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COHERENCE EFFECTS AND AVERAGE MULTIPLICITY
IN DEEP INELASTIC SCATTERING AT SMALL x

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ABSTRACT

The average hadron multiplicity in deep inelastic scattering at small x is calculated in this paper. Its relationship with the average multiplicity in e^+e^- annihilation is established. As shown the results do not depend on a choice of the gauge vector. The important rôle of coherence effects in both space-like and time-like jet evolution is clarified.

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1. Deep inelastic scattering (DIS) is a classical example of a process proceeding at small distances, where quarks and gluons are relevant objects. Perturbative QCD has been successfully applied to the description of the structure function of this process. It is clear, however, that the study of final hadron states in DIS promises to give much more information on both nucleon structure and multiple hadron production at small distances.

Nevertheless, up to now the techniques of perturbative QCD were mainly used for the calculation of inclusive spectra and multiplicities in another hard process, namely e^+e^- annihilation.

Multiple production in DIS was considered for the first time in Refs. [1]. Later on [2] we studied in more detail final hadron states in DIS at fixed values of the Bjorken variable x . As has been shown in Ref. [2] there is a transparent relation between average hadron multiplicity in DIS and that in e^+e^- annihilation.

The aim of the present paper is to study multiple production in DIS at small x . According to Ref. [2], the average multiplicity in DIS $\langle N \rangle^{\text{DIS}}$ is formed from the average multiplicity of the particles produced in parton subprocess $\langle \hat{N} \rangle$ and the average number of nucleon remnant fragments $\langle N_0 \rangle$ (provided colour compensation is soft). For fixed x the quantity $\hat{N}$ has been calculated within the framework of perturbative QCD in Refs. [1]. In the present paper we extend the relation between $\hat{N}$ and $N_{e^+e^-}$ to the range of small x . The transition from $\hat{N}$ to $\langle N \rangle^{\text{DIS}}$ is also discussed.

Note that in perturbative QCD it is not the average number of observed hadrons which is directly calculated but the multiplicity of virtual colour partons. Nothing is known a priori about the gauge dependence of such a quantity. That is why in the paper we pay careful attention to the gauge properties of $\hat{N}$.

In our treatment we exploit coherence effects in both time-like and space-like jet evolution to a considerable extent.

2. The inclusive distributions of soft partons is mostly studied for time-like evolution, say, the fragmentation of massive quark and gluon jets in e^+e^- annihilation.

The account of interference diagrams in the double logarithmic approximation leads to non-trivial results. Namely, interference graphs give a destructive contribution to the inclusive spectrum at small $\bar{x}$, which completely compensates the contribution from the greater part of the phase space of the ladder graphs [3]. In

order to obtain a correct result at small $\bar{x}$ it is sufficient to sum only ladder diagrams, where the angles of the successive emissions $\bar{\theta}_i$ strongly decrease (Fig. 1)

$$\bar{\theta}_1 \gg \bar{\theta}_2 \gg \dots \quad (1)$$

The angular ordering (1) reflects the fact that large angle emission becomes coherent [3].

Let us consider in more detail the evolution of the initial quark with momentum k in e^+e^- annihilation (Fig. 1). In multiple production its rôle is to "trigger" massive gluon jets. Let $n_g(r^2, \bar{Q}_0^2)$ be the number density for the parton with virtuality $\bar{Q}_0^2$ in the gluon jet with fixed virtuality r^2 . n_g depends on invariant variables. In order to take into account the coherence effects in an equation for n_g one should recast the restriction (1) in terms of virtualities. Let $z = \ell_+/r_+$, then the condition (1) is equivalent to the following one

$$z \gg r^2/k^2. \quad (2)$$

Inequality (2) is in addition to the kinematical one

$$z \geq \ell^2/r^2. \quad (3)$$

The average multiplicity in the gluon jet, whose virtuality varies up to k^2 , is connected with n_g through the relation

$$N_g(k^2, \bar{Q}_0^2) = \int_{\bar{Q}_0^2}^{k^2} \frac{dr^2}{r^2} n_g(r^2, \bar{Q}_0^2). \quad (4)$$

It is N_g that defines the average multiplicity in e^+e^- annihilation (see below).

