THE TRANSVERSE HADRONIC ENERGY IN W AND Z^0 PRODUCTION

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ABSTRACT

A thorough analysis is presented of the hadronic transverse energy produced in association with W^+ and Z^0 weak bosons at $p\bar{p}$ colliders. The complete rapidity and transverse energy distributions are derived, including the perturbative annihilation and Compton terms and the resummation of soft gluons. The first non-leading corrections to the corresponding Sudakov exponent are included in full and their effects are discussed. To make contact with the actual data a number of non-perturbative inputs is also necessary. We extract this information from minimum bias measurements. A detailed comparison of the resulting theoretical predictions with the available experimental data is finally presented.

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1. - Introduction

In the last few years the QCD description of lepton pair production in hadron-hadron collisions has been further developed in connection with W and Z^0 production experiments⁽¹⁾. A complete treatment of the lepton pair transverse momentum ($\vec{q}_T$) dependence has been given⁽²⁾ which includes the whole body of the available theoretical information⁽³⁻⁷⁾. The resulting $\vec{q}_T$ and rapidity (y) double distribution correctly reproduces the perturbative behaviour at large q_T ⁽³⁾ and incorporates the soft gluon exponentiation at leading⁽⁴⁾ and next-to-leading logarithmic accuracy.^(5,6) Upon integration over $\vec{q}_T$ the known perturbative results⁽⁷⁾ for the total cross section and the rapidity distribution $d\sigma/dy$ are exactly recovered (including the terms of order α_s which give rise to the "K-factor"). The phenomenological implications for W/ Z^0 production at the CERN $p\bar{p}$ Collider and for the physics at future multi-TeV Colliders have been discussed in detail in Ref.8. Also, by comparison with the ISR data, it was shown⁽⁹⁾ that the considerable spread of the q_T distribution in going from ISR to Collider energies is remarkably well reproduced without need of introducing any appreciable amount of intrinsic q_T or other similar non perturbative parameters.

In the present paper we extend the previous analysis by considering the transverse energy (E_T) distribution of the hadrons produced in association with the lepton pair. A direct comparison with the actual data is more difficult in this case. In fact, by perturbative QCD one can only predict the E_T distribution at the parton level (radiation from initial partons). But the contribution of the proton-anti-proton debris and of the final parton hadronization products is substantially more important for E_T than for q_T distributions and can only be taken into account by a purely phenomenological description. In spite of this not only the study

of E_T distributions is important per se but it also leads to interesting results in connection with the available data.

We improve over the existing analysis⁽¹⁰⁻¹¹⁾ of this problem on several aspects. We work out the complete y and E_T double distribution including finite terms from both $q\bar{q}$ annihilation and $q\bar{q} \rightarrow g$ Compton scattering. The complete non leading corrections⁽⁵⁾ to the soft gluon exponent are also included and their effect is discussed in detail. We study the sensitivity of the results to the value of α_s and to small changes in the Sudakov exponent. We display the energy dependence of the E_T distribution from Collider to Tevatron energies. Some approximate distributions (perturbative, leading double logarithmic) are compared to our results and their limitations are made evident. Some subtle points related to the positive definite sign of E_T which in intermediate stages leads to integrals which are complex numbers are brought up and clarified. Finally, we discuss the comparison of theory and experiment.

2. - Analytic Results

At order α_s the perturbative E_T and q_T distributions are identical. Schematically, for the $q\bar{q}$ term the perturbative E_T distribution at order α_s is of the form:

$$\begin{aligned} \frac{d\sigma}{dE_T} = & \sigma_0 (1+A) \delta(E_T) + 4B \left(\frac{\ln Q/E_T}{E_T} \right)_+ + \\ & + \frac{2C}{(E_T)_+} + 2E_T Y(E_T^2). \end{aligned} \quad (1)$$

where σ_0 is the lowest order cross section and A , B and C are of order α_s and independent of E_T in the limit of a fixed coupling α_s . $Y(E_T^2)$, also of order α_s , is a regular function of E_T at $E_T=0$. The definition of the "plus" distri-

bution is with respect to the kinematic upper limit of the E_T integration

$$\int_0^{A_T} g(x) f_+(x) dx = \int_0^{A_T} (g(x) - g(0)) f(x) dx \quad (2)$$

where

$$A_T = (E_T)_{\text{MAX}} = \frac{(S+Q^2)^2}{4S\chi^2 y} - Q^2 \quad (3)$$

The total cross section is given by

$$\sigma = \sigma_0(1+A) + \int_0^{A_T^2} Y(E_T^2) dE_T^2 \quad (4)$$

Note that the integral of the B and C terms over the whole E_T range vanishes because of eq.2. These terms become large when $E_T \ll Q$ and have to be resummed to all orders. The related exponentiation proceeds according to:

$$\begin{aligned} \frac{1}{\sigma_0} \frac{d\sigma}{dE_T} &= \sum_{n=0}^{\infty} \frac{1}{n!} \left[\pi \int_0^{A_T} f_+(E_{T_i}) dE_{T_i} \right] \delta(E_T - \sum_{i=1}^n E_{T_i}) = \\ &= \int_{-\infty}^{\infty} \frac{db}{2\pi} e^{ibE_T} \sum_{n=0}^{\infty} \frac{1}{n!} \left[\int_0^{A_T} f_+(E) e^{-ibE} dE \right]^n = \\ &= \int_{-\infty}^{\infty} \frac{db}{2\pi} e^{ibE_T} \exp \int_0^{A_T} f(E) (e^{-ibE} - 1) dE \end{aligned} \quad (5)$$

Note that two consequences of $E_{T_i} \geq 0$ are that the integrals over E_{T_i} in eq.(5) are complex numbers and that the corresponding upper limits of integration can be replaced with any number not smaller than E_T without affecting the final result for $1/\sigma_0 d\sigma/dE_T$ in eq.(5). These are important dif-

ferences with respect to the $q\bar{q}$ case.

The final result is of the form:

$$\frac{d\sigma}{dE_T} = 2E_T Y(E_T^2) + \int_{-\infty}^{\infty} \frac{db}{2\pi} e^{ibE_T} \sigma_0(1+A) \exp S(b) \quad (6)$$

where $S(b)$ can be read from eq.(5).

Following this general outline we now consider the different contributions to the E_T and y double distribution. The emphasis will be on those aspects which are specific of the E_T distribution and new.

Starting from eqs.(42, 43) of ref.2, we first consider the perturbative expression of the $q\bar{q}$ annihilation term at order α_s :

$$\frac{d\sigma}{dE_T dy} = 2E_T \left[X_q(E_T^2, Q^2, y) + Y_q(E_T^2, Q^2, y) \right] \quad (7)$$

where Q and y are the W/Z^0 (or lepton pair) mass and rapidity, respectively, and X_q , Y_q are the singular and non singular parts, respectively, of the E_T distribution. X_q and Y_q are specified in eqs.(42) and (63) of ref.2. In particular:

$$\begin{aligned} 2E_T X_q(E_T^2, Q^2, y) = & \sigma_0 \left\{ H(x_1^0, x_2^0) \left[\delta(E_T) + \frac{\alpha_s}{2\pi} C_F \left[8 \left(\frac{2Q/E_T}{E_T} \right)_+ - \right. \right. \right. \\ & \left. \left. \left. - \frac{6}{(E_T)_+} - \left(6 \ln \frac{A_T}{Q} + 4 \ln^2 \frac{A_T}{Q} \right) \delta(E_T) \right] \right\} + \right. \\ & + \frac{\alpha_s}{2\pi} C_F \left[\frac{2}{(E_T)_+} + \delta(E_T) \left(2 \ln \frac{A_T}{\mu} - \frac{1}{\epsilon} \right) \right] \left(H_1 \otimes P_{q\bar{q}} + H_2 \otimes P_{q\bar{q}} \right)^{(8)}_+ \\ & \left. + \frac{\alpha_s}{2\pi} C_F \left[H_1 \otimes C_q + H_2 \otimes C_q \right] \right\} \end{aligned}$$

where, as in ref.2, $x_{1,2}^0 = \sqrt{\tau} e^{\pm y}$, $\tau = Q^2/S$, $H(x_1^0, x_2^0) \sim [q(x_1^0, \mu^2) \cdot \bar{q}(x_2^0, \mu^2) + 1 \leftrightarrow 2]$, P_{qq} is the splitting function⁽¹²⁾, $C_F = 4/3$,

$$H_1 \otimes f \equiv \int_{x_1^0}^1 \frac{dz}{z} H\left(\frac{x_1^0}{z}, x_2^0\right) f(z) \quad (9)$$

and similarly for $H_2 \otimes f$. The coefficient of the $H_{1,2} \otimes P_{qq}$ terms is singular, because of the $1/\hat{\epsilon}$ left over from dimensional regularization. This singularity will eventually be reabsorbed in the definition of parton densities by shifting at the same time their scale from μ to a value of order Q . The $H_{1,2} \otimes c_q$ terms are instead finite and are fixed once a precise definition of the quark densities is given.

