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Модель с логарифмически нарушенным скейлингом и лептон-нуклонные взаимодействия при высоких энергиях

Предложена кварк-партонная модель с логарифмически нарушенным скейлингом. Принципы, лежащие в ее основе, не противоречат квантовой хромодинамике. Получены явные выражения для кварковых и глюонных распределений в нуклоне. На их основе рассчитаны сечения глубоконеупругого ep-, ed- и $\nu(\bar{\nu})N$-рассеяния. Результаты расчетов сравниваются с экспериментальными данными.

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A Model with Logarithmic Scaling Violation and High-Energy Lepton-Hadron Interactions

INTRODUCTION

Due to an intensive development of Quantum Chromodynamics (QCD) a notable progress has been achieved recently in understanding the structure of elementary particles and the nature of approximate scaling and its violation. Now we can formulate rigorous QCD predictions for the short-distance phenomena. Unfortunately, these predictions are expressed only either in terms of the moments of deep-inelastic structure functions or as the Q^2 evolution equation [1] for those ones. To define the very structure functions in the whole range of the x and Q^2 values we have either to invert the moments or to solve the Q^2 evolution equation. As is well known, in both cases initial information on x-dependence of the structure functions at a fixed Q_0^2 value is necessary. Usually it is extracted from the relevant experimental data. Thus, we can see that up to now the deep inelastic structure functions have not been clear theoretically (without any empiricism) calculated within QCD. Evidently, the main reason for this situation is a strong dependence of the short-distance phenomena on the long-range nature of theory. We can factorize the long and short distances [2], but not separate them completely. It is clear, for self-consistent solution of the short distances problem in the framework of QCD the long-range structure of theory should be established. Recently there has been developed a series of phenomenological models [3], which imply the QCD scaling violation. The structure-function parametrizations which have been provided with these models are very useful for applications and permit to check particular QCD properties.

In this paper we propose a quark-parton model, based on the Kuti, Weiskopf ideas [4], that is used here in more radical ways, and some principles of QCD. The model contains two free parameters and agrees well with the experimental data.

The paper is organized as follows:

In Sec. I main assumptions of our model are formulated and explicit expressions for quark and gluon distribution functions in a nucleon are obtained.

In Sec. II the cross-sections of deep inelastic eN-, ed- and $\nu(\bar{\nu})$ N-scattering and the asymmetry of polarized ep-scattering are calculated. The ed-scattering should be calculated taking into account the nucleon Fermi motion in the deuteron.

Section III is dedicated to the comparison of model predictions with the experimental ep, ed and $\nu(\bar{\nu})$N data.

1. QUARK-GLUON DISTRIBUTIONS

Here follows the list of our model assumptions.

1. As usual, we shall assume that a nucleon consists of three valence quarks which define its quantum numbers, and of a singlet "sea" of quark-antiquark pairs and neutral vector gluons.

2. Single-particle distribution of valence quarks is selected in the form

$$f = \frac{x^{1-\alpha(0)}}{\sqrt{x^2 + \frac{\mu^2}{\vec{p}^2}}}, \qquad /1.1/$$

where $\alpha(0)=1/2$ is the intercept of A_2-meson Regge trajectory; x, the fraction of a nucleon longitudinal momentum $\vec{p}$ carried by a parton of mass μ.

We have taken these assumptions without alterations from ref.[4]. It provides a conventional Regge asymptotic of a nondiffractive component of the forward virtual Compton amplitude.

The following two assumptions are decisive for the description of scaling violation and present a generalization of the Kuti, Weisskopf ideas[4].

3. We will assume quark-gluon interaction to be renormalizable (QCD as an example) and the scaling violation arises mainly due to the coupling constant renormalization[5], which results in the change of the "bare" coupling constant to the running one in all expressions for the structure functions. Of course, there are more ingenious effects of scaling violation (for example, because of the wave function renormalization) but in the framework of our phenomenological consideration we will not take them into account.

4. Finally, let us assume that the nature of a nucleon "sea" is similar to the one of an equivalent photon system of quantum electrodynamics. The present assumption seems to be justified at least for the case of large momentum transfers and small coupling constant[6]. Thus it looks quite natural from the point of view of asymptotically free theories.