In what follows we will use axial gauge

$$n_\mu A_\mu = 0. \quad (5)$$

In such a gauge we can regard quarks and gluons in perturbative QCD as partons and hard processes (e^+e^- annihilation, DIS) admit a probabilistic (parton) interpretation.

Now let us show that the multiplicity in a gluon jet emitted in a physical process (say, off an initial quark in e^+e^- annihilation) does not depend on a choice of gauge vector n_μ .

Consider the rather general case

$$n_\mu = (n_+, n_-, 0), \quad n_\pm \neq 0, \infty. \quad (6)$$

Taking account of the inequalities (2) and (3) leads to the equation [4]

$$n_g(r^2, \bar{Q}_0^2) = r^2 \delta(r^2 - \bar{Q}_0^2) + \frac{d_s(r^2)}{2\pi} \int_{\bar{Q}_0^2}^{r^2} \frac{d\ell^2}{\ell^2} \times \int_{\max(\ell^2/r^2, r^2/k^2)}^1 dz \hat{P}_g^{gg}(z, n) n_g(\ell^2, \bar{Q}_0^2). \quad (7)$$

The non-regularized splitting function $\hat{P}_g^{gg}$ has a singularity which depends on the gauge vector

$$\hat{P}_g^{gg}(z, n) \Big|_{z \rightarrow 0} \simeq 2C_A / (z + \Delta), \quad (8)$$

where

$$\Delta = r^2 n_+ / k^2 n_-. \quad (9)$$

By comparing (8) and (9) with (2) we convince ourselves that the quantity Δ can be omitted in the expression for $\hat{P}_g^{gg}$ (8). In other words, the n_- dependence of n_g disappears provided relevant coherence effects in time-like jet evolution are taken into account.

The multiplicity in the gluon jet emitted by the initial antiquark in e^+e^- annihilation is treated analogously (after changing $n_\pm \rightarrow n_\mp$). It can be shown that the expression for n_g remains the same in light-cone gauge ($n_\pm = 0$) as well [4].

Finally, the relation between average multiplicity in e^+e^- annihilation and that in the gluon jet looks like

$$N_{e^+e^-}(Q^2, \bar{Q}_0^2) = 2C_F \int_{\bar{Q}_0^2}^{Q^2} \frac{dk^2}{k^2} \int_{\bar{Q}_0^2}^{k^2} \frac{d\tilde{k}^2}{\tilde{k}^2} \frac{d_s(\tilde{k}^2)}{2\pi} N_g(\tilde{k}^2, \bar{Q}_0^2). \quad (10)$$

We see that $N_{e^+e^-}$ does not depend on n_μ . The formula (10) will be used below.

3. A few words should be said about the structure function of DIS. At small x it is defined by a gluon cascade. In the approximation

$$1 \ll \ln 1/x, \quad \alpha_s(Q^2) \ln 1/x \lesssim 1 \quad (11)$$

a ladder-type diagram with a strong $k_{\perp}$ -ordering contributes to D_g , where $k_{\perp}$ means transverse momenta along the ladder [5]. Notice that the $k_{\perp}$ -ordering is different from angular ordering due to the occurrence of small fractions of space-like momenta.

