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SCALE VIOLATIONS IN DEEP INELASTIC LEPTON-HADRON PROCESSES AND THE POWER LAW IN LARGE P_T HADRON COLLISIONS

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Нарушение масштабности в глубоконеупругих лептон-адронных процессах и степенной закон в адронных столкновениях при больших P_T

Рассмотрена связь нарушений бьеркеновского скейлинга в глубоконеупругих лептон-адронных процессах и характера степенного поведения сечений образования адронов при больших поперечных импульсах. В рамках предположения о кварковой структуре адронов получена аналитическая зависимость инклюзивных сечений от параметров нарушения в глубоконеупругой области. Анализируются условия согласования наблюдаемой Q^2 -зависимости структурных функций адронов и поведения инклюзивных сечений, определяемого размерным кварковым анализом в различных областях изменения x_T, θ -переменных.

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Matveev V.A., Slepchenko L.A., Tavkhelidze A.N. **E2 - 11580**

Scale Violations in Deep Inelastic Lepton-Hadron Processes and the Power Law in Large P_T Hadron Collisions

Connection between violation of Bjorken scaling in deep inelastic lepton scattering and power law behaviour in hadronic high P_T processes is considered. In the framework of quark-parton model an analytic expression of inclusive cross section in terms of the scale breaking parameters is obtained. The consistency conditions of the scaling deviations and the power behaviour of hadronic cross sections given by the quark counting rules are analyzed in different regions of x_T, θ -variables.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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1. INTRODUCTION

As is known, the dimensional quark counting (DQC) gives the following asymptotics of the exclusive binary processes and e.m. form factors [1]:

$$\frac{d\sigma}{dt}(AB\to CD)\sim s^{-(n_A+n_B+n_C+n_D-2)}\cdot f(t/s);\; F_A(t)\sim|t|^{-n_A+1}, \quad (1.1)$$

$-n_A, n_B, \dots,$ are the minimal number of valence constituents of the hadrons $A, B, \dots$

Derivation of the corresponding quark counting rules for the high P_T inclusive particle spectra is related to consideration of some additional assumptions about the origin of hadronic constituents and dynamical mechanisms of their interaction.

Assuming the validity of DQC, the inclusive cross section for large P_T production has the following form [2-4]

$$E\,d\sigma/d^3p(AB\to C+X) = \sum_{A'B'C'D'} P_T^{-N(A'B'C'D')} f(x_T, \theta), \quad (1.2)$$

where according to eq. (1.1) exponent $N(A'B'C'D')$ depends on the number of hadron constituents taking part in the large angle exclusive scattering. When $A'B'C'D'$ are quark states DQC gives P_T^{-4} law in accordance with parton model result [5,6]. In the case when $A'B'C'D'$ permit the existence of states with hadronic quantum numbers, DQC leads to the values $N\geq 8$, ref. [3,7]. Thus, the dynamical assumptions play the crucial role in deriving the power laws for the cross sections at large transverse momenta.

It is well known that due to the hadron structure the effective hadronic interaction is not a local one. This leads to the difference in the description of deep inelastic lepton-hadron and hadron-hadron processes. As was shown in ref.[8], however, in the asymptotic regime of large s and $t(Q^2)$ the hadron-hadron interaction can be effectively local and inclusive cross sections can be expressed in terms of the hadron structure functions $F_H(x)$, similar to those observed in deep inelastic lepton processes, i.e., $F_H(x) \sim F_2(x) \equiv \nu W_2(x)$ and

$$E d\sigma/d^3 p(AB \to C+X) = \frac{1}{\pi} \sum \iint_{\min} F_A(x,Q^2) F_B(y,Q^2) \times$$

$$\times \frac{d\hat{\sigma}}{dt}(\hat{s},\hat{t}) z^{-1} D_C(z,Q^2) \frac{dxdy}{xy}, \tag{1.3}$$

where

$$z = x_1/x + x_2/y, \quad x_1 = -u/s = -x_T/2 \operatorname{ctg}\theta/2, \quad x_2 = -t/s = -x_T/2 \operatorname{tg}\theta/2$$

$$x_T = 2P_T/\sqrt{s}, \quad s = P_T^2 x_1 \cdot x_2, \quad Q^2 = 2x/x_T z \cdot P_T^2.$$

