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MULTIHADRON PRODUCTION IN DEEP INELASTIC SCATTERING
AND ITS RELATION WITH THAT IN e^+e^- ANNIHILATION

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ABSTRACT

The similarity and difference of particle production in deep inelastic scattering and e^+e^- annihilation is discussed. Comparison with the EMC data is given and predictions for higher energies are made.

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The aim of this paper is to discuss the relation between particle production in e^+e^- annihilation and that in deep inelastic scattering (DIS). The problem is not new; nevertheless, to date there does not exist any definite answer. In what follows, we give, on the basis of QCD evolution theory (augmented with unavoidable "plausible" assumptions), a formula which relates in a quite simple manner the multiplicities in both processes. As a by-product we obtain an explanation for the Q^2 -dependence of the hadron multiplicity, which was recently discovered by EMC.

To begin with, let us consider a spacelike photon which moves with high enough momentum. Its partonic (quark and gluon) fluctuations are pictured in Fig. 1. The first way to "materialize" these fluctuations in the form of hadronic final states (assuming some hadronization mechanism) is to continue the photon four-momentum, q , to the timelike region. After that the photon decays into off-mass-shell partons and then into hadron final states. This process, which is certainly familiar e^+e^- annihilation into hadrons, is pictured in Fig. 2. However, there exists another way. One can add to the spacelike photon some external particle with finite momentum (a proton at rest, say). Partons from the photon fluctuations will interact with this (target) particle, while the remaining partons will easily materialize, being timelike (see Fig. 3). It is customary in the parton model that partons which interact with the target have presumably low energies, so it is quite natural to expect that the final states in this process (which is nothing but DIS) will be almost the same as in e^+e^- , except for the energy carried away by the parton which triggers the process.

The above picture for hadron production looks most natural in the laboratory frame, where the virtual photon moves with very high ($\sim Q^2 = -q^2$) energy, and the partonic description of its quantum fluctuations is justified.

Let us now write down formulae. From the light-cone operator expansion, it follows that the structure function for DIS can be cast in the following factorized form [we assume that the kinematical variables Q^2 and $W^2 = (q+p)^2$ lie in the region where higher twists are negligible]

$$F(x, Q^2) = \sum_i \int_x^1 \frac{dz}{z} F_i(Q^2, Q_0^2, \frac{x}{z}) f_i(Q_0^2, z) \quad (1)$$

where $x = Q^2/(Q^2+W^2)$, i labels the triggering parton. In the operator product expansion, Q_0^2 is the renormalization scale for the corresponding composite operator. In our parton picture, it may be interpreted as the size scale ($\sim Q_0^{-1}$) of the

triggering parton, while z is related to its energy-momentum, k_μ . In the laboratory frame $k_0 \approx Q_0^2/2Mz - zM/2$ and M is the target mass. F_i in this picture describes the fragmentation of the virtual photon, while f_i is related to the interaction of the parton i with the target. F_i can be calculated in perturbative QCD, but f_i cannot, as we choose Q_0^2 to lie in the intermediate region between perturbative and non-perturbative phases. However, we still demand that $\alpha_s(Q_0^2)/2\pi \lesssim 1$.

According to Eq. (1), the deep inelastic process continues as follows in the laboratory frame. The photon decays into a $q\bar{q}$ pair with high virtualities $\sim Q^2$ (one negative and one positive). After this, both partons evolve with decreasing virtualities and momenta down to the stage of hadronization. It is worth noting that, at not very low x ($\gtrsim 0.1$), the triggering parton is dominantly $\bar{u}$, and to a lesser degree $\bar{d}$, while u , d and $s\bar{s}$, $c\bar{c}$, $b\bar{b}$ pairs practically do not contribute to the process. From this we arrive at the following expression for the average multiplicity in DIS:

$$\begin{aligned} \langle n \rangle^{DIS} &= \sum_i \int dz w_i(Q^2, x, z) (n_i + n_{oi}) \\ &= \langle n \rangle + \langle n_o \rangle \end{aligned} \quad (2)$$

where $\langle \dots \rangle$ denotes the average with weight $w_i = F_i f_i / zF$, n_i is the multiplicity in the photon fragmentation, while n_{oi} is the multiplicity in the antiquark-proton interaction. In Eq. (2) we implicitly assume that the colour compensation process is so soft that it does not affect the kinematics of independent fragmentation.