In the range (11) leading powers of both $\ln Q^2/Q_0^2$ and $\ln 1/x$ should be kept. This leads to the following formula (the coupling constant α_s is fixed) [5]

$$\begin{aligned} x D_g(Q^2, Q_0^2, x) &= \delta(1-x) + \sum_{m=1}^{\infty} \bar{\alpha}_s^m \\ &\times \sum_{k=1}^{m+1} \frac{1}{(2m-k)!} (\ln 1/x)^{2m-k} \frac{1}{(k-1)!} (\ln Q^2/Q_0^2)^{k-1}. \end{aligned} \quad (12)$$

where $\bar{\alpha}_s \equiv \alpha_s C_A/\pi$. The expression in the right-hand side of (12) can be presented in the form

$$\begin{aligned} x D_g(Q^2, Q_0^2, x) &= \delta(1-x) + \left(\frac{\ln 1/x}{\bar{\alpha}_s \ln Q^2/Q_0^2} \right)^{-1/2} \\ &\times \left[I_1 \left(2 \sqrt{\bar{\alpha}_s \ln Q^2/Q_0^2 \ln 1/x} \right) + \sum_{m=1}^{\infty} \left(\frac{\ln 1/x}{\ln Q^2/Q_0^2} \right)^m \right. \\ &\left. \times I_{2m-1} \left(2 \sqrt{\bar{\alpha}_s \ln Q^2/Q_0^2 \ln 1/x} \right) \right], \end{aligned} \quad (13)$$

where I_n is a modified Bessel function. Thus the following asymptotic relation takes place at small x

$$D_g \simeq D_g^{\text{LLA}} \sum_{m=0}^{\infty} \left(\frac{\ln 1/x}{\ln Q^2/Q_0^2} \right)^m, \quad (14)$$

where D^{LLA} means the gluon structure function obtained in leading logarithmic approximation (LLA).

4. Recently, the parton inclusive distribution in DIS at small x has been studied in Refs. [6a]^{*}). It appears to be given by a "jet-calculus" diagram [7] with $k_{\perp}$ -ordering in the propagators and ϕ -ordering at the vertices, where ϕ is the angle of the emitted soft parton with respect to the jet axis (Fig. 2) [6a]. Let us underline that ϕ -ordering

$$\phi_1 \gg \phi_2 \gg \dots \quad (15)$$

and θ -ordering in space-like parton evolution differ.

For angular distributions an additional vertex factor related to virtual corrections arises. It seems to disappear in integrated y -distribution [6a].

So, there are two types of coherence in multiple production in DIS at small x . Firstly, the emission of jets by a space-like parton becomes coherent at large ϕ_i . Secondly, large $\bar{\theta}_i$ -angle emission in the evolution of these time-like jets shows coherence as well (see Fig. 2)

$$\begin{aligned} \phi_1 &\gg \bar{\theta}_1^1 \gg \bar{\theta}_2^1 \gg \dots \\ \phi_2 &\gg \bar{\theta}_1^2 \gg \bar{\theta}_2^2 \gg \dots \\ &\dots \dots \dots \end{aligned} \quad (16)$$

ϕ -ordering is equivalent to the restriction

$$1-y \gg \ell^2/k^2, \quad (17)$$

while the kinematical restriction looks like

$$1-y \geq \bar{Q}_0^2/|\ell^2|. \quad (18)$$

In order to calculate $\hat{N}$ we use the light-cone gauge

$$n_{\mu} = (1, 0, 0). \quad (19)$$

In this gauge the emission off a struck quark is suppressed. At $1 \ll \ln 1/x \ll \ln Q^2/Q_0^2$ the leading contribution to $\hat{N}$ is given by the "jet-calculus" diagram pictured in Fig. 3a. The corresponding analytical expression for $\hat{N}_a$ is of the form

^{*}) This point has been briefly discussed also in Ref. [6b].

$$D_g^q(Q^2, Q_0^2, x) \hat{N}_a = \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \int_x^1 \frac{dz}{z} \frac{\alpha_s(k^2(1-z))}{2\pi} \times \hat{P}_q^{qg}(z) D_g^q(k^2/z, Q_0^2, x/z) N_g(|k^2|(1-z)). \quad (20)$$