Therefore, if the origin of the high P_T hadrons is the hard scattering of quarks and Bjorken scale breaking is defined in a satisfactory way* (one can extract this information from deep inelastic lepton experiments[11-13]), the observed experimental results on high transverse momentum scattering can be explained as a certain screening of the canonical P_T^{-4} quark behaviour. See, particularly refs.[14-20]. Notice, that in this case P_T^{-8} behaviour does not play a role of the fundamental law[3,21] like as it was discussed in CIM[7]. Moreover, going beyond F.F. assumptions, the power law exponent N is not necessarily 8 for meson, and 12* for baryon production, but depends crucially[22] on the experimental conditions, i.e., on the regions of the measured $P_T(x_T)$ and θ-range and on the corresponding nonscaling regions in Q^2, x_{Bj} -variables. Furthermore, nonscaling behaviour in lepton processes is enhanced in going to the hadron high P_T scattering. Taking into account the connection $Q^2 \to P_T^2$, a proper defining of the leading contributions in integral (3) is of great importance. We stress this point in connection with some nonphysical extrapolations made in refs.[14,15].

The present article is organized as follows. In Sec. 2 we derive some analytical expressions for the high P_T-hadron production cross section in terms of the scale breaking parameters. In Secs. 3 and 4 we give a brief discussion of some asymptotic results in different regions of x_T. The detailed analysis of the x_T, θ -dependence of exponent $N-4 \approx \epsilon(x_T, \theta)$, as well as asymptotic predictions will be published elsewhere.

2. Assuming the validity of equation (1.3), suppose that the main dynamical mechanism leading to the high P_T hadrons is the hard constituent elastic scattering. In this case, most general form of the quark cross section is

$$\frac{d\hat{\sigma}}{dt}(ab \to cd) \approx c_\sigma \overline{\alpha_s^2} \cdot \hat{s}^{-N} f(x_1, x_2) =$$

$$= \bar{\alpha}_s^2 c_\sigma P_T^{-4} (x_1/xz)^{N-\sigma_1} (x_2/yz)^{N-\sigma_2}. \tag{2.1}$$

Inserting (2.1) into (1.3) and making the suitable change of variables, we get:

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*A scheme for the scaling violation $F_2(x,Q^2)$ can be obtained within AF gauge theories[9,10].

* Note, that this number has a meaning only in constituent interchange model[2], see in this connection ref.[22].

$$\Sigma^{AB\to C}(P_T,x_T,\theta)\equiv E\frac{d\sigma}{d^3p}(AB\to C) =$$

$$=\frac{c_\sigma\bar{\alpha}_s^2}{\pi}P_T^{-4}\int_{x_1}^{1-x_2}dz_1\int_{x_2}^{1-z_1}dz_2\; z_1^{1-\sigma_1}\cdot z_2^{1-\sigma_2}\times$$

$$\times F_A\Big(\frac{x_1}{z_1},\frac{2x_1P_T^2}{x_Tz_1(z_1+z_2)}\Big)F_B\Big(\frac{x_2}{z_2},\frac{2x_1P_T^2}{x_Tz_2(z_1+z_2)}\Big)\times$$

$$\times D_C\Big(z_1+z_2,\frac{2x_1p_T^2}{x_Tz_1(z_1+z_2)}\Big)(z_1+z_2)^{-1}. \qquad (2.2)$$

We shall use the following parametrization for the quark distribution and decay functions

$$F_H(x,Q^2)=F_H(x,Q_0^2)\cdot\Phi_H(Q^2/Q_0^2,x),\; H=A,B$$

$$D_C(z,Q^2)=D_C(z,Q_0^2)\cdot\Phi_C(Q^2/Q_0^2,z), \qquad (2.3)$$

where the functions Φ ϕ are specified by the scaling deviations in deep inelastic lepton scattering and

$$F_H(x,Q_0^2)=c_Hx^h(1-x)^H, \quad H=A,B$$

$$D_C(z,Q_0^2)=d_Cz^\alpha(1-z)^C \qquad h=\alpha,\beta \qquad (2.4)$$

satisfying the sum rules

$$c_H^{-1}=\int_0^1F_H(x,Q_0^2)dx, \qquad d_C^{-1}=\int_0^1D_C(z,Q_0^2)dz. \qquad (2.5)$$