With the mentioned assumptions the "bare" quark and gluon distributions can be obtained. The "bare" distribution (see ref.[7]) is that one which does not reflect the fact that quarks and gluons are constituents of a concrete hadron.

The spectrum $f(E_\gamma)$ of "equivalent" photons has the form:

$$f(E_\gamma) \sim \alpha / E_\gamma, \qquad /1.2/$$

where α is an electromagnetic coupling constant; E_γ, photon energy. Exp. /1.2/ has the form of a single-particle phase space distribution.

Taking the 4th assumption we will choose the following form for the distribution functions of the "sea" quark ϕ and gluons ψ

$$\phi = a'(g)\frac{1/3}{\sqrt{x^2 + \frac{\mu^2}{\vec{p}^2}}}, \qquad \psi = a''(g)\frac{e^{-\beta x}}{\sqrt{x^2 + \frac{\mu^2}{\vec{p}^2}}} \qquad /1.3/$$

Here g is the quark gluon coupling constant; $a'(g)$ and $a''(g)$ are some unknown functions of g.

It should be noted that in a contrast to the distribution function ϕ the ψ-function has an additional Boltzmann factor $e^{-\beta x}$. The introduction of this factor is supported by the fact that the gluon gas in a nucleon has, evidently, a rather high density which exceeds significantly the density of the quark gas. (It is known at least that more than 60% of the nucleon momentum is carried by neutral gluons). Thus, it is reasonable to describe it with an equilibrating statistical distribution like the Boltzmann distribution. Taking into account assumption 3 we have to change in /1.3./ g to the running coupling constant $\bar{g}(Q^2)$. In the case of a renormalized interaction $\bar{g}(Q^2)$ depends on Q^2 logarithmically. Thus it follows:

$$\phi(x,Q^2) = a'(\bar{g}(Q^2))\frac{1/3}{\sqrt{x^2 + \frac{\mu^2}{p^2}}}, \qquad \psi(x,Q^2) = a''(\bar{g}(Q^2))\frac{e^{-\beta x}}{\sqrt{x^2 + \frac{\mu^2}{p^2}}} \qquad /1.4/$$

It means that within our assumptions x and Q^2 -dependence of the "bare" distribution functions are factorized.

Now, when we know the "bare" distributions /1.4/, we can obtain the distribution functions of quarks and gluons inside a nucleon. To do this let us consider all sorts of multi-particle configurations corresponding to the nucleon in the infinite momentum frame (IMF). The probability of an N -particle state can be written down in the form:

$$dP_N(x_1,x_2,\dots,x_N)=Z\cdot\frac{1}{k_1!\,k_2!\,k_3!\,\ell!}\times$$

$$\delta(1-\sum_{i=1}^{N}x_i)\prod_{i=1}^{3}f(x_i)dx_i\prod_{j=1}^{k_1+k_2+k_3}\phi(x_j)dx_j\times \qquad /1.5/$$

$$\prod_{k=1}^{\ell}\psi(x_k)dx_k\,.$$

Here Z is the normalization constant; $N=3+\sum_i k_i+\ell$, $k_i = 0,2,4,\dots$ is the number of "sea" quarks and antiquarks, ℓ is the number of gluons in a given N -particle state, i is the flavour. Integrating over all x -variables except one and summing over all possible configurations we find the sought distribution functions (see refs. [4,8]). After a series of transformations we get the following expressions:

$$G^p_{1V}(x,Q^2)=2G_{2V}=2Z\lim_{p\to\infty}f\cdot\hat{A}^{(2)}[f,\phi,\psi].$$

$$G^p_{1C}(x,Q^2)=G^p_{2C}=G^p_{3C}= \qquad /1.6/$$

$$Z\lim_{p\to\infty}\phi\,\hat{A}^{(3)}[f,\phi,\psi],$$

$$G^p_0(x,Q^2)=Z\lim_{p\to\infty}\psi\,\hat{A}^{(3)}[f,\phi,\psi],$$

where G^p_{iV}, G^p_{iC}, G^p_0 are the distribution functions in a proton for valence quarks, "sea" (anti-) quarks, and gluons, respectively; i is the flavour. For neutron $G^N_{1V}=G^p_{2V}$, $G^N_{2V}=G^p_{1V}$.

