SOFT COMPONENT OF HARD REACTIONS
AND NUCLEAR SHADOWING
(DIS, Drell-Yan reaction, heavy quark production)∗

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Abstract

We consider deep-inelastic lepton scattering, Drell-Yan lepton pair and heavy quark production in the reference frame of the target. These reactions traditionally treated as hard have, however, a substantial leading twist soft contribution. One of the manifestations of such a soft component, nuclear shadowing, is overviewed.

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1 Introduction

A parton-model interpretation of a high-energy reaction depends upon a reference frame. For instance, the shadowing in the nuclear structure function measured in deep-inelastic scattering (DIS) at small $x$ looks in the infinite momentum frame of the nucleus like a fusion of the parton clouds of the bound nucleons [1]. On the other hand, one can view this shadowing as inelastic corrections [2] in the rest frame of the nucleus (see for example [3]). To be convinced that they are the same, one can compare these mechanisms of shadowing either with the Feynman graphs, or with the Reggeon diagrams. Only observables may be Lorentz invariant, but not a space-time interpretation.

Another example is the Drell-Yan (DY) reaction of lepton pair production [4]. The very structure of the DY formula led to a wide spread opinion that this mechanism corresponds to instantaneous lepton pair production in quark-antiquark annihilation independently of a reference frame. According to the factorization theorem the cross section of $\ell\bar{\ell}$ pair production is proportional to the subprocess cross section, $q\bar{q} \rightarrow \ell\bar{\ell}$, times the hadronic distribution functions which are the probabilities to find a quark and an antiquark with definite momenta and virtualities in the colliding hadrons. From this point of view the nuclear suppression of lepton pair production can be treated as a shadowing in the nuclear structure function.

However, even the very statement that the annihilating quark and antiquark belong to the beam and to the target respectively (or vise versa) is not Lorentz invariant. For instance, in the target rest frame both $q$ and $\bar{q}$, as well as the lepton pair, should be considered as a beam hadron fluctuation, related to the beam distribution function, and vise versa in the beam rest frame. Correspondingly, the DY process can be treated as an electromagnetic bremsstrahlung of heavy photons by the beam or the target, depending on a reference frame. Different partonic interpretations of DY reaction in different reference frames correspond, nevertheless, to the same Feynman graphs.

Heavy quark production is also treated as a hard reaction, subject to perturbative QCD [5]. However, similar to the DY case, one is free to choose any reference frame depending
on convenience. Correspondingly, the partonic interpretation varies: the subprocess looks like $q\bar{q}$ annihilation or gluon fusion in the rest frame of the $Q\bar{Q}$ pair, however, it should be treated as a heavy gluon bremsstrahlung, looking from the infinite momentum frame of the $Q\bar{Q}$.

In this paper we consider all these processes in the rest frame of the target and use the light-cone representation for the projectile wave function, which allows to take explicitly into account the color-coherence effects, missed in the standard perturbative consideration. We come to the conclusion that these processes have much in common, and all of them have a substantial soft component, which scales in $Q^2$ and manifests itself particularly in nuclear shadowing. Bjorken and Kogut were first who claimed existence of such a component in DIS and its importance for the $Q^2$ scaling \[6\]. A hadronic fluctuation of the virtual photon have to have a low intrinsic transverse momentum to interact softly. Such a fluctuation produces jets aligned along the photon momentum. This idea was discussed in the framework of QCD in \[7\] and got its explicit formulation in the line-cone wave functions formalism \[8\]. Analogous small-$k_T$ component of projectile fluctuations was found in DY reaction and heavy quark production in \[9\], and a light-cone formalism was developed in \[10, 11\].

We overview briefly in what follows the light-cone wave function formalism for the above hard reactions, emphasizing manifestations of the soft component.

