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TRANSVERSE HADRONIC ENERGY IN $p\bar{p}$ AND pp HIGH ENERGY COLLISIONS

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The production of transverse hadronic energy in high energy $p\bar{p}$ and pp collisions is discussed as the effect of multigluon soft emission produced in the hard scattering of the proton (antiproton) constituents. The mechanism proposed is in agreement with the observed features at SPS and $p\bar{p}$ collider energies.

Hadron–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, reported by the NA5 group [1] at the CERN SPS and the E557 experiment [2] at Fermilab, has been even more copiously observed by the UA1 [3] and UA2 [4] collaborations at the CERN $p\bar{p}$ collider. The measured cross sections, which change strongly with energy, are much larger than expected from QCD four jet predictions [6]. In addition, the event structure does not indicate a sensible contribution from high p_T jets which originated in a hard scattering of the constituents.

On the other hand, our present understanding [7] of the transverse momentum properties of the large mass lepton pairs produced in hadron collisions, in the framework of perturbative QCD beyond lowest order [8], naturally leads us to expect the production of hadrons at large k_T, accompanying the lepton pair. The reason for that simply follows from the fact that the transverse momentum p_T distribution of a pair of mass M originates from multigluon bremsstrahlung from the initial quark legs, when $p_T \ll M$, which is the most likely case. This mechanism has been shown [9] to successfully describe the transverse momentum properties of Drell–Yan pairs in the various experiments performed so far. One expects [10] therefore to observe the corresponding excess of transverse hadronic energy associated with Drell–Yan pairs and also weak bosons.

In the present letter we suggest that the same mechanism is also responsible for the large production of transverse hadronic energy observed in pp and $p\bar{p}$ collisions at SPS and collider energies. The idea is quite simple. In any hard scattering process among the constituents, which for simplicity we will take to be the valence quarks (antiquarks), a fraction of the initial quark's sub-energy $\sqrt{\hat{s}}$ is released in the form of soft QCD radiation, whose spectrum can be calculated to all orders in α_s.

Furthermore, the emitted radiation being soft, the corresponding spectrum factorizes and, to leading order, is independent of the particular hard scattering process. Summing on the final states and integrating over the initial quark momenta, one obtains a transverse momentum distribution and a corresponding transverse energy flow which should be observable in all high-energy experiments with larger transverse energy triggers. Of course, the very tail of the spectrum might well be modified by the detailed dynamics of a particular hard scattering process [e.g. 11], similar to what happens in the lepton pair production when $p_T \sim O(M)$. Our results support this interpretation of the experimental data.

The techniques to calculate the soft transverse momentum spectrum in the double logarithmic approximation (DLLA) are well known [12]. Then, in a $qq(\bar{q})$ hard collision with centre-of-mass energy $Q \equiv \sqrt{\hat{s}}$, the emission from the initial legs of QCD radiation which sums to a total transverse momentum K_T is given by

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$$\frac{1}{\sigma_0^h}\frac{d\sigma^h(Q)}{d^2K_T} = \sum_{n=0}^{\infty}\frac{1}{n!}\prod_{j=1}^{n}\int d^2k_{Tj} f_+(k_{Tj}) \times \delta^{(2)}\left(\sum_{i=1}^{n} k_{Ti} - K_T\right) = \frac{1}{2\pi}\int_0^{\infty} b\, db\, J_0(bK_T)\exp[\Delta(b,Q)]\,, \quad (1)$$

with

$$\Delta(b,Q) = \pi\int dq_T^2\, f_+(q_T)\, J_0(bq_T) \equiv \frac{4c_F}{\pi}\int_0^{Q/2}\frac{dq_T}{q_T}\ln\left(\frac{Q}{q_T}\right)\alpha(q_T)[J_0(bq_T)-1]\,. \quad (2)$$

In eq. (1), σ_0^h is the Born cross section for the hard process $qq(\bar{q}) \to h$. The upper bound of the transverse momentum phase space of the emitted radiation has been put to $Q/2$.

Similarly, the transverse energy distribution is given by

$$\frac{1}{\sigma_0^h}\frac{d\sigma^h(Q)}{dE_T} = \int d^2K_T\left(\frac{1}{\sigma_0^h}\frac{d\sigma^h(Q)}{d^2K_T\, dE_T}\right)\,, \quad (3)$$

with

$$\frac{1}{\sigma_0^h}\frac{d\sigma^h(Q)}{d^2K_T\, dE_T} = \sum_{n=0}^{\infty}\frac{1}{n!}\prod_{j=1}^{n}\int d^2k_{Tj}\, f_+(k_{Tj}) \times \delta^{(2)}\left(\sum_{i=1}^{n} k_{Ti} - K_T\right)\delta\left(\sum_{i=1}^{n} k_{Ti} - E_T\right)\,. \quad (4)$$