The $\hat{A}^{(k)}$ -operator is defined by the relation

$$\hat{A}^{(k)}[f,\phi,\psi]=\int_{-\infty}^{\infty}dy\,e^{i(1-x)y}\,(\int_0^{\infty}dz\,e^{-iyz}f(z))^k\times \qquad /1.7/$$

$$\exp\{\int_0^{\infty}dt\cdot e^{-iyt}(3\phi(t)+\psi(t))\}.$$

Substituting /1.4/ into /1.6/ and fixing the normalization we get the final expressions:

$$G^p_{2V}(x,Q^2)=\frac{x^{-1/2}(1-x)^{\tau(Q^2)}}{B(\frac{1}{2},\tau(Q^2)+1)}\cdot\frac{\Phi(a''(Q^2),\tau(Q^2)+1;-\beta(1-x))}{\Phi(a''(Q^2),\tau(Q^2)+3/2;-\beta)}$$

$$G^p_{1C}(x,Q^2)=\frac{a'(Q^2)}{3x}(1-x)^{\tau(Q^2)+\frac{1}{2}}\frac{\Phi(a''(Q^2),\tau(Q^2)+\frac{3}{2};-\beta(1-x))}{\Phi(a''(Q^2),\tau(Q^2)+3/2;-\beta)} \qquad /1.8/$$

$$G^p_0(x,Q^2)=3\cdot\frac{a''(Q^2)}{a'(Q^2)}\cdot e^{-\beta x}\cdot G^p_{1C}(x,Q^2)$$

$\Phi(\alpha,\beta;z)$ is a degenerated hypergeometric function, $\tau(Q^2)=a'(Q^2)+a''(Q^2)$.

The distributions of concrete-type quarks in a nucleon in the case of SU(3)-symmetric "sea" are expressed through G_V and G_C in an ordinary way. In the case of broken SU(4)-symmetry of the nucleon "sea" (charm is suppressed) we can write

$$u(x,Q^2)=2G^p_{2V}(x,Q^2)+\frac{1}{2}(1-\alpha_c)G^p_{1C}(x,Q^2)$$

$$d(x,Q^2)=G^p_{2V}(x,Q^2)+\frac{1}{2}(1-\alpha_c)G^p_{1C}(x,Q^2) \qquad /1.9/$$

$$\bar{u}(x,Q^2)=\bar{d}(x,Q^2)=\bar{S}(x,Q^2)=S(x,Q^2)=\frac{1}{2}(1-\alpha_c)G^p_{1C}(x,Q^2)$$

$$C(x,Q^2)=\bar{C}(x,Q^2)=\frac{3}{2}\alpha_c G^p_{1C}(x,Q^2)\,.$$

where $0 \le a_c \le 1/4$ is the parameter of charm suppression in a nucleon. It is not our model parameter, but clearly a kinematical one. The suppression mechanism we have used is a simple replacement of a' by $(1-a_c)a'$ and of a'' by $(1-a_c)a''$ in the "bare" distributions /1.4/ of light and charmed quarks, respectively.

It would be more natural to accept a_c-parameter to depend on x and Q^2. However, it is not our goal to clarify the properties of $a_c(x, Q^2)$-function and below we will consider the a_c-parameter as some averaged value of this function over the kinematic region we are interested in.

