To be submitted to
Nuclear Physics

ISTITUTO NAZIONALE DI FISICA NUCLEARE
Laboratori Nazionali di Frascati

LNF-83/51(P)
27 Luglio 1983

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HIGH E_T PRODUCTION IN QCD

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ABSTRACT

The production of large transverse energy is discussed as the effect of multigluon emission produced in the hard scattering of the hadron constituents. Analytical formulae are given, based on resumming techniques to all orders in α_s, and succesfully compared with recent experiments at SPS, ISR and $p\bar{p}$ collider energies.

1. - INTRODUCTION

Hadron production of high transverse energy-momentum (E_T-p_T) is normally believed to result from hard scattering among hadron constituents and to provide a direct test of the short-distance behaviour of strong interactions. A characteristic jet structure is then expected to dominate the final states.

Motivated by these ideas hadron collisions at very high energies have recently been investigated by various experiments(1-5) with large geometrical acceptance. The abundant production of events with large values of transverse energy, first reported by the NA5 group(1) at the CERN SPS, has been confirmed at Fermilab(2) and ISR(3) energies, and even more copiously observed(4-5) at the CERN $p\bar{p}$ collider. However the measured cross sections, which change strongly with energy, are much larger than excepted(6) from QCD four jet production. In addition, up to resonably large E_T, the event structure does not indicate a sensible contribution from high p_T jets originated in a hard scattering of the constituents. Only very recently calorimeter-triggered experiments at the CERN ISR(7) and $p\bar{p}$ collider(5,8) have finally observed an important contribution from jet production at very large transverse energy.

The question arises therefore as to whether there exist two physically distinct mechanisms operating in a hadron-hadron collision at large E_T, appearing with a different configuration of the final states as well as with quite different cross sections. While there is little doubt that the jetty events, which are seen at higher E_T, reflect a hard scattering of the constituents and can henceforth compared with the predictions of perturbative QCD, the origin of the softer component has not been definitively clarified yet and various interpretations have been given so far. We only mention here models(9,10) with sizeable effects of fragmentation into hadrons of the constituents participating into the hard scattering and of the spectator partons, soft models(11) without hard scattering which generate events isotropically and gluon bremsstrahlung from the constituents(12).

In particular in a recent letter(12) we have suggested that a mechanism similar to the one responsible for the transverse momentum properties(13) of Drell-Yan pairs and weak bosons produced in hadron collisions is also operating here. The idea is quite simple. In any hard scattering process among the constituents a fraction of the initial parton subenergy is released in the form of soft QCD radiation, whose spectrum can be calculated to all orders in α_s. Furthermore the corresponding spectrum factorizes and, to leading order, is independent of the particular hard scattering process. Of course the very tail of the spectrum is modified by the detailed dynamics of a particular hard scattering process which takes over at appropriate large E_T, similarly to what happens in lepton pair production when $p_T \sim 0(M)$. In spite of the naîveté of the model used to illustrate these ideas - to simplify the calculation only quark-quark interactions were considered, with energy scaling cross sections - the results reproduced the qualitative behaviour of the data.

In the present paper we investigate further this scheme, providing a quantitative study of the idea of a unique dynamics - the hard scattering of the hadron constituents accompanied by emission of QCD radiation - operating in high energy hadron collisions at large E_T(14,15). We adress in particular to various questions of both theoretical and experimental interest: the role played by gluons as a source of radiation and the corresponding observable implications in comparison with quark initiated processes, the dependence on the beam type (π, p, $\bar{p}$), the effect of scaling violations, the relative role played by the hard and the soft component, the dependence on experimental kinematical acceptance (Δy–, $\Delta\phi$ – cuts), and so on.

A remarkable feature of our results lies on the fact that they are all analytic. Based on certain improved leading logarithmic approximations, which have been carefully studied and phenomenologically tested(16) in electroweak pairs production and e^+e^- annihilation, they can be taken with comfortable confidence not only for direct comparison with data, but also for testing various Montecarlo models, inspired by QCD but usually including ad hoc assumptions, which are of common use for data analysis.

The paper is organized as follows. In Section 2. we report the basic formulae for hard jet production, in Section 3. we discuss the soft gluon contributions and finally in Section 4. we discuss our results and give our final conclusions.