By going to b -space one obtains:

$$\begin{aligned} \int_0^{A_T} 2EX_q(E^2, Q^2, y) e^{-ibE} dE = \alpha_s \left\{ H(x_1^0, x_2^0) \left[1 + \right. \right. \\ \left. + \frac{\alpha_s}{2\pi} C_F \int_0^{A_T} \left(8 \frac{\ln Q/E}{E} - \frac{6}{E} \right) (e^{-ibE} - 1) dE - 6 \ln \frac{A_T}{Q} - \right. \\ \left. - 4 \ln^2 \frac{A_T}{Q} \right] \left. + \frac{\alpha_s}{2\pi} C_F \left[\ln \frac{A_T^2}{\mu^2} - \frac{1}{\hat{\epsilon}} + 2 \int_0^{A_T} \frac{dE}{E} (e^{-ibE} - 1) \right] \right\} \quad (10) \\ \cdot \left(H_1 \otimes P_{qq} + H_2 \otimes P_{qq} \right) + \frac{\alpha_s}{2\pi} C_F \left(H_1 \otimes c_q + H_2 \otimes c_q \right) \end{aligned}$$

Note that the $H_{1,2} \otimes P_{qq}$ terms cannot be directly reabsorbed into the quark densities by a change of scale because the final scale would be complex. Rather we proceed as follows. After exponentiation the $H \otimes P_{qq}$ terms take the form $e^{P(b)} \otimes H$ with

$$P(b) = \frac{\alpha_s}{2\pi} C_F \left[\ln \frac{A_T^2}{\mu^2} - \frac{1}{\hat{\epsilon}} + 2 \int_0^{A_T} \frac{dE}{E} (e^{-ibE} - 1) \right] P_{qq} \quad (11)$$

We define:

$$\frac{dK}{dE_T} = \int_{-\infty}^{\infty} \frac{db}{2\pi} e^{ibE_T} e^{P(b)} \otimes H \quad (12)$$

Thus

$$K(E_T) = \int_0^{E_T} dE' \frac{dK}{dE'} \quad (13)$$

By using a standard technique⁽¹³⁾ the quantity $K(E_T)$ can exactly be evaluated for $0 \leq E_T \leq A_T$. The result (at α_s fixed) is:

$$K(E_T)_{E_T \leq A_T} = \exp \left[\frac{\alpha_s}{2\pi} C_F P_{qq} \left(\ln \frac{E_T^2}{\mu^2} - \frac{1}{\epsilon} \right) \right] \otimes H \quad (14)$$

It is simple to recognize that $K(E_T) \equiv K(x_1^q, x_2^q, E_T)$ is the product of quark and antiquark densities at the scale E_T . The inclusion of gluon terms essentially does not change this result, but simply adds the P_{qg} contribution to the evolution from μ to E of quark densities. From eq.(12) we have:

$$\exp P(b) \otimes H = \int_{-\infty}^{\infty} dE e^{-ibE} \frac{dK}{dE} \quad (15)$$

From eq.(14) we only know dK/dE inside the range $0 \leq E \leq A_T$. Fortunately the integral in eq.(15) can in fact be restricted to this interval, because the difference does not contribute when inserted in the final expression of the form (5). We can thus make the replacement:

$$\exp P(b) \otimes H \rightarrow \int_0^{A_T} dE e^{-ibE} \frac{dk}{dE} = \quad (16)$$

$$= ib \int_0^{A_T} dE e^{-ibE} k(E) + k(A_T) e^{-ibA_T}$$

The surface term is crucial to reproduce the correct total cross section and average value of E_T (which both coincide with the corresponding perturbative results)(10).

After some final manipulations the complete result, also including the Compton contribution, becomes (see eq.(6)):

$$\frac{d\sigma}{dE_T dy} = 2E_T Y(E_T^2, Q^2, y) + \sigma_0 \left\{ \int_{-\infty}^{\infty} \frac{db}{2\pi} e^{ibE_T} \cdot R(b, y) \exp S(b, y) \right\} \quad (17)$$