The a' and a''-functions we have introduced earlier define the form of scaling violation. In the present paper we will use the following simple phenomenological representation for a' and a'' /7/:

$$a'(Q^2) = a''(Q^2) = a(Q^2) = \frac{a}{\bar{g}^2(Q^2)}. \qquad /1.10/$$

In the case of QCD the running coupling constant $\bar{g}(Q^2)$ is:

$$\bar{g}^2(Q^2)/4\pi = \frac{12\pi}{25\,\ln Q^2/\Lambda^2}, \qquad /1.11/$$

where $\Lambda = 0.5$ (GeV/c) (the conventional value). The rigorous calculation of the $a'(Q^2)$ and $a''(Q^2)$-functions in the framework of QCD may be a subject for a subsequent paper. Here we notice that the representation /1.4/-/1.5/ leads, according to /1.8/, to the increase of the "sea" contribution at large Q^2 in the region $x \to 0$. Moreover, /1.4/-/1.5/ and /1.8/ automatically ensure the decrease of structure functions with the growth of Q^2 in the large-x region. As is well known both phenomena are observed experimentally and predicted by QCD.

A peculiar prediction of our model arises from /1.8/ and /1.10/-/1.11/. Let us find the separation point $x_s(Q^2)$ which satisfies the equation

$$\frac{\partial F(x_s, Q^2)}{\partial Q^2} = 0. \qquad /1.12/$$

Here F is one of the structure functions /1.9/. For large Q^2 we can obtain

$$x_s(Q^2) \sim \frac{1}{a'(Q^2) + a''(Q^2)} = \frac{1}{2a(Q^2)}. \qquad /1.13/$$

Thus, $x_s(Q^2)$ is a decreasing function of Q^2 since $a(Q^2)$ is an increasing one. It leads to the contraction with the growth of Q^2 of the region $x < x_s(Q^2)$ where the structure functions are increasing functions of Q^2 and to the disappearance of this region in the limit $Q^2 \to \infty$.

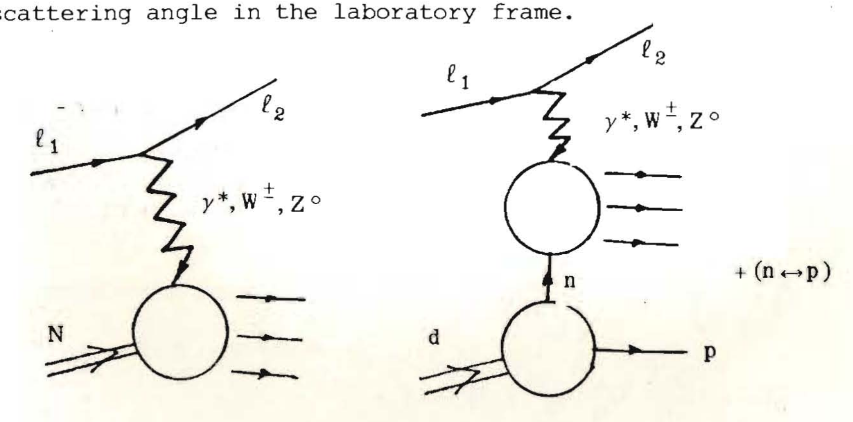
2. THE CALCULATION OF THE CROSS-SECTIONS

In the case of deep inelastic electromagnetic interactions the differential cross section of scattering has the following form in a single-photon approximation

$$\frac{d^2\sigma}{dE'\,d\Omega} = \frac{4\alpha^2 E'^2}{Q^4}\left(2W_1 \sin^2\frac{\theta}{2} + W_2 \cos^2\frac{\theta}{2}\right) \times \qquad /2.1/$$

$$(1 + \zeta_\ell \zeta_\parallel \Delta_\parallel + \zeta_\ell \zeta_\perp \Delta_\perp),$$

where as usual ζ_ℓ is the degree of longitudinal electron (muon) polarization (with respect to the electron momentum $\vec{K}$ in the laboratory frame); $\zeta_\perp$ and $\zeta_\parallel$ are components of target polarization on the lepton scattering plane, perpendicular and parallel to $\vec{K}$, respectively; θ is the lepton scattering angle in the laboratory frame.



a) b)

Fig.1.

