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LEPTON PAIR PRODUCTION IN TWO-DIMENSIONAL QCD

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ABSTRACT

Large q^2 lepton pair production is calculated in leading order in $1/N_c$ in two-dimensional quantum chromodynamics (QCD). From this confinement model the (Drell-Yan) parton model result is recovered.

Ref.TH.2266-CERN
12 January 1977

1. INTRODUCTION

The gauge theory of coloured quarks and gluons is considered as a promising candidate for a theory of strong interactions. Its light-cone behaviour is apparently confirmed by experimental findings. Not so much is known, however, on its long-distance properties. There is much hope that coloured objects like quarks and gluons are confined and will not appear in the spectrum of the theory. In order to prove this one should go beyond ordinary perturbation theory -- a difficult task. The best one can do up to now in 3+1 dimensions is to try to calculate with confinement enforced by a lattice^{1,2)} or bag^{3,4)} prescription. The crucial question⁴⁾ then remains -- not answered satisfactorily until now -- if in some limit (lattice constant $a \rightarrow 0$, bag constant $R \rightarrow \infty$) one returns to the original gauge theory.

A method particularly suited to study the confinement problem seems to be the $1/N_c$ ($U(N_c)$ colour group) expansion of 't Hooft⁵⁾. In the limit $N_c \rightarrow \infty$, $g^2 N_c$ fixed, the leading behaviour of a given amplitude is determined by the sum of the infinite set of planar diagrams (without fermion loops). This expansion resembles the dual loop expansion⁶⁾.

In 1+1 dimensions, the $1/N_c$ expansion has been successfully used to explore the spectrum of quantum chromodynamics (QCD) and the behaviour of pure hadronic and current amplitudes^{7,8,9,10)}. Since this model is super-renormalizable, it exhibits Bjorken scaling. Confinement is almost trivial in two dimensions due to the linearly rising gluon potential. Hence we do not learn much about the confinement in the 4-dimensional case. If the latter is supposed, however, 2-dimensional QCD might be a useful laboratory to explore its consequences.

Only a few quantities like $\sigma_{\text{tot } e^+e^-}$ or the Q dependence of the moments of structure functions can be traced back to the light-cone structure of the theory. Other processes like inclusive production in e^+e^- annihilation or lepton pair production in hadron collisions cannot be related to the light-cone behaviour in a simple way but still can be described by parton model ideas. The parton model, however, is incomplete since it does not explain the fate of partons in the final state. It is of tremendous interest to see to what extent parton model ideas survive in a confinement model which evidently has proper final states.

In this paper we study lepton pair production in inclusive hadron scattering in the framework of 2-dimensional QCD. Section 2 contains our calculation. In Section 3 we discuss our result which actually is of the Drell-Yan¹¹⁾ type. We also mention non-leading bremsstrahlung contributions.

2. CALCULATION OF THE LEADING CONTRIBUTION TO MASSIVE PHOTON PRODUCTION

In leading order in $1/N_c$ the production of a lepton pair in meson-meson high-energy scattering is accompanied by a single heavy resonance in the final state. The transition probability for this process is of the order $1/N_c$, like in the parton model.

The type of diagrams we are going to consider in this section is shown in Fig. 1. Other diagrams of the same order in $1/N_c$ are shown in Fig. 2. As will be argued in Section 3 these bremsstrahlung-type contributions are dynamically suppressed by a factor s^{-4} against the first ones. All planar gluon interactions have to be included in the frame of the (quark) graph of Fig. 1. (The graph as drawn in Fig. 1 contains a twist which, of course, has to be also in the gluons.) Practical calculations in 2-dimensional QCD have been carried out⁷⁻¹⁰⁾ in the light-cone gauge ($A_- = A_0 - A_1 = 0$). In this paper we follow that approach. We should mention, however, that the consideration of more general axial gauges led to some inconsistencies which are not resolved so far¹²⁾.

