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INCLUSIVE ELECTROPRODUCTION IN QUANTUM FIELD THEORY

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INCLUSIVE ELECTROPRODUCTION IN QUANTUM FIELD THEORY

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Инклюзивное электророждение адронов в квантовой теории поля

Показано, что на основе теории спинорного поля и псевдоскалярного глюонного поля с псевдоскалярной связью диаграммный подход приводит к партонной картине для инклюзивного процесса $e+p\to e+h+X$. Скейлинг структурных функций получается в различных кинематических областях.

Работа выполнена в Лаборатории теоретической физики ОИЯИ.

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Gàlfi L. E2 - 11223

Inclusive Electroproduction in Quantum Field Theory

It is shown that starting from the quark-pseudoscalar gluon interacting Lagrangian the diagrammatic approach leads to parton picture for the process $e+p\to e+h+X$. The scaling behaviour of certain structure functions is investigated in different kinematic regions.

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

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I. Photon induced inclusive reactions are usefull tools for probing the hadron structure. Theoretical attempt to explain the experimental results of deep inelastic electron-nucleon scattering led to the parton picture of hadrons. Different versions of the parton model were used to describe leptoproduction of hadrons on nucleon. The operator product expansion (OPE) on the other hand proved to be a nice method to investigate the deep inelastic scattering. It has, however, little use for those processes which are not light cone dominated. The diagrammatic approach which is based on the investigation of the asymptotics of Feynman diagrams proved to be equivalent with the OPE, where the latter can be used, but gives wider applicability [1].

In this paper we apply the diagrammatic approach to inclusive electroproduction $e + p \to e + h +$ anything. We show that starting from a quark-pseudoscalar gluon interacting Lagrangian the diagrammatic approach leads to parton picture. (We believe that our predictions remain valid for the standard quark-vector gluon model of hadrons, too). The scaling behaviour of certain structure functions was investigated earlier by OPE in the target fragmentation region [2], and by parton model calculations [3,4]. We deduce an expression which is a substitute of the OPE in the current fragmentation region and gives the same scaling results as the parton models.

The essential kinematics is presented in Sect. II. In Sect. III the asymptotics of the one-particle distribution function is investigated in different kinematic regions. This is a slight extension of the results of [5]. Assuming finite charge renormalization for the quark-gluon theory, the current fragmentation region is discussed in detail in Sect. IV. We do not construct bound states from quarks. We consider the physical hadrons as collinear beams of their constituents. We hope this simplification does not influence the asymptotic predictions.

II. The scattering process is shown in *Fig. 1*. The following notation is used:

$$\nu = \frac{pq}{m},\ \nu_1 = \frac{p'q}{m_1},\ \kappa = \frac{pp'}{m},\ \omega = -\frac{2pq}{q^2},\ Q^2 = -q^2.$$

In the laboratory frame

$$k = E(1, \sin\psi, 0, \cos\psi),$$

$$k' = E'(1, \sin(\psi+\theta), 0, \cos(\psi+\theta)),$$

$$q = (q_0, 0, 0, q_3),$$

$$p' = (p'_0, p_T\cos\phi, p_T\sin\phi, p_L),$$

m and m_1 are the masses of the proton and the detected hadron, respectively. The lepton mass is neglected. In the lab. frame the differential cross section is

$$p'_0 \frac{d\sigma}{dQ^2 d\nu\, d^3p'} = \frac{4\pi\alpha^2}{Q^4 E^2} \ell^{\mu\nu} W_{\mu\nu},$$