The longitudinal $\Lambda_{\parallel}$ and transversal $\Lambda_{\perp}$ asymmetries are

$$\Lambda_{\parallel} = \frac{(E + E'\cos\theta) M g_1 + Q^2 g_2}{W_1 + \frac{1}{2}\,\mathrm{ctg}^2(\frac{\theta}{2}) W_2}, \qquad /2.2/$$

$$\Lambda_{\perp} = \frac{E'\sin\theta\,(M g_1 + 2E g_2)}{W_1 + \frac{1}{2}\,\mathrm{ctg}^2(\frac{\theta}{2}) W_2}.$$

For the case of weak deep inelastic interactions the differential cross section in the lowest order in the week coupling constant G is defined as follows:

$$\frac{d^2\sigma^{\pm}}{dx\,dy} = \frac{G^2 ME}{\pi}\left(\frac{M_{ex}^2}{M_{ex}^2 + Q^2}\right)^2 [(1 - y - \frac{Mxy}{E}) F_2^{\pm} + \qquad /2.3/$$

$$\frac{xy^2}{2} F_1^{\pm} \mp (y - \frac{y^2}{2}) x F_3^{\pm}],$$

where M_{ex} corresponds to $W^{\pm}$ or Z° bozon masses; $x = Q^2/2M\nu$; $y = \nu/E$. E is the energy of the neutrino beam.

In the framework of a parton model one may obtain correspondence rules [9] between quark distributions of a nucleon and structure functions $F_i^{\pm}$ for the interactions with the hadron V-A current of the general form:

$$J_{\mu} = \sum_{i,j} \bar{q}_i \gamma_{\mu} (C_{ij}^V + C_{ij}^A \gamma_5) q_j, \qquad /2.4/$$

where q_i is the field of an i-th quark; C_{ij}^V and C_{ij}^A are the vector and axial coupling constants. Taking into account the kinematic mass corrections to scaling and the threshold effects due to the heavy quarks production we can find these rules in the form:

$$F_k^{\pm} = \sum_{i,j} (F_{ji(k)}^{\pm} + \tilde{F}_{ij(k)}^{\pm}) \qquad /2.5/$$

$$F_{ji(1)}^{\pm} = (C_{ij}^{V^2} + C_{ij}^{A^2}) \phi_j^{\pm} (\xi_i) \theta (1 - \xi_i) \qquad /2.6/$$

$$\tilde{F}_{ij(1)}^{\pm} = F_{ij(1)}^{\mp} \qquad /2.7/$$

$$F_{ji(3)}^{\pm} = \mp 2 C_{ij}^V C_{ij}^A \phi_j^{\pm}(\xi_i) \theta(1 - \xi_i), \quad \tilde{F}_{ij(3)}^{\pm} = F_{ij(3)}^{\mp} \qquad /2.8/$$

$$F_{ji(2)}^{\pm} = \xi_i F_{ji(1)}^{\pm}, \quad \tilde{F}_{ij(2)}^{\pm} = F_{ij(2)}^{\mp}. \qquad /2.9/$$

Here $\phi^+(z,Q^2)$ and $\phi^-(z,Q^2)$ correspond to the momentum distribution functions of i-th type quarks and antiquarks. $\theta(1-\xi_i)$ factors define the thresholds of production of quarks with masses m_i. $F_{ij}^{\pm}$, $\tilde{F}_{ij}^{\pm}$ are partial structure functions of the quark transitions: F_{ij}^{+} and $\tilde{F}_{ij}^{-}$ for $q_i \to q_j$, F_{ij}^{-} and $\tilde{F}_{ij}^{+}$ for $\bar{q}_i \to \bar{q}_j$. Using the reduced ξ-scaling variable $\xi_i = (Q^2 + m_i^2)/2M\nu$, we have taken into account the most important mass corrections to scaling for the processes considered below. ξ_i coincides with the conventional ξ-scaling variable [10] to the second order in $(M/Q)^2$. The summation in /2.5/ is performed over such a set of quark transitions which leads to the final hadronic state of the considered reaction.