2 Scaling variable for nuclear shadowing

In spite of the fact that DIS probes only the quark distribution function, there is wide spread believe that shadowing in the gluon distribution function is approximately the same as observed in DIS. This could be true, if electromagnetic probe of gluon distribution were really hard. However this is not the case. Indeed, a virtual photon participates in strong interaction through its hadronic fluctuations. Assuming that the transverse separation $\rho$ of the $q\bar{q}$ fluctuation of a highly virtual photon is small, one can use the property of color transparency of the dipole cross section \[12\] $\sigma(\rho) \propto \rho^2$. If this is the case, the total photoabsorption cross
section § is related to the gluon distribution function \( xG_N(x, Q^2) = g_N(x, Q^2) \). 

\[
\sigma_{\text{tot}}^N(x, Q^2) = \int_0^1 d\alpha \int d^2\rho |\Psi_{qq}(\rho, \alpha)|^2 \sigma(\rho, x) \approx \frac{\pi^2}{3} \alpha_s(\rho) \langle \rho^2 \rangle g_N(x, Q^2) 
\]

(1)

Here \( \alpha \) is the relative fraction of the photon light-cone momentum carried by the quark. Hereafter we neglect the small contribution of the longitudinal component of the photon wave function and the quark mass unless it is important. The wave function reads \[8\],

\[
|\Psi_{qq}(\rho, \alpha)|^2 \approx \frac{6\alpha_{em}}{(2\pi)^2} \sum_i^{N_f} Z_i^2 [1 - 2\alpha(1 - \alpha)] \epsilon^2 K_1^2(\epsilon \rho),
\]

(2)

where \( K_1(z) \) is modified Bessel function,

\[
\epsilon^2 = \alpha(1 - \alpha)Q^2 + m_q^2
\]

(3)

According to (2), (3) the mean \( q - \bar{q} \) transverse separation squared, \( \rho^2 \propto 1/\epsilon^2 \), is of the order of \( 1/Q^2 \), except the edges of the kinematical region \( \alpha \) or \( 1 - \alpha \sim m_q^2/Q^2 \), where the \( q\bar{q} \) fluctuation acquires a large transverse separation, \( \rho^2 \sim 1/m_q^2 \) (actually, the quark mass should be replaced by an effective cut off of the order of \( \Lambda_{QCD} \)). The presence of this soft component in the \( q\bar{q} \) fluctuation of the photon, makes questionable the validity of the electromagnetic probe for the gluon distribution. Indeed, if one needs, for example, to figure out how many nucleons are in a nucleus, and one uses, say, a pion-nucleus inelastic interaction as a probe, one obviously gets a wrong result, namely, the number of nucleons \( \sim A^{2/3} \). The source of the trouble is the softness of the probe, \( \sigma_{\text{tot}}(\pi N) \) is too large. To get the correct answer one should use a probe with a sufficiently small cross section, \( \sigma \ll 1/R_A\rho_A \), where \( R_A \) and \( \rho_A \) are the radius and the density of the nucleus.

The same problem arises for probing the gluon distribution by a photon, due to asymmetric \( q\bar{q} \) fluctuations with large transverse separations, which experience shadowing, interacting with a cloud of gluons. Such shadowing looks like a suppression of the photoabsorption cross section on the nucleus, but it still does not mean that the gluon density at small \( x \) is reduced in the nucleus. One should disentangle a nuclear suppression of gluon density at small \( x \) and a regular, Glauber-like shadowing of the photon fluctuations propagating
through the nucleus. Both, however, depend on the density of gluons which originate from different nucleons and overlap.

Let us test universality of the relation between shadowing and the gluon density \cite{15} on a sample of available data on DIS on nuclei. We start with the general expression for the photoabsorption cross section on a nucleus \cite{16, 12, 8}, which takes into account all the inelastic shadowing corrections \cite{2}.

\[
\sigma_{\gamma^*A}^{\gamma A}(x, Q^2) = 2 \int d^2b \left\langle 1 - \left[ 1 - \frac{\sigma(\rho, x) T(b)}{2A} \right]^A \right\rangle
\]

(4)

The partial amplitude of $q\bar{q}$ elastic scattering is averaged over the eigenstates of interaction, which have a definite transverse sizes. In the case of photoabsorption the averaging is weighted with the photon wave function squared (2), in analogy to (1). Nuclear thickness $T(b) \approx \int_{-\infty}^{\infty} dz \rho_A(b, z)$.

