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INCLUSIVE SCATTERING AT LARGE TRANSVERSE MOMENTA IN MASSIVE QUANTUM ELECTRODYNAMICS

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1. INTRODUCTION

There is great interest in the study of hadron-hadron inclusive scattering at large transverse momentum p_T because of their possible relation to basic processes at small distances. The interest has arisen from the discovery at the CERN, ISR that the production of particles with large momenta transverse to the beam direction is larger than expected. Berman, Bjorken and Kogut [1], Blankenbecler, Bridsky and Gunion [2] consider the basic process to be an interaction of fundamental particles (partons) within the hadrons and these parton pictures predict the power law behaviour of the form

$$E \frac{d^3\sigma}{d^3p} \underset{\substack{s\to\infty \\ \frac{p_T}{\sqrt{s}}\ \text{fixed},}}{\longrightarrow} p_T^{-n} f\left(\frac{p_T}{\sqrt{s}}, \theta_{cm}\right), \tag{1.1}$$

where $\sqrt{s}$ and θ_{cm} are the centre of mass energy and scattering angle. The data analysis suggests that for $\sqrt{s} > 5$ *GeV* and $p_T > 0.5$ *GeV* the spectrum for $pp \to \pi^0 + X$ has a power law behaviour rather than falling exponentially. NAL experiment data [3] for $0.5 < p_T < 4$ *GeV* and for incident proton momentum between 50 and 4000 *GeV* is well fitted by

$$E \frac{d^3\sigma}{d^3p} = N(p_T^2 + 0.86)^{-4.5} \left(1 - \frac{2\frac{p_T}{\sqrt{s}}}{\sin\theta_{cm}}\right)^4, \tag{1.2}$$

where $N \simeq 5\ mb/GeV^2$. Again as $\sqrt{s}$ increase for fixed p the spectrum rises. In addition, the large p_T inclusive cross section somewhat increases with energy in violation of Feynman scaling. The presemt status of experimental data [4] for hadronic inclusive reactions is shown in *fig. 1*.

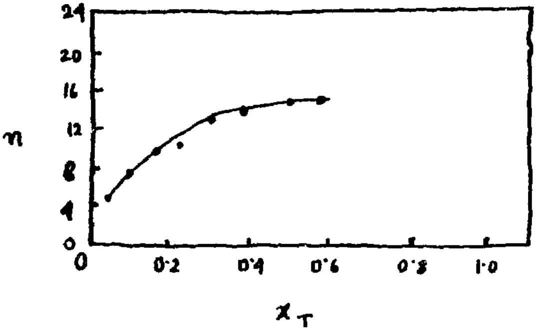


Fig. 1. The local value of n *at different values of* x_T *for proton.*

The experimental discovery of power law behaviour in inclusive hadron scattering at large transverse momentum leads to two basic theoretical attitudes. One is the parton picture [5] where high p_T particles are due to parton-parton wide angle scattering, such a model predicts $n=4$ the scattered partons are quarks and higher value of n if they are diquarks, mesons, etc. Again from the point of view of quantum field theory the experimental power law behaviour in inclusive reactions is quite interesting, because Feynman Graphs consist of integrals of rational functions and one would expect from perturbation theory a power law behaviour in addition to lns factors. Efremov [6] considered large p_T onclusive processes in field theory with scalar or pseudoscalar Yukawa coupling. By summing logarithmic terms of all diagrams, assuming finite charge renormalization he found

$$E\frac{d^3\sigma}{d^3p} \sim \frac{1}{(p_T^2)^n}\left(\frac{1}{x_T}\right)^{\alpha_a(0)+\alpha_b(0)-2}, \qquad (1.3)$$

where $n = 2+\epsilon_c(g_0^2)+2\epsilon(g_0^2)-[k(\alpha_a,g_0^2)k(\alpha_b,g_0^2)-1]$, ϵ_c, ϵ are anomalous dimensions, α_a, α_b are the leading Regge singularities of elastic amplitude, and $k(\alpha_a, g_0^2)$, $k(\alpha_b, g_0^2)$ are some functions of invariant charge.