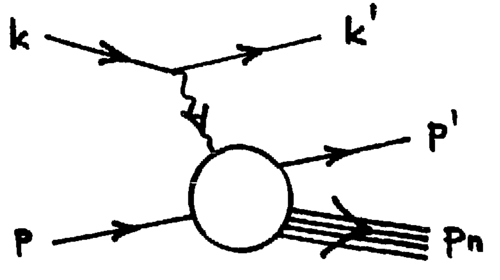

Fig. 1

where $\ell_{\mu\nu}$ is the lepton tensor,

$$W_{\mu\nu} = \sum_n \langle p | J_\mu (0) | p', n \rangle \times$$

$$\times \langle p', n | J_\nu (0) | p \rangle (2\pi)^4 \delta(p + q - p' - p_n)$$

is the hadron tensor. Summation over the polarization of the scattered electron and that of the detected hadron and average over the polarization of the incoming electron and proton is understood. Expanding in terms of invariant amplitudes, we have

$$W_{\mu\nu} = (-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}) W_1 + \frac{1}{m^2}(p_\mu - \frac{pq}{q^2} q_\mu)(p_\nu - \frac{pq}{q^2} q_\nu) W_2 +$$

$$+ \frac{1}{m_1^2}(p'_\mu - \frac{p'q}{q^2} q_\mu)(p'_\nu - \frac{p'q}{q^2} q_\nu) W_3 +$$

$$+ \frac{1}{mm_1} [(p_\mu - \frac{pq}{q^2} q_\mu)(p'_\nu - \frac{p'q}{q^2} q_\nu) + \mu \leftrightarrow \nu] W_4 .$$

Integrating over the azimuthal angle ϕ, one can define the integrated structure functions by

$$\tilde{W}_{\mu\nu} = \int \frac{d^3 p'}{p'_0} W_{\mu\nu}\, \delta(\kappa - \frac{pp'}{m}) \delta(\nu_1 - \frac{p'q}{m_1}) =$$

$$= (-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}) \tilde{W}_1 + \frac{1}{m^2}(p_\mu - \frac{pq}{q^2} q_\mu)(p_\nu - \frac{pq}{q^2} q_\nu) \tilde{W}_2 ,$$

where $\tilde{W}_i = \tilde{W}_i(q^2, \nu, \nu_1, \kappa)$,

$$\frac{d\sigma}{dQ^2 d\nu\, d\kappa\, d\nu_1} = \frac{4\pi\alpha^2}{Q^4 E^2} \ell^{\mu\nu} \tilde{W}_{\mu\nu} =$$

$$= \frac{4\pi\alpha^2}{Q^4} \frac{E'}{E} (\cos^2 \frac{\theta}{2} \tilde{W}_2 + 2 \sin^2 \frac{\theta}{2} \tilde{W}_1) .$$

We define the cross section for the initial proton to absorb a virtual photon (transverse or longitudinal) and produce the detected hadron plus missing mass by

$$p'_0 \frac{d\sigma^{T,L}}{d^3 p'} = \frac{4\pi\alpha^2 m}{m\nu + q^2/2} \epsilon_\mu^{T,L} \epsilon_\nu^{T,L*} W_{\mu\nu},$$

where $\epsilon_\mu^{T\pm} = \frac{1}{\sqrt{2}}(0,1,\pm i,0)$, $\epsilon_\mu^{L} = \frac{1}{\sqrt{Q^2}}(q_3,0,0,q_0)$ in the lab. frame. Then we have for the cross sections averaged over the azimuthal angle:

$$p'_0 \frac{d\sigma^T}{p_T dp_T dp_L} = \frac{4\pi\alpha^2 m}{m\nu + q^2/2} \frac{\sqrt{\nu^2 + Q^2}}{m_1} \tilde{W}_1 ,$$

$$p'_0 \frac{d\sigma^L}{p_T dp_T dp_L} = \frac{4\pi\alpha^2 m}{m\nu + q^2/2} \frac{\sqrt{\nu^2 + Q^2}}{m_1} [\nu \tilde{W}_2 (\frac{\omega}{2m} + \frac{1}{\nu}) - \tilde{W}_1].$$

The deep inelastic region is defined by $Q^2 \to \infty$, $\nu \to \infty$ with ω finite. The definitions of the target and current fragmentation region are the following: pp' finite, $qp' \to \infty$ with qp'/qp finite: target fragmentation. $pp' \to \infty$, $qp' \to -\infty$ with qp'/qp and pp'/qp finite: current fragmentation.