Using (2.5)-(2.9) it is easy to calculate the structure functions for interactions with electromagnetic current, with charged and neutral currents (in concrete calculations we use the standard 4-quark model WS-GIM) and for the case of charmed hadron state [9] production. Assuming (like it has been made in ref. [4]) that the main contribution into the polarized scattering is made by valence quarks and that [10]

$$\frac{1}{2}[G_{1V}^{p\uparrow} - G_{1V}^{p\downarrow}] = \frac{2}{3} G_{1V}^{p} \qquad /2.10/$$

$$\frac{1}{2}[G_{2V}^{p\uparrow} - G_{2V}^{p\downarrow}] = -\frac{1}{3} G_{2V}^{p}$$

we obtain the polarized structure functions

$$M^2 \nu g_1 = \frac{1}{4}(\frac{4}{9}(G_{1V}^{p\uparrow} - G_{1V}^{p\downarrow}) + \frac{1}{9}(G_{2V}^{p\uparrow} - G_{2V}^{p\downarrow})) = \frac{5}{18} G_{2V}^{p} \qquad /2.11/$$

$$g_2 = 0.$$

If the $W^p_{1,2}$ and $W^n_{1,2}$ nucleon structure functions are known it is possible to proceed to the calculation of the deuteron structure functions. There is a large number of approaches to this problem. In the present paper we use the conventional Atwood-West approach /12/ . As long as we intend to consider only the $x > 1$ region, the difficulties of this formalism, which have been pointed out in ref. /13/ , are not emerged. The upper blocks of the diagram (Fig. 1b) are described by the nucleon structure functions and depend on their Fermi-motion momentum inside of a deuteron.

With the help of the relativistic spherically-symmetrical deuteron wave function ψ it becomes possible to perform the averaging over the nucleon Fermi-motion momenta and to obtain the deuteron structure functions:

$$W^d_{\mu\nu} = \int \frac{d^3\vec{p}_s}{E_s/M_s} \, |\psi(|\vec{p}_s|)|^2 (W^n_{\mu\nu} + W^p_{\mu\nu}) \equiv W^{n(s)}_{\mu\nu} + W^{p(s)}_{\mu\nu} , \qquad /2.12/$$

where $W^{n,p}_{\mu\nu}$ are the structure functions of neutron and proton, respectively; $\vec{p}_s$, E_s , M_s are the 3-momentum energy and mass of a free spectator nucleon emitted from the bottom block of the diagram in Fig. 1b. In our calculations we have used the Reid "hard core" wave function /14/ extended to the relativistic region in accordance with Atwood-West /12/ .

3. COMPARISON WITH EXPERIMENT

We have compared the above obtained predictions of our model for deutron and proton targets with the experimental data on the deep inelastic ed-, ep- /15/ and $\nu(\bar{\nu})N$ -scattering /16,17/ . To test the prediction of our model in more detail we have carried out a joint analysis of these data, using a unique set of free parameters of the model. The best agreement of theoretical curves with experimental points (see for illustration Figs. 2 and 3 presenting a part of data analysed) is obtained at the following values of free parameters: a =5.2, β =-3.5, a_c =0.23. In this case $\chi^2/\bar{\chi}^2$=370/378. Using the found values of the varied parameters we have calculated the curves for the longitudinal asymmetry coefficient of electron scattering on proton. The comparison of the theoretical curves with the known experi-

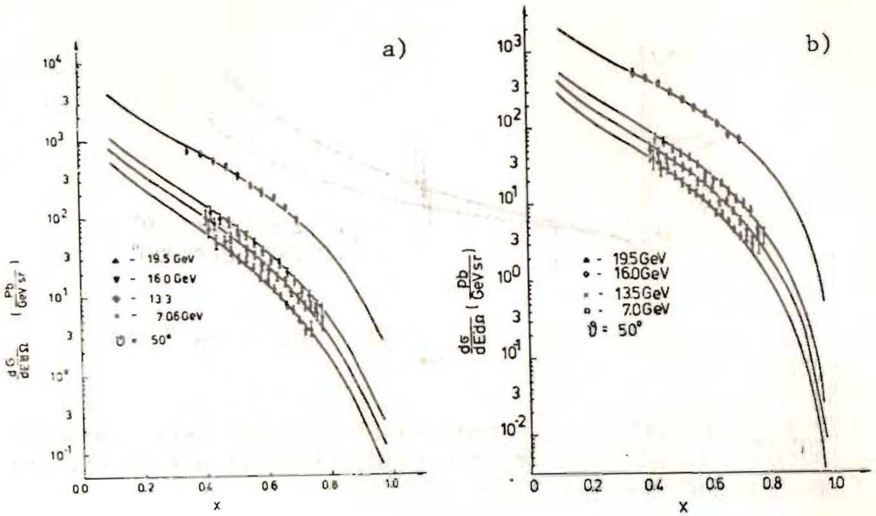


Fig. 2. Differential cross section of deep inelastic a) ep-, b) ed-scattering. Solid lines are the predictions of our model.