Alternatively, one can define a transverse energy flow, corresponding to the total transverse momentum K_T, as [10]

$$\frac{1}{\sigma_0^h}\frac{d\Sigma^h(Q)}{d^2K_T} \equiv \int\left(\frac{1}{\sigma_0^h}\frac{d\sigma^h(Q)}{d^2K_T\, dE_T}\right)E_T\, dE_T = \sum_{n=1}^{\infty}\frac{1}{n!}\prod_{j=1}^{n}\int d^2k_{Tj}\, f_+(k_{Tj})\left(\sum_{i=1}^{n} k_{Ti}\right) \times \delta^{(2)}\left(\sum_{i=1}^{n} k_{Ti} - K_T\right) = \frac{1}{2\pi}\int_0^{\infty} b\, db\, J_0(bK_T)\exp[\Delta(b,Q)]\, g(b,Q)\,, \quad (5)$$

where

$$g(b,Q) = \frac{4c_F}{\pi}\int_0^{Q/2}\left(dq_T \ln\right)\frac{Q}{q_T}\;\alpha(q_T)\, J_0(bq_T)\,. \quad (6)$$

From eqs. (1) and (5) the average transverse energy $E_T(K_T)$ is then calculated as a function of the transverse momentum K_T as

$$E_T(K_T) = \left(\frac{1}{\sigma_0}\frac{d\Sigma}{d^2K_T}\right)\Bigg/\left(\frac{1}{\sigma_0}\frac{d\sigma}{d^2K_T}\right)\,. \quad (7)$$

The quantities (5) and (7) are particularly useful in the study of the transverse energy distributions accompanying electroweak pairs produced in hadronic collisions, along with the K_T distribution of the pair. In addition, they can also be used as a measure of the hadronic transverse energy released in purely hadronic processes. In fact, the two distributions (3) and (5) become very similar if the transverse energies are not too small. In the following we will concentrate on eqs. (5) and (7).

To obtain the final distributions one has only to take into account the initial parton densities, namely

$$\frac{d\sigma}{dK_T^2} \sim \sum_{ij,h}\int dx_1 dx_2\, q_i(x_1)\, q_j(x_2)\frac{d\sigma_{ij}^h}{dK_T^2}(Q^2 = x_1x_2S) \quad (8)$$

and similarly for $d\Sigma/dK_T^2$, where S is the energy squared of the hadronic $pp(p\bar{p})$ system.

A few comments are in order here. For eq. (8) to be valid, one has to have $Q > \sqrt{s_0}$, $\sqrt{s_0}$ being some energy threshold large enough to justify (i) the factorization of the K_T distributions in eqs. (1) and (5) and (ii) the assumption that a hard scattering process is taking place. Furthermore, if $\sqrt{s_0}$ is large, we can, as a first approximation, restrict the sum over the initial partons to the valence quarks and therefore neglect the sea and gluon contributions, which are expected to play a role for $x \to 0$.

With this caveat, eq. (8) becomes

$$\frac{1}{\sigma_0}\frac{d\sigma}{dK_T^2} \sim \int\frac{dx_1 dx_2}{x_1x_2}\, q_v(x_1)\, q_v(x_2) \times \int b\, db\, J_0(bK_T)\exp[\Delta(b,(Sx_1x_2)^{1/2})] \quad (9)$$

and similarly

$$\frac{1}{\sigma_0}\frac{d\Sigma}{dK_T^2} \sim \int \frac{dx_1 dx_2}{x_1 x_2} q_v(x_1)\, q_v(x_2)$$
$$\times \int b\, db\, J_0(bK_T) \exp[\Delta(b, (Sx_1x_2)^{1/2})]$$
$$\times\ g(b, (Sx_1x_2)^{1/2}), \qquad (10)$$

where we have taken into account a factor $1/(x_1x_2)$ to scale the hard cross sections. Taking for simplicity $q_v(x) \sim x^\alpha(1-x)^\beta$ with $\alpha = -1/2$ and $\beta = 3$, we finally obtain

$$\frac{1}{\sigma_0}\frac{d\sigma}{dK_T^2} \sim \int_{p_0}^{1} \frac{dp}{p} F(p) \int_0^\infty b\, db\, J_0(bK_T) \exp[\Delta(b, \sqrt{Sp})] \qquad (11)$$

and

$$\frac{1}{\sigma_0}\frac{d\Sigma}{dK_T^2} \sim \int_{p_0}^{1} \frac{dp}{p} F(p) \int_0^\infty b\, db\, J_0(bK_T)$$
$$\times \exp[\Delta(b, \sqrt{Sp})]\ g(b, \sqrt{Sp}), \qquad (12)$$

where

$$F(p) = p^{-1/2}\{(1+p)(1+p^2+8p)\ln p^{-1/2}$$
$$-\tfrac{1}{6}(1-p)[11(1+p)^2 + 16p]\}, \qquad (13)$$

$p = (x_1x_2)$ and $p_0 = \max(s_0/S, 4k_T^2/S)$.