III. We investigate now the asymptotic behaviour of the one-particle distribution function $f_{\gamma^* p \to hX}$ with the help of the diagrammatic approach. According to the generalized optical theorem $f_{\gamma^* p \to hX} = \frac{1}{\sigma_t s} \mathrm{disc}_{M^2} T$,

where σ_t is the total photoabsorption cross section, T is the $\gamma^* p\bar{h} \to \gamma^* p\bar{h}$ forward scattering amplitude, $s = (p+q)^2$, $M^2 = (p+q-p')^2$. We consider a quark-pseudoscalar gluon theory with $L_I = g\bar{\psi}\gamma_5 \psi\phi + h\phi^4$. The method is based on the investigation of the asymptotic behaviour of Feynman diagrams. The results can be formulated by three "rules of game": 1) The behaviour of diagrams of high energy processes in the region where some of the variables are large is connected with the scale regime of such subgraphs which being contracted into a point make the initial diagram weakly connected and independent of the large variables. 2) The scale regime for such subgraphs generates simple pole in the sum of the Mellin parameters corresponding to those variables the dependence on which disappears as a result of the contraction at the point equal to the scale dimension of the subgraph. 3) The coefficient of the pole is a product of the contribution from the contracted part, which depends on the bare coupling constants g_0, h_0 and those from the weakly connected part of the diagram resulting from the contraction and depending on the renormalized coupling constants. This factorization property can be influenced by the choice of the large variables. We assume a finite charge renormalization: $g^2(s) \approx g_0^2 + (g^2 - g_0^2)(s/m^2)^{-\kappa}$ with $g_0^2 \ll 1$.

1) Let us consider first the target fragmentation region (*Fig. 2*). Asymptotic variables we choose, $2pq$, $-2p'q$, q^2. $u=(p-p')^2$ is small. For the forward scattering amplitude we can write the Mellin transform with respect to the large variables:

$$T = \frac{1}{(2\pi i)^3} \int_{-i\infty}^{i\infty} dj_1 dj_2 dj_3 \left(\frac{2pq}{q}\right)^{j_1+j_2} \left(-\frac{p'q}{qp}\right)^{j_2} (q^2)^{j_1+j_2+j_3} \times$$

$$\times F(j_1, j_2, j_3, u).$$

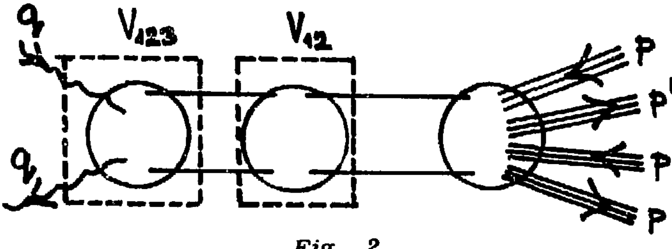

Fig. 2

The contraction V_{12} kills the pq- and p'q- dependence and generates the factor $[w^{-1}(j_1+j_2) - B_{\gamma\gamma}(0)]^{-1}$ in the function $F(j_1, j_2, j_3, u)$. The functions w and B can be calculated by perturbation theory. If the leading singularity is the pole at $j_1 + j_2 = \alpha(0)$ (Pomeranchuk-pole), then one gets $\omega^{\alpha(o)}$ Regge behaviour for $\omega \gg 1$. The contraction V_{123} kills the dependence of all large variables and in the weak bear coupling constant (WBCC) approximation generates the factor $(j_1+j_2+j_3-\eta_4-w(j_1+j_2))^{-1}$ in F, where the dimension of the contracted subgraph $\eta_4 \sim O(g_o^2)$, $w \sim O(g_o^2)$, so we get $(q^2)^{O(g_o^2)}$ asymptotics. In this way we get that in the ordinary Regge limit with $\alpha(0) = 1$ $f_{\gamma^* p \to hX} \sim \frac{1}{q^2} K(x_F, u)$, where x_F is the Feynman variable. The triple Regge region can be examined exactly in the same way as in "ordinary" inclusive processes.