In the same theory Roth[7] considered the process and found by summation of leading logarithmic diagrams:

$$E\frac{d^3\sigma}{d^3p} = \frac{1}{p_T^4} f\left(\frac{p_T}{\sqrt{s}}, \theta_{cm}, \xi_s\right),$$

where

$$\xi_s = -\frac{1}{5}\ln\left[1-\frac{5g^2}{16\pi^2}\ln s\right].$$

Roth has found a strong relation to deep inelastic scattering and light cone physics with the inclusive hadronic interactions in pseudoscalar Yukawa. This feature has already occured in many models. In this model there appeared an explicit dependence on the electroproduction structure function νW_2. In all cases he considered the spectrum falls for large fixed p_T and increasing energy.

In the present paper large p_T inclusive processes, e.g., $p+p \to p+X$, will be studied in renormalizable field theory of a vector field V_μ of mass μ interacting with a spin 1/2 field of mass m.

It appears, however, that the interest to this mechanism is connected with the well known fact that in the diffractive region $s \gg t$ it leads to the nonscaling behaviour $\sigma \sim$ constant in contrast with scalar gluon mechanism $\sigma \sim 1/s$.

The low order Feynman diagrams considered by us give the following result for large s

$$E\frac{d^3\sigma}{d^3p} \sim \frac{1}{s^2}\ln^2 s\, f(x_1, y_1),$$

which gives

$$E\frac{d^3\sigma}{d^3p} \simeq \frac{1}{p_T^4}\ln^2 s\, f(x_1, y_1),$$

i.e., the same power law behaviour as scalar gluon-exchange. This power law hopes to work in the region of small x_T as it seen from experimental data (*Fig. 1*).

We have not studied the relation between deep inelastic scattering and light cone method with inclusive scattering at large p_T. We hope to study this in a future paper.

II. KINEMATICS AND LOW ORDER DIAGRAMS

We will consider the inclusive processes, where all scalar products are large

$$2ab = s$$

$$2bp_1 = sx_1$$

$$2ap_1 = sy_1$$

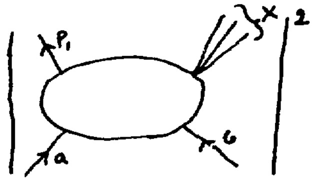
and $s \to \infty$.

We consider the limit $s \to \infty$, x_1, y_1 fixed which corresponds to large transverse momentum $p_T^2 \sim x_1 y_1 s$ and to $\cot^2(\theta/2) = x_1/y_1$, where θ is the angle at which p_1 emerges in the centre of mass frame. We should note that x_1, and y_1 are kinematically bounded by $p_T^2/s < x_1$, $y_1 < 1 - p_T^2/s$.

We are interested in calculating the Lorentz invariant inclusive differential cross section $E\, d^3\sigma/d^3p$ in a region where the transverse momentum is large. There are two basic ways in which this can be done. One way to calculate is to evaluate the contribution due to all exclusive channels and then to sum all them up (*fig. 2*).

But there is also another way to calculate the spectrum which points out an interesting symmetry and has been made plausible [8] that the inclusive differential cross section for the process $a + b \to p_1 +$ anything is related to discontinuity in $M^2=(a+b-p_1)^2=s(1-x_1-y_1)$ for the forward three-to-three amplitude as is shown in *fig. 2*.

We shall calculate all the diagrams necessary to determine the structure of leading diagrams. Once the structure of the behaviour is realised we shall assume that this structure will appear in all orders of perturbation theory.

$$E_1 \frac{d^3\sigma}{d^3p_1} = \frac{1}{s} \sum_{\substack{\text{final states} \\ X}}$$  $=$

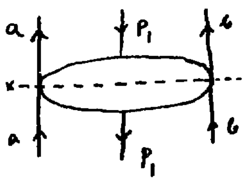
$= \frac{1}{s}$ 

Fig. 2. The inclusive spectrum either as a sum over exclusive processes or as the discontinuty in the forward three to three amplitude.

For the theory which we have considered to calculate $p+p \to p+X$ the following diagrams in the lowest order Feynman diagrams are possible.

There will be four diagrams similar to (i)-(iv) and those can be obtained from 4 (i)-(iv) interchanging p_2 and a by p_3 and b.