Making use of expansion in r.h.s of (4) one can represent the nuclear shadowing effects in the form,

\[
R_{A/N}^{DIS}(x, Q^2) = \frac{\sigma_{\gamma^*A}^{\gamma A}(x, Q^2)}{A \sigma_{\gamma^*N}^{\gamma N}(x, Q^2)} = 1 - \frac{1}{4} \frac{\langle \sigma^2(\rho, x) \rangle}{\langle \sigma(\rho, x) \rangle} \langle T(b) \rangle + ...,
\]

(5)

were $\langle T(b) \rangle = (1 - 1/A) / A \int d^2b T^2(b)$. The cross section of interaction of a $q\bar{q}$ fluctuation with a nucleon in can be represented in the form \cite{17}, $\langle \sigma(\rho, x) \rangle \approx \langle \sigma(\rho) \rangle / x^{\Delta_P(Q^2)}$, where $\Delta_P(Q^2)$ measured at HERA \cite{22, 23} reaches a large value $0.3 \div 0.4$ at high $Q^2$. On the other hand, $\langle \sigma^2(\rho, x) \rangle$ is dominated by the soft asymmetric fluctuations, and its $x$-dependence is governed by the soft Pomeron, $\langle \sigma^2(\rho, x) \rangle \approx \langle \sigma^2(\rho) \rangle / x^{2\Delta_{sof}(Q^2)}$, where $\Delta_{sof} \approx 0.1$. Correspondingly, $\langle \sigma^2(\rho) \rangle$ has a leading twist behaviour $\propto 1/Q^2$. Therefore, one can write

\[
\frac{\langle \sigma^2(\rho, x) \rangle}{\langle \sigma(\rho, x) \rangle} = \frac{N}{F_2^p(x, Q^2)} \left( \frac{1}{x} \right)^{2\Delta_{sof}},
\]

(6)

where N is a free parameter.

An important ingredient of formula (4) is the assumption that $x$ is sufficiently small, $x \ll 1/m_N R_A$, which guarantees that all the parton clouds of nucleons with the same impact parameter overlap in the antilaboratory frame. This is the same as to say that the lifetime
of the photon fluctuation is much longer than the nuclear radius. However, most of available data are in the region of $x$, where this condition is broken. Finiteness of the lifetime, which is usually called coherence time, can be taken into account taking into account the phase shifts between $q\bar{q}$ wave packets produced at different longitudinal coordinates, which suppress the effective nuclear thickness in (3),

$$\langle \tilde{T}(b) \rangle = \frac{1}{A} \int d^2 b \left[ \int_{-\infty}^{\infty} dz \, \rho_A(b, z) \, e^{i q z} \right]^2 \quad (7)$$

Here $q = (Q^2 + M^2) / 2 \nu$ is the longitudinal momentum transfer in the diffractive photo-production of a hadronic state with mass $M$.

Available data on nuclear shadowing are taken at different values of $x$ and $Q^2$, even within the same experiment, what makes it difficult to compare with theoretical calculations. Since we expect the shadowing effects to be dependent only on the amount of overlapping gluons, the data should be plotted against a new variable $n(x, Q^2)$,

$$n(x, Q^2) = \frac{N}{4 F^2_F(x, Q^2) \langle \tilde{T}(b) \rangle \left( \frac{1}{x} \right)^{2 \Delta_{soft}}} \quad (8)$$

which has to be calculated for each experimental point.

Although one may have a plausible guess about the value of parameter $N$, calculation is quite ambiguous. Instead, we intend to test the scaling of nuclear shadowing versus new variable $n(x, Q^2)$, which is $N$-independent.