2) In the central region (*Fig. 3*) the convenient choice of asymptotic variables is $2pq$, $-2p'q$, $-2pp'$, q^2. For the forward scattering amplitude we have the following representation:

$$T = \frac{1}{(2\pi i)^4} \int_{-i\infty}^{i\infty} \prod_{i=1}^{4} dj_i \left(\frac{pq \cdot m_o^2}{pp' \cdot p'q}\right)^{j_1} \left(-\frac{2p'q}{q^2}\right)^{j_1+j_2} \times$$

$$\times \left(-\frac{2pp'}{m_o^2}\right)^{j_1+j_3} (q^2)^{j_1+j_2+j_4} F(j_i),$$

m_o is a scale parameter. The contraction V_{12} kills the pq- and p'q - dependence and generates the factor $[w^{-1}(j_1+j_2) - B_{\gamma\gamma}(0)]^{-1}$ in F. With the Pomeranchuk-pole one gets $(-2pq/q^2)^{\alpha(o)}$ Regge behaviour. The contraction V_{13} kills the pq- and pp' -dependence and generates the factor $[w^{-1}(j_1+j_3) - B_{p\bar{p}}(0)]^{-1}$ in F. With the Pomeranchuk pole one gets $(-2pp'/m_0^2)^{\alpha(o)}$ Regge behaviour. The contraction V_{124} kills the $pq-$, $p'q$ and q^2-dependence and in the WBCC approximation generates the factor $[j_1+j_2+j_4-\eta_4-w(j_1+j_2)]^{-1}$ in F, so we get $(q^2)^{O(g_0^2)}$ asymptotics. In this way we get that in the conventional Regge limit $f_{\gamma^* p\to hX} \sim \frac{1}{q^2}\frac{ut}{s\cdot s_o}\int dj_1 \left(\frac{ss_o}{ut}\right)^{j_1} \chi(j_1, j_2, j_3)$. We have assumed that t/q^2 is large. In this case $\frac{ut}{s} = \frac{s^2-q^4}{s^2}p_T^2 - \frac{q^2(s-q^2)}{s\sqrt{s}}p_T \sim p_T^2$, and the large p_T^2-behaviour can be examined in the same way as in pure hadronic processes.

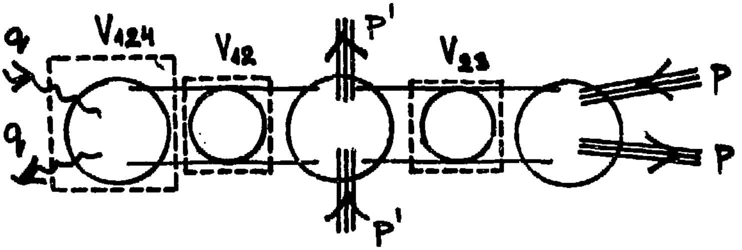

Fig. 3

3) We consider now the current fragmentation region (*Fig. 4*). The set of asymptotic variables is $2pq$, $-2pp'$, $-2p'q$, q^2. We write the forward scattering amplitude in the following form (a detailed analysis will be given in the next section):

$$T=\frac{1}{(2\pi i)^4}\int_{-i\infty}^{i\infty}\prod_{i=1}^{4} dj_i \left(\frac{2pq}{q^2}\right)^{j_1+j_2}\left(-\frac{pp'}{pq}\right)^{j_3}\times$$