We study the inclusive differential cross section as a function of s, x_1, y_1

$$E_1 \frac{d^3\sigma}{d^3 p_1} = \left\{ \frac{g^8 (-1)^{2 s_{p1}}}{4(2\pi)^8 (2s_a+1)(2s_b+1)} \right\} \frac{1}{s} \int |M_{fi}|^2 \, \delta(a+b-p_1+c-p_2-p_3) \times$$

$$\times d^4c \, d^4p_2 d^4p_3 \, \delta(p_2^2-m^2)\delta(p_3^2-m^2) \, \delta(c^2-m^2), \tag{2.1}$$

where s_a , s_b , s_{p_1} are the spins of the particles carrying momentum a , b , p_1 , respectively.

First we will calculate *fig. 3.*

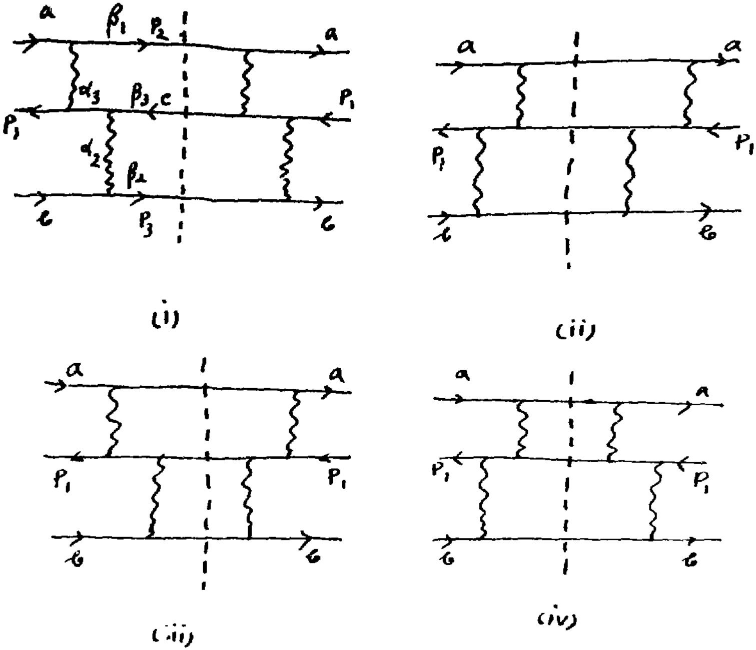


Fig. 3. Wigth order graphs for $p(a)+p(b)\rightarrow p(p_1)+X$ *at large* p_T.

From *fig. 3(i)-(iv)*

$$M_{fi} = \frac{\bar{W}(p_2)\gamma^{\mu} W(a)}{\{(a-p_2)^2-\mu^2\}}[\bar{W}(p_1)]\gamma_{\mu}\frac{\hat{b}-\hat{p}_3+\hat{c}+m}{(b-p_3+c)^2-m^2}\gamma_{\nu} + \tag{2.2}$$

$$+\gamma_{\nu}\frac{\hat{a}-\hat{p}_2+\hat{c}+m}{(a-p_2+c)^2-m^2}\gamma_{\mu}[W(c)]\frac{\bar{u}(p_3)\gamma_{\nu}u(b)}{\{(b-p_3)^2-\mu^2\}}.$$

We can write (2.1) as

$$E_1 \frac{d^3\sigma}{d^3p_1} = \frac{C_0}{s} \int |M_{fi}|^2 \, \delta^4(a+b+c-p_1-p_2-p_3)\,\delta(p_2^2-m^2)\,\delta(p^2-m^2)\times$$

$$\times\delta(c^2-m^2)\, d^4p_2\, d^4p_3\, d^4c \ , \tag{2.3}$$

where

$$C_0 = \frac{g^8 (-1)^{2s_{p_1}}}{4(2\pi)^\delta (2s_a+1)(2s_b+1)} . \tag{2.4}$$