The main contribution to the integral in z comes from large z . This circumstance enables one to simplify the expression (20). Taking into account the relation between N_g and $N_{e^+e^-}$ (10) we obtain

$$D_g^q(Q^2, Q_0^2, x) \hat{N}_a \simeq \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} D_g^q(k^2, Q_0^2, x) \times \frac{\partial}{\partial \ln k^2} N_{e^+e^-}(|k^2|). \quad (21)$$

Now let us turn to the contributions which rise with decreasing x . One of them, connected with quark jet fragments, is presented by the diagram in Fig. 3b. The treatment is the same as in the previous case. Omitting the details of the calculation we show the final formula

$$D_g^q(Q^2, Q_0^2, x) \hat{N}_b \simeq \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \frac{\partial}{\partial \ln k^2} D_g^q(k^2, Q_0^2, x) \times N_{e^+e^-}(|k^2|). \quad (22)$$

It immediately follows from (21) and (22) that

$$\hat{N}_a + \hat{N}_b = N_{e^+e^-}(W^2). \quad (23)$$

The equality (23) holds up to the corrections $O(\sqrt{\alpha_s(Q^2)})$ and $O((\ln 1/x)^{-\frac{1}{2}})$.

There is another contribution to $\hat{N}$ rising with decreasing x (see Fig. 3c). The multiplicity of gluon jet fragments in Fig. 3c is given by the formula

$$\begin{aligned}
 D_g^q(Q^2, Q_0^2, x) \hat{N}_c^1 &= \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} \int_{x/z}^1 \frac{dz}{z} P_g^{q\bar{q}}(z) \\
 &\times \int_{Q_0^2}^{1/k^2} \frac{dl^2}{l^2} \int_{x/z}^1 \frac{dy}{y} \frac{\alpha_s(l^2(1-y))}{2\pi} \hat{P}_g^{gg}(y) D_g^g(l^2/y, Q_0^2, x/zy) \\
 &\times N_g(l^2(1-y)). \quad (24)
 \end{aligned}$$

The restrictions (17) and (18) are assumed to be imposed on the integration over the variable y .

It is convenient to subdivide the splitting function $\hat{P}_g^{gg}(z)$ into three parts

$$\hat{P}_g^{gg} = 2C_A (P_1 + P_2 + P_3), \quad (25)$$

where $P_1(z) = 1/(1-z)$, $P_2(z) = 1/z$, $P_3(z) = z(1-z)-2$. Let $\hat{N}_c^{i1}$ ($i = 1, 2, 3$) denote a contribution to $\hat{N}_c^1$ corresponding to P_i . Then one finds [rejecting the terms proportional to $\sqrt{\alpha_s(Q^2)}$ and $(\ln 1/x)^{-\frac{1}{2}}$]

$$\begin{aligned}
 D_g^q(Q^2, Q_0^2, x) \hat{N}_c^{11} &= \frac{C_A}{2C_F} \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} \frac{\partial}{\partial \xi_k} D_g^q(k^2, Q_0^2, x) \\
 &\times N_{e^+e^-}(1k^2) - \\
 &- \frac{3C_A}{4C_F} \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} \int_{Q_0^2}^{1/k^2} \frac{dl^2}{l^2} \frac{\alpha_s(l^2)}{2\pi} \frac{\partial^2}{\partial \xi_l^2} D_g^q(l^2, Q_0^2, x) \\
 &\times N_{e^+e^-}(1l^2), \quad (26)
 \end{aligned}$$

$$\begin{aligned}
 D_g^q(Q^2, Q_0^2, x) \hat{N}_c^{21} &= \frac{C_A}{2C_F} \int_{Q_0^2}^{W^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} \int_{Q_0^2}^{1/k^2} \frac{dl^2}{l^2} \frac{\alpha_s(l^2)}{2\pi} \\
 &\times \frac{\partial^2}{\partial \xi_l^2} D_g^q(l^2, Q_0^2, x) N_{e^+e^-}(1l^2).
 \end{aligned}$$

Here

$$\bar{\xi}_Q = \int_{Q_0^2}^{Q^2} \frac{dk^2}{k^2} \frac{d_3(k^2)}{2\pi} \quad (27)$$

is a QCD evolution parameter.