Then, the invariant cross section (2.2) can be represented in the form

$$\Sigma^{AB\to C}(P_T,x_T,\theta)=C\,P_T^{-4}\;I(P_T,x_T,\theta)x_1^\alpha x_2^\beta,$$

$$C=\bar{\alpha}_s^2/\pi\,c_\sigma c_A c_B d_C$$

$$I(P_T,x_T,\theta)=\int_{x_1}^{1-x_2}dz_1\int_{x_2}^{1-z_1}dz_2\,z_1^{1-\sigma_1-\alpha}z_2^{1-\sigma_2-\beta}(z_1-x_1)^A\times$$

$$\times(z_2-x_2)^B(1-z_1-z_2)^C(z_1+z_2)^{\sigma_1+\sigma_2+\gamma-1}\mathcal{P}(z_1z_2), \qquad (2.6)$$

and the function $\mathcal{P}(Q^2,z_1,z_2)$ defines the screening of the P_T^{-4} law due to the nonscaling behaviour of $F(x,Q^2)$, $D(z,Q^2)$. We have:

$$\mathcal{P}(Q^2/Q_0^2,z_1,z_2)=\Phi_A\Big(\frac{2x_1P_T^2/Q_0^2}{x_Tz_1(z_1+z_2)},\frac{x_1}{z_1}\Big)\Phi_B\Big(\frac{2x_1P_T^2/Q_0^2}{x_Tz_1(z_1+z_2)},\frac{x_2}{z_2}\Big)\times$$

$$\times\phi_C\Big(\frac{2x_1P_T^2/Q_0^2}{x_Tz_1(z_1+z_2)},z_1+z_2\Big), \qquad \mathcal{P}(1,z_1z_2)=1. \qquad (2.7)$$

Hence, shifting the variables

$$z_i=x_i+(1-\bar{x})y_i, \quad i=1,2$$

$$\bar{x}=x_1+x_2=x_T/\sin\theta\simeq x_R \qquad (2.8)$$

from eqs. (2.6) and (2.7), we obtain

$$\Sigma^{AB\to C}(P_T,x_T,\theta)=C\,P_T^{-4}(1-\bar{x})^{A+B+C+2}x_1^\alpha x_2^\beta J^{AB\to C}(P_T,x_T,\theta)$$

$$J^{AB\to C}=\int_0^1dy_1\int_0^{1-y_1}dy_2\,y_1^Ay_2^B\frac{(1-y_1-y_2)^C[\bar{x}+(1-\bar{x})(y_1+y_2)]^D}{(x_1+(1-\bar{x})y_1)^{\mathcal{A}}(x_2+(1-\bar{x})y_2)^{\mathcal{B}}}\times$$

$$\times\;\mathcal{P}(Q_1^2,y_1,y_2) \qquad (2.9)$$

$$\mathcal{A}=A+\alpha+\sigma_1-1,\quad \mathcal{B}=B+\beta+\sigma_2-1,\quad D=\sigma_1+\sigma_2+\gamma-1.$$

In the framework of the assumptions made equations (2.8) and (2.9) define the relation between the high P_T hadron inclusive cross section and scaling violations in deep inelastic lepton scattering. Integral $J^{AB\to C}$ depends parametrically on the external variables x_T, θ, and, as will be discussed below, can be reduced for the physically interesting cases of small and threshold values of x_T.

3. We proceed now with the consideration of some limiting forms of the $\Sigma(\bar{x})$ in the regions of $\bar{x}$-variable where eqs. (2.8-9) permit the analytical calculations.