2. - THE HARD JET YIELD

We start with the well known formula(17) for the invariant jet cross section for hadron-hadron collisions at c.m. energy squared S.

$$p_o \frac{d\sigma}{d^3p} = \sum_{ijkl} \iint dx_1 dx_2 F_i(x_1,q^2) F_j(x_2,q^2) \frac{s}{\pi} \delta(s+t+u) \frac{d\sigma^{ijkl}}{dt}(s,t,u) \ , \tag{1}$$

in terms of the parton subprocesses $p_i+p_j \rightarrow p_k+p_l$, where $s=x_1x_2S$, t and u are the usual Mandelstam variables, and the indices i, j, k, l refer to q, $\bar{q}$ and g. The variable q^2 appearing in the structure functions $F_i(x,q^2)$ represents some characteristic scale of order p_T^2, when scaling violations are incorporated in the calculation. The expressions for $d\sigma/dt$ have been reported many times earlier(17) and are summarized below for the reader's convenience:

$$\frac{d\sigma}{dt} = \pi\alpha_s(q^2) \frac{|A(s,t,u)|^2}{s^2} \ , \tag{2}$$

with $|A|^2$ indicated below for the various QCD subprocesses

$$q_iq_j \rightarrow q_iq_j,\ q_i\bar{q}_j \rightarrow q_i\bar{q}_j (i\neq j) \qquad |A|^2 = \frac{4}{9}(s^2+u^2)/t^2 \tag{3a}$$

$$q_iq_i \rightarrow q_iq_i \qquad |A|^2 = \frac{4}{9}\left(\frac{s^2+u^2}{t^2} + \frac{s^2+t^2}{u^2}\right) - \frac{8}{27}\frac{s^2}{ut} \tag{3b}$$

$$q_i\bar{q}_i \rightarrow q_i\bar{q}_i \qquad |A|^2 = \frac{4}{9}\left(\frac{s^2+u^2}{t^2} + \frac{t^2+u^2}{s^2}\right) - \frac{8}{27}\frac{u^2}{st} \tag{3c}$$

$$q_i\bar{q}_i \rightarrow gg \qquad |A|^2 = \frac{32}{27}\left(\frac{u^2+t^2}{ut}\right) - \frac{8}{3}\left(\frac{u^2+t^2}{s^2}\right) \tag{3d}$$

$$gg \rightarrow q_i\bar{q}_i \qquad |A|^2 = \frac{1}{6}\left(\frac{u^2+t^2}{ut}\right) - \frac{3}{8}\left(\frac{u^2+t^2}{s^2}\right) \tag{3e}$$

$$q_ig \rightarrow q_ig \qquad |A|^2 = -\frac{4}{9}\left(\frac{u^2+s^2}{us}\right) + \left(\frac{u^2+s^2}{t^2}\right) \tag{3f}$$

$$gg \rightarrow gg \qquad |A|^2 = \frac{9}{2}\left(3 - \frac{ut}{s^2} - \frac{us}{t^2} - \frac{st}{u^2}\right) \tag{3g}$$

As well known, and from direct inspection of Eqs. (3), the dominant contributions to the various cross sections come from terms $\sim 1/t^2$, $1/u^2$, reflecting the exchange of vector gluons. Furthermore the colour factors present in the different $|A|^2$ clearly favor gluon initiated subprocesses, as already emphasyzed in the literature(17). These two observations will be of importance also for the soft gluon effects which are discussed in the next section.

From Eqs. (1) and (2) one obtains

$$\frac{d\sigma^{hard}}{dE_T} = \sum_{ijkl} \iint dx_1 dx_2\, F_i(x_1,q^2)\, F_j(x_2,q^2) \frac{2\,p_T}{\sqrt{1-4p_T^2/s}}\, \pi \frac{\alpha_s^2}{s^2} |A|^2_{ijkl} \ , \tag{4}$$

where $E_T=2\,p_T$ is the transverse energy corresponding to a 2-jet final state with transverse momentum p_T.

It is useful to introduce a scaled c.m. parton subenergy squared $p \equiv s/S = x_1 x_2$, so that the hard contribution to the transverse energy distribution finally becomes

$$\frac{d\sigma^{hard}}{dE_T} = \frac{\pi E_T}{S^2} \sum_{ijkl} \int \frac{dp}{p^2} \int_p^1 \frac{dx_1}{x_1} F_i(x_1,q^2)\, F_j\left(\frac{p}{x_1},q^2\right) \frac{\alpha_s^2}{\sqrt{1-E_T^2/Sp}} \left|A(p,E_T)\right|^2_{ijkl} \tag{5}$$

This expression can be used with some confidence only for sufficiently large of E_T to describe the tail of the distributions experimentally observed. In fact soft terms neglected in this lowest order calculation dominate and completely modify the E_T behaviour at medium E_T values. Even at large E_T, higher order corrections of $O(\alpha_s^3)$ might be of importance, as indicated by preliminary estimates[18]. For the present purpose Eq. (5) is sufficient to give us a resonable estimate of the role played by the hard component relative to the soft one, because of the lack of a precise knowledge of the gluon structure functions leading to a theoretical uncertainty of order of a factor of two.

