PHOTON DIFFRACTIVE DISSOCIATION IN DEEP INELASTIC SCATTERING

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I. INTRODUCTION

The new ep-collider HERA gives us not only the possibility to measure the structure function at very small x (up to $x \sim 10^{-4}$) where the vacuum singularity-pomeron is dominated but also to study the diffractive dissociation of virtual photon in deep inelastic ep-collision. The last problem is interesting due to two reasons.

From the reggeon phenomenology point of view the process of photon dissociation in deep inelastic scattering is the most direct way to measure the value of triple-pomeron vertex G_{3P}. The absorptive corrections are negligible here while in the case of hadron-hadron interactions the absorptive pomeron cuts contributions change crucially the inclusive cross sections in the triple-reggeon region ($x \to 1$) [1-3]. In particular, it was shown in refs. [2,3] that the value of the correct bare vertex G_{3P} may more than 4 times exceeds its effective value measuring in the triple-reggeon region and reaches the value of about 40-50% of the elastic pp-pomeron vertex. On the contrary in deep inelastic processes the perpendicular momenta q_t of the secondary particles are large enough. Due to this fact the absorptive corrections which proportional to $1/q_t^2$ are suppressed. Thus in deep inelastic reactions one can measure the absolute value of G_{3P} vertex in the most direct way and compare its value and q_t dependence with the leading log QCD predictions. For example in the perturbative QCD one expects the transverse energy jet distribution of the form dE_{tj}^2/E_{tjet}^4 in the G_{3P} rapidity reggeon ($y_{jet} \sim y_{3P}$) for photon dissociation events on the contrary to the conventional deep inelastic logarithmic distribution dE_t^2/E_t^2 for $\sigma_{in}(\gamma * p)$.

At last one can estimate the absorptive cut corrections in the small x region using the AGK cutting rules [4] and the diffractive dissociation cross section σ^D.

For simplicity in this paper we use the leading log approximation (LLA) keeping in our formulas only the maximum power of perpendicular momenta ($\ln Q^2$) and longitudinal ($\ln 1/x$) logs.

II. THE CROSS SECTION OF PHOTON DISSOCIATION (BORN APPROXIMATION)

Let us consider the process of photon diffractive dissociation into a three jets (quark, antiquark and gluon) as a typical and realistic example of a triple pomeron event in the deep inelastic collision. In the Born approximation this reaction is described by the diagrams of the fig.1 type. The generalization to the LLA case is evident and will be done below.

It is convenient to use the Sudakov variables [5]

$$q = \alpha Q' + \beta p + q_t; \quad d^4q = s/2 \; d\alpha d\beta d^2 q_t; \quad s = 2pQ'; \quad p^2 = Q'^2 = 0$$
$$Q' = Q + \beta_Q p; \quad \beta_Q = x_B = |Q^2|/s \tag{1}$$

and axial gauge $A_\mu p_\mu = 0$. The propagator of gluon with momentum k takes the form

$$d_{\mu\nu}(k)/k^2 = 1/k^2 \, [g_{\mu\nu} - (k_\mu p_\nu + p_\mu k_\nu)/(pk)] \tag{2}$$

The following equations are fulfilled in this gauge

$$p_\mu d_{\mu\nu} = 0; \quad Q'_\mu d_{\mu\nu}(k) = -k_{t\nu}/\alpha_k - 2\beta_k p_\nu/\alpha_k \simeq -k_{t\nu}/\alpha_k \tag{3}$$

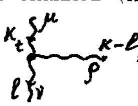
and at $k^2 = 0$, $k_\mu d_{\mu\nu}(k) = 0$, $d_{\mu\mu} = 2$, $d_{\mu\nu}(k) d_{\nu\rho}(k) = d_{\mu\rho}(k)$.

The logarithmic kinematics which essential for our process fig.1 satisfies the inequalities

$$\begin{gathered} |Q^2| \gg q_t^2 \gg k_t^2, m_t^2, \; l_t^2 \gg \mu^2 \\ \alpha_1 \simeq 1 \gg \alpha_2 \gg \alpha_k \gg \alpha_l, \alpha_m \\ \beta_l, \beta_m \gg \beta_Q \sim \beta_2 \gg \beta_1 \;, \end{gathered} \tag{4}$$

where the μ is some characteristic mass or the inverse size of a target-proton.