From equation (2.2) one has

$$|M_{fi}|^2 = \frac{1}{\{(a-p_2)^2-\mu^2\}^2\{(b-p_3)^2-\mu^2\}} \mathrm{Sp}\,[(\hat{p}_2+m)\gamma_\mu (\hat{a}+m)\gamma_{\mu'}]\times$$

$$\times\,\mathrm{Sp}[(\hat{p}_3+m)\gamma_\nu(\hat{b}+m)\gamma_{\nu'}]\,\mathrm{Sp}[(\hat{p}_1+m)M(\hat{c}+m)\bar{M}] , \tag{2.5}$$

where

$$M = \gamma_\mu \frac{\hat{b}-\hat{p}_3+\hat{c}+m}{(b-p_3+c)^2-m^2}\gamma_{\nu'} + \gamma_{\nu'} \frac{\hat{a}-\hat{p}_2+\hat{c}+m}{(a-p_2+c)^2-m^2}\gamma_\mu . \tag{2.6}$$

We shall neglect mass terms in the numerator because these do not contribute to the leading term. Equation (2.5) can be written as

$$|M_{fi}|^2 = \frac{1}{\{(a-p_2)^2-\mu^2\}^2\{(b-p_3)^2-\mu^2\}^2} \mathrm{Sp}[\hat{p}_2\gamma_\mu \hat{a}\gamma_{\mu'}]\,\mathrm{Sp}[\hat{p}_3\gamma_\nu \hat{b}\gamma_{\nu'}]\times$$

$$\times\,\mathrm{Sp}\Big[\frac{\hat{p}_1\gamma_\mu\hat{s}_1\gamma_\nu\hat{c}\gamma_{\nu'}\hat{s}\gamma_{\mu'}}{G_1^2} + \frac{\hat{p}_1\gamma_\nu\hat{s}_2\gamma_\mu\hat{c}\gamma_{\mu'}\hat{s}\gamma_{\nu'}}{G_2^2} + \tag{2.7}$$

$$+\frac{\hat{p}_1\gamma_\mu\hat{s}_1\gamma_\nu\hat{c}\gamma_{\mu'}\hat{s}_2\gamma_{\nu'}}{G_1G_2} + \frac{\hat{p}_1\gamma_\nu\hat{s}_2\gamma_\mu\hat{c}\gamma_{\nu'}\hat{s}_1\gamma_{\mu'}}{G_1G_2}\Big] = A_1+B_1+C_1+D_1 ,$$

where

$$s_1=(b-p_3+c),\quad s_2=p_1+p_3-b,\quad G_1=(b-p_3+c)^2-m^2,$$

$$G_2=(a-p_2+c)^2-m^2.$$

For the delta function and propagator we shall use α representation, e.g.,

$$\frac{1}{(a-p_2)^2-\mu^2+i\epsilon}=\frac{1}{i}\int_0^\infty d\alpha_1\, e^{i\alpha_1\{(a-p_2)^2-\mu^2+i\epsilon\}}$$

$$\delta(c^2-m^2)=\frac{1}{2\pi}\left(\int_0^\infty e^{i\beta_3(c^2-m^2)}d\beta_3+\int_0^\infty e^{-i\beta_3(c^2-m^2)}d\beta_3\right).$$

To make computation easier before calculating spur we integrate over c, then we will have

$$E_1\frac{d^3\sigma}{d^3p_1}=\frac{C_0}{s(2\pi)^3}\int |M_{fi}|^2\delta(p_2^2-m^2)\delta(p_3^2-m^2)\delta(c^2-m^2)d^4p_2\times \tag{2.8}$$

$$\times d^4p_3\Big|_{c=p_1+p_2+p_3-a-b}$$

Therefore

$$\int A_1\,\delta(p_2^2-m^2)\delta(p_3^2-m^2)\delta(c^2-m^2)d^4p_2\,d^4p_3\Big|_{c=p_1+p_2+p_3-a-b}$$

$$=-\frac{C_0}{s(2\pi)^3}\int \alpha_1\alpha_2\alpha_3\,d\alpha_1 d\alpha_2 d\alpha_3\, d\beta_1 d\beta_2 d\beta_3\,\times \tag{2.9}$$

$$\times e^{i[(\alpha_2+\beta_2+\beta_3)p_3^2+2p_3\{(p_2+p_1-a-b)\beta_3-\alpha_2 b\}]}\times(\hat{A}_1)\times$$

$$\times e^{i[(\alpha_1+\alpha_3+\beta_1+\beta_3)p_2^2+2p_2\{\beta_3(p_1-a-b)-\alpha_1 a+p_1\alpha_3-a\alpha_3\}]}\times$$

$$\times e^{i\beta_3 s(1-x_1-y_1)}\,e^{-i\alpha_3 sy_1}\times$$

$$\times e^{-i\{(\beta_1+\beta_2+\beta_3)m^2+(\alpha_1+\alpha_2)(\mu^2-i\epsilon)+\alpha_3(m^2-i\epsilon)\}}.$$