A small change ($\lesssim 20\%$) in β does not significantly affect our results. We have parametrized $\alpha(k^2)$ as $\alpha(k^2) = 12\pi/25 \ln[(k^2+\lambda^2)/\Lambda^2]$, with $\lambda \approx 1$ GeV and $\Lambda \approx 0.3$, in agreement with previous phenomenology at low k^2 [9].

The average transverse energy E_T is plotted versus K_T in fig. 1 at the $p\bar{p}$ collider energy $\sqrt{S} = 540$ GeV and for $\sqrt{s_0} \approx 60$ GeV (see later). The dependence is simply $E_T \approx \sqrt{2}K_T + c$, as a good approximation, and the constant c also clearly indicates a release of E_T when all transverse momenta k_{Ti} balance to zero.

Then eq. (12) is shown in fig. 2, versus E_T, and compared with the experimental data of the UA1 collaboration [3]. The theoretical predictions, for various thresholds $\sqrt{s_0}$, are normalized to the data by a common factor. For the range of values shown in the figure, all curves converge at $E_T \gtrsim 50$ GeV, corresponding to $\bar{K}_T \approx 30$ GeV, and therefore become absolute predictions independent of s_0. The value $\sqrt{s_0} = 2\bar{K}_T \approx 60$ GeV has been used in fig. 1 for $E_T \lesssim 50$ GeV.

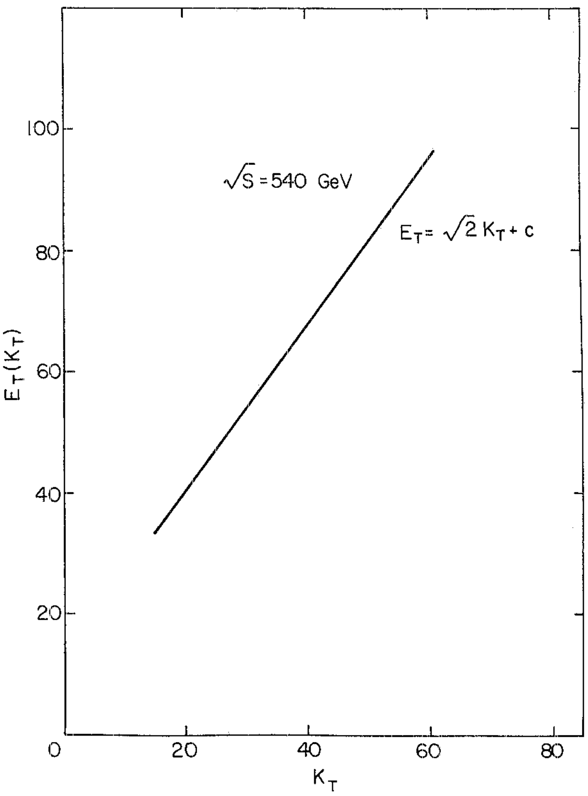

Fig. 1. Average transverse energy $E_T(K_T)$ as a function of K_T at $\sqrt{S} = 540$ GeV.

There is a simple argument which easily explains this behaviour. Neglecting the constraint of transverse momentum conservation in eq. (1), eq. (11) can then be approximately written as

$$\frac{d\sigma}{dK_T^2} \sim \int dx_1 dx_2\, q_v(x_1)\, q_v(x_2) \frac{d\sigma}{dK_T^2}(Sx_1x_2)$$
$$\sim \int_{p_0}^{1} \frac{dp}{Sp} F(p) \frac{1}{K_T^2} \ln\left(\frac{Sp}{K_T^2}\right)$$
$$\times \exp\left[-\frac{\alpha c_F}{2\pi}\ln^2\left(\frac{Sp}{K_T^2}\right)\right], \qquad (14)$$

with $F(p)$ given by eq. (13) and the exponential corresponds to an effective quark form factor of DLLA [12]. This form factor essentially cuts the integral at a value $p \sim p_0$, which depends on s_0 versus K_T^2. Then, for large enough K_T, eq. (12) becomes

$$\mathrm{d}\sigma/\mathrm{d}K_T^2 \sim K_T^{-4}\, F(4K_T^2/S) \qquad (15)$$

and similarly

$$\mathrm{d}\Sigma/\mathrm{d}K_T^2 \sim K_T^{-3}\, F(4K_T^2/S) \qquad (16)$$

which essentially describes the decrease observed at large E_T in fig. 2.