It is well known that the only longitudinal polarizations of t-channel (vertical in fig.1) gluons give the cross section not disappearing at small $x = |Q^2|/s \ll 1$. Due to this fact we multiply the upper end of the propagators $d_{\mu'\nu}$ by the $g_{\mu\mu'}$-tensor and retain in $g_{\mu\mu'} = (p_\mu Q'_{\mu'} + Q'_\mu p_{\mu'})/(pQ') + g^t_{\mu\mu'}$ the first term $p_\mu Q'_{\mu'}/(pQ')$ only. So the polarizations of the t-channel gluons are in proportion to p_μ at the upper ends and to $Q'_{\mu'} d_{\mu'\nu}(k)/(pQ') = -k_{\nu t}/(\alpha_k s/2)$ at the lower ends of the propagators. In this case (and gauge) the vertex emitting s-channel gluon (horizontal in fig. 1) is equal to



$$= \frac{-k_{t\mu}}{\alpha_k \cdot s/2} p_\nu \Gamma_{\mu\rho\nu} = -2k_{t\mu} \tag{5}$$

and the vertex of s-channel gluon 'k' - t-channel gluon 'l' interaction

$$\Gamma_{\rho\delta\sigma} p_\delta/(pQ') = 2\alpha_k g_{\rho\sigma}$$

coincides with the vertex of the classical current soft gluon ($\alpha_l \ll \alpha_k$) emission $j_\mu = 2k_\mu/(k-l')^2$. To fulfil the gauge invariance conditions simultaneously with the diagram fig.1a one needs to consider a group of analogous graphs (fig. 1b, 1c) and so on, where the upper end of any t-channel gluon joints to any quark line q_1 or q_2. The soft gluon l' (or l) emission by the quark line is described by the same classical current $J_\mu = 2q_{i\mu}/(q_i-l')^2$. Let us pick out the term $M_{\mu\sigma}$ in the photon dissociation amplitude, which is changed with the permutation of the gluons (l,l') upper ends. In the case of the graph of fig.1a we get $M^a_{\mu\sigma} = k_{\mu t}k_{\sigma t}/k^2$, and for other diagrams $M^b = -(k+l')_{\mu t}(k+l')_{t\sigma}/(k+l')^2$, $M^c = -(k-l)_{t\mu}(k-l)_{t\sigma}/(k-l)^2$ and

$M^d = k_{\mu t}k_{t\sigma}/k^2$. The square of the whole amplitude is equal to

$|M^\Sigma|^2 = |M^a + M^b + M^c + M^d|^2 = M^\Sigma_{\mu\sigma}(k,l)\, M^\Sigma_{\mu\sigma}(k,m) = 2R(k,l)R(k,m)\pi^2_,$
where we integrate over the angles of transverse momenta vectors $\vec{l}_t$
and $\vec{m}_t$ and introduce the definition

$$R(k,l) = \theta(l^2 - k^2) + l^2/k^2\ \theta(k^2-l^2) \qquad (6)$$

Finally, the photon dissociation cross section takes the form

$$\left.\frac{d\sigma^D(\gamma* \longrightarrow q\bar{q}g)}{dt}\right|_{t=0} = 2\Sigma_F e_F^2 \frac{4\pi\alpha_{e.m.}}{|Q^2|}\left[\frac{\alpha_s}{\pi}\right]^5 \phi(l)\phi(m)\frac{2[z^2+(1-z)^2]}{l^4 m^4 q_t^2} \qquad (7)$$

$$\frac{\alpha\beta_k}{\beta_k} dl^2 dm^2 d^2q_t d^2k_t dz R(k,l)R(k,m)$$

The factor 2 reflects the fact that not only the quark q_1 (as in fig.1, $\alpha_1 \simeq 1$) but the antiquark q_2 also can carry away almost the whole momentum of virtual photon. We sum over this two configuration here. e_F is the charge of the quark in the electron charge units; $\alpha_{e.m.} = 1/137$. The function $\phi(l)$ describes the emission of the two gluons l and l' by the proton. At $l_t >> \mu$ each individual valence quark emits its own pair of gluons l,l' and $\phi(l) \simeq 3$, but at small $l_t < \mu$ the function $\phi(l) \propto l_t^2 R_N^2$, where R_N is the radius of a nucleon.