As for $\hat{N}_c^3$ it appears to be an $O(\sqrt{\alpha_s(Q^2)})$ correction with respect to $\hat{N}_c^{1,2}$ and, therefore, can be ignored in the sum $\hat{N}_c = \sum_i \hat{N}_c^i$.

In order to estimate $\hat{N}_c$ at $1 \ll \ln 1/x \ll \ln Q^2/Q_0^2$, we use the analytical expressions for D_g^q and $N_{e^+e^-}$ obtained in perturbative QCD. As follows from the results of Section 3 (14), it is sufficient to choose the LLA formula for D_g^q . For N_g we need the expression derived by taking into account coherence effects in gluon jet evolution (see, for instance, Ref. [8]).

Everything said above results in the relation

$$\hat{N}_c = N_{e^+e^-}(W^2) \frac{C_A}{2C_F} \left[\varepsilon - \frac{3}{2} \varepsilon^2 + O(\varepsilon^3) \right], \quad (28)$$

where

$$\varepsilon = \sqrt{\frac{d_3(Q^2)}{4\pi} \frac{\ln 1/x}{\bar{\xi}_Q}}. \quad (29)$$

As for particle emissions off the lower part of the gluon ladder their total contributions to $\hat{N}$ is suppressed at least by the factor $\sqrt{\alpha_s(Q^2)}$. So, we arrive at the following formula for the average multiplicity in DIS off a parton at small x

$$\begin{aligned} \hat{N} \equiv \hat{N}_a + \hat{N}_b + \hat{N}_c \simeq N_{e^+e^-}(W^2) \left[1 + \frac{C_A}{2C_F} \left(\varepsilon - \right. \right. \\ \left. \left. - \frac{3}{2} \varepsilon^2 \right) \right]. \end{aligned} \quad (30)$$

5. The expression (30) has been derived in light-cone gauge. Let us show that it does not depend on a choice of gauge vector. Consider, say, temporal gauge

$$\eta_\mu = (1, 1, 0). \quad (31)$$

In it the following relations are valid

$$\begin{aligned} \hat{N}^{s,q} &= 1/2 N_{e^+e^-}(W^2), \\ \hat{N}_a + \hat{N}_b &= 1/2 N_{e^+e^-}(W^2), \end{aligned} \quad (32)$$

where $\hat{N}^{s,q}$ denotes the average multiplicity of struck quark fragments.

Now let us check that $\hat{N}_c$ remains the same when we pass from light-cone gauge to temporal gauge. In the gauge (31) the equation for N_c has the form (24) where $\hat{P}_g^{gg}(y)$ is replaced by $\hat{P}_g^{gg}(y, \tilde{\Delta})$. It is important that N_g in (24) does not depend on n_μ (see Section 2).

The function $\hat{P}_g^{gg}(y, \tilde{\Delta})$ has the following singularity

$$\left. \hat{P}_g^{gg}(y, \tilde{\Delta}) \right|_{y \rightarrow 1} \simeq 2C_A / (1 - y + \tilde{\Delta}), \quad (33)$$

where

$$\tilde{\Delta} = |l^2|/W^2. \quad (34)$$

In light-cone gauge $\tilde{\Delta}$ was equal to 0. Now remember the dynamical restriction (17). It is seen to overlap the singularity of $\hat{P}_g^{gg}$ (33). Thus, the quantity $\tilde{\Delta}$ can be omitted in (33). This in turn means that equation for $\hat{N}_c$ remains the same in temporal gauge and so does the total multiplicity $\hat{N}$ (30). Analogous arguments can be given in the more general case (6).