The rules for calculating skeleton diagrams of the type of Fig. 1 have been extensively discussed in the literature^{8,9)}. Introducing light-cone co-ordinates the + component of loop integrations can easily be performed. The γ algebra becomes trivial. It results in a kind of "time" ordered perturbation theory. All quark propagators are cancelled by the vertex functions. At the vertices the 't Hooft bound state wave function $\phi_n^{ab}(k_-/p_-)$ (Fig. 3) remains. The gluon propagator is $P(1/k_-^2)$ and is attached to quark lines which are turned around in "time". The $q\bar{q}$ scattering Green's function is

$$G^{ab}(x, y, \mu^2) = \sum_n \frac{\phi_n^{ab}(x) \phi_n^{ab}(y)}{\mu^2 - \mu_n^2} \quad (1)$$

where $\mu^2 = m_n^2 \pi / g^2 N$ is given in the natural unit $g^2 N / \pi$ (as all other masses with letter μ in the following). For current amplitudes we have to substitute (with $y = \ell_- / q_-$ in the notation of Fig. 4)

$$\begin{aligned} \phi_n^{aa}(y) &\rightarrow \frac{\pi}{g^2 N_c} \frac{q_-}{\sqrt{2}} \int_0^1 dx G^{aa}(x, y, \frac{\pi q^2}{g^2 N_c}) g_i^a(x, q) \\ &= \frac{\pi}{g^2 N_c} \frac{q_-}{\sqrt{2}} G_i^a(y, q) \end{aligned} \quad (2)$$

with

$$g_s^a(x, q) = \frac{-1}{\sqrt{2} q_-} \left(\frac{m_a}{x} - \frac{m_a}{1-x} \right) \quad (2a)$$

for scalar currents and

$$g_\mu^a(x, q) = g_{\mu+} - \frac{m_a^2}{q_-^2} \left(\frac{1}{x} + \frac{1}{1-x} \right) g_{\mu-} \quad (2b)$$

for vector currents if $0 \leq y \leq 1$. m_a is the quark mass. For momenta outside this range we also have a "direct" photon contribution.

The current amplitude considered here is similar to the large mass production amplitude studied in Ref. 9) as a candidate for a Pomeron in two dimensions (if squared and summed over the heavy particles). Again there are 11 diagrams -- now of "vector dominance" type, shown in Figs. 5a-5d if we choose a reference frame where^{*)}

$$\begin{aligned} p_{1-} &\sim \sqrt{s}, \quad p_{2-} \sim \frac{m_n^2}{\sqrt{s}} \\ q_- &= x\sqrt{s}, \quad q_+ = z\sqrt{s}, \quad q^2 = xz s \end{aligned} \quad (3)$$

There are two further diagrams not present in the pure hadronic case since they contain a "direct" photon coupling (Fig. 6). As shown in Ref. 9) the four graphs of the type 5d and also those of Fig. 5e cancel separately in a pairwise manner. This remains true for the current amplitude. The cancellation of 5a,b,c, which led to a vanishing Pomeron in Ref. 9, is however incomplete in our case, as shown in the following.

For large n the bound state wave function $\phi_n^{ab}(x)$ scales⁸⁾

$$\lim_{n \rightarrow \infty} \phi_n^{ab} \left(\frac{\xi}{\mu_n^2} \right) = \phi^a(\xi) \quad (4)$$

*) In the rest of the paper we will suppress the - index.

The integral equation

$$\left(1 - \frac{\mu_a^2 - 1}{\xi}\right) \phi^a(\xi) + P \int_0^\infty \frac{d\eta \phi^a(\eta)}{(\eta - \xi)^2} = 0 \quad (5)$$

can be used for the missing-mass vertex to obtain a cancellation of graph 5c with part of graph 5a. However, the corresponding scaling functions for the current vertex following from Eqs. (2)

$$\lim_{q^2 \rightarrow \infty} \frac{q^2}{g^2 N_c} G_i^a \left(\frac{g^2 N_c}{q^2}, q \right) = h_i^a(\xi, q) \quad (6)$$

fulfil integral equations⁸⁾ slightly different from Eq. (5)

$$\left(1 - \frac{\mu_a^2 - 1}{\xi}\right) h_i^a(\xi, q) + P \int_0^\infty \frac{d\eta h_i^a(\eta, q)}{(\eta - \xi)^2} = f_i^a(\xi, q) \quad (7)$$

with

$$f_s^a(\xi, q) = - \left(g^2 \frac{N_c}{\pi}\right)^{-1/2} \frac{q_+}{\sqrt{2}} \frac{\mu_a}{\xi} \quad (7a)$$

and

$$f_\mu^a(\xi, q) = \left(g_{\mu+} - g_{\mu-} \mu_a^2 \frac{q_+}{q_-} \frac{1}{\xi}\right) \quad (7b)$$

for scalar and vector currents, respectively.