Note, that the P_T-dependence of the $\mathcal{P}(Q^2, y_1 y_2)$ function can be represented in the following way:

$$\mathcal{P}(Q^2, y_1, y_2) \sim (P_T^2)^{-f(\bar{x}, y_1 y_2)} \quad (3.1)$$

$$f(\bar{x}, y_1 y_2) = f(1 + \frac{x_1 x_2 - (1-\bar{x})^2 y_1 y_2}{x_1 x_2 + (1-\bar{x})(x_1 y_2 + x_2 y_1) + (1-\bar{x})^2 y_1 y_2},$$

$$\bar{x} + (1-\bar{x})(y_1 + y_2)). \quad (3.2)$$

A. Weak Screening $(\bar{x} \to 0)$

This limit corresponds to the small $x_T(x)$ range and is specified as approach to some boundary quark transverse momenta $P_T \geq P_T^*$ at the infinite energies $\sqrt{s} \to \infty$ *. In this case eqs. (2.8-9) and (3.1-2) give

$$\Sigma_0^{AB\to C}(P_T, x \to 0) = c_0 P_T^{-4(1+\epsilon_0)}, \quad (3.3)$$

where

$$\epsilon_0 = 1/2\, f(0, \overline{y_1 + y_2}),$$

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*Compare the discussion of the energy-angular constraints for jet production in the subsequent section (see also ref. /25/).

i.e., are defined by the scale breaking effects at the origin and correspond to the small deviations from the P_T^{-4} behaviour.

B. Strong Screening $(\bar{x} \to \bar{x}_{thr.})$

In this limit characterized by the threshold region of $x_T = \sin\theta$ variable or $x_R \to 1$ (neglecting the masses), anomalous P_T-dependence from nonscaling effects can provide maximal deviation from the canonical value $N=4$, seen now in the available FNAL-ISR range, i.e.,

$$\Sigma_1^{AB\to C}(P_T, \bar{x} \to 1) = c_1 P_T^{-(4+\epsilon_1)}$$

$$\epsilon_1 = 2f(2, 1). \quad (3.4)$$

Thus, eq. (3.4) corresponds to an increase in power exponent $N-4 = \epsilon_1$. Generally speaking, the nonscaling behaviour of the quark distributions (and decay) functions*, which define the precise form of the (3.3) and (3.4) limits could be obtained within AF gauge theories as the effect of different gluon corrections. In the framework of QCD, however, there is no selfconsistent method for the perturbative calculations of higher orders, and we restrict ourselves to the phenomenological ansatz. Particularly, for the sake of simplicity, we use a parametrization

$$\nu W_2^{\ell N}(x, Q^2) = F_2(x, Q^2) = F_H(x, Q_0^2)(Q^2/Q_0^2)^{-f_H(x)}, \quad H = A, B$$

$$f_H(x) = a + bx, \quad f_C = cz, \quad a = -0.25, \quad b = 1.0, \quad c = 0.5. \quad (3.5)$$

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*We do not consider here high quark k_T-effects which could give rise to some enhancement in our consideration. See refs. /26,27/ .

This pattern of empirical scale breaking contradicts, however, to the AF predictions of logarithmic Q^2-dependence of the $\nu W_2(x)$ moments, but is still applicable (numerically) as a quide to the recent deep inelastic experimental data. These power behaved breaking terms, in principle, can play a role of transient phenomena, and the right asymptotic region can be considered as the change of regime $(Q^2)^{-a} \to (\ln Q^2)^{-\alpha}$.

Thus, for the two simple cases considered we obtain the following P_T-behaviour of single particle distributions*

A. $\Sigma_0^{AB\to\pi} = c_0 P_T^{-4(1+\epsilon_0)}$, $\epsilon_0 = a + 3/8c$ — meson production $N \approx 4.0$ (3.6)

$\Sigma_0^{AB\to p} = c_0' P_T^{-4(1+\epsilon_0')}$, $\epsilon_0' = a + 3/8b$ — baryon production $N \approx 4.5$ (3.7)

and

B. $\Sigma_1^{AB\to\pi} = c_1 P_T^{-(4+\epsilon_1)}$, $\epsilon_1 = 4(a+b+c/2)$ — meson production $N \approx 8.0$ (3.8)

$\Sigma_1^{AB\to p} = c_1' P_T^{-(4+\epsilon_1')}$, $\epsilon_1' = 4a + 6b$ — baryon production $N \approx 9.0$ (3.9)

respectively.

The numerical constants $c_{0,1}$ contain the information about spin (angular) dependence of the quark cross sections, as well as the precise threshold behaviour of quark distribution (decay) functions. For the concrete type of QCD Born graph contributions to the $d\hat{\sigma}/d\hat{t}$ and $F(x, Q^2)$ the latter are fully calculable quantities.