3. - THE SOFT GLUON YIELD

The relevant role played by the QCD phenomena of multiple gluon emission in hard processes, and particularly for dilepton production in hadron collisions, has been recognized by now by many authors[16]. The high degree of sophistication reached by resumming the perturbation theory to all orders in α_s in certain improved double leading approximations has matched the impressive accuracy of the experimental data, leading to a quite detailed undestanding of the Drell-Yan processes and weak boson production[19]. In particular, from the successful analysis of the p_T effects one naturally expects some peculiar features in the p_T, E_T, $\langle n \rangle$-distributions of the hadron produced in association to the lepton pairs and weak bosons as well[20]. In absence of a correspondingly detailed observation of such hadronic properties[21] it is natural to look for similar effects in other hadronic processes. Then a qualitative study of multigluon emission effects in purely hadronic hard collisions was carried out in Ref. (12). Here we give a more quantitative analysis of this phenomenon.

We start with the soft transverse momentum K_T distribution of the QCD radiation emitted by the initial legs, when there is no detection of the final jets resulting from the hard collision (minimum bias events).

$$\frac{dP}{d^2K_T} \equiv \frac{1}{\sigma_o^h} \frac{d\sigma^h(Q)}{d^2K_T} = \sum_{n=0}^{\infty} \frac{1}{n!} \prod_{j=1}^{n} \int d^2k_{Tj}\, f_+(k_{Tj})\, \delta^{(2)}\left(\sum_{i=1}^{n} \vec{k}_{Ti} - \vec{K}_T\right) = \frac{1}{2\pi} \int_o^{\infty} bdb\, J_o(b K_T) \exp\left[\Delta(b, Q)\right] , \tag{6}$$

with

$$\Delta(b, Q) = \pi \int dq_T^2\, f_+(q_T)\, J_o(bq_T) \equiv \frac{2(c_1 + c_2)}{\pi} \int_o^{Q/2} \frac{dq_T}{q_T} \ln\left(\frac{Q}{q_T}\right) \alpha(q_T) \left[J_o(bq_T) - 1\right] . \tag{7}$$

In Eqs. (6-7), σ_o^h is the Born cross section for the hard process $p_1 + p_2 \rightarrow h$, and c_1 and c_2 are the appropriate colour factors for quarks and gluons, i.e. $c_q \equiv c_F = 4/3$ and $c_g \equiv c_A = 3$. The upper bound of the transverse momentum phase space of the emitted radiation has been put to Q/2.

Similarly the soft E_T spectrum, where E_T is total transverse energy of the QCD radiation emitted by the initial legs, is given[20] by

$$\begin{aligned}\frac{dP}{dE_T} &\equiv \frac{1}{\sigma_o^h} \frac{d\sigma^h(Q)}{dE_T} = \sum_{n=0}^{\infty} \frac{1}{n!} \prod_{j=1}^{n} \int d^2k_{Tj}\, f_+(k_{Tj})\, \delta\left(\sum_{i=1}^{n} k_{Ti} - E_T\right) \\ &= \frac{1}{2\pi}\int_{-\infty}^{\infty} dx\, e^{-ix E_T} \exp\left\{\int d^2q_T\, f(q_T)\left[e^{i x q_T} - 1\right]\right\} \\ &= \frac{1}{\pi}\int_{0}^{\infty} dx \cos\left[x E_T - \frac{2(c_1+c_2)}{\pi}\int_0^{Q/2} \frac{dq_T}{q_T} \ln\left(\frac{Q}{q_T}\right) \alpha(q_T) \sin(q_T x)\right] \cdot \\ &\quad \cdot \exp\left\{\frac{2(c_1+c_2)}{\pi}\int_0^{Q/2} \frac{dq_T}{q_T} \ln\left(\frac{Q}{q_T}\right) (q_T) \left[\cos(q_T x) - 1\right]\right\}\end{aligned} \tag{8}$$

It could be of interest, in some particular processes, to look at the distribution of transverve energy associated to the production of a transverse momentum K_T. This is the case, for example, in the study of the hadronic system accompanying Drell-Yan pairs on weak bosons. Then, the cross section weighted by gluonic transverse energy is[20]

$$\begin{aligned}\frac{1}{\sigma_o^h} \frac{d\Sigma^h(Q)}{d^2K_T} &= \sum_{n=1}^{\infty} \frac{1}{n!} \prod_{j=1}^{n} \int d^2k_{Tj}\, f_+(k_{Tj}) \left(\sum_{i=1}^{n} k_{Ti}\right) \delta^{(2)}\left(\sum_{i=1}^{n} \vec{k}_{Ti} - \vec{K}_T\right) \\ &= \frac{1}{2\pi}\int_0^{\infty} bdb\, J_o(b K_T) \exp\left[\Delta(b, Q)\right] \cdot g(b, Q)\end{aligned} \tag{9}$$

where

$$g(b, Q) = \frac{2(c_1+c_2)}{\pi}\int_0^{Q/2} dq_T \ln\left(\frac{Q}{q_T}\right) \alpha(q_T)\, J_o(b q_T)\,.$$

From Eqs. (6) and (9) the average transverse energy $\bar{E}_T(K_T)$ is then calculated as a function of the transverse momentum K_T as

$$\bar{E}_T(K_T) = \left(\frac{1}{\sigma_o} \frac{d\Sigma}{d^2 K_T}\right) / \left(\frac{1}{\sigma_o} \frac{d\sigma}{d^2 K_T}\right) \tag{11}$$

It is worth of noting that the distributions (8) and (9), the latter expressed in terms of $\bar{E}_T$, become very similar if the transverse energies are not too small.

