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MASSIVE LEPTON PAIR PRODUCTION IN HADRON-HADRON COLLISIONS

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ABSTRACT

We study the production of massive lepton pairs in hadron-hadron collisions in the framework of the massive quark model, the general field theoretical approach to confined quarks. Like for the other aspects of deep inelastic physics, our description is seen to agree remarkably well with presently available data.

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1. INTRODUCTION

Massive lepton pair production in hadron-hadron collisions has, no doubt, had a rôle of paramount importance in the subnuclear physics of the last decade. Pioneered at CERN fifteen years ago¹⁾, this line of research has scored the remarkable successes of the discovery of the J/ψ particle²⁾ and of the T, T' particles³⁾, in addition to the revelation of very interesting aspects of the lepton pair continuum, which is now being explored in a systematic way^{*)}. As for the future prospects of this field there is a good chance that the first Z^0 's, if they exist in the predicted mass range, will be produced in hadronic collisions and detected in their lepton pair decay.

On the theoretical front, the first analyses^{4), 5)} of this process are at the origin of a number of theoretical developments whose impact in the physics of the last decade cannot be underestimated: the foundation of light cone physics⁵⁾ and of the Mueller-Regge⁶⁾ approach to high energy inclusive reactions.

It was, in fact, the attempt to extend the notions of Bjorken scaling to the production of a massive (large Q^2) timelike current, that led to an operator formulation of the product of two currents when their separation x is close to the light cone $x^2 = 0$. The importance of the light cone operator product expansion in linking the asymptotic freedom property of perturbative QCD to experimental observations in deep inelastic scattering (DIS) is too well known to need any further elaboration⁷⁾.

As for the massive lepton pair production it has been stated⁸⁾, and is widely believed, that this process is not dominated by the light cone singularities of the product of two electromagnetic currents. We would like to point out that this statement is inaccurate. What is correct, is that the leading singularity due to the peculiarities of strong interactions at high energy (Regge behaviour, strong p_T cut-offs, etc.) does not dominate for large Q^2 ; while it is the next-to-leading singularity^{**)} that yields the asymptotic behaviour. This is why the first theoretical analysis^{4), 5)} of massive lepton pair production did not succeed in accounting for the experimental observations, while the simple parton model⁹⁾ constituted a valid, first approximation approach to its physics. The former analysis dealt only with the leading singularity, while the latter picked up the next-to-leading singularity in a very simple (and somewhat unrealistic) theoretical framework.

*) The latter aspect owes a great deal to the work of Leon Lederman.

**) i.e., the operator of "twist"⁴⁾.

That the parton model result could have a validity well beyond its limited scope was first shown in the context of the Massive Quark Model (MQM) several years ago¹⁰⁾. In this paper we wish to take up again the MQM analysis, and apply it, together with our previous results on DIS^{11),12),13)}, to the prediction of massive lepton pairs' continuum production in different collisions. This paper concludes a series^{11),12),13)}, devoted to the detailed application of the MQM to deep inelastic phenomena; the success of our analysis is, in our opinion, a further confirmation of the validity of the general approach of the MQM and of its more detailed version which has been called Quark Geometrodynamics¹⁴⁾ (QGD).

2. THE MQM ANALYSIS OF LEPTON PAIR PRODUCTION

The process we wish to study and its kinematics are depicted in Fig. 1. h and h' denote two hadrons, $(\ell, \bar{\ell})$ is a lepton pair and $(\vec{k}_T + \vec{k}'_T)^2 \ll Q^2$ is assumed throughout the paper. In the limit $(m_\ell^2/Q^2) \rightarrow 0$, we can write the cross-section for this process:

$$\frac{d\sigma}{dQ^2 d^4q} = \frac{\alpha^2}{12\pi^3} \frac{1}{Q^2 s} S(q^2 - Q^2) w_\mu{}^\mu; \quad (2.1)$$