We calculated $n(x, Q^2)$ for each point of data [18, 19] for $R^{DIS}_{A/N}(x, Q^2)$ and plotted them against this variable in fig. 1. The data demonstrate excellent scaling in $n(x, Q^2)$. The absolute value of shadowing fixes the parameter $N \approx 3 \text{ GeV}^{-2}$.

Notice that the data with $Q^2 \ll m^2_\rho$ is subject to vector dominance model, they expose the same nuclear shadowing as real photons, independently on $x$. This is the reason of saturation of nuclear shadowing at low $x$, claimed in [20, 13]. We excluded the data points [19] with $Q^2 < 0.5 \text{ GeV}^2$ from the analyses.
Figure 1: Normalized by $1/A$ nucleus-to-nucleon ratio of the structure functions plotted versus scaling variable $n(x, Q^2)$, defined in (8). The data are from [18, 19].

3 Nuclear shadowing of Drell-Yan lepton pairs

The Born diagrams of perturbative QCD describing the quark (antiquark)-nucleon interaction with radiation of a heavy photon, converting into a lepton pair of mass $M$, are shown in fig. 2.

Figure 2: Born diagrams for lepton pair production.
We are interested in moderate values of \( x_1 \) and small \( x_2 \ll 1 \), where \( x_{1,2} = \pm x_F/2 + \sqrt{x_F^2/4 + M^2}/s \), what corresponds to gluonic exchanges. The calculation of these diagrams in impact parameter representation \[10\] leads to the following expression for the cross section of lepton pair production in \( hN \) interaction,

\[
M^2 \frac{d\sigma_{hN}^{hN}}{dM^2dx_1} = \int_{x_1}^1 \alpha F_q^h x_1 x_1 \int d^2 \rho |\Psi_{q\bar{l}}(\alpha, \rho)|^2 \sigma_{q\bar{l}}(q \rightarrow q\bar{l}) .
\]

(9)

Here \( F_q^h(z) \) is the quark distribution function of the hadron.

\[
|\Psi_{q\bar{l}}(\alpha, \rho)|^2 \approx \frac{(Z_q\alpha_{em})^2}{\pi^2} \text{Im} \Pi(M^2) \alpha^2 K_1(M^2) \rho^2 \tau^2,
\]

(10)

is the light-cone wave function squared of the \( q\bar{l} \) Fock component of the projectile quark in the mixed \( \rho - \alpha \) representation, where \( \rho \) is the transverse separation between \( q \) and the center of gravity of the \( l\bar{l}\)-pair, and \( \alpha \) the fraction of the projectile quark light-cone momentum carried by the \( l\bar{l} \) pair. The value of \( \Pi(M^2) \approx 1/12\pi \) was calculated in \[9\].

There is a close similarity between wave function \[10\] and that for the \( q\bar{q} \) fluctuation of a photon \[2\]. The parameter \( \tau \) is also close to \( \epsilon \), defined in \[3\], \( \tau^2 = (1 - \alpha)M^2 + \alpha^2 m_q^2 \).

Most surprisingly, the cross section of freeing of the \( l\bar{l} \)-pair turns out to be the same dipole interaction cross section of a colorless \( q\bar{q} \) pair with a separation \( \alpha\rho \), \( \sigma_{q\bar{l}}(q \rightarrow q\bar{l}) = \sigma(\alpha\rho, x_2) \).

This has a natural interpretation \[10\]. The color screening factor \( 1 - \exp(i\vec{q}\vec{r}) \) in the dipole interaction cross section \( \sigma(\rho) \) \[12\] originates from the difference in the impact parameters of gluon attachments to \( q \) or \( \bar{q} \). The same is valid for the graphs in fig. 2. The quark trajectories before and after the photon radiation have different impact parameters. The \( q\gamma^* \) fluctuation with transverse separation \( \rho \) has a center of gravity, which coincides with the impact parameter of the parent quark. It is separated by \( (1 - \alpha)\rho \) from the \( \gamma^* \) and \( \alpha\rho \) from the \( q \). Thus, the diagrams in fig. 2 have relative phase factor \( \exp(i\alpha\rho) \) and different signs, what leads to the dipole form of the cross section.

