22, 23}, $\Delta_{\text{eff}} \approx 0.07 - 0.08$. Our results show that this distinction remains substantial down to quite low $Q^2 \sim 2 \text{ GeV}^2$. Note that the observed $Q^2$-dependence of $\Delta_{\text{eff}}$ contradicts the Pomeron factorization. This is not surprising since perturbative QCD calculations \cite{24, 25} show that the Pomeron is a more complicated singularity, a cut or a sequence of poles.

The rising $Q^2$-dependence of $\Delta_{\text{eff}}(Q^2)$ means that the $x$-dependence of the cross section $\sigma(\rho, x)$ is steeper at smaller $\rho$. Indeed, the larger $Q^2$ is, the smaller is $\langle \rho^2 \rangle \propto \ln(Q^2/\lambda^2)/Q^2$.

Despite the smallness of the mean transverse size of the photon fluctuations participating in DIS, of the order of $\sim 1/Q^2$, the shadowing terms in the expansion eq. (3) are dominated by large transverse separations in $q\bar{q}$ fluctuations. This can be argued using the relation, $\langle \rho^1 \rangle/\langle \rho^2 \rangle = 2.4/\lambda^2 \ln(Q^2/\lambda^2)$, which is obtained by the same averaging procedure as is defined in eq. (2). This relation is the manifestation of a salient feature of the hadronic fluctuations of the photon, namely, a huge dispersion of the transverse size

\[ \Delta_{\text{eff}}(x, Q^2) \]
distribution, $\langle \rho^4 \rangle \gg \langle \rho^2 \rangle^2$. Although we used the perturbative photon wave function, this conclusion has a rather general character. It is a result of the interplay of perturbative and nonperturbative contributions in the deep-inelastic cross section. The former results from the small-size fluctuations corresponding to "symmetric pairs", $\alpha \sim 1 - \alpha$. They are presented with large probability in the photon wave function, but have a small, $\sim 1/Q^2$ interaction cross section. On the contrary, the highly asymmetric fluctuations with $\alpha$ or $(1 - \alpha) \sim \lambda^2/Q^2$, have a very small, $\sim 1/Q^2$ weight in the photon wave function, but a large interaction cross section, typical for hadrons. In the case of double scattering in the nucleus the perturbative contribution turns out to be very much suppressed by the factor of $1/Q^4$, while the nonperturbative part is suppressed only once by the weight factor of $1/Q^2$. Thus, the soft interaction dominates nuclear shadowing.

Actually, just this effect is responsible for the scaling behavior of nuclear shadowing [2] and of unitarity corrections to the photoabsorption cross section on a nucleon [7].

Once the interaction responsible for shadowing is essentially soft, its $x$-dependence is governed by the soft $\Delta_{eff}(\lambda^2) \approx 0.1$, rather than the hard one. This is confirmed by the recent study of diffractive dissociation by the H1 collaboration [27], which claimed $\Delta_P = 0.1 \pm 0.03 \pm 0.04$.

To proceed further with the calculation of $n(x, Q^2, A)$, note that in eqs. (2), (3) we temporarily used an assumption that the photon energy $\nu = Q^2/2m_Nx$ in the nuclear rest frame is sufficiently high to make the lifetime of the photon fluctuation long compared with the nuclear size, so that it propagates through the whole nucleus with a frozen intrinsic separation $\rho$. However, most of the data available are in the transition region of $x$, where the lifetime, usually called coherence time, is comparable with the nuclear radius. The finite coherence time can be taken into account by introducing a phase shift between $q\bar{q}$ wave packets produced at different longitudinal coordinates, in the same way as for inelastic corrections [28], or in the vector dominance model [1]. The mean nuclear thickness function of eq. (3) should be replaced by an effective one,

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This conclusion is in variance with the statement in [26] that at high $Q^2$ the unitarity corrections vanish as $\sim 1/Q^2$ and one sees the single Pomeron exchange.
\[ \langle \tilde{T}(b) \rangle = \frac{A - 1}{A^2} \int d^2 b \left[ \int_{-\infty}^{\infty} dz \, \rho_A(b, z) \, e^{iqz} \right]^2 \approx \langle T(b) \rangle \, F_A^2(q_L) . \]  

Here \( F_A(q_L) = \exp(-q_L^2 R_A^2 / 6) \) is the nuclear longitudinal formfactor and \( R_A \) is the mean nuclear radius. For the sake of simplicity we use the Gaussian form for the nuclear density which is quite precise for the calculation of \( F_A(q_L) \). Calculating \( \langle T \rangle \) we use the realistic parameterization of nuclear density [29].

The decrease of the effective nuclear thickness function \( \langle \tilde{T}(b) \rangle \) at large \( q_L \) can be interpreted as a result of shortness of the hadronic fluctuation path in the nucleus, if we are in the nuclear rest frame, or as an incomplete overlap of the gluon clouds of the nucleons which have the same impact parameter in the infinite momentum frame of the nucleus [30, 31].

In order to calculate the longitudinal momentum transfer in DIS, \( q_L = (Q^2 + M^2) / 2\nu \), one needs to know the effective mass of the produced \( q\bar{q} \) wave packet. However, a \( q\bar{q} \) state with definite separation \( \rho \) does not have a definite mass. This is a typical problem for those who work in the eigenstate basis of interaction. We evaluate \( q_L \approx 2xm_N \) assuming \( M^2 \sim Q^2 \). Thus, the parameter which controls the value of \( \langle \tilde{T}(b) \rangle \) is \( x \).