At least $2[z^2+(1-z)^2]$ is the ordinary kernel of the Gribov-Lipatov-Altarelli-Parisi [6] evolution equation corresponding to the gluon (k) into a quark-antiquark pair (q_1q_2) transition; $z = \beta_Q/(\beta_Q+\beta_2)$.

The integration over l_t (or m_t) converges at $l_t \gg k_t$ ($m_t \gg k_t$) and has a logarithmic character at $l_t < k_t$, as the function $R(k,l) = l_t^2/k_t^2$ at $l_t < k_t$. Due to this reason the transverse momentum distribution of the slowest (in the proton-target rest frame) jet (k) takes the form dk_t^2/k_t^4 and the mean value $\langle k_{t,jet}\rangle$ is controlled by the lower limit of integration $-\mu$. The only possibility to have rather large $k_{t,jet}$ is to select the event with large k_t experimentally As one can see the k_t distribution in photon dissociation events is more soft than in conventional deep inelastic collision corresponding to diagrams fig.2. The inclusive spectra in the last case are proportional to dk_t^2/k_t^2 The longitudinal momentum distribution has the usual logarithmic form $d\beta_k/\beta_k = dE_k/E_{k,jet}$. It means that the distribution is $d\sigma^D/dM^2 \sim 1/M^2$. In the same way one can calculate the cross section of virtual photon dissociation into the two jets (quark and antiquark jets, see fig.3)

$$\left.\frac{d\sigma^D(\gamma^*\to q\tilde{q})}{dt}\right|_{t=0} = 2\sum_F e_F^2\frac{\alpha_{e.m.}}{|Q^2|}\left(\frac{\alpha_s}{2\pi}\right)^4 \Phi(l)\Phi(m)\,\frac{2[z^2+(1-z)^2]}{l_t^2 m_t^2 q_t^4}\times \tag{8}$$

$$\times\, dz(1-z)^2 d^2l_t d^2m_t d^2q_t\cdot 4/27$$

The quark-antiquark system (q_1q_2) is in the colour singlet state here and the coefficient 4/27 is the colour factor. The integration over l_t (or m_t) is logarithmic in the region $\mu \ll l_t$, $m_t \ll q_t$, analogously to the previous case (eq. (7)), and the quark jet transverse momentum spectrum is of the type of dq_t^2/q_t^4. Therefore one is needed to select artificially the events with large $q_t \gg \mu$ in order to have the possibility to use the perturbative QCD formulas.

III. THE LLA GENERALIZATION

To take into account the higher order α_s corrections in the LLA approximation it is enough to emit the arbitrary number of gluons (or quark-antiquark pairs) between the lines ll', mm' or t-channel gluons k,k^*. These additional partons are shown by the dash lines in fig. 4. The graph of fig.4 represents now the decay of the central "pomeron" (the ladder "k" described the photon dissociation into a group of quark and gluon jets) into the two ladders "ll'" and "mm". To get the diffractive dissociation cross section one has to change the factors $(4\alpha_s/3\pi)\int\phi(l)[dl_t^2/l_t^2]$ and $4\alpha_s/3\pi\int\phi(m)[dm_t^2/m_t^2]$ in eqs. (7),(8) by the gluon structure functions $xD_N^G(x=\beta_l,k^2,\mu^2)$ and $xD_N^G(x = \beta_m,k^2,\mu^2)$ correspondingly (the function $xD_N{}^G(x,k^2,\mu^2)$ describes the gluon distribution inside the target proton). The upper block of the virtual photon dissociation takes the form $\bar{x}D_{\gamma*}(Q^2,k^2, \bar{x} = \alpha_k = k_t^2/\beta_k s)$. In the functions D^G and $D_{\gamma*}$ we have written down the initial and the final virtualities k^2 and μ^2 or Q^2 and k^2. It seems that in leading log approximation where the photon dissociation block is described by the structure function $D_{\gamma*}$ containing a lot of partons the form of the amplitude $M^{(1)}$(with any permutations of the gluons (l,l') upper ends) changes crucially. But this is not the case. For deep inelastic scattering one has the usual LLA ordering conditions

$$|Q^2| \gg q_{2t}^2 \gg \ldots \gg k_t''^2 \gg k_t'^2 \gg k_t^2 , l_t^2, m_t^2. \qquad (9)$$