Equation (6) is defined for positive q^2 only via an averaging procedure. A critical discussion in Section 3 will justify this step.

Because of the difference between Eqs. (5) and (7) on the right-hand side, the cancellation of graph 5(b) with the rest of graph (5a) is incomplete and we are left with a term

$$M_a + M_b + M_c = \frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \left(\frac{g^2 N_c}{s} \right)^{1/2} \mu_{n_2}^2 \phi_{n_1}^{ab}(x) \int_0^1 dy \phi_{n_2}^{ac}(y) \phi^c((1-y)(1-z)\mu_{n_2}^2) f_i^a(yz\mu_{n_2}^2, q) \quad (8)$$

The amplitudes M_a, M_b, M_c (averaged over the missing-mass spectrum) are given in Appendix A. The graphs of Fig. 6 (see there for notation) are calculated as

$$\begin{aligned} \tilde{M}_a + \tilde{M}_b &= \frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \left(\frac{g^2 N_c}{s} \right)^{1/2} \int dl dl' d\tau' \frac{1}{(\tau' + l')^2} \\ &\quad G^{ab} \left(\frac{l}{(1-x)\sqrt{s}}, \frac{l'}{(1-x)\sqrt{s}}, -z(1-x) \frac{s}{g^2 N_c/\pi} \right) h_i^a \left(\frac{l+q}{q-}, q \right) \\ &\quad \phi_{n_1}^{ab} \left(x - \frac{l}{\sqrt{s}} \right) \phi_{n_2}^{ac} \left(\frac{\tau'}{m_{n_2}^2/\sqrt{s}} \right) \left\{ \phi_n^{cb} \left(\frac{m_{n_2}^2/\sqrt{s} - \tau'}{(1-x)\sqrt{s}} \right) - \phi_n^{cb} \left(\frac{m_{n_2}^2/\sqrt{s} + l'}{(1-x)\sqrt{s}} \right) \right\} \end{aligned} \quad (9)$$

Introducing again scaling functions $\phi(\xi)$ we finally obtain

$$\begin{aligned} \left(l = \xi \frac{g^2 N_c/\pi}{\sqrt{s}}, l' = \xi' \frac{m_{n_2}^2}{\sqrt{s}}, \tau' = y \frac{m_{n_2}^2}{\sqrt{s}} \right) \\ \tilde{M}_a + \tilde{M}_b = \frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \left(\frac{g^2 N_c/\pi}{s} \right)^{1/2} \phi_{n_1}^{ab}(x) \int_0^\infty d\xi \int_0^\infty d\xi' \int_0^1 dy \\ \frac{1}{(\xi' + y)^2} \frac{1}{\pi^2} \int_0^\infty \frac{d\lambda}{\lambda+1} \phi^a(\lambda z \xi) \phi^a(\lambda z \xi' \mu_{n_2}^2) f_i^a(z \xi, q) \\ \phi_{n_2}^{ac}(y) \left\{ \phi^c((1-y)(1-z)\mu_{n_2}^2) - \phi^c((1-z)(1+\xi')\mu_{n_2}^2) \right\} \end{aligned} \quad (10)$$