$$\times\left(-\frac{2p'q}{q^2}\right)^{j_3}(q^2)^{j_1+j_2+j_3+j_4}F(j_i).$$

The contraction V_{12} kills the pq- and pp'-dependence and produces the factor $[w^{-1}(j_1+j_2)-B_{p\bar{p}}(0)]^{-1}$ in F. If $\omega \gg 1$ one gets $\omega^{\alpha(0)}$ Regge behaviour with the Pomeranchuk pole $j_1+j_2=\alpha(0)$. The contraction V_{1234} kills the dependence of all large variables, and in the WBCC approximation generates the factor $(j_1+j_2+j_3+j_4-\eta-[k(j_1+j_2)k(j_2+j_3)-1])^{-1}$

in F, where the dimension of the contracted subgraph $\eta_6=-1+O(g_0^2)$, $k(a)=w(a)+1$, so we get $(q^2)^{-1+O(g_0^2)}$ asymptotics.

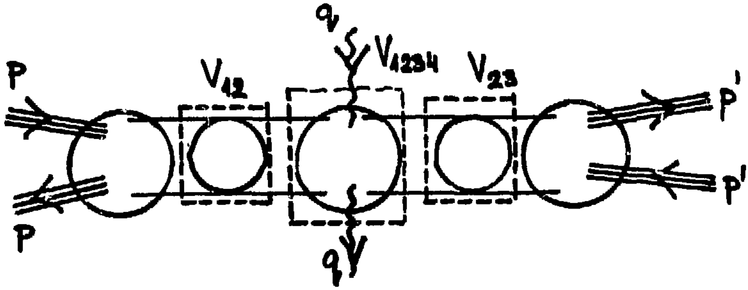

Fig. 4

In this way we get that in the ordinary Regge limit

with $\alpha(0)=1$ $\quad f_{\gamma^* p\to hX} \to (\frac{1}{q^2})^2 \tilde{K}(x_F, \frac{t}{q^2})$.

In the triple Regge region, where s, M^2 and $\frac{s}{M^2}$ are all large, one can get a more detailed representation/5/. The q^2-behaviour will be much more complicated because the dimension of the contracted subgraph strongly depends on the quark content of the detected hadron.

IV. In this section we discuss the current fragmentation region in detail. The forward scattering amplitude in α-representation has the following form:

$$T = \sum_{\text{diagr.}} H(g) \int_0^\infty \frac{\prod_\sigma d\alpha_\sigma}{D^2(\alpha)} C(\alpha, r_i) \exp i \left[\sum_{i=1}^4 \frac{r_i A_i(\alpha)}{D(\alpha)} \right],$$

where the functions $D(\alpha)$, $C(\alpha, r_i)$, $A_i(\alpha)$ are determined by the topology of the diagram, $H(g)$ is the product of the coupling constants, $r_1 = 2pq$, $r_2 = -2p'p$, $r_3 = -2p'q$, $r_4 = q^2$, the masses being neglected. We take the Mellin transform with respect to the large variables:

$$T = \frac{1}{(2\pi i)^4} \int_{-i\infty}^{i\infty} \prod_{i=1}^4 dj_i \left(\frac{2pq}{q^2}\right)^{j_1 + j_2} \left(-\frac{pp'}{pq}\right)^{j_2} \times$$

$$\times \left(-\frac{2p'q}{q^2}\right)^{j_3} (q^2)^{j_1 + j_2 + j_3 + j_4} F(j_i),$$

where

$$F(j_i) = \sum_{\text{diagr.}} H(g) \int_0^\infty \frac{\prod_\sigma d\alpha_\sigma}{D^2(\alpha)} \sum_{n_i} g(n_i, \alpha) \times$$

$$\times \prod_{i=1}^4 \left(i \frac{A_i(\alpha)}{D(\alpha)}\right)^{j_i - n_i} \Gamma(n_i - j_i) \frac{1 + e^{i\pi j_2}}{2}.$$