Due to this fact a soft gluon (l or l') interacts with the system of partons $q_1,q_2,\ldots,k''.k'$ which transverse size is of the order of $1/k_t'$ and is much smaller than the gluon Compton wave length $\lambda \sim 1/l_t$ as a whole. Within the accuracy of the order of $\sim l_t/k_t \ll 1$ such an interaction is equivalent to the soft gluon interaction with the individual particle with the same colour charge (which is equal to the minus colour charge of gluon "k"). In this sense the contributions of the graph fig.1 b,c and so on repeat exactly the amplitudes of gluon l (or l') - quark-antiquark (q_1q_2) interaction that have been discussed in the previous section. Thus the whole amplitude M^Σ retains its form: $|M^\Sigma|^2 = 2R(k,l)R(k,m)\pi^2$ (see eq.(6)).

IV. THE COMPARISON WITH THE HADRON DISSOCIATION

Let us consider the process of proton diffractive dissociation into a system of particles which includes a large k_t jet. Such a kinematics corresponds to UA-8 collab.experiment [7]. In this case we have an inversed ordering

$$\mu \ll q_{2t} \ll \ldots \ll k'_t \ll k_1 \qquad (10)$$

and the hard gluon l with the wave length $\lambda \sim 1/l_t \ll r \sim 1/k'_t$ (where r is the size of a system of partons $q_1 q_2, \ldots \tilde{k}'', \tilde{k}'$) really emits by the last parton with the largest transverse momentum k_t only. In another case the large momentum $l_t \sim k_t \gg k'_t$ destroyes the logarithmic integrations over $d^2k''_t/k''^2$ and so on inside the structure function $D_h(x=\alpha_k, k^2, \mu^2)$ playing the role of a central pomeron (function D_{γ_*}) here. As a result the amplitudes $M^{(i)}$ changes in the following way. Gluon "k" emits by the previous parton "k'" and this vertex does not depend on transverse momentum $k_{t\mu}$. The vector dependence of amplitudes $M^{(i)}$ is caused by the s-channel gluon k emission vertex (s) only. This leads to

$$\tilde{M}^{\Sigma}_{\mu}(k,l) = 2\frac{k_{t\mu}}{k_t^2} - \frac{(k+l)_{t\mu}}{(k+l)_t^2} - \frac{(k-l)_{t\mu}}{(k-l)_t^2}$$

and after the angle (l_t or m_t) integrations one gets the effective triple-pomeron vertex

$$G_{3P}^{(h)} \leftrightarrow |\tilde{M}^{\Sigma}|^2 = 4\theta(l_t^2 - k_t^2)\theta(m_t^2 - k_t^2)/k_t^2 \qquad (11)$$

In the kinematics described above the regions of logarithmic integrations $[\mu \ll l, m \ll k]$ are absent. The integrals over l_t, m_t converge at $l_t, m_t > k_t$ and are equal to $\int\theta(l_t^2-k_t^2)\theta(m_t^2-k_t^2)\, dl_t^2 dm_t^2/l_t^4 m_t^4 = 1/k_t^4$.

V. THE t-DEPENDENCE

The t-dependence of photon dissociation amplitude reveals itself in two factors. First of all the transferred momentum $\tilde{Q}^2=t$ ($\tilde{Q}=l-l'=m-m'$) restricts the interval of logarithmic integration over l_t (or m_t) inside the structure functions $D^G(x,k^2)$. At $t \neq 0$ the initial virtuality is equal to $\tilde{Q}^2=t$ instead of μ^2. Besides that the dissociation amplitude is multiplied by the proton-target form factors $F_N(\tilde{Q}^2)$ as the valence quark emitting the ladder $l, l'=l-\tilde{Q}$ (or $m, m' = m-\tilde{Q}$) gets the momentum $\tilde{Q}$. Thus the cross section

$$d\sigma^D/dt = [d\sigma^D(t=0)/dt]\,[F_N(t)\;D^G(x,k^2,t)/D^G(x,k^2,\mu^2)]^2.$$