For symmetry reasons the structure of the total amplitude -- the sum of Eqs. (8) and (10) -- is very simple. Using the parity invariance of the theory

$$\begin{aligned} M_S &\xrightarrow{P} M_S \\ M_{\pm} &\xrightarrow{P} M_{\mp} \end{aligned} \quad (11)$$

and recalculating the amplitudes in the new Lorentz frame we obtain relations similar to Eqs. (8) and (10) but with x and z interchanged. Thus the total amplitude is of the form

$$M_S = \frac{2\sqrt{\pi} c_s}{\sqrt{N_c}} \phi_{n_1}^{ab}(x) \phi_{n_2}^{ac}(z) \quad (12)$$

for the scalar current and

$$M_{\mu} = \frac{2\sqrt{\pi} c_v}{\sqrt{N_c}} \frac{m_a}{\sqrt{s}} \left(\frac{g_{\mu+}}{z} - \frac{g_{\mu-}}{x} \right) \phi_{n_1}^{ab}(x) \phi_{n_2}^{ac}(z) \quad (13)$$

for the vector current. Remark that the latter is divergenceless as it should be. The constants c_s, c_v cannot depend on n_1, n_2, m_b, m_c since these dependences are completely described by the ϕ as is easily seen from Eqs. (8) and (10) plus the symmetry argument. To fulfil (12) and (13) the following remarkable relations have to be true

$$\begin{aligned} 2\sqrt{\pi} c_s \phi_{n_2}^{ac}(z) &= -\frac{2}{\sqrt{\pi}} \int_0^{\infty} d\xi' \int_0^1 dy \frac{1}{\pi} \int_0^{\infty} \frac{d\lambda}{\lambda+1} \frac{1}{(\xi'+y)^2} \\ &\quad \phi^a(\lambda z \xi' / \mu_{n_2}^2) \phi_{n_2}^{ac}(y) \left\{ \phi^c((1-z)(1-y) \mu_{n_1}^2) - \phi^c((1-z)(1+\xi') \mu_{n_2}^2) \right\} \\ &+ \frac{2\mu_a}{\sqrt{\pi}} \int_0^1 \frac{dy}{y} \phi_{n_2}^{ac}(y) \phi^c((1-y)(1-z) \mu_{n_2}^2) \end{aligned} \quad (14)$$

and

$$\begin{aligned}
 2\pi C_V \phi_{n_2}^{ac}(z) &= \frac{4}{\sqrt{2\pi}} \int_0^\infty d\xi' \int_0^1 dy \frac{1}{\pi} \int_0^\infty \frac{d\lambda}{\lambda(\lambda+1)} \frac{1}{(\xi'+y)^2} \\
 &\quad \phi^a(\lambda z \xi' \mu_{n_2}^2) \phi_{n_2}^{ac}(y) \{ \phi^c((1-y)(1-z) \mu_{n_2}^2) \\
 &\quad - \phi^c((1-z)(1+y) \mu_{n_2}^2) \} \\
 &\quad - \frac{4}{\sqrt{2\pi}} \frac{\mu_{n_2}^2}{\mu_a} z \int_0^1 dy \phi_{n_2}^{ac}(y) \phi^c((1-y)(1-z) \mu_{n_2}^2)
 \end{aligned} \tag{15}$$

Here we used the relations ^{*)}

$$\int_0^\infty \frac{d\xi}{\xi} \phi^a(\xi) = \frac{1}{\mu_a^2} \int_0^\infty d\xi \phi^a(\xi) = \frac{\pi}{\mu_a} \tag{16}$$

discussed in Ref. 8).

Relations (14) and (15) do not follow trivially from 't Hooft's bound state equation. Like in previous calculations ^{8,9)} it would be nice to see explicitly (with some powerful identities) the covariant character of the calculations. We were not able to prove them in general and hence checked them with computer calculations. The behaviour for $z \rightarrow 1$ is obviously the same on both sides due to the behaviour

$$\begin{aligned}
 \phi^c(\xi) &\sim C^c \xi^{\beta_c} \text{ for } \xi \rightarrow 0 \\
 \phi_{n_2}^{ac}(z) &\sim C^{ca} (1-z)^{\beta_c} \text{ for } z \rightarrow 1
 \end{aligned} \tag{17}$$

For $z \rightarrow 0$ in Eq. (15) only the first term contributes in leading order and is again consistent with the left-hand-side. In Eq. (14) the integral over λ can be written as