Here $Y=Y_q+Y_g$ is given in eqs.(62),(63),(64) of ref.2,

$$R(b, y) = \int_0^{A_T} dE e^{-ibE} \frac{dk}{dE} \left[1 + \frac{\alpha_s}{2\pi} \frac{4}{3} \int_Q^{A_T} \left(8 \ln \frac{Q}{E} - \frac{6}{E} \right) e^{-ibE} dE \right] +$$

$$+ \frac{\alpha_s}{2\pi} \frac{4}{3} \left[\int_{x_1^0}^1 \frac{dz}{z} f_q(z) H\left(\frac{x_1^0}{z}, x_2^0\right) + \int_{x_2^0}^1 \frac{dz}{z} f_q(z) H\left(x_1^0, \frac{x_2^0}{z}\right) \right] +$$

$$+ \frac{\alpha_s}{2\pi} \frac{1}{2} \left[\int_{x_1^0}^1 \frac{dz}{z} f_g(z) K_2\left(\frac{x_1^0}{z}, x_2^0\right) + \int_{x_2^0}^1 \frac{dz}{z} f_g(z) K_1\left(x_1^0, \frac{x_2^0}{z}\right) \right] \quad (18)$$

where f_q and f_g are the appropriate kernels, given in eq.(59) of ref.2, for quark densities defined from the structure function F_2 of leptonproduction, $K_1 \equiv (q(x_1) + \bar{q}(x_1)) \cdot g(x_2)$, $K_2 \equiv (q(x_2) + \bar{q}(x_2)) \cdot g(x_1)$. In the terms of $O(\alpha_s)$ the scale P^2 of α_s and of parton densities is not specified. It

is natural to expect $O(E_T^2) \lesssim P^2 \lesssim O(Q^2)$. Note that the last four terms in R could be omitted provided that in the lowest order term the corresponding replacement:

$$H \rightarrow H + \frac{\alpha_s}{2\pi} \frac{4}{3} [f_q \otimes H_1 + f_q \otimes H_2] + \frac{\alpha_s}{2\pi} \frac{1}{2} [f_g \otimes k_1 + f_g \otimes k_2] \quad (19)$$

is made. Finally $S(b, y)$ in eq.(17) is given by

$$S = \int_0^Q \frac{dE}{E} (e^{-ibE} - 1) (4B \ln \frac{Q}{E} - 2C) \quad (20)$$

The contribution of the integration range above Q does not lead to large logarithms. It is convenient to only treat this term at first order and include it in R because it partially cancels the logs which appear in eq.(10). B and C in eq.(20) are known to a better accuracy than to order α_s :⁽⁵⁾

$$B = \frac{4}{3} \frac{\alpha_s(E^2)}{2\pi} [1 + D \alpha_s(E^2) + \dots] \quad (21)$$

$$C = -2 \frac{\alpha_s(E^2)}{\pi} [1 + F \alpha_s(E^2) + \dots] \quad (22)$$

where, in the $\overline{MS}$ scheme:

$$D = \frac{1}{2\pi} \left(\frac{67}{6} - \frac{\pi^2}{2} - \frac{5}{9} n_f \right) \quad (23)$$

The value of F is also known^(f1) and was included in our numerical calculations, but it is less important because the whole C term is suppressed by a large logarithm with respect to the B term. We argue, in agreement with ref.11, that D and F are the same as for the q_T distribution although the E_T and q_T distributions are different when two gluons are emitted. D and F are in fact process independent (although they depend on the renormalization scheme). For example, D is related to the coefficient of $((1+z^2)/(1-z))_+$ in the two loop corrections to $P_{qq}(z)$ (as defined and computed in ref.15), whose Mellin transform is a two loop anomalous dimension.

By using the fact that the real (imaginary) parts of R and S are even (odd) in b , eq.(17) can be cast into the form

$$\frac{d\sigma}{dE_T dy} = 2E_T Y(E_T^2, Q^2, y) + \sigma_0 \int_0^\infty \frac{db}{\pi} \left[\text{Re } R \cos(bE_T + \text{Im } S) - \text{Im } R \sin(bE_T + \text{Im } S) \right] \exp \text{Re } S \quad (24)$$

The physical implications of this final result will be discussed in the next section.

3. - Phenomenological Implications

In this section we derive numerical results, discuss their physical implications and compare them with the available experimental information. Some preliminary results were presented in ref.16.

f1) See the second article in ref.5, where the expressions of D and F were derived by using the results established in ref.14.

All computations in the following correspond to parton densities given by the set 1 of Duke-Owens.⁽¹⁷⁾ In general, technical details are handled here exactly as in ref.2 which can therefore be consulted by the interested reader. A difference which is worth discussing has to do with the freezing of $\alpha_s(E^2)$ at small E^2 . In the integral for the exponent $S(b,y)$, given in eq.20, the integration range in E extends down to $E=0$. However, the asymptotic form of $\alpha_s(E^2)$ becomes meaningless at $E \lesssim \Lambda$. A simple modification of $\alpha_s(E^2)$ at small E is given by (at one loop accuracy):

$$\alpha_s^{-1}(E^2) = b \ln \frac{E^2 + a\Lambda^2}{\Lambda^2}$$

and is used here as well as in ref.2.