So far we have considered just the energy momentum properties of the partonic system. To obtain the final distributions in hadron collisions, one has to take into account the parton densities, as in Eqs. (4,5) namely

$$\frac{d\sigma^{soft}}{d E_T} = \sum_{ij,h} \iint dx_1\, dx_2\, F_i(x_1, q^2)\, F_j(x_2, q^2)\, \frac{d\sigma^{h}}{dE}(x_1 x_2 S) = \\ = \sum_{ij,h} \int_{p_{min}}^{1} dp \int_{p}^{1} \frac{dx_1}{x_1}\, F_i(x_1, q^2)\, F_j\left(\frac{p}{x_1}, q^2\right) \frac{d\sigma^{ij-h}}{dE_T}(pS) \quad , \tag{12}$$

and similarly for $d\sigma/d^2K_T$ and $d\Sigma/d^2K_T$.

A few comments are in order here. First, one could be interested is studying the properties of the hadronic system produced in association to a particular hard final state, namely for fixed h. This is the case, for example, in studying the reaction $H_1 + H_2 \to J_1 + J_2 + X$, where J_1 and J_2 are two hard jects. Then the energy-momentum properties of the accompanying hadrons will give us informations on the QCD dynamics in the initial parton system. The different colour structure of quarks and gluons will affect the properties of the emitted radiation, as it explicitly appears in Eqs. (7-8-10) through the C_F, C_A factors. A measurement[22] therefore of the K_T (or E_T) distribution of such hadronic system produced in p-$\bar{p}$ collisions at the collider energies, where the gluon contribution is overwhelming, as discussed in the next section, whould be of particular interest, as compared for example to the analogous distributions for Drell-Yan pairs, which is dominated by quark-antiquark annihilation.

On the other hand, let us consider the case where one does not single out a particular hard final state, as for the so called minimum bias events. Then one has to sum over h and integrate over the final parton momenta. The radiation from the final legs is authomatically taken into account, because the process of dressing the final partons corresponds to a unit probability. This can be explicitly seen from the distributions (6-9) which are indeed normalized to one. Of course a measured value of E_T differs from the radiative-E_T of Eq. (8), by an amount $\langle E_{T\,fin} \rangle$, which can be directly estimated from the final partonic state, as discussed explicitly below. Then, for fixed i and j, the initial parton states, integration over the final states gives, using Eqs. (2-3-8),

$$\sum_h \frac{d\sigma^{ij\to h}}{dE_T} = \frac{\pi \alpha_s^2}{|t_{min}|}\, c_{ij}\, \frac{dP}{dE_T} \quad , \tag{13}$$

where the factor c_{ij} are obtained from Eqs. (3) and $|t_{min}| \simeq p^2_{Tmin}$ is the minimum value of the transfer momentum in the subprocesses $p_i + p_j \to h$. Eq. (13) is a consequence of the factorizazion property of the soft spectrum, as indicated explicitly in Eq. (8). Therefore the quantity $(p_{Tmin})^{-2}$ fixes the scale of the soft component to d $/dE_T$. From our understanding of deep inelastic phenomena one would expect that a value $(P_{Tmin})^{-2} \sim 1$ GeV^{-2} gives the right order of magnitude, and indeed the analysis of the experimental data for $\sqrt{S} \sim 20 \div 540$ GeV, as shown in the next section, confirms this naive expectation.

One finally obtains from Eqs. (12-13) the result

$$\frac{d\sigma^{soft}}{dE_T} = \sum_{ij} c_{ij} \int_{p_{min}}^{1} dp \int_{p}^{1} \frac{dx_1}{x_1}\, F_i(x_1\,, q^2)\, F_j\left(\frac{p}{x_1}, q^2\right) . \frac{\pi\alpha_s^2}{|t_{min}|}\, \frac{dP}{dE_T}(Sp)\,, \tag{14}$$

and similarly for $d\sigma/dK_T$ and $d\Sigma/dK_T$.