VI. THE NUMERICAL ESTIMATIONS

1. At the HERA energies $s \sim 10^5$ GeV^2 a possibility arises to measure the processes with sufficient large $Q^2 \sim 10-100$ GeV^2 and very small $x = 10^{-4}-10^{-3}$. It is known that at so small x the structure function $D^G(x,l^2,\mu^2)$ increases rapidly due to a strong scaling violation. The maximum parton density is restricted by the unitarity constraint and is proportional to l^2/μ^2. This question had been discussed in detail in review [8] where it was shown that the $D^G(x,l^2) \propto l^2/\mu^2$ behaviour is correct in the region $l^2<q_0^2(x)$. At larger $l_t \gg q_0(x)$ the factor $D^G(x,l^2)/l^2 \propto dl^2/l^4$ and the integration over l_t in the expressions (7),(8) for diffractive dissociation cross section converges at $l_t \sim q_0(x)$ (or $m_t \sim q_0(x)$ in the case of the integration over m_t). Consequently the characteristic transverse momenta of a jet k_t is of the order of $q_0(x=\beta_1)$. The value of $q_0(x)$ calculated in the LLA increases with log 1/x accordingly to the regulation in $q_0^2(x) = 3.56\,\sqrt{\ln 1/x}$. A phenomenological analysis of semihard processes experimental data permits us to find the preasymptotic

corrections [9]

$$q_0^2(x) = Q_0^2 + \Lambda^2 \exp\,(3.56\ \sqrt{\ln 1/3x}) \qquad (12)$$

(where the $Q_0^2 = 2\ GeV^2$ and $\Lambda = 52 MeV$). The momentum $q_0(x)$ reaches the value of **3.2** and **4.8 GeV** at $x=2\cdot 10^{-3}$ and $6\cdot 10^{-4}$ respectively. It means that the typical transverse momentum of a quark jet is of the order of 3.2-4.8 GeV for the case of photon into a two quark jets dissociation at $|Q^2| = 100\text{-}30\ GeV^2$. The inclusive cross section $d\sigma^D(\gamma^* \to q\tilde{q})/dtdq_{1t}^2$ starts to decrease as dq_{t1}^2/q_{t1}^4 only at $q_{1t} > q_0(x)$, i.e. $q_{t1} > 3.2$ GeV for $Q^2 = 100\ GeV^2$ and $s = 10^5\ GeV^2$.

The experimental study of q_t-distribution in the virtual photon dissociation events is one of the best way to observe the predicted in the framework of perturbative QCD effect of parton density saturation [8] in the small "x" region (the word saturation means that the function $xD^G(x,l^2,\mu^2) \sim l^2/\mu^2$ at $l^2 < q_0^2(x)$ and does not increase faster than $\ln^2 x$ for further decrease of x). This is also a good and direct way to check the expression (12) for $q_0(x)$ values.

In the case of photon into a three jet dissociation ($\gamma^* \to q\tilde{q}g$) the characteristic value of $\bar{k}_{t\,jet} \simeq q_0(x = \beta_k)$ is expected to diminish with the decrease of the jet energy ($E_j = \alpha_k Q'$) (in the proton rest frame). For example, at $\alpha_k \simeq 10^{-3}$ the estimated value $\bar{k}_t \simeq 1.65$ GeV ($\beta_k \simeq 0.028$) and for $\alpha_k = 10^{-2}$ it is $k_t \simeq 2.4$ GeV ($\beta_k = 0.056$).

2. To estimate the number of intermediate partons inside the structure functions D^G and $D_{\gamma*}$ (i.e. inside the ladders ll' or "k") it is convenient to use the formula

$$\bar{n} = \alpha_s[\partial D(x,k^2)/\partial\alpha_s]/D(x,k^2).$$

In the double log approximation (DLA) the function $D^G(x,k^2) \sim I_1(v)/v$, where I_1 is the modified Bessel function

$v = \sqrt{(16N_c/b)\ln[\alpha_s(\mu^2)/\alpha_s(k^2)]\ln 1/x} \simeq \sqrt{(4N_c\alpha_s/\pi)\ln k^2/\mu^2 \ln 1/x}$, and $\alpha_s = 4\pi/b\ln k^2/\lambda^2$.