*) We differ from Ref. 8) (second reference) by a factor $\sqrt{2}$ due to an incorrect counting of contributing resonances in that paper. We have checked Eq. (16) on the computer.

$$\int_0^{\infty} \frac{d\lambda}{\lambda+1} \phi^e(\lambda \pm \xi' \mu_{n_2}^2) = - \int_0^{\infty} \frac{d\lambda}{\lambda(\lambda+1)} \phi^e(\lambda \pm \xi' \mu_{n_2}^2) + \frac{\pi}{\mu_a} \quad (18)$$

The term corresponding to the second term in Eq. (18) has to cancel with the second term in Eq. (14). This cancellation has been checked in a computer calculation. The rest then gives the same behaviour for $z \rightarrow 0$ as the left-hand-side. In the limit $z \rightarrow 0$ we obtain the equation

$$2 \sqrt{\pi} c c^{ac} = \frac{2}{\pi^{3/2}} (\mu_{n_2}^2)^{\beta_a} \int_0^{\infty} d\xi \int_0^1 dy c^a \xi^{\beta_a} \phi_{n_2}^{ac}(y) \frac{1}{(\xi+y)^2} \left(\frac{\pi}{\sin \pi \beta_a} \right) \{ \phi^c((1-y)\mu_{n_2}^2) - \phi^c((1+\xi)\mu_{n_2}^2) \} \quad (19)$$

also obtained in the derivation of the imaginary part of an amplitude in the Regge limit in Ref. 9). We evaluated the constant $c = c_v/\sqrt{2} = c_s$ on the computer. It turns out to be independent also of index a and compatible with 1.

3. DISCUSSION

Since the wave functions $\phi_n(x)$ are normalized as $\int_0^1 \phi_n^2(x) dx = 1$, $\phi_n^2(x)$ can be interpreted as a probability distribution of quarks in the meson. This is in agreement with the calculation of the deep inelastic structure function. Equations (12) and (13) with $c_s = c_v/\sqrt{2} = 1$ then give just the parton model (Drell-Yan) result that the cross-section for lepton pair production is given by the product of probabilities to produce a quark and an antiquark out of bound states 1 and 2 times the cross-section for the annihilation of this $q\bar{q}$ system with a lepton pair (Fig. 7).

Like in Ref. 8) one could argue that the leading $1/N_c$ contribution is obtained from the region of $G(x, y, q^2)$ where the resonance width is important since the smoothing of a propagator square is required in the cross-section calculation. However, in the case considered here, all the high-mass "vector mesons" decouple

from the amplitude in leading order in s because of the cancellation of the purely hadronic amplitudes following from relations of type (5). The right-hand side of Eq. (7) is obtained from the integrated completeness relation for the ϕ_n $[\sum_n \phi_n(x) \phi_n(y) = \delta(x-y)]$ and hence the collective effect of all resonances coupling to the current is important. One can take out the resonances with mass near q^2 -- these are the poles for which width corrections are important, without effecting the completeness sum in leading order since⁸⁾

$$\int_0^1 dx \phi_n(x) \sim \frac{1}{\mu_n^2} \quad (20)$$

for large n .

The question if experiments support the factor $1/N_c$ in front of the Drell-Yan result (without colour) is not completely settled although recent experiments can be interpreted in favour of such a colour factor¹³⁾. It is interesting to speculate what might happen if transverse dimensions could be taken into account in a calculation not completely different from the one presented here in two dimensions. Actually there exist attempts to first solve the two-dimensional world and to take into account transverse mixing afterwards²⁾. It is conceivable that the cancellation leading to the decoupling of high-mass resonances becomes incomplete, and that we obtain an additional piece analogue to the Pomeron which is now of order unity in the $1/N_c$ expansion if we follow the arguments about width effects in the averaged propagator square presented in Ref. 10).