* Note, that we do not consider here the leading particle production mechanism.

4. Now we would like to discuss briefly a very interesting connection between our results and the general asymptotic bounds on high-energy inclusive spectra which follow from analytic considerations of quantum field theory [30,31]. More precisely, we shall consider the low energetic particle production, i.e., the limit $x_R \to 0$, via the hard constituent scattering mechanism. As we have seen above, the corresponding cross section takes the form

$$(E d\sigma/d^3 p)_{jet} \sim P_T^{-4} \cdot f(x_R \approx 0). \qquad (4.1)$$

This means that all the dimensional factors which determine the behaviour of quark-parton distribution functions, etc., with all nonscaling effects cancel somehow in that low-energy region. One can guess that this is not an accident but the fundamental fact which does not depend significantly on the details of dynamical models.

Bearing this in mind, we compare eq. (4.1) with the low-energy theorem in the theory of inclusive reactions [31]. This theorem derived by A.A.Logunov and M.A.Mestvirishvili on the general grounds of analytic properties of multiparticle scattering amplitudes, says that

$$\lim_{x\to 0} \int_0^{x\sqrt{s}/2} \frac{d\sigma}{d\vec{k}}\, d\vec{k}_c \le \frac{\pi}{s} \log^2 s \ . \qquad (4.2)$$

To compare eq. (4.2) with our result, we have to integrate eq. (4.1) up to some lowest value $(P_T^*)_{min}$ which is determined by the requirement that the effective region of quark k_T-distribution in the initial hadrons will not be overlaped and by proper treatment of the jet production (cf. ref. [25]). The QCD calculations give as estimation [32]

$$(P_T^*)_{min} \ge (\overline{k_T})_{QCD} = \alpha_s \sqrt{Q^2} \sim \alpha_s \sqrt{s}$$

$\alpha_s = g_s^2/4\pi$ - the running coupling constant of strong interactions.

Thus, we have

$$\lim_{x\to 0} \left(\frac{d\sigma}{dx}\right)_{jet} = \frac{1}{s}\left(\frac{f(x=0)}{a_s^2}\right) \qquad (4.3)$$

what is consistent with the low-energy theorem in the form (4.2).

SOME CONCLUDING REMARKS

i) As was previously mentioned, the case of the weak screening (3.6-7) and the corresponding eq. (4.3) should be valid at rather high energies of colliding hadrons and fixed final $P_T \geq P_T^* \sim 1\div 2$ *GeV/c*, i.e., $x_T \to 0$. This interval of the transverse momenta could be associated with the passing to the high P_T regime in production cross sections. This phenomenon is already seen in cosmic ray data[33] and has support in the available accelerator energy range. For instance, there is some evidence from data analysis on the basis of radial scaling[34] at $x_R \to 0$ and world data analysis[22]. A similar transition region manifests itself also in the fixed angle elastic cross section[35] at the small $|t|$ values (corresp. $P_T^* \sim$ ~ 1.7 *GeV/c*) in accordance with the (3.6-7) limit and the exclusive-inclusive connection[36].

ii) In the case of the medium x_T range, averaging between the strong and weak screening limits provides the resulting values of power law exponent N which are less than 8 for meson, and 10 for baryon productions, respectively. It should be noted that the results of the world data analysis support the absence of the stable values $N=8\,(12)$ for $\pi, K\,(p^{\pm})$-production cross sections. Besides, some evidence of the decreasing character of exponent N comes from ISR experiments[23,24] and preliminary data of CCHK, CCOR, CSZ groups[37,38], where $N \simeq 6.3$ for mesons.

iii) We should like to emphasize that the weak decrease in $N(\theta)$ dependence for the small values of production angle θ in sense of (3.1-2)* does not contradict to the high P_T experimental data[22,24] and is in accordance with the corresponding elastic scattering data showing the smaller values of $N_{excl.}$ $(d\sigma/dt \sim s^{-N_{excl.}} \cdot \tilde{f})$ for smaller fixed angles[39].

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* Note, that the effective power analysis N_{eff} could mix all these effects, and the corresponding results can be considered only in the sense of average values[7,17].

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