We would like to make a few comments. First, our results cannot be trusted at very large E_T ($2E_T \lesssim \sqrt{S}$), where both the DLLA break down and genuine hard scattering effects are expected to take over. Indeed, evidence for jet structures at very large E_T have very recently been suggested [4]. Furthermore, the agreement with experiments could seem somewhat accidental because of our simplified assumption of taking only valence quarks into account. Indeed, for $\sqrt{s_0} \approx 60$ GeV and $\sqrt{S} = 540$ GeV one still finds $x \gtrsim 0.1$ for the fractional quark momenta, which are rather small, and one would therefore expect this simple picture to be spoiled by the occurrence of other effects which come over for $x \to 0$ [13]. The most obvious one comes from hard qg and gg interactions which might modify our results, particularly at low E_T. We have checked however that the shape of the distribution obtained for large E_T from qg scattering, with $g(x) \sim (1-x)^5/x$, is steeper than that shown in fig. 2. It is even steeper for gg scattering. A more detailed analysis of the various contributions will be given elsewhere.

At very low E_T the absence in the theoretical curves of the experimentally observed dip reflects simply the failure of the proportionality between eqs. (3) and (5). For the reason explained above, and at the level of the present analysis, this region is not accessible anyway.

As a further test of the model, we have compared our predictions with the data of the E557 experiment [2] at Fermilab, in pp collisions at $\sqrt{S} = 27.4$ GeV. The results are shown in fig. 3. Again, the curves for various $\sqrt{s_0}$ converge at $E_T \approx 9$ GeV, corresponding to $\bar{K}_T' \approx 5$ GeV. The decrease at large E_T is faster than the data, as expected both for kinematical considerations and genuine hard scattering effects. Note, however, that the absolute normalization of the theoretical curves in fig. 3 is simply provided by that for p$\bar{\text{p}}$ in fig. 2, scaled by a factor $(s_0/s_0') = (\bar{K}_T/\bar{K}_T')^2$. This indicates that our interpretation of the data is essentially correct. It is also worth noticing that the strong S dependence shown by the data in the interval $\sqrt{S} \approx 27$–540 GeV is well reproduced by the theory.

As a final remark, the mechanism proposed here naturally leads us to expect large multiplicities associated with large values of E_T, as well as an increase in the average transverse momentum per particle, similar to what was observed in e^+e^- annihilation. These features are borne out by the data [1–5].

In conclusion, we have suggested that the production of hadronic events with large values of transverse energy observed in pp and p$\bar{\text{p}}$ collisions, with CM energies varying by a factor of twenty, can be under-

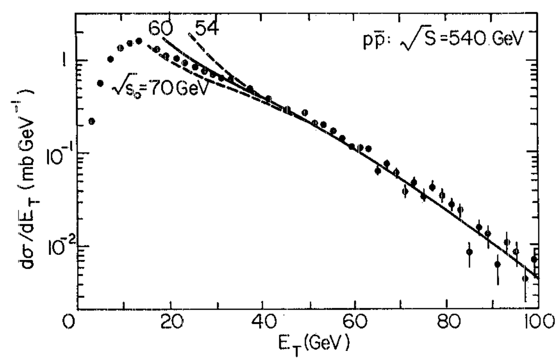

Fig. 2. Transverse hadronic energy distributions in p$\bar{\text{p}}$ collisions at $\sqrt{S} = 540$ GeV. The data are from ref. [3]. The curves show eq. (12) as a function of $E_T(K_T)$ [eq. (7)], for various values of $\sqrt{s_0}$.

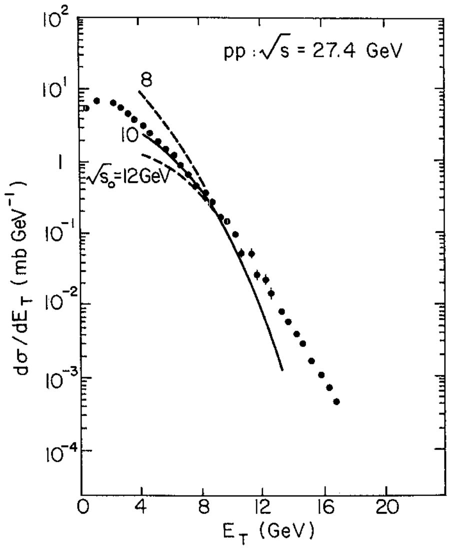

Fig. 3. The same as in fig. 2 for pp collisions at $\sqrt{S} = 27.4$ GeV. The data are from ref. [2].

stood as a standard effect of multigluon soft emission produced in the hard scattering processes of the constituents.

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