So far the lower bound p_{min} of the parton-parton subenergy has not been indicated explicitly. Being the distribution dP/dE_T appropriately normalized to unity, the quantity

$$\sigma^{inel} \equiv \sum_{ij} c_{ij} \int_{\bar{p}_{min}}^{1} dp \int_{p}^{1} \frac{dx_1}{x_1} F_i(x_1) F_j\left(\frac{p}{x_1}\right) \frac{\pi\alpha_s^2}{|t_{min}|} \tag{15}$$

represents approximately the inelastic cross section, the hard contribution being negligible. Eq. (15) then fixed the absolute lower bound $\bar{p}_{min}$, as a function of the parton densities which contribute to the actual reaction. It is however clear that for a fixed and sufficiently large E_T in Eq. (14), simple kinematical considerations demand $pS \gtrsim 2E_T^2$. Then $p_{min}=\max(\bar{p}_{min}, 2E_T^2/S)$.

Let us finally comment on the validity of our expression (14). It is clear that this result cannot be trusted at very large E_T ($2E_T \gtrsim \sqrt{S}$), when both the DLLA breaks down and genuine hard scattering effects, discussed in the previous section, are expected to take over. Indeed, evidence for jet structure at very large E_T has been clearly observed at ISR and $p\bar{p}$ collider[5,7,8]. Eq. (14) is therefore expected to hold in a range of medium-large values of E_T and the relative role the two components will indicate more precisely the transition region for each process. On the other hand Eq. (14) cannot be valid at too small values of E_T, because the contributions to E_T from initial and final states cannot be simply factorized as in Eq. (13). This is clear from the following observation. Let us write, for the observed E_T-distribution,

$$\frac{d\sigma}{dE_{Tobs}} = \iint dE_{Trad}\, dE_{Tfin} \left(\frac{d^2\sigma}{dE_{Trad}\, dE_{Tfin}}\right) \delta\left(E_{Tobs} - E_{Trad} - E_{Tfin}\right) , \tag{16}$$

where E_{Trad} and E_{Tfin} denote the contributions of from the initial and final states respectively. In compact notation, and within the usual soft approximation, we have

$$\frac{d^2\sigma}{dE_{Trad}\, dE_{Tfin}} \sim \sum_{ij,h} \int dp\, F_{ij}(p) \frac{dP(p)}{dE_{Trad}} \frac{d\sigma^h(p)}{dE_{Tfin}} \tag{17}$$

Then for $E_{Trad} \gg \langle E_{Tfin} \rangle$, one obtains

$$\frac{d\sigma}{dE_{Tobs}} \simeq \sum_{ij} \int dp\, F_{ij}(p) \left[\frac{dP}{dE_{Trad}}\right]_{E_{Tobs}=E_{Trad}+\langle E_{Tfin}\rangle} \sum_h \int dE_{Tfin} \left(\frac{d\sigma^h}{dE_{Tfin}}\right) \tag{18}$$

which reduces to Eq. (14) with a simple shift of the transverse energy scale. In the actual calculations $\langle E_{Tfin} \rangle$ is of order of a few GeV. Notice that for integral over E_{Tfin} in the r.h.s. of (18) automatically includes the bremsstrahlung from the final legs. Furthermore, the radiation emitted from the initial legs, up to a value $\vec{K}_T$ (or E_{Trad}), does not affect appreciably the average $\langle E_{Tfin} \rangle$, after $\vec{K}_T$ is summed to the back-to-back final parton transverse momenta. The situation here is different from the analogous case in deep inealstic scattering, where the knocked parton has $\langle E_{Tfin} \rangle \sim 0$ in absence of bremsstrahlung and acquires then an average $\langle E_{Tfin} \rangle \sim \langle E_{Trad} \rangle$, just because of balancing the $\vec{K}_T$ of the radiation.

On the other hand, for small E_T, one can write Eqs. (16-17) as

$$\frac{d\sigma}{dE_{Tobs}} \sim \sum_{ij,h} \int dp\, F_{ij}(p) \int dE_{Tfin} \left(\frac{d\sigma}{dE_{Tfin}}\right)^h \left[\frac{dP}{dE_{Trad}}\right]_{E_{Trad}=E_{Tobs}-E_{Tfin}} , \qquad (19)$$

which gives a better answer for this region. Of course a more accurate calculation could be done numerically, but it will further complicate the numerical analysis, and is out of the purpose of this paper.

4. - RESULTS AND DISCUSSION

In the previous sections we have given the main formulae, and for phenomenological applications we still have to specify the parton densities. We will first use scaling quark and gluon distributions, which are simpler for numerical use and enough to describe the main features of our results. The effect of scaling violations will be discussed later. Then the valence p and π structure functions are

$$x\, u^P(x) = A_u\, x^{0.5} (1-x)^3 \;; \qquad x\, d^P(x) = A_d\, x^{0.5} (1-x)^4 \;; \qquad x\, V^{\pi}(x) = A_{\pi}\, x^{0.5} (1-x) \;, \qquad (20)$$

where A_u, A_d and A_{π} are determined from the normalized conditions $\int u^P(x)\, dx=2$, $\int d^P(x)\, dx=1$ and $\int V^{\pi}(x)\, dx=1$.