where

$$w_{\mu\nu} = - \int d^4x e^{-iqx} \langle h(p_1) h'(p_2) | J_\mu(x) J_\nu(0) | h(p_1) h'(p_2) \rangle \quad (2.2)$$

is the hadronic tensor relevant for such process. Its MQM diagram is drawn in Fig. 2. In the limit $s \rightarrow \infty$, $s^* = s[1 + \tau(1 - 1/x_1 - 1/x_2)] \rightarrow \infty$, the diagram in Fig. 2 is dominated by the Mueller-Regge expansion of Fig. 3, where the sum runs over the leading Regge trajectories of hadronic physics (Pomeron, vector mesons, ...).

In the region where experimental data are presently available¹⁵⁾⁻¹⁹⁾ ($\sqrt{s} \approx 20$ GeV, $\div 60$ GeV) Regge poles with $\alpha(0) = \frac{1}{2}$ (vector mesons, tensor mesons) give a contribution proportional to $1/\sqrt{s^*}$, i.e., a small subasymptotic correction which we neglect in view of the present quality of the experimental data. Thus in the following only the Regge singularity with $\alpha(0) = 1$ (the Pomeron) shall be considered.

As for the meaning of the diagrams in Fig. 3, in order to clear up a frequent misunderstanding, we would like to stress that they should be interpreted in a way very different from the final state interaction diagrams of the parton model²⁰⁾, which are known to give no contribution to the eight point function discontinuity $W_{\mu\nu}$ due to the AGK cancellation²¹⁾. For the MQM there can be no cancellation due to the fact that no cut can be performed across quark legs, which by the confinement assumption cannot carry any singularity. Thus in the diagrams of Fig. 3 only one cut is possible, whose positive definiteness is obvious.

In the Mueller-Regge expansion of Fig. 3 the $q\bar{q}$ - $h\bar{h}$ -Pomeron vertex appears (Fig. 4) which has already been encountered in the DIS case. It should be noted that, as far as the p_T dependence of F_q^h is concerned, we need only assume a fast decrease (as usual in the Mueller-Regge analysis) to obtain scaling and the asymptotic validity of the Callan-Gross relation^{*)}. From our recent analysis of DIS¹¹⁾ it turns out that, if we define

$$\int F_{q\bar{q}}^h(x, p_1^2) d^2 p_1 = \frac{(2\pi)^4 b}{\beta} f_{q\bar{q}}^h(x), \quad (2.3)$$

where β is the residue of the $q\bar{q}P$ vertex and b is the exponential slope of the forward peak, the DIS structure function $F_1^{eh}(x)$ is given by:

$$F_1^{eh}(x) = \text{Tr} [Q^2 (f_q^h(x) + f_{\bar{q}}^h(x))] \quad (2.4)$$

We are now ready to compute the hadronic tensor $W_{\mu\nu}$. Using the MQM rules laid down in Ref. 10)**), we obtain

*) In Ref. 10) the following parametrization was adopted:

$$F_q^h(x, p_T) = 2(2\pi)^3 \beta_h \frac{f_q^h(x)}{x} e^{-\mu p_1^2} \frac{\mu}{\pi}$$

**) The quark propagator, assuming $m_q \approx 0$, is approximated by $(2\pi) \not{p} \delta(p^2)$.

$$\begin{aligned}
 W_{\mu}^{\mu} = & -\frac{1}{3} \int \frac{d^4 \ell}{(2\pi)^4} \int \frac{d^4 \ell'}{(2\pi)^4} \text{Tr} \left[\gamma^{\mu} \left(-\frac{q}{2} + \ell \right) \gamma^{\sigma} \left(-\frac{q}{2} + \ell' \right) \gamma_{\mu} \left(\frac{q}{2} + \ell' \right) \right. \\
 & \left. \gamma_{\sigma} \left(\frac{q}{2} + \ell \right) \right] e^{b(\ell-\ell')^2} (c^*)^{\alpha_p(0)-1} (2\pi)^4 \delta \left[\left(\frac{q}{2} + \ell' \right)^2 \right] \delta \left[\left(\frac{q}{2} + \ell \right)^2 \right] \\
 & \delta \left[\left(-\frac{q}{2} + \ell \right)^2 \right] \delta \left[\left(-\frac{q}{2} + \ell' \right)^2 \right] \sum_q \left\{ Q_q^2 F_q^h(x_1, p_1^2) F_q^{h'}(x_2, p_1'^2) + (q \leftrightarrow \bar{q}) \right\}
 \end{aligned} \quad (2.5)$$