The cross section of DY pair production off a nuclear target can be written in analogy to \[3\],

\[8\]
\[ M^2 \frac{d\sigma_{\text{DY}}^{bA}}{dM^2 dx_1} = 2 \int d^2 b \left\langle 1 - \left[ 1 - \sigma(q \rightarrow q\bar{l}\bar{l}) \frac{T(b)}{2A} \right]^A \right\rangle, \]  

where averaging over \( \alpha \) and \( \rho \) is weighted in the same way as in (9).

The corrections for finiteness of the lifetime of the \( q\bar{l} \) fluctuation can be done in the same way as in the case of DIS. The coherence time turns out to be about the same \([\text{10}]\), \( l_c = 1/q \approx 1/2x_2m_N \).

Keeping in (11) the first-order shadowing correction we find for the nucleus-to-nucleon ratio of DY cross sections

\[ R_{A/N}^{\text{DY}}(x_F, M^2) \approx 1 - n[x_2, M^2(1 - x_1)] \]  

where \( n(x, Q^2) \) is defined in (8).

Comparing (12) with (5) we see that the factorization theorem should be modified: DIS and DY structure function are to be compared at \( Q^2 = M^2(1 - x_1) \). The reason for this correction is clear: the average \( q\gamma^* \) fluctuation becomes more and more asymmetric at \( x_1 \rightarrow 1 \), since integration over \( \alpha \) in (9) goes from \( x_1 \) to 1. Thus, the projectile fluctuations in DY reaction are more asymmetric, i.e. are softer, than in DIS. This provides a deviation from factorization and should be taken into account.

As soon as parameter \( N \) is fixed by the data on nuclear shadowing in DIS, one can predict the nuclear suppression of DY lepton pairs. Comparison with the results of the E772 collaboration \([\text{24}]\) at 800 GeV shown in fig. 3 demonstrates a good agreement.

Note that the increase of nuclear shadowing towards \( x_F = 1 \) is mainly due to the increase of the fluctuation lifetime \( t_c \), similar to DIS.

4 Heavy quark production

Analogous to DY reaction the standard approach to heavy quark production is based on the factorization theorem and uses the quark and gluon distribution functions of the colliding hadrons. Heavy quark pair production may also be treated, similar to the DY lepton
pairs, as a bremsstrahlung in the infinite momentum frame. In this case we have the same diagrams as in fig. 2, with replacement of $\gamma^* \rightarrow l\bar{l}$ by $g^* \rightarrow Q\bar{Q}$, and on top of that a few new ones shown in fig. 4. These new Feynman graphs correspond to direct interaction of the virtual gluon and heavy quarks with the target. Comparing with the standard approach one may say that those diagrams on fig. 2 correspond to $q\bar{q}$ annihilation, while ones in fig. 4 correspond to gluon fusion mechanism. The vertex $gg \rightarrow Q\bar{Q}$ in fig. 4a is decomposed explicitly in fig. 4b.

The important result of the previous section, the cross section of lepton pair production in form (9), can be naturally extended to a general case of a virtual decay of a parton $a \rightarrow b + c$, where $a, b, c$ may be colored or colorless [11]. The cross section of freeing the fluctuation $bc$, which has a transverse separation $\vec{\rho}$ and relative fraction $\alpha$ of the longitudinal momentum carried by $b$, equals to the dipole total cross section of interaction of a colorless system $\bar{a}bc$ with relative transverse separations, $\vec{\rho}$ ($bc$), $(1 - \alpha)\vec{\rho}$ ($b\bar{a}$) and $\alpha\vec{\rho}$ ($\bar{a}c$): $\sigma(a \rightarrow bc) = \sigma_{\bar{a}bc}[\rho, (1 - \alpha)\rho, \alpha\rho]$

The color screening provides the infra-red stability of the freeing cross section, even if $a$
Figure 4: The additional to Fig. 2 Born diagrams for heavy quark production.

is colored. The cases of $\gamma^* \rightarrow q\bar{q}$ and $q \rightarrow q\gamma^*$ were discussed above. Below we encounter new examples, like $q \rightarrow qg^*$ and $g \rightarrow Q\bar{Q}$ fluctuations.