The term $\langle n_o \rangle$ is defined by the available invariant mass in $\bar{q}p$ interaction, $M_o^2 = (k+p)^2 < Q_0^2(1-x)/x + M^2$. At $x \gtrsim 0.1$ and $Q_0 = 0.7 \div 1.0$ GeV this mass varies between 1 and 3 GeV. So, taking into account that a baryon must be produced and that the above energy is typical for resonance production, it seems justified to assume $\langle n_o \rangle$ to be constant. Its value will be given below. As for what concerns the term $\langle n \rangle$, it was shown in Ref. 1) that $n_i(Q^2, z/x)$ coincides with the contribution of the triggering quark to $n_{e^+e^-}$ at energy $\tilde{W}^2 = (z-x/1-x)W^2$. This seems almost evident when one considers gluon bremsstrahlung from the initial $q\bar{q}$ pair, and realizes that it is necessary to subtract the energy-momentum which corresponds to the triggering $\bar{q}$. From Eq. (2) one can obtain a more transparent relation:

$$\langle n \rangle^{DIS} = \langle n \rangle_{l.q.}^{e^+e^-} (\kappa W^2) + \langle n_o \rangle \quad (3)$$

where $\kappa = \kappa(Q^2)$ is some number between 0 and 1 and l.q. refers to the light quark contribution.

In order to estimate the size and Q^2 -dependence of κ , it is useful to pass to the frame where the target moves rapidly. As $z = k^+/p^+$, it is invariant under the longitudinal boosts, and so is the function $f_i(z, Q_0^2)$. Now $f_i(z, Q_0^2)$ is interpreted as the distribution function of the parton inside the target hadron. From what we have already mentioned, it follows that f_i is essentially $f_{u,d}$, i.e., the valence quark distribution. For the latter, we apply the quark counting rule

$$f_i(Q_0^2, z) \sim (1-z)^{2n-1} \quad (4)$$

where n is the number of spectator quarks. For the proton, $n = 2$. Phenomenologically, n appears to be slightly higher, which can be related both with a softer distribution for d-quarks²⁾ and some Q_0^2 -dependence. As to the function $F_i(Q^2, z/x, Q_0^2)$, its form is defined by the QCD formula for valence quarks:

$$F_i(Q^2, Q_0^2, z/x) \sim (1-x/z)_+^{2C_F\xi-1} / \Gamma(2C_F\xi) \quad (5)$$

where $C_F = (N_c^2 - 1)/2N_c = 4/3$,

$$\xi = \int_{Q_0^2}^{Q^2} \frac{dk^2}{k^2} \frac{\alpha_s(k^2)}{2\pi} = \frac{6}{33-2n_f} \ln \frac{\ln Q^2/\Lambda^2}{\ln Q_0^2/\Lambda^2} \quad (6)$$

and n_f is the number of active flavours.

In principle one must use Eq. (2) to integrate the corresponding contributions in $\langle n \rangle^{e^+e^-}$. Already at this level one can see that, due to Eq. (5), the main contribution in the integral in z comes from growing z when Q^2 increases. However, Eq. (4) shows in turn that bound-state effects oppose this growth of z . So the efficiency factor $(z-x)/(1-x)$, which stands in front of W^2 in the expression for $\tilde{W}^2$, is much less than 1 and grows with Q^2 very slowly. In order to give

a more transparent expression, we simply substitute instead of κ in Eq. (3) the average value for $(z-x)/(1-x)$, which appears to be

$$\kappa = C_F \xi / (C_F \xi + n) \quad (7)$$

Note that κ , as defined in Eq. (7), changes from ≈ 0.12 at $Q^2 = 100 \text{ GeV}^2$ to only 0.19 at $Q^2 = 10^7 \text{ GeV}^2$!

Now let us turn to data. The latter show³⁾ that up to $W \approx 20 \text{ GeV}$, $\langle n \rangle^{\text{DIS}}$ is fitted accurately with the logarithmic dependence $A+B \ln W^2$. Contrary to this, in e^+e^- annihilation, a simple logarithmic fit fails to describe the data beyond $W \approx 10 \text{ GeV}$, where $\langle n \rangle^{e^+e^-}$ begins to grow quite rapidly up to $W \approx 50 \text{ GeV}$. Part of this growth is due to heavy quark hadronic decay, but what is more important is the very form of the $\langle n \rangle^{e^+e^-}(W)$, which appears to correspond to the QCD expression $\langle n \rangle \sim \exp \sqrt{c \ln W}$, as was obtained in Refs. 4).