Now we are in a position to estimate the nuclear shadowing correction per nucleon \( n(x, Q^2, A) \) in eq. (3), which reads

\[ n(x, Q^2, A) = \frac{1}{4} \frac{N}{F_2^p(x, Q^2)} \langle T(b) \rangle F_A^2(q_L) \left( \frac{1}{x} \right)^{2\Delta_{eff}(x^2)} . \]  

The scaling variable \( n(x, Q^2, A) \) can be interpreted as a measure of the amount of those gluons which take part in the interaction with the \( q\bar{q} \) fluctuation during its lifetime. Eq. (5) is a model-dependent realization of the general expression (4).

3. Comparison with the data

The variable \( n(x, Q^2, A) \) has been calculated from eq. (5) using the values of \( x, Q^2 \) and \( A \) corresponding to data from the NMC experiment [32, 33] as well as the results of our fit to \( F_2^p(x, Q^2) \). The expected scaling dependence of the nuclear shadowing on \( n(x, Q^2, A) \) is not affected by the overall normalization factor \( N \). However, in order to have a correct slope of \( n \)-dependence corresponding to eq. (3) we choose \( N = 3 \, GeV^{-2} \). The data on the ratio of the nuclear and nucleonic photoabsorption cross sections \( R_{A/N}(x, Q^2) \) is plotted against
$n(x, Q^2, A)$ in Fig.2. They demonstrate a good scaling in $n(x, Q^2, A)$ within a few percent accuracy.

![Graph showing data on nuclear shadowing at small $x$ from the NMC experiments versus the scaling variable $n(x, Q^2, A)$ as defined in eq. (5). The slope of the straight line corresponds to the normalization factor in eq. (5) $N = 3 \text{ GeV}^{-2}$.]

We should comment more on the procedure of the calculation of $n(x, Q^2, A)$:

(i) Our considerations are valid only for small $x$, so we limit the $x$-region to $x < 0.07$. At larger $x$, the nuclear structure functions show a small enhancement of a few percent relative to the proton one, which results in $R_{A/N}(x, Q^2) > 1$ for $n(x, Q^2, A)B_0$. A plausible assumption is that about the same antishadowing correction extends down to smaller $x$, where it is compensated by stronger shadowing effects. Such a behavior, for instance, is expected in the model of swelling bound nucleons. The antishadowing effect may have some $A$ dependence, what would cause a small, a few percent relative shift of the data in fig. 2 corresponding to different nuclei, but will not change the slope of $n$-dependence. Since the physics of antishadowing is beyond the scope of our present consideration, and
the effect is numerically very small, we do not try to incorporate with it, but just have renormalized the solid line \( R_{A/N}(n) = 1 - n \) in fig. 2 by 3% up to make the comparison easier.

(ii) The data points [33] for \( Q^2 < 0.5 \text{ GeV}^2 \) were excluded from the analysis because they are in the realm of the vector dominance model, rather than DIS. They should correspond to the same nuclear shadowing as is experienced by the \( \rho \)-meson. This is the reason for the saturation of nuclear shadowing at small \( x \), claimed in [35, 33].

(iii) A further important observation is that \( R_{A/N}(x, Q^2) \) depends to a good accuracy linearly upon \( n(x, Q^2, A) \) at least for \( n < 0.2 \). The higher order terms in the expansion in eq. (3) are expected to violate the linearity in \( n \). This implies that those terms are small. A model-dependent evaluation of the next shadowing correction shows that it is small indeed.

(iv) There are two contributions to the shadowing in DIS [12, 13, 14], one comes from the suppression of the gluon density as a consequence of gluon fusion \( gg \beta g \), which corresponds to the triple Pomeron graph in the framework of standard Regge phenomenology. Another contribution to the shadowing comes from the Glauber-like rescattering of the \( q\bar{q} \) fluctuation off gluons. The latter mechanism, which was mostly under consideration above, can also be viewed upon as a parton fusion, but as a fusion of gluons into a \( q\bar{q} \) pair, \( gg\beta q\bar{q} \). In the Regge-model this process corresponds to the Pomeron-Pomeron-Reggeon graph. Both mechanisms lead to the same form of the variable \( n(x, Q^2, A) \) in eq. (3). However, the formfactor \( F_A(q_L) \) has a different form for the triple-Pomeron mechanism due to the contribution of heavier hadronic fluctuations of the photon, \( F_A(q_L) \propto Ei(-x^2m_N^2R_A^2/6) \). Here \( Ei \) is the integral exponential function. We checked that in the \( x \) and \( A \) domain investigated, the admixture of the triple-Pomeron mechanism does not affect the \( n(x, Q^2, A) \)-scaling within the error bars of the data available. It may, however, cause a deviation from the scaling for heavy nuclei.

We hope that forthcoming high-statistic data on heavy nuclei from the NMC Collaboration may help to disentangle these two mechanisms of shadowing. This is important if one wants to predict the unitarity corrections to the proton structure function at small \( x \) or the photon diffractive dissociation cross section, because the admixture of the triple-Pomeron affects the normalization constant \( N \) in eq. (5).

4. Summary
Starting from the QCD dynamics of deep-inelastic scattering at small $x$ we have found a new variable $n(x, Q^2, A)$ which scales all available data on nuclear shadowing in DIS at small-$x$. This variable measures the number of gluons which a hadronic fluctuation of the virtual photon interacts with during its lifetime.

An important observation is also that shadowing corrections at small $x$ and large $Q^2$ grow less steeply than $[F^p_2(x, Q^2)]^2$. This is because nuclear shadowing is a subject to soft rather than hard physics.

The observed scaling of nuclear shadowing as function of $n(x, Q^2, A)$ supports our assumptions on the dynamics of nuclear shadowing.

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