Similarly there will be seven other terms interchanging $\beta_1, \beta_2, \beta_3$ by $-\beta_1, -\beta_2, -\beta_3$.

After inspection from the spur calculation we see that the maximum contribution will come from the expression evaluated from the spur calculation of (2.7)

$$
\begin{aligned}
&512[(p_1p_2)(ap_1)\{(bp_1+bp_2-ab)(p_3p_1-ap_3)+(bp_1-ab)(p_3p_1+p_3p_2-ap_3)+\\
&+p_3p_1(bp_1+bp_2-ab)+(p_3p_2)(bp_1+bp_2-ab)+(bp_1)(p_3p_1+p_3p_2-ap_3)+\\
&+(bp_2)(p_3p_1+p_3p_2-ap_3)\}-(p_1p_2)^2(ab)(p_3p_1+p_3p_2-ap_3)-\\
&-(p_1p_2)^2(ap_3)(bp_1+bp_2-ab)+(ap_1)^2(bp_1+bp_2-ab)(p_2p_3)+\\
&+(ap_1)^2(bp_2)(p_3p_1+p_3p_2-ap_3)\}]
\end{aligned}
\tag{2.10}
$$

and this expression will give us

$$
E_1\frac{d^3\sigma}{d^3p_1} \simeq \text{constant.} \tag{2.11}
$$

But we shall show constant terms are cancelled. We will assume $s \gg m^2$, $s \gg \mu^2$ and first scaling $(\beta_1, \alpha_1, \alpha_3, \beta_3) = \lambda(\beta_1, \alpha_1, \alpha_3, \beta_3)$ and letting $\lambda \to 0$, we have from (2.9) after applying the Mellin transform

$$
\begin{aligned}
&\frac{128\pi^4 s^4}{(2\pi)^3}x_1 y_1^2 C_0 \int \frac{\alpha_1^2\alpha_2^2\alpha_3}{K_0^3X_0^3}d\alpha_1 d\beta_1 d\alpha_2 d\beta_2 d\alpha_3 d\beta_3\left(x_1-1+\frac{\alpha_1}{X_0}\right)\times\\
&\times\frac{i}{2}\int_{\delta-i\infty}^{\delta+i\infty}\frac{dj(-isF)^j}{\Gamma(j+1)\sin\pi j(j+4)}\times\\
&\times e^{-i\{m^2(\beta_1+\beta_2+\beta_3)+\alpha_3(m^2-i\epsilon)+(\alpha_1+\alpha_2)(\mu^2-i\epsilon)\}}\,\delta(\alpha_1+\beta_1+\alpha_3+\beta_3-1),
\end{aligned}
\tag{2.12}
$$

where $X_0 = \alpha_1 + \beta_1 + \alpha_3 + \beta_3 - \frac{\beta_3^2}{K_0}$, $K_0 = \alpha_2 + \beta_2 + \beta_3$.

Again scaling $(\alpha_2, \beta_2, \alpha_3, \beta_3) = \lambda_1(\alpha_2, \beta_2, \alpha_3, \beta_3)$ and letting again $\lambda_1 \to 0$, we have

$$\frac{128\pi^4 s^4 x_1 y_1^2 C_0}{(2\pi)^3} \int \frac{\alpha_1^2 \alpha_2^2 \alpha_3 \, d\alpha_1 \, d\alpha_2 d\alpha_3 \, d\beta_1 \, d\beta_2 d\beta_3}{(\alpha_1 + \beta_1)^3 (\alpha_2 + \beta_2)^3} (x_1 - \frac{\beta}{\alpha_1 + \beta_1}) \times$$

$$\times \frac{i}{2} \int \frac{dj\,(-isF)^j}{\Gamma(j+1) \sin \pi j \, (j+4)^2} \times \tag{2.13}$$

$$\times e^{-i\{m^2(\beta_1 + \beta_2 + \beta_3) + \alpha_3(m^2 - i\epsilon) + (\alpha_1 + \alpha_2)(\mu^2 - i\epsilon)\}} \delta(\alpha_1 + \beta_1 - 1)\delta(\alpha_2 + \beta_2 + \alpha_3 + \beta_3 - 1),$$