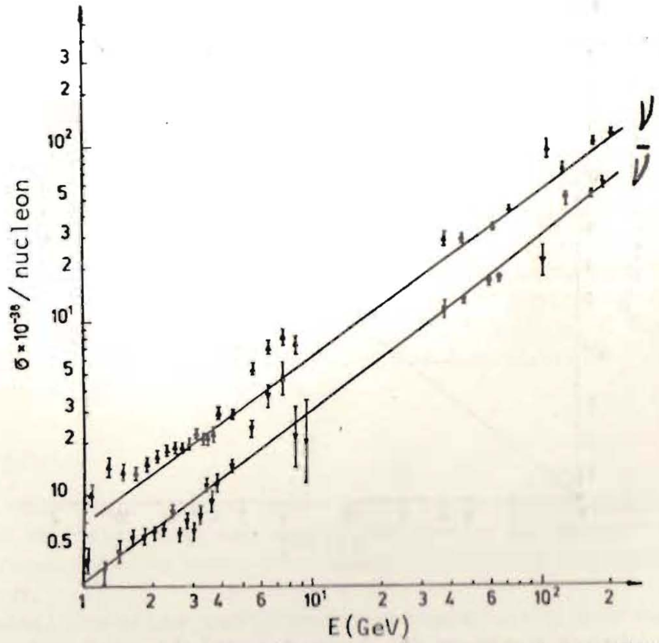


Fig. 3. Total cross section of deep inelastic $\nu(\bar{\nu})N$ -scattering. Solid lines are the predictions of our model.

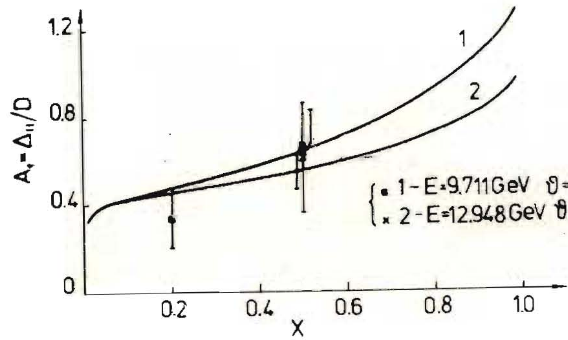


Fig. 4. Longitudinal asymmetry of deep inelastic ep-scattering. Solid lines are the predictions of our model. The points are the experiment [16].

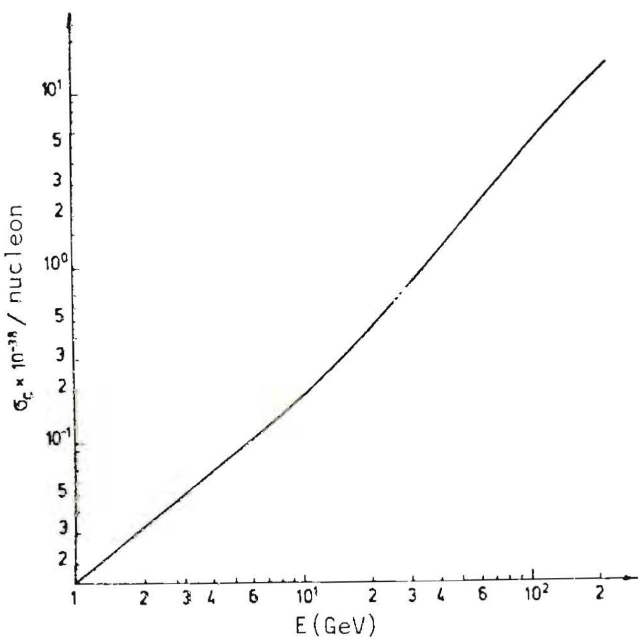


Fig. 5. Cross section of charmed-particle production in deep inelastic $\nu(\bar{\nu})N$-scattering. The prediction of our model.