This mnemonic rule applied to the whole set of diagrams in fig. 3, leads to the $Q\bar{Q}$ production cross section, which has a structure similar to that for DY reaction (9),

$$
M^2 \frac{d\sigma_{hN}}{dM^2dx_1} = \int_{x_1}^1 d\alpha F_q^b \frac{x_1}{\alpha} \int d^2\rho \int d^2r |\Psi_{qg^*}(\alpha, \rho)|^2 \times
\int_0^{\beta_{\text{max}}} d\beta |\Psi_{Q\bar{Q}}(\beta, r)|^2 \left[ \sigma(q \rightarrow qg) + \sigma(g \rightarrow Q\bar{Q}) \right],
$$

(13)

$$
\sigma(q \rightarrow qg) = \sigma_{qq}[\rho, (1 - \alpha)\rho, \alpha\rho],
$$

$$
\sigma(q \rightarrow Q\bar{Q}) = \sigma_{qq}[\beta r, (1 - \beta)r, r].
$$

(14)

Here, $r$ is the transverse separation of the $Q\bar{Q}$ pair; $\beta$ is a share of the gluon longitudinal momentum carried by $Q$. The total cross section for the interaction of a colorless $gqq$ system with a nucleon can be expressed in terms of $q\bar{q}$ dipole cross sections $[25]$, $\sigma_{gqq}(\rho_1, \rho_2, r) = 9/8 \{\sigma_{qq}(\rho_1) + \sigma_{qq}[(1 - y)\rho_2]\} - 1/8\sigma(r)$.

The $qg^*$ Fock state wave function squared, $|\Psi_{qg^*}(\tau, \rho)|^2$, has the same form as eq. (14), except for a replacement in the normalization, $(Z_q\alpha_{em})^2 \rightarrow (\alpha_s/2N_c)^2$. We denote by $g^*$
a point-like $Q\bar{Q}$ pair in the color-octet point-like state, which is indistinguishable from a gluon. The interaction can resolve such a pair only if it has a separation, i.e., a color dipole moment.

The squared wave function $|\Psi_{Q\bar{Q}}(\beta, r)|^2$ of the $Q\bar{Q}$ fluctuation of a gluon is equal to the photon’s one (2) up to a simple change of the normalization.

Interaction with the spectator light quark has a substantial soft component at $1 - \alpha \sim m_q^2/M^2$. This gives a contribution to (13) which is $\propto F_{gh}^q(x_1)/M^2 \ln[(1 - x_1)M^2/m_q^2]$.

The contribution of the direct interaction was considered also in [26]. It is suppressed by a factor $1/M^2_{Q\bar{Q}}$, since $\langle r^2 \rangle \sim 1/m_{Q\bar{Q}}^2$. This contribution to cross section (13) is $\propto (1 - x_1)F_{gh}^q(x_1)/M^2$ and is small compared with the spectator quark interaction.

The soft interaction component under discussion provides a substantial nuclear suppression, similar to that in DY reaction. If the lifetime of the $qQ\bar{Q}$ fluctuation is long in comparison with the nuclear radius, the nuclear cross section in frozen approximation reads similar to (11)

$$M^2 \frac{d\sigma^{hA}_{Q\bar{Q}}}{dM^2 dx_1} = 2 \int d^2b \left(1 - \left\{1 - \left[\sigma(q \rightarrow qg) + \sigma(g \rightarrow Q\bar{Q})\right] \frac{T(b)}{2A}\right\}\right)^A,$$

The averaging corresponds to the integration over $\rho, r, \alpha$ and $\beta$ in (13). Analogous to DIS and DY reaction this nuclear shadowing of heavy quark production scales with $M^2$.