The gluon distributions are given by

$$x\, G^P(x) = 3(1-x)^5 \;; \qquad x\, G^{\pi}(x) = 2(1-x)^3 \qquad (21)$$

The effect of scaling violations will be calculated using the parametrization of Ref. (23)

$$x\, v(x, Q^2) = \frac{NC}{B(D+1,\, A/C)}\, x^A (1-x^C)^D \qquad (22)$$

with N=2, $A=0.421-0.0412\,\bar{s}$, $C=2-0.6223\,\bar{s}^{\,0.8}$ and $D=3.37+0.4319\,\bar{s}$ for $v=u^P$; N=1, $A=0.364-0.0368\,\bar{s}$, $C=2-0.5414\,\bar{s}^{\,0.8}$ and $D=5.09+0.3463\,\bar{s}$ for $v=d^P$, where, as usual, $\bar{s}=\ln\left[\ln(Q^2/\Lambda^2)/\ln(Q_o^2/\Lambda^2)\right]$ with $Q_o^2=4\ \mathrm{GeV}^2$ and $\Lambda=0.4$ GeV. Furthermore the gluons are parametrized as(23)

$$x\, G(x, Q^2) = A(1+Bx+C\,x^2)(1-x)^D + E\, e^{-Fx} \qquad (23)$$

with

$A = 0.9243 + 2.51\,\bar{s}^{\,0.5}$ $\qquad$ $D = 6 + 1.454\,\bar{s}$

$B = 8.558 - 9.227\,\bar{s}^{\,0.3} - 0.655\,\bar{s}^{\,1.5}$ $\qquad$ $E = 11.29\,\bar{s}^{\,2}$

$C = 53.57 - 68.78\,\bar{s}^{\,0.3} + 19.3\,\bar{s}$ $\qquad$ $F = 41.24 + 5071\,\bar{s}$.

In the following our results will be based on Eqs. (20-21) unless stated explicitly. Finally we have parametrized $\alpha(k_\perp^2)$ as $\alpha(k_\perp^2) = 12\,\pi/25\, \ln\left[(k_\perp^2+\lambda^2)/\Lambda^2\right]$ with $\lambda=1$ GeV and $\Lambda\sim 0.3$ GeV, in agreement with previous phenomenology of Drell-Yan processes[16].

We show first in Fig. 1 our results for p collisions at 300 GeV compared with the NA5 results[1]. The full curve refers to the soft contribution of Eqs. (14-18). The normalization is absolute and, as stated above, depends on $|t_{min}|^{-1}$ which is taken to be 1 GeV^2. An inelastic cross section $\sigma^{inel} \approx 22$ mb fixed p_{min}, through Eq. (15), to the value $\bar{p}_{min} \simeq 10^{-2}$. Furthermore the scale of the observed E_T differs from E_{Trad} by the amount $\langle E_{Tfin} \rangle \simeq 2$ GeV, which is estimated from the hard cross section. The latter is also shown in Fig. 1 (dashed line). Finally the approximation for small E_T given by Eq. (19) is also

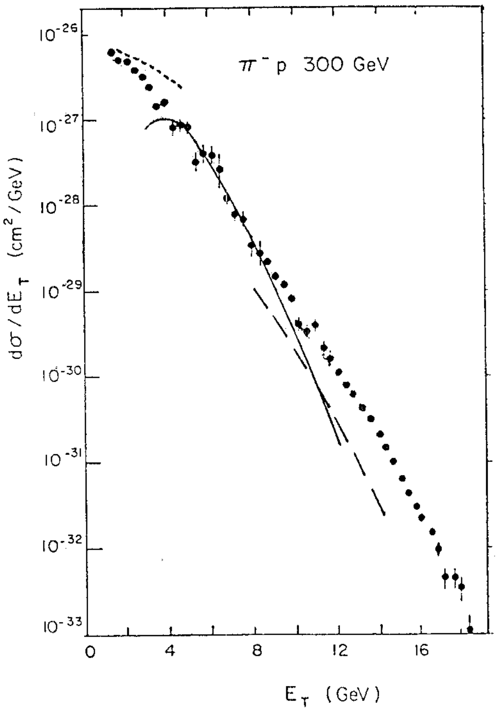

FIG. 1 - Transverse energy distribution in πp collisions at $\sqrt{s} \simeq 24$ GeV. The data are from Ref. (1). See text for the theoretical curves.

indicated (dotted line). From inspection of this figure it is clear that the soft mechanism describes quite well the observed spectrum up to reasonally large E_T, until the hard component takes over. Notice that the falloff of the soft term is steeper when scaling violations are taken into account. Furthermore a renormalization of the hard term by a sort of k-factor, of order of two, as indicated by preliminary estimates, is also suggested by the figure. We would'nt like to push too much the comparison with data at very small E_T, because of possible effects from experimental cuts and fragmentation, as well as for the theoretical uncertainty of our approach to this very soft region.