where $\alpha_p(0) = 1$, and

$$\begin{aligned}
 p_1^2 &= \left(-\frac{q}{2} + \ell \right)_1^2 \\
 p_1'^2 &= \left(\frac{q}{2} + \ell \right)_1^2
 \end{aligned} \quad (2.6)$$

and the colour factor $1/3$ has been written down explicitly.

From (2.1) the double differential cross-section $d^2\sigma/dM dx_F$ ($\sqrt{Q^2} = M$) can be immediately obtained as

$$\frac{d^2\sigma}{dM dx_F} = \frac{1}{4M} \frac{s^2}{3\pi^3} \frac{1}{s} \frac{1}{(4\tau + x_F^2)^{1/2}} \int d^2q_{\perp} W_{\mu}^{\mu}, \quad (2.7)$$

and performing the integrations in (2.5) and (2.7) we obtain the following formula:

$$M^3 \frac{d^2\sigma}{dM dx_F} = \frac{32\pi\alpha^2\tau}{9(x_1 x_2)} \frac{\pi b}{\beta^2} \text{Tr} \left[Q_q^2 f_q^h(x_1) f_q^{h'}(x_2) + (q \leftrightarrow \bar{q}) \right] \quad (2.8)$$

A few comments are now in order. No explicit assumption about the p_T^2 dependence of $F(x, p_T^2)$ has been necessary to perform the integrations over ℓ_T^2 and q_T^2 . The double differential cross-section, according to (2.8), is given by a convolution of q and $\bar{q}$ distribution functions whose structure is similar to the Drell-Yan formula of the parton model⁹⁾. Also the normalization of (2.8) turns out to be completely fixed; in fact, setting $\beta^2/\pi b = \lambda$, the famous ratio R of e^+e^- annihilation is given by

$$R = \frac{\sigma(e^+e^- \rightarrow \text{hadrons})}{\sigma(e^+e^- \rightarrow \mu^+\mu^-)} = \lambda \cdot 3 \sum_{i=1}^{N_f} Q_i^2. \quad (2.9)$$

In the parton model $\lambda = 1$, whereas experimentally λ appears to be somewhat bigger than 1. As suggested by a quark geometrodynamical calculation¹⁴⁾, we shall take

$$\lambda = \frac{\beta^2}{\pi b} = \frac{4}{\pi}. \quad (2.10)$$

Equations (2.8) and (2.10) summarize the "scaling result" for the differential cross-section $M^3 d\sigma/dM dx_F$; which follows from the exchange of the "Pomeron".

The q and $\bar{q}$ distribution functions in the hadrons h and h' are precisely those appearing in the DIS cross-sections in the Bjorken limit, which have been found to scale^{10),11)}. We list here the results for such distribution functions, following from our previous analysis of DIS on a nucleon target¹¹⁾:

$$f_u^p(x) = f_d^n(x) = \frac{1}{2} \left\{ \frac{5}{3} \alpha_0 (1-x)^3 + \frac{6\theta_{20}}{\sqrt{x}} (1-x)^5 + \frac{3}{2} \tilde{\theta}_0 \frac{(1-x)^9}{x} \right\} \quad (2.11)$$