For missing masses $M_x^2 \sim O(1)$ we have $x = 1 - O(1/s)$ or $z = 1 - (1/s)$. The cross-section is suppressed by factors $s^{-2\alpha_b-1}$ and $s^{-2\beta_c-1}$, respectively. The bremsstrahlung contributions of Fig. 2 can also be calculated in 2-dimensional QCD. In a frame with $p_1 \sim \sqrt{s}$ (see Eq. (3)) we have nine graphs for Fig. 2a drawn in Fig. 7. Since they all contain at least one large momentum transfer (through a gluon or the G) they are suppressed by a factor s^{-2} . Hence we have a suppression factor s^{-4} in the cross-section. The graphs corresponding to Fig. 2b are separately gauge invariants and are also suppressed by s^{-2} for symmetry reasons^{*)}.

In summary, we have calculated massive photon production in a (two-dimensional) quark-gluon theory which explicitly shows confinement. We recover the general structure as well as the normalization of the corresponding parton model (Drell-Yan) result. This gives an explicit example that final state interactions necessary to prevent quarks from appearing in the final state may not destroy the simple quark-parton results.

*) Although the cancellation of the "direct" photon graphs proportional to $s^{-\beta}$ appearing now in ref. system (3) is non-trivial.

Of course, the ultimate goal is to perform calculations in four dimensions. Calculations of the type presented here may serve as a laboratory for this difficult project.

Details on the numerical results and the available programs may be obtained from the authors.

Acknowledgements

We would like to thank D. Amati, J. Ellis and J. Ranft for discussions, Eva Edberg for her help with the computer calculation, and D. Amati for reading the manuscript.

APPENDIX

Here we list the expressions corresponding to the graphs of Fig. 5a,b,c.

$$M_a = \frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \frac{q_-}{\sqrt{g^2 N_c \pi}} \int_0^1 dy \phi_{n_1}^{ab}(x) \phi_{n_2}^{ac}(y) \\ \phi^c((1-y)(1-z)\mu_{n_2}^2) G_i^a(yz\mu_{n_2}^2) \left\{ \frac{\mu_a^2-1}{y} + \frac{\mu_c^2-1}{1-y} - \mu_{n_2}^2 \right\}$$

$$M_b = -\frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \frac{q_-}{\sqrt{g^2 N_c \pi}} \int_0^1 dy \int_0^\infty d\eta_b \phi_{n_1}^{ab}(x) \phi_{n_2}^{ac}(y) \\ \phi^c((1-y)(1-z)\mu_{n_2}^2) G_i^a(\eta_b z \mu_{n_2}^2) / (y - \eta_b)^2$$

$$M_c = -\frac{1}{\sqrt{N_c}} \frac{4}{\sqrt{2\pi}} \frac{q_-}{\sqrt{g^2 N_c \pi}} \int_0^1 dy \int_0^\infty d\eta_c \phi_{n_1}^{ab}(x) \phi_{n_2}^{ac}(y) \\ \phi^c(\eta_c(1-z)\mu_{n_2}^2) G_i^a(yz\mu_{n_2}^2) / (1-y-\eta_c)^2$$

in the frame (3) with

$$y = \frac{\ell - q_- - p_{2-}}{p_{2-}} \quad \text{in (a) and}$$

$$y = \frac{\ell'}{p_{2-}} \quad \text{in (b) and (c)}$$

and

$$\eta_b = \ell \frac{\sqrt{s}}{m_{n_2}^2} \quad \eta_c = (\ell + p_{2-} - q_{1-}) \frac{\sqrt{s}}{m_{n_2}^2}$$

A factor of $(1/g^2 N_c \pi)^{\frac{1}{2}}$ for the density of state normalization in the missing-mass channel is included.

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Figure captions

- Fig. 1 : Leading quark graph for inclusive virtual photon production.
- Fig. 2 : Bremsstrahlung quark graphs.
- Fig. 3 : Bound state vertex.
- Fig. 4 : Pole dominated photon vertex substituting the bound state vertex.
- Fig. 5 : Leading quark-gluon graphs with pole dominated photon vertex contributing to inclusive virtual photon production. (d) and (e) represent four diagrams each with different single gluon interactions.
- Fig. 6 : Leading quark gluon graphs with "direct" photon coupling.
- Fig. 7 : Quark-gluon bremsstrahlung diagrams.

FIG.1

FIG.2

FIG.3

FIG.4

FIG.5

FIG. 6

FIG. 7