Scaling again $(\alpha_3, \beta_3) = \lambda_2(\lambda_3, \beta_3)$ and letting again $\lambda_2 \to 0$ we have

$$\frac{128\pi^4 s^4 x_1 y_1^2 C_0}{(2\pi)^3} \int \frac{\alpha_1^2 \alpha_2^2 \alpha_3 \, d\alpha_1 d\alpha_2 d\alpha_3 d\beta_1 d\beta_2 \, d\beta_3}{(\alpha_1 + \beta_1)^3 (\alpha_2 + \beta_2)^3} (x_1 - \frac{\beta_1}{\alpha_1 + \beta_1}) \times$$

$$\times \frac{i}{2} \int \frac{dj\,(-isF)^j}{\Gamma(j+1) \sin \pi j \,(j+3)(j+4)^2} \times \tag{2.14}$$

$$\times e^{-i\{m^2(\beta_1 + \beta_2 + \beta_3) - i\epsilon + \alpha_3(m^2 - i\epsilon) + (\alpha_1 + \alpha_2)(\mu^2 - i\epsilon)\}} \times$$

$$\times \delta(\alpha_1 + \beta_1 - 1)\delta(\alpha_2 + \beta_2 - 1)\delta(\alpha_3 + \beta_3 - 1),$$

where

$$F = \frac{\beta_1 \beta_2 \beta_3}{(\alpha_1+\beta_1)(\alpha_2+\beta_2)} - \frac{y_1 \beta_1 (\alpha_3+\beta_3)}{(\alpha_1+\beta_1)} - \frac{x_1 \beta_2 \beta_3}{(\alpha_2+\beta_2)}. \tag{2.15}$$

We see that the constant terms which expected by leading singularity for $E_1 d^3\sigma/d^3p_1$ cancell in massive electrodynamics when we add all the diagrams of *fig. 3*. This is due to gauge invariance and renormalizability of the theory.

Now we have to consider the next leading terms from the diagrams 3(i)-(iv). We have seen the next maximum contribution $j = -4$ is also cancelled. The next term which will contribute to the inclusive differential cross section is found to be of the form

$$E_1 \frac{d^3\sigma}{d^3p_1} \propto \frac{1}{s^2} \ell n^2 \frac{s}{\mu^2} f(x_1, y_1). \tag{2.16}$$

Similarly, we have calculated $E_1 \frac{d^3\sigma}{d^3p_1}$ for the diagrams 4(iv)-(iv) and it is found to be of the form

$$E_1 \frac{d^3\sigma}{d^3p_1} \propto \frac{1}{s^2} \ell n \frac{s}{\mu^2} f(x_1, y_1). \tag{2.17}$$

As we have expected, the interchange p_1 by b or a will give less contribution. Similarly by interchanging a by b in *fig. 4(i)-(iv)* we obtain the same result but in place of x_1 we will have y_1, etc.

The analogous calculations for $e^+ + e^- \to e^+ + e^- + e^+ + e^-$ in quantum electrodynamics were done by Budnev et al. /9/, Baier and Fadin /10/, Kuraev and Lipatov /11/ in eight order and the total cross section is found to be $\sigma \underset{s\to\infty}{\sim} \ell n^3 s$.