Choosing the values λ = 100 MeV, k_t= 3 GeV, μ = 1 GeV and b=9 (three sorts of light quarks - u,d,s) one gets $\bar{n}\simeq 1/4$ (for $s=10^5 GeV^2$ and β_k=0.03). It means that from the DLA point of view the expressions (7),(8) calculated in the Born approximation give the reasonable estimations for photon dissociation cross sections. Nevertheless the leading log 1/x corrections (of the type of $(\alpha_s \ln 1/x)^n$) are significant. If one uses the result of refs. [10,11] (where the term of the type of $(\alpha_s \ln 1/x)^n$ has been summed up) after differentiation of the corresponding amplitude Imf = $e^{\xi}/\sqrt{\xi}$ (here the $\xi = \ln(1/x)12\cdot \alpha_s \cdot \ln 2/\pi$) one gets $\bar{n}_k \simeq 0.9$ and $\bar{n}_l = \bar{n}_m \simeq 1.1$. Thus the leading log 1/x corrections (including the reggeization of the gluons [10]) are needed to be taken into account. I hope it will be done elsewhere.

3. But before this let us estimate the scale of photon diffractive dissociation cross sections using the Born approximation (eqs. (7) and (8)).

$$\left.\frac{d\sigma(\gamma^* \to q\tilde{q})}{dt\, dq_t^2}\right|_{t=0} = \frac{\alpha_s^4}{|Q^2|\, q_t^4} \ln^2(q_t^2/\mu^2)\cdot 1.2\ 10^{-4} \simeq \frac{6\cdot 10^{-6}}{|Q^2|q_t^4}$$

at $q_t^2=10\mu^2$ and α_s=0.31

$$\left.\frac{E_k d\sigma(\gamma^* \to q\tilde{q}g)}{dE_k dt\, dk_t^2}\right|_{t=0} = \frac{\alpha_s^5}{|Q^2|k_t^4}\ln(Q^2/k^2)\ln^2(k_t^2/\mu^2)1.5\cdot 10^{-3} \simeq \frac{5\cdot 10^{-5}}{|Q^2|k_t^4}. \quad (13)$$

This is not the small cross sections. Moreover we expect these cross sections should increase by the factor 10 -40 due to the increasing of structure functions D^G and $D_{\gamma*}$ in the region of small x if one takes into account the corrections of the order of $\alpha_s \ln 1/x$ as have been discussed just above.

The triple pomeron vertex G_{3P} in comparison with the elastic nucleon scattering vertex is equal to

$$G_{3P} = g_{el} \frac{2N_c}{3C_2} \frac{\alpha_s(k^2)}{\alpha_s(\mu^2)} \frac{\mu^2}{k^2}$$

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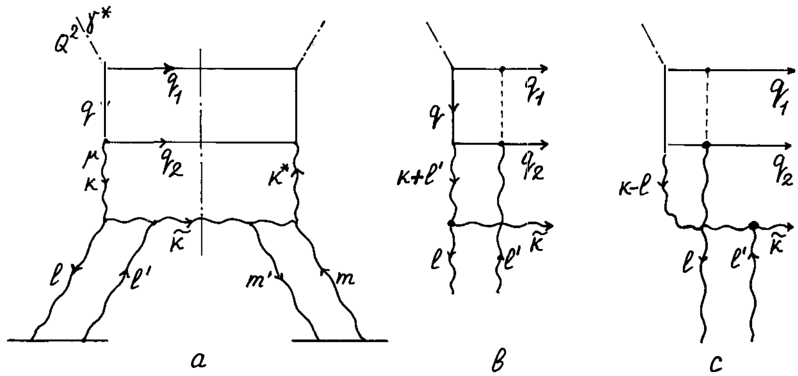

Fig. 1.

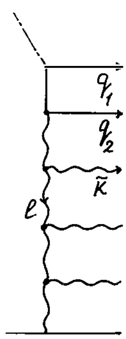

Fig. 2.

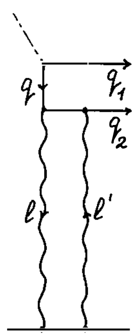

Fig. 3.

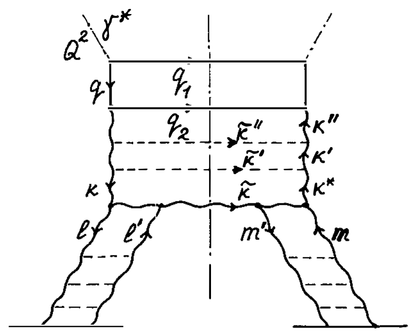

Fig. 4.