In Fig. 2a our results are compared with the UA1 data[4] at the CERN $p\bar{p}$ collider. The inelastic

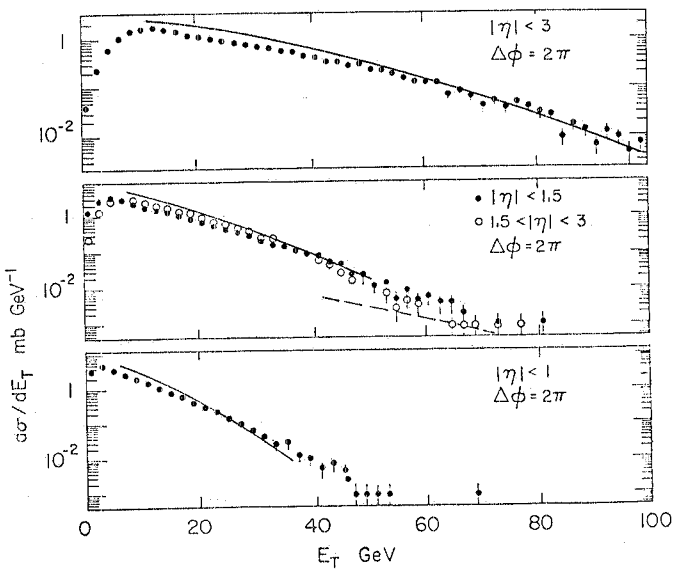

FIG. 2 - Transverse energy distributions in $p\bar{p}$ collisions at $\sqrt{s} = 540$ GeV. The data are from Ref. (4). See text for the theoretical curves.

cross section has been fixed to the value $\sigma^{inel} \simeq 40$ mb. The normalization is the same given above, i.e. $|t_{min}|^{-1}=1$ GeV^{-2}. The theoretical curve is obtained with no restriction on the phase space of the emitted radiation. In order study the rapidity (y) and the azimuthal (ϕ)-dependence we have restricted the full Δy or $\Delta\phi$ gluon ranges by a factor of two and three, and compared with data in Figs. (2b, c). From the inspection of these figures follows that all dependences on the kinematical variables S, y and ϕ are well reproduced. Furthermore in Fig. (2b) the region of transition between the soft and the hard components, and the latter is shown in the figure, is around E_T=50 GeV, which is also consistent with the data.

So far the gluon contribution, which is dominant at the collider energies, has been calculated with a scaling density $g_p(x)$ given in Eq. (21). The effect of scaling violations, as well as the use of gluon density of Eq. (23), referred as g', is indicated in Fig. 3, and amounts to a correction factor of order 2-3 to the scaling result.

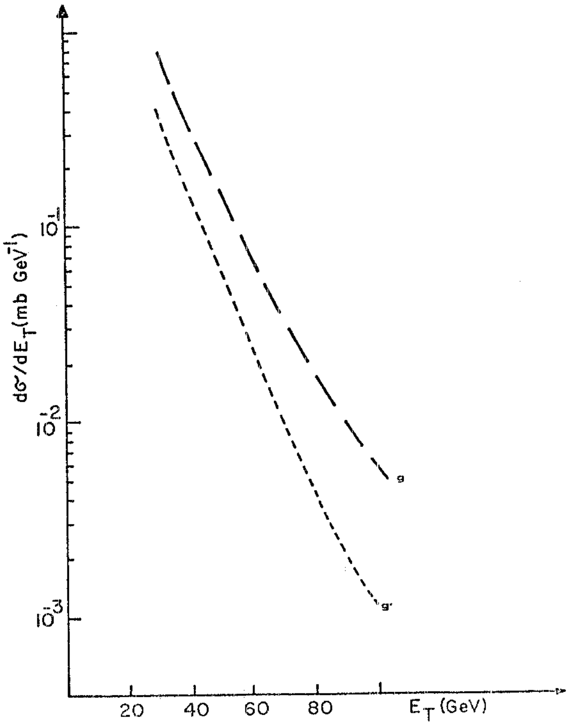

FIG. 3 - Transverse energy distribution in $p\bar{p}$ collisions at $\sqrt{s}$= 540 GeV, for two different gluon distributions as defined in the text.