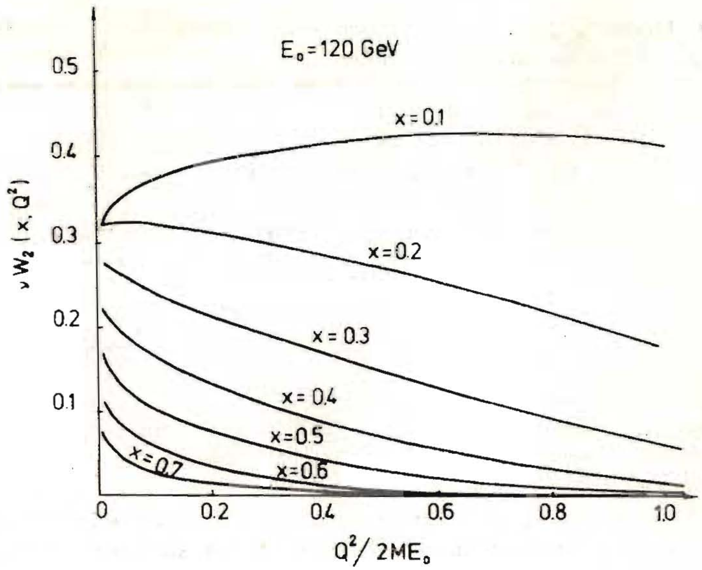


Fig. 6. The model predictions for behaviour of νW_2^p -structure function in the range of large momentum transfers.

mental points[18] is given in Fig. 4. The cross section of the charm production $\sigma^c(\Delta c = \pm 1)$ has been also calculated (see Fig. 5).

The considered experimental data cover a rather kinematic region: $4 \leq Q^2 \leq 30$ GeV, $0.34 \leq x \leq 0.97$ and are well described by our model. Thus it seems quite natural to extend the model predictions to the large Q^2 region which is not well investigated yet. In Fig. 6 we show the Q^2-dependence of the electromagnetic structure function νW_2^p of the proton predicted by our model.

4. CONCLUSION

Thus, the quark parton model with logarithmic scaling violation we have proposed describes quite successfully the experimental data on ep-, ed-, and $\nu(\bar{\nu})N$ -deep inelastic scattering.

The model contains two free parameters which are defined from the experimental data. The basis ideas of our model do not contradict QCD. Moreover it may happen that they follow from QCD in a natural way. Clarification of such a possibility is in fact the substantiation problem for our model in the context of this theory.

We would like to summarize some specific predictions of the model. It predicts that the separation point $x_s(Q^2)$ /1.13/ tends to zero when Q^2 tends to infinity. This means that the region $x \le x_s(Q^2)$, where structure functions are increasing functions of Q^2, contracts and disappears in the limit $Q^2 \to \infty$.

Gluons inside a nucleon are treated as a gas with the Boltzmann-type distribution functions. It is necessary to point out that the negative value of β-parameter in /1.3/ obtained from the comparison with experimental data does not lead to any contradictions. A self-consistent thermodynamical description requires negative values of the Boltzmann exponential power (i.e., $\beta > 0$) only in the large momentum limit that ensures the finiteness of the free energy integral. So, we consider the β-parameter as an average value of a certain function $\beta(x)$ (in $\exp(-x\beta(x))$) over the interval $0.34 \le x \le 0.97$ (experimental region) and $\beta(x) > 0$ at $x \sim 1$.

In principle, there may be another point of view on the negative value of the β-parameter. It can be regarded as the case of negative temperatures (in the $e^{-E/T}$ -distribution) encountered in the statistical systems at the inverse population of energy levels (e.g., laser). The latter possibility can be realized only in the weakly self-interacting systems. The gluon gas obeying QCD satisfies this condition at small distances ($Q^2 \to \infty$). Certainly such an interpretation is not well grounded on and it has a very preliminary character.

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