$$f_d^p(x) = f_u^n(x) = \frac{1}{2} \left\{ \frac{1}{3} \alpha_0 (1-x)^3 + \frac{3\theta_{20}}{\sqrt{x}} (1-x)^5 + \frac{3}{2} \tilde{\theta}_0 \frac{(1-x)^9}{x} \right\} \quad (2.12)$$

$$f_K^p(x) = f_S^p(x) = f_K^n(x) = f_S^n(x) = f_S^p(x) = f_S^n(x) = \frac{1}{2} \left\{ \frac{3}{2} \tilde{\theta}_0 \frac{(1-x)^9}{x} + \theta_0 \frac{(1-x)^7}{\sqrt{x}} \right\} \quad (2.13)$$

$$f_S^p(x) = f_S^n(x) = \frac{3}{4} \tilde{\theta}_0 (1-x)^9 \quad (2.14)$$

where the values of the constants are¹¹⁾:

$$\alpha_0 = 0.85, \quad \theta_{10} = 0.10, \quad \theta_{20} = 0.41, \quad \tilde{\theta}_0 = 0.13. \quad (2.15)$$

As for mesonic targets, in a recent paper¹³⁾ we have shown that in the MQM framework the q and $\bar{q}$ distribution functions can be expressed in terms of only one parameter. We found that such distribution functions are $SU(3)$ symmetric and are given by:

$$f_u^{\pi^+}(x) = A \frac{(1-x)^3}{3x} F(2,1,4;1-x) + \frac{B}{54x} (1-x)^{3/2} F(1,2,7/2;1-x) + C(1-x) \quad (2.16)$$

$$f_{\bar{u}}^{\pi^+} = A \frac{(1-x)^3}{3x} F(2,1,4;1-x) , \quad (2.17)$$

where $F(a,b,c;z)$ are hypergeometric functions and the constants A, B and C satisfy the following relations:

$$A = \frac{3}{4} \frac{\tilde{\theta}_0}{1.7} \simeq 0.06, \quad (2.18)$$

$$C + 1.06 B = 1 . \quad (2.19)$$

A preliminary analysis of the π structure functions in massive μ pairs production¹³⁾ has given the following value for C

$$C \simeq 0.36 \quad (2.20)$$

These formulae complete our analysis. But before we are ready to compare our results with the available experimental information, we must look in a more detailed way into the problem of the normalization of the cross-section.

3. THE FINAL STATE INTERACTION ENHANCEMENT

In the preceding section our calculation has been carried out by supposing that the $q\bar{q}$ pair with large mass, once it emerges from the hadronic crucible, arranges itself in a massive photon without further ado. This assumption is certainly true in the parton model and in theories where the quarks are free at asymptotic energies, but it cannot be true in a world, like the one conceived by QGD, where a $q\bar{q}$ pair at large mass has a non-negligible probability amplitude to make up those massive mesons, which are the original inhabitants of the

"Primitive World"*) . Thus the amplitude A , to produce a massive photon through the annihilation of a $q\bar{q}$ pair, receives the contributions from the diagrams of Fig. 5. Note that qualitatively these "final state" contributions are similar to the ones computed in perturbative QCD²²⁾, whose origin lies with gluon radiation in the final state. However, we should recall that our approach, contrary to perturbative QCD, does not suffer from the usual troubles of a perturbative calculation, connected with the unjustified hypothesis that the confining forces do not affect the quark behaviour in the asymptotic region.

To return to the amplitude A , we must compute the contribution A_b . According to the MQM approach A_b is, in the high Q^2 region, well represented by Pomeron exchange, and comparing with the $(e^+e^- \rightarrow \text{hadrons})$ amplitude we easily get:

$$\text{Im } A_b \rightarrow \lambda A_a \quad (3.1)$$

where λ is given in (2.10). By means of a dispersion relation we arrive at the result summarized in Fig. 6. Thus the net result of this mechanism of final state enhancement is to multiply Eq. (2.8) by the following factor

$$f(Q^2) = \lambda^2 + \left[1 - \frac{\lambda}{\pi} \log \frac{Q^2}{Q_0^2} \right]^2. \quad (3.2)$$

Needless to say $f(Q^2)$ violates "scaling"; however, as it can be seen from Fig. 7, in the regions of $\sqrt{s}$ and $\sqrt{\tau}$ where the experimental data are presently available the violation is fairly smooth and gentle. It will certainly be very interesting to check whether our peculiar predictions will accurately describe the energy regions that the Collider and ISABELLE will open to observation in the next few years.