Notice that our results do not depend on any assumption on the fragmentation into hadrons of quarks and gluons. This is not unconceivable since we look at totally inclusive distributions and therefore hadronization effects should be almost completely whashed out. Furthermore, because of a strong dominance of the gluons contribution, one should observe differences in the structure of the final states, as compared, for example, with what it is observed in purely quark initiated processes, like the production of Drell-Yan pairs or weak bosons. In particular the E_T behaviour of the hadron multiplicities or the transverse energy-momentum properties should reflect the effect of the different colour factors C_F and C_A associated to the QCD radiation emitted from the initial partons. A typical reaction to study would be, for example, $p\bar{p} \to J_1+J_2+X$, where J_1 and J_2 are two hard jets observed in the final state. Then the properties of the associated hadrons in X should show appropriate similarities with those observed for the minimum bias events and at the same time reveal differences with the case of the production of electro-weak pairs.

We also shown in Fig. 4 the comparison with the very recent ISR data of the AFS collaboration[24]. The theoretical curve, from Eq. (14), which includes scaling violations, is calculated with no restriction on the phase space of the emitted radiation, with an inelastic cross section $\sigma^{inel} \simeq 30$ mb. Notice here a large contribution from the quark-quark subprocess.

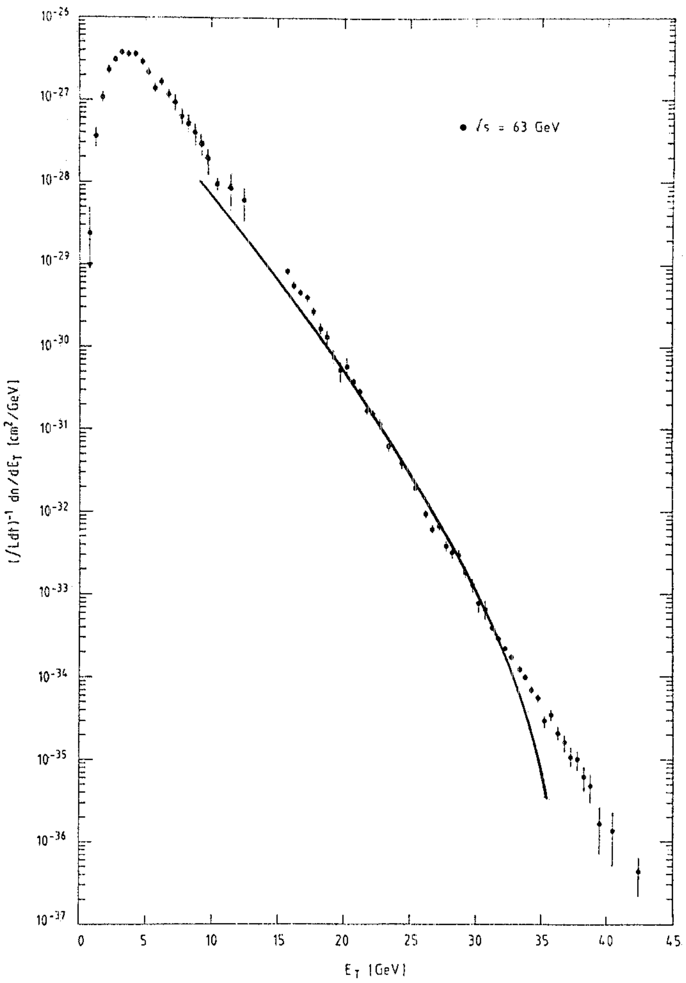

FIG. 4 - Transverse energy distribution in pp collisions at $\sqrt{s}$= 63 GeV. The data are from Ref. (24). See text for the theoretical curve.

From the above discussion it follows that the soft E_T distributions obey the following scaling law

$$E_T \frac{d\sigma^{soft}}{dE_T} \simeq f\left(x_T = \frac{2E_T}{\sqrt{S}}\right) \tag{24}$$

with the scale of σ^{soft} given by $1/|t_{min}|$. The approximate validity of (24) is based on neglecting logarithmic corrections from scaling violations and on the assumption $E_T = E_{Trad} + \langle E_{Tfin} \rangle \simeq E_{Trad}$, which is reasonable for E_T not very small. Notice that, on the contrary, the hard component obeys a different scaling law, namely

$$E_T^3 \frac{d\sigma^{hard}}{dE_T} \simeq f(x_T) \ . \tag{25}$$

In conclusion we have shown that most of the events which are produced at reasonably large E_T in high energy hadron-hadron collisions are the manifestation of gluon bremsstrahlung associated to the hard parton scattering. The application to this class of processes of the resumming techniques to all orders in α_s, which were found quite successful in describing the p_T properties of electro-weak Drell-Yan pairs, led us to a simple, analytical description of the observed phenomena, independent of any assumption on the hadronization of the partons.The resulting formulae describe, with no free parameters, the experimental findings for a quite large range of energies, and with various beams. More informations could be extracted from the data, in addition to the obvious study of back-to-back jets. They can give complementary informations on QCD dynamics, and compared with the observed features in Drell-Yan processes, could reveal important differences leading to dyscriminate between quarks and gluons.

The author gratefully aknowledges the kind hospitality of the CERN Theory Division, where part of this work was done.

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