Typical value of the predicted effect for charm production is about $20 - 30\%$ for heavy nuclei. We avoid here comparison with available data on open charm hadron production, since it needs to include into consideration the hadronization stage. This is expected to rearrange substantially the produced hadron momenta [11].

Summarizing, there is a substantial soft contribution to DIS, DY pair and heavy quark production, which is a leading twist effect. It has a common origin in all reactions, namely, highly asymmetric in longitudinal momentum projectile hadronic fluctuations are subject to color opacity. These soft components lead to nuclear shadowing, diffractive production. It causes a deviation from the factorization. In this mini-review we treated all these reactions in the same approach and found a close similarity between them.
References

[1] O.V. Kancheli, Pisma ZHETF, **18** (1973) 469

[2] V.N. Gribov, Sov. Phys. JETP **29** (1969) 483

[3] Yu.L. Dokshitzer, V.A. Khoze, A.H. Mueller and S.I. Troyan, 'Basics of Perturbative QCD', Editions Frontiéres, 1991

[4] S.D. Drell and T.M. Yan, Phys. Rev. Lett. **25** (1970) 316

[5] P. Nason, S. Dawson and R.K. Ellis, Nucl. Phys. **B303** (1988) 607; **B327** (1989) 49

[6] J.D. Bjorken and J. Kogut, Phys. Rev. **D8** (1973) 1341

[7] L.L. Frankfurt and M.I. Strikman, Phys. Rept. **160** (1988) 235

[8] N.N. Nikolaev and B.G. Zakharov, Z. Phys. **C49** (1991) 607

[9] S.J. Brodsky et al., Nucl. Phys. B **369** (1992) 519

[10] O. Benhar, B.Z. Kopeliovich and A. Zieminski, 'Soft component of Drell-Yan mechanism of lepton pair production', paper in preparation.

[11] J. Hüfner and B.Z. Kopeliovich, 'Color coherence in hadroproduction of heavy quarks from nucleons and nuclei’, paper in preparation.

[12] Al.B. Zamolodchikov, B.Z. Kopeliovich and L.I. Lapidus, JETP Lett. **33** (1981) 612.

[13] V. Barone et al., Z. Phys. **C58** (1993) 541

[14] B. Blättel et al., Phys. Rev. Lett. **71** (1993) 896

[15] B.Z. Kopeliovich and B. Povh, 'What can we learn from nuclear shadowing about the proton structure function at small x?', paper in preparation

[16] B.Z. Kopeliovich and L.I. Lapidus, Sov. Phys. JETP Lett. (1978)
[17] E.A. Kuraev, L.N. Lipatov and V.S. Fadin, Sov. Phys. JETP 44 (1976) 443; 45 (1977) 199; Ya.Ya. Balitskii and L.I. Lipatov, Sov. J. Nucl. Phys. 28 (1978) 822; L.N. Lipatov, Sov. Phys. JETP 63 (1986) 904

[18] CERN NMC, P. Amaudruz et al., Z. Phys. C51 (1991) 387

[19] CERN NMC, M. Arneodo et al., submitted to Nucl. Phys.

[20] FNAL E665, M.R. Adams et al., Phys. Lett. B287 (1992) 375

[21] B.Z. Kopeliovich and B.G. Zakharov, Phys. Rev. D44 (1991) 3466

[22] DESY ZEUS, M. Derrick et al., DESY 94-143, August 1994

[23] DESY H1, G. Rädel, H1-10/94-390, 1994

[24] D. M. Alde et al, Phys. Rev. Lett. 64 (1990) 2479.

[25] N.N. Nikolaev and B.G. Zakharov, Jülich preprint, KFA-IKP(Th)-1993-17

[26] N.N. Nikolaev, G. Piller and B.G. Zakharov, Jülich preprint, KFA-IKP(TH)-1994-43