4. RESULTS AND COMPARISON WITH THE EXPERIMENTS

All the presently available data have uncertainties in the normalization of the order of 25%, so it is appropriate to look first at the normalization independent cross-section ratios such as $d\sigma(\pi^+N \rightarrow \ell^+\ell^- + X)/d\sigma(\pi^-N \rightarrow \ell^+\ell^- + X)$, $d\sigma(K^+N \rightarrow \ell^+\ell^- + X)/d\sigma(K^-N \rightarrow \ell^+\ell^- + X)$ and $d\sigma(\bar{p}N \rightarrow \ell^+\ell^- + X)/d\sigma(pN \rightarrow \ell^+\ell^- + X)$.

*) For a detailed discussion of the "Primitive World", as an important step to describe the "Real World", see G. Preparata, Rapporteur talk at the EPS conference on High Energy Physics, Geneva 1979.

In Fig. 8 we present our predictions for the ratios $d\sigma(K^+p \rightarrow \mu^+\mu^-X)/d\sigma(K^-p \rightarrow \mu^+\mu^-X)$ and $d\sigma(pp \rightarrow \mu^+\mu^-X)/d\sigma(p\bar{p} \rightarrow \mu^+\mu^-X)$. Preliminary results for these quantities are already available from the NA3 experiment⁶⁾, but we could not report them due to their lack of acceptance corrections. However, the results of Fig. 2 of Ref. 6) do agree qualitatively with our predictions. Figure 9 reports a comparison between our prediction for $d\sigma(\pi^+N \rightarrow \mu^+\mu^-X)/d\sigma(\pi N \rightarrow \mu^+\mu^-X)$ and the experimental data of the CIP collaboration⁵⁾; both the experimental results and our prediction are integrated in the region $x_F \geq 0$. Figure 10 shows a comparison between our predictions for the charge asymmetry π^+/π^- for a platinum (upper curve) and a proton (lower curve) target and the experimental data of the NA3 experiment⁶⁾. In Fig. 11 the differential cross-section $M^3 d\sigma(\pi^-Pt \rightarrow \mu^+\mu^-X)/dM$ is compared with the experimental data from NA3 collaboration⁶⁾ at two different energies. The scaling violations coming from the renormalization factor $f(Q^2)$ (with $Q_0 = 4$ GeV) are very small as one can see from the comparison between the two theoretical curves (calculated at $p = 200$ GeV/c and $p = 280$ GeV/c). An over-all uncertainty in the experimental points of 15% has not been included. Figure 12 compares our prediction (with $Q_0 = 4$ GeV) for $M^3 d\sigma(pN \rightarrow \mu^+\mu^-X)/dM$ with the experimental data from the CFS collaboration⁷⁾. The agreement in shape is excellent, but the normalization is not; however, we note that the experimental points are affected by large systematic errors, which can easily account for the discrepancies between data and prediction. From Figs 11 and 12 we cannot draw any firm conclusion as to the value of the parameter Q_0 appearing in Eq. (23). $Q_0 = 4$ GeV seems to work well, however, the large systematic errors in the normalization of the experimental data imply a large uncertainty on Q_0 also.

We can thus conclude that the MQM affords us a good, quantitative understanding of the production of lepton pairs in hadronic collisions.