Now it is clear why the DIS data grow like the low-energy part of $\langle n \rangle^{e^+e^-}$: the κ factor is too low in Eq. (3). However, at higher Q^2 , κ becomes larger, and finally at some Q^2 it will be possible to observe the same QCD growth of the multiplicity as in e^+e^- .

The history of the Q^2 -dependence started in 1975, when the Cornell group published its results⁵⁾, which showed that at $1.4 \text{ GeV}^2 < Q^2 < 8 \text{ GeV}^2$ and $1.7 \text{ GeV} < W < 4 \text{ GeV}$, $\langle n \rangle^{\text{DIS}}$ decreases with Q^2 like $-\ln Q^2$. Furthermore, the neutrino experiments at FNAL⁶⁾ ($1 \text{ GeV}^2 < Q^2 < 30 \text{ GeV}^2$, $1 \text{ GeV} < W < 10 \text{ GeV}$) and CERN⁷⁾ ($1 \text{ GeV}^2 < Q^2 < 80 \text{ GeV}^2$, $3 \text{ GeV} < W < 13 \text{ GeV}$) did not show any significant Q^2 -dependence of the multiplicity. Finally, the EMC³⁾ ($4 \text{ GeV}^2 < Q^2 < 100 \text{ GeV}^2$, $4 \text{ GeV} < W < 20 \text{ GeV}$) discovered a weak but statistically significant growth of the multiplicity of charged hadrons with Q^2 . Also in Ref. 3) it was demonstrated that the Lund model leads to a decrease of the multiplicity with Q^2 .

In Fig. 4, we compare the EMC data to the fit based on Eqs. (3) and (6), with the following choice of parameters: $n = 2.3$, $Q_0 = 0.7 \text{ GeV}$, $\langle n_0 \rangle = 2.2$, $\Lambda = 0.25 \text{ GeV}$. Reference 3) gave a phenomenological fit which takes into account the growth with Q^2 :

$$\langle n \rangle^{\text{DIS}} = A + B \ln W^2 + C \ln Q^2, \quad B, C > 0 \quad (8)$$

Figure 5 shows our predictions in comparison with the EMC fit (8) at $x = 0.1$ and $W \geq 10$ GeV, and with the QCD fit for $\langle n \rangle^{e^+e^-}$. One can see that, according to our approach, $\langle n \rangle^{\text{DIS}}$ at high enough W will reveal the typical QCD growth, which is delayed in DIS due to the bound-state effects and slow QCD evolution of the structure functions.

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REFERENCES

- 1) G.Sh. Dzhaparidze, A.V. Kiselev and V.A. Petrov, Proc. IV Intern. Workshop on High Energy Physics and Field Theory (Protvino, July 1981), Vol. 1 (1982) 12; Sov. J. Nucl. Phys. 35 (1982) 927;
A. Bassetto, Nucl. Phys. B202 (1982) 493.
- 2) J. Feltesse, Proc. Intern. Conf. on High Energy Physics (Bari, July 1985), p. 979.
- 3) EMC, M. Arneodo et al., Phys. Lett. 165B (1985) 222.
- 4) A. Bassetto, M. Ciafaloni and G. Marchesini, Nucl. Phys. B163 (1980) 477.
- 5) B. Gibbard et al., Phys. Rev. D11 (1975) 2367.
- 6) M. Derrick et al., Phys. Rev. D25 (1982) 624.
- 7) N. Schmitz, Proc. XV Int. Symp. on Multiparticle Dynamics (Lund, June 1984) 764.

FIGURE CAPTIONS

- Fig. 1 : General picture of quark-gluon quantum fluctuations of the spacelike photon.
- Fig. 2 : Hadronic final states from the timelike photon.
- Fig. 3 : Hadronic final states from the spacelike photon.
- Fig. 4 : Q^2 -dependence of the charged hadron multiplicity in μp DIS. Solid lines correspond to the fit based on Eq. (3).
- Fig. 5 : Comparative display of the high-energy predictions for $\langle n_{ch} \rangle$ in e^+e^- and DIS.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5