We have seen before that asymptotics in q^2 is determined by the leading singularity in the sum Σj_i, and in the WBCC approximation it is the pole at $\eta_6 + [k(j_1 + j_2) k(j_2 + j_3) - 1] = -1 + O(g_o^2)$. Then we get

$$F_{\text{lead}} = \left(\frac{1}{\mu^2}\right)^{\Sigma j_i + 1} \sum_{a,b} \tilde{f}_p^a (j_1 + j_2, \mu^2) \times$$

$$\times \tilde{T}_{ab} (j_i, g_o^2) \tilde{g}_b^h (j_2 + j_3, \mu^2),$$

where $\left(\frac{1}{\mu^2}\right)^{\Sigma j_i + 1} \tilde{T}_{ab}$ is the result of integration over small α-parameters ($\Sigma \alpha_\sigma < 1/\mu^2$, $\sigma \in V_{1234}$), that means the scale regime. The indices a, b stand for different intermediate states. To get this formula we have made the following procedure. Taking $\alpha_\sigma = \lambda_V \alpha'_\sigma$, $\Sigma \alpha'_\sigma = 1$ for the parameters of lines belonging to the subgraph V_{1234}, we have integrated for λ_V from 0 to $1/\mu^2$. The functions $A_i(\alpha)$, $D(\alpha)$ can be written in the smallest order of λ_V in the following form:

$$A_i(\alpha) = A_i(L)A_i(V)A_i(R), \quad D(\alpha) = D(L)D(V)D(R),$$

where L = left, R = right to the contracted subgraph V. This factorization property (rule 3) makes possible to get equation (1), which is the substitute of the operator product expansion in the current fragmentation region. We can now define the parton distribution and fragmentation functions /1,6/:

$$f^a_p(x,\mu^2) = \frac{1}{2\pi i}\int_{-i\infty}^{i\infty} \tilde{f}^a_p(j_1+j_2,\mu^2)\, x^{-(j_1+j_2)} d(j_1+j_2),$$

$$g^h_b(y,\mu^2) = \frac{1}{2\pi i}\int_{-i\infty}^{i\infty} \tilde{g}^h_b(j_2+j_3,\mu^2)\, y^{(j_2+j_3)} d(j_2+j_3) .$$

The inverse Mellin transformation gives the formulas:

$$\tilde{f}^a_p(j_1+j_2,\mu^2) = \int_0^1 \frac{dx}{x} x^{j_1+j_2} f^a_p(x,\mu^2),$$

$$\tilde{g}^h_b(j_2+j_3,\mu^2) = \int_0^1 \frac{dy}{y} y^{-(j_2+j_3)} g^h_b(y,\mu^2).$$

Putting these expressions into (1), we get the following result for the forward scattering amplitude:

$$T = \int_0^1\int_0^1 \frac{dxdy}{xy} \sum_{a,b} f^a_p(x,\mu^2)\, g^h_b(y,\mu^2) \times$$

$$\times \int_{-i\infty}^{i\infty} dj_1\, dj_2\, dj_3 \left(x\frac{2pq}{q^2}\right)^{j_1+j_2} \left(-\frac{pp'}{y\cdot pq}\right)^{j_2} \left(-\frac{2p'q}{y\cdot q^2}\right)^{j_3} \frac{1}{q^2} T_{ab} =$$

$$= \int_0^1\int_0^1 \sum_{a,b} \frac{dxdy}{xy} f_p^a(x,\mu^2)\, g_b^h(y,\mu^2)\, T_{a\gamma^* b\to a\gamma^* b}(p\to xp, p'\to \frac{p'}{y}),$$