III. DISCUSSIONS

From our calculations, we see that the differential cross section for the inclusive scattering $p+p \to \binom{p}{\pi} + X$ would behave as

$$E_1 \frac{d^3\sigma}{d^3p_1} \simeq \frac{1}{p_T^4} f(x_1, y_1).$$

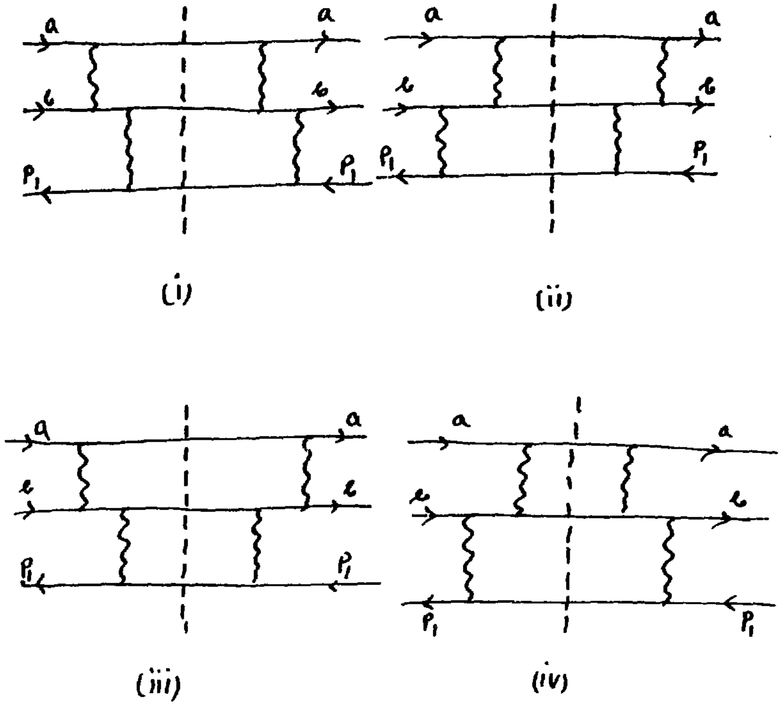


Fig. 4. Eigth order graphs by interchange $b \to p_1$.

This result cannot explain the experimental data which shows in the whole x_T region more quark decreasing dependence on p_T for the differential cross section in inclusive scattering at large transverse momenta (*fig. 1*). Only the region $x_T \ll 1$ give the hope that such sort of mechanism would work there.

We see from our calculation that we have maximum contributions from *fig. 3* with the parameters β_1 , β_2 , β_3 , a_1 , a_2 , a_3 are small. This strongly suggests that large p_T inclusive scattering samples a region where every particle is turned somehow close to every other particle.

So, the vector gluon exchange results for the high p_T inclusive process in the same behaviour as scalar gluon in spite of special "nondimensional" behaviour σ- constant in diffractive region $s \gg t$ [12].

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REFERENCES

1. *S.Berman, J.D.Bjorken and J.Kogut. Phys. Rev., D4, 3388 /1971/.*
2. *R.Blankenbecler, S.Brodsky and J.Gunion. Phys. Lett., 42B, 461 /1972/.*
3. *D.C.Carey et al. Phys. Rev. Lett., 33, 327, 330 /1974/.*
4. *J.W.Cronin et al. Phys. Rev., D11, 3105 /1975/.*
5. *S.D.Ellis and M.B.Kisslinger. Phys. Rev., D9, 2027 /1974/. S.Brodsky and G.Ferrar. Phys. Rev. Lett., 31, 1152 /1973/.*
   A.S.Kanofsky and K.F.Klenk. Phys. Rev. Lett., 31, 1323 /1973/. P.V.Landshoff and J.C.Polkinghoene. Phys. Rev., D8, 927 /1973/.
6. *A.V.Efremov. Yad. Fiz., 19, 196 /1974/; Fortschr. d. Phys., 22, 575 /1974/.*
7. *M.W.Roth. Phys. Rev., D12, 219 .1975/.*
8. *J.C.Polkinghorne. Nuovo Cimento, 7A, 555 /1972/.*
9. *V.M.Budnev, I.F.Ginzburg, G.V.Meledin and V.G.Serbo. Phys. Lett., 39B, 526 /1972/. Nucl. Phys.m 63B, 519 /1973/. Phys. Rep., 15C, 181 /1975/.*
10. *V.N.Baier and B.C.Fadin. JETP, 61, 477 /1971/.*
11. *E.A.Kuraev and L.N.Lipatov. Yad. Fiz., 16, 1060 /1972/.*
12. *H.Cheng and T.T.Wu. Phys. Rev., D1, 1775 /1970/.*

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