For $a=0$ the integral in eq.20 would diverge at $E=\Lambda$. The most reasonable range of values for a is found by a compromise between two requisites. On one hand a must be sufficiently large that the steep dependence of the results on a (induced by the singularity at $a=0$) is quenched down. On the other hand a must be small enough so that the asymptotic form of $\alpha_s(E^2)$ is approached as fast as possible in E . The best value of a , found by trial and error, will be larger if the results are more sensitive to the infrared region in E . In Ref.2 for the q_T distribution, we set $a=2$. In the present case a larger value is necessary to stabilize the results: $a=5$. This increased infrared sensitivity is completely due to $\text{Im } S(b,y)$. The complex factor $e^{-ibE}-1$ leads to $\cos bE-1$ and $-\sin bE$ in the real and imaginary parts, respectively. At small values of bE , the sine is damped only linearly in E and this makes most of the difference.

In fig. 1 we start by a comparison of the transverse energy and transverse momentum distributions. Their common first order perturbative limit is also shown in the same

figure. The plots refer to $\sqrt{S} = 630$ GeV and $\Lambda_{\overline{MS}} = 200$ MeV. We recall that the complete E_T and q_T distributions include the subleading contributions in the exponent $S(b,y)$ given in eqs.(20) and (23) which are obtained in the $\overline{MS}$ prescription⁽¹⁸⁾. Correspondingly α_s must be evaluated at two loop accuracy in $S(b,y)$ and therefore the label $\overline{MS}$ on Λ is justified. It is immediately seen from fig. 1 that the peak or the average of the E_T distribution are located at much larger values than for the q_T distribution.

Note the slow rise of $d\sigma/dE_T$ at small E_T . Recall that $d\sigma/dq_T^2$ at $q_T = 0$ is large because of the contribution of several gluons with vanishing vector sum of transverse momentum⁽⁴⁾⁽⁶⁾. This mechanism clearly does not apply to the E_T distribution, and the importance of multiple gluon emission is reflected in a suppression of $d\sigma/dE_T$ at small E_T . It is very important to observe that the perturbative limit for the E_T distribution is only approached at very large values of E_T i.e. for $E_T \sim M_W$. Thus the practical use of lowest order perturbative calculations for E_T distributions is drastically reduced with respect to the q_T case. Recall that the integral under the E_T and q_T curves is the same (but for terms of order $\alpha_s^2(M_W)$) and corresponds to the total cross section up and including terms of order α_s (for the perturbative q_T distribution the same integral is also obtained only if the contribution of the $\delta(q_T^2)$ terms is taken into account).

The energy dependence of the E_T distribution in the range of interest for the CERN $p\bar{p}$ collider and the FNAL Tevatron is shown in fig.2. The difference between the two distributions at $\sqrt{S}=540$ and 630 GeV is much smaller than the existing theoretical ambiguities on either one.

The E_T distribution is rather sensitive to relatively small changes in the exponent $S(b,y)$. In fig. 3 we display the effect on the E_T distribution for $\sqrt{S}=630$ GeV of the non leading terms D and F defined in eqs.(20) to (23). As expec-

ted, most of the difference is induced by the Kodaira-Tren-tadue term, but the additional variation (of the same sign) due to F is also appreciable. The curves refer to the same value of $\Lambda_{\overline{MS}}=200$ MeV. Note, however, that α_s is computed at two loop accuracy only when D (with or without F) is included. We find that the quantitative importance of the two loop corrections to $S(b,y)$ can be reabsorbed for practical purposes by a tolerable change of Λ in the one loop expression for S . To illustrate this point we also included in fig. 3 a plot of the uncorrected E_T distribution with $\Lambda=300$ MeV. Our results may appear to be in contrast with the statement of ref. 11 that the uncorrected and fully corrected exponents lead to similar results. One important point is that they consider the whole term C in eq.(20) as a correction, while in our case the uncorrected term includes the lowest order expressions of both B and C . Furthermore their result refers to the real part of $S(b,y)$, while we find that the complete result is much sensitive to the imaginary part of $S(b,y)$. In conclusion we find that by omitting the non-leading corrections to $S(b,y)$ one would overestimate Λ by a factor ~ 1.5 .

Finally the dependence of the E_T distribution on the value of $\Lambda_{\overline{MS}}$ can be appreciated from fig.4.

After this purely theoretical description of the characteristic features of the partonic E_T distribution, we now consider the relation of our calculations