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FIGURE CAPTIONS

- Fig. 1 : The process $hh' \rightarrow \gamma_V X$ and its kinematics.
 $\begin{array}{c} \gamma_V X \\ \searrow \\ \ell \bar{\ell} \end{array}$
- Fig. 2 : The MQM diagram for the hadronic tensor $W_{\mu\nu}$.
- Fig. 3 : The Mueller-Regge expansion of the tensor $W_{\mu\nu}$, in the limit $s \rightarrow \infty$, $s^* = s[1 + \tau(1 - 1/x_1 - 1/x_2)] \rightarrow \infty$.
- Fig. 4 : The $q\bar{q}$ - $h\bar{h}$ -P vertex.
- Fig. 5 : The coupling of $q\bar{q}$ to the virtual photon in the timelike region.
 (a) the direct coupling; (b) the vector meson coupling.
- Fig. 6 : Connection between the full $q\bar{q}\gamma_V$ vertex in the timelike region and the point-like coupling.
- Fig. 7 : The factor $f(Q^2)$ plotted as a function of $\sqrt{\tau}$ for $\sqrt{s} = 19.4$ (-), 27.4 (---) and 62 GeV (.....), for two values of the parameter Q_0 (2 and 4 GeV).
- Fig. 8 : Theoretical predictions for the ratios $d\sigma(K^+p \rightarrow \mu^+\mu^-X)/d\sigma(K^-p \rightarrow \mu^+\mu^-X)$ and $d\sigma(pp \rightarrow \mu^+\mu^-X)/d\sigma(p\bar{p} \rightarrow \mu^+\mu^-X)$ as a function of $\sqrt{\tau}$.
- Fig. 9 : The ratio $R = d\sigma(\pi^+N \rightarrow \mu^+\mu^-X)/d\sigma(\pi^-N \rightarrow \mu^+\mu^-X)$ for an isoscalar target. The solid curve is our prediction, compared with the experimental results of the CIP collaboration¹⁵⁾ on a carbon target at the lab momentum $p = 225$ GeV/c and $x_F \geq 0$.
- Fig. 10 : The ratio $d\sigma(\pi^+N \rightarrow \mu^+\mu^-X)/d\sigma(\pi^-N \rightarrow \mu^+\mu^-X)$ plotted vs $M_{\mu^+\mu^-}$. The upper curve is our prediction for a Pt target, while the lower curve refers to a proton target. The experimental points are from the NA3 experiment¹⁶⁾ at $\sqrt{s} = 19.4$ GeV.

Fig. 11 : The differential cross-section $M^3 d\sigma(\pi^- Pt \rightarrow \mu^+ \mu^- X)/dM$ plotted vs $\sqrt{s}$. Experimental points are from the NA3 collaboration⁶⁾ at two different energies: open points at $p = 200$ GeV/c, full points at $p = 280$ GeV/c. The dashed and the full curves refer to our predictions ($Q_0 = 4$ GeV) for $p = 200$ GeV/c and $p = 280$ GeV/c respectively.

Fig. 12 : The differential cross-section $M^3 d\sigma(pN \rightarrow \mu^+ \mu^- X)/dM$ plotted as a function of $\sqrt{s}$. Experimental points are from the CFS collaboration at $p = 400$ GeV/c and are adapted from Ref. 15). No systematic error has been included. Both experimental points and the theoretical predictions are integrated for $x_F \geq 0$.

FIG. 1

$$\begin{aligned}
 (p_1 + p_2)^2 &= s \\
 (k + k')^2 &= Q^2 \\
 x_1 &= \frac{q^2}{2 p_1 q} \\
 x_2 &= \frac{q^2}{2 p_2 q} \\
 x_F &= x_1 - x_2 \\
 \tau &= x_1 x_2 = \frac{q^2}{s}
 \end{aligned}$$

FIG. 2

FIG. 3

$$= (\gamma_\sigma)_\alpha^\beta F_q^h(x_1, p_\perp^2)$$

FIG. 4

$$= A_a + A_b$$

FIG. 5

FIG. 6

FIG. 7

FIG.8

FIG.9

FIG. 10

FIG. 11

FIG. 12