We underline that in our treatment we exploited coherence effects in both space-like evolution [ϕ -ordering in gluon jet emission (15)] and time-like evolution [$\bar{\theta}$ -ordering in gluon jet fragmentation (1)].

The average multiplicity of partons moving along the current ($\hat{N}_F$) and along the target ($\hat{N}_B$) are also gauge invariant quantities (for fixed x this has been shown in Ref. [9]). Nevertheless, the picture of final state production in DIS

depends crucially on a gauge. For instance, in light-cone gauge (19) one can no longer interpret $\hat{N}_F$ as a result of the struck quark fragmentation.

Up to now we were considering multiple production of virtual partons. It is natural to suppose that a parton transition into hadrons is the same for both DIS and e^+e^- annihilation. Then the relation between $\hat{N}$ and $N_{e^+e^-}$ obtained above (30) can be regarded as the relation which connects hadron multiplicities in two hard processes.

Recently, an attempt to calculate $\hat{N}$ has been made in Ref. [10]. The authors started from an undressed gluon ladder diagram for the structure function and considered all possible insets of a real soft gluon in such a diagram. Later on this gluon was regarded to be a massive gluon jet. This approach leads to an overestimation of $\hat{N}$. Namely, the $O(\epsilon^2)$ correction obtained in Ref. [10] was positive, while the correct $O(\epsilon^2)$ term has the opposite sign and differs significantly in absolute value (28).

Average hadron multiplicity in leptonproduction events has been also considered in a Monte Carlo parton shower model [11]. Note that the model does not incorporate the coherence in space-like cascade.

6. In conclusion, we will discuss real lepton-hadron scattering. Everything mentioned above concerned mainly multiple hadron production in DIS off a parton (gluon). In the case of lepton scattering off a hadron (at small x) the corresponding formula for average hadron multiplicity looks like [2]

$$\begin{aligned} \langle N \rangle^{\text{DIS}}(W^2, Q^2) &= \int_x^1 \frac{dz}{z} w(Q^2, x, z) \\ &\times \left[\hat{N}(\hat{W}^2) + N_0(M_0^2) \right] = \langle \hat{N} \rangle + \langle N_0 \rangle, \end{aligned} \quad (35)$$

where $\langle \dots \rangle$ denotes the average with weight

$$\begin{aligned} w(Q^2, x, z) &= D_g^q(Q^2, Q_0^2, x/z) f_g(Q_0^2, z) \\ &\times \left[\int_x^1 \frac{dz}{z} D_g^q(Q^2, Q_0^2, x/z) f_g(Q_0^2, z) \right]^{-1} \end{aligned} \quad (36)$$

and f_g is the distribution function of a gluon with virtuality Q_0^2 inside the target

hadron.

The quantity $\hat{N}$ has been calculated in Section 4. It depends on the effective energy W available for partial production in parton subprocess

$$\hat{W}^2 = (z-x)/(1-x) W^2 \equiv k W^2. \quad (37)$$

The average multiplicity of hadron remnant fragments N_0 is some phenomenological function of its invariant mass $M_0^2 = (1-z)(Q_0^2/z + M_h^2)$ where M_h is a target mass.

$w(z)$ tends rapidly to 0 at $z \rightarrow 1$ and the average value of z is close to x . Therefore, the contributions to $\langle N \rangle^{\text{DIS}}$ which rise with decreasing x become smaller after taking into account the target hadron [$\ln 1/x$ is replaced by $\ln \tilde{z}/x$ in the expression (29)]. The efficiency factor k which stands in front of W^2 in (37) is much less than 1. That is why the growth of $\langle N \rangle^{\text{DIS}}(W, Q^2)$ in W at fixed Q^2 is expected to be delayed in comparison with e^+e^- annihilation multiplicity growth.

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Fig. 1

Fig. 2

Fig. 3a

Fig. 3b

Fig. 3c