This result is just the parton picture of the process in the current fragmentation region. The proton emits a parton of the type a with momentum xp. Absorbing the virtual photon it converts into parton of the type b, among the fragments of which is the hadron h. Taking the imaginary part, we get the following hard-scattering like formula:

$$p_0' \frac{d\sigma^{T,L}}{d^3p'} = \frac{1}{\text{flux}} \int_0^1\int_0^1 \sum_{a,b} \frac{dxdy}{xy} f_p^a(x,\mu^2) g_b^h(y,\mu^2) \times$$

$$\times \operatorname{Im} T_{a\gamma^{*T,L} b\to a\gamma^{*T,L} b}(xp, \frac{p'}{y}).$$

The amplitude $T_{a\gamma^* b\to a\gamma^* b}$ describes the small distance subprocess. It can be calculated by renormalization group improved perturbation theory using the fact, that the cross section is independent both of the splitting parameter μ and of the renormalization point of an R-operation /7/. In smallest order it is determined by the diagrams in *Fig. 5*. In the elementary subprocess the parton can get large transverse momentum. Using the parametrization $p' = \alpha(\beta p + q) + p_T$ where p_T is spacelike and orthogonal to p and q, and taking $\mu^2 = Q^2$, we get the following expression for the integrated structure functions:

$$\nu \tilde{W}_1 \sim \int_0^1\int_0^1 \frac{dxdy}{xy} [\sum_a S(a+\gamma^* \xrightarrow{T} a+\text{gluon}) f_p^a(x,Q^2)(g_a^h(y,Q^2) + g_{gl}^h(y,Q^2)) +$$

$$+ \sum_b S(\text{gluon}+\gamma^*\to b+\bar{b}) f_p^{gl}(x,Q^2)(g_b^h(y,Q^2) + g_{\bar{b}}^h(y,Q^2)).$$

and similarly for $\nu^2 W_2$ with $\gamma^* \xrightarrow{T} \gamma^{*T+L}$, where

$$S \sim \frac{\bar{g}^2(Q^2)}{Q^2} \delta((\frac{\alpha\beta}{y} - x)(\frac{\alpha}{y} - 1)\omega - (\frac{\alpha}{y} - 1)^2 + \frac{p_T^2}{q^2 y^2}).$$

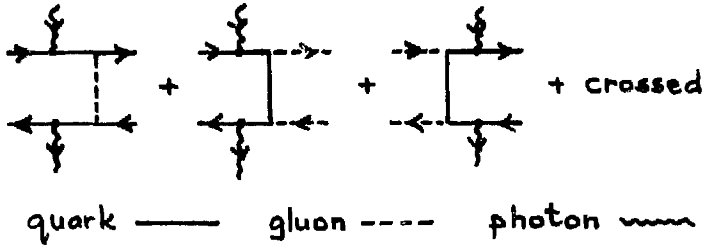

Fig. 5

Now it is easy to see that in the current fragmentation region $\nu^2 \tilde{W}_1$ and $\nu^3 \tilde{W}_2$ scale and are functions of α, β and p_T^2/q^2. If p_T^2/q^2 is small, then $\alpha = x_F - \frac{1}{(\omega-1)x_F} \frac{p_T^2}{q^2} + \dots$, $\beta = \frac{1}{\omega} - \frac{1}{\omega x_F^2} \frac{p_T^2}{q^2} + \dots$. In the case of $y \sim \alpha$, $p_T^2 \sim 0$, however, the argument of the δ-function gives an additional power of ν and with the proper scaling variables x_F, p_T^2 $\nu \tilde{W}_1$ and $\nu^2 \tilde{W}_2$ scales [3,8]. So, we can predict that there is a definite change in scaling somewhere between $p_T^2 \sim 0$ and $p_T^2 \sim Q^2$.

In the target fragmentation region we cannot get such a simple parton picture. Here $\nu \tilde{W}_1$ and $\nu^2 \tilde{W}_2$ scale and are functions of ω, $p'q/pq$ and pp' only. One can see that the reason of the difference of the scaling behaviours is the fact that the dimension of the contracted subgraph is -1 in the current and 0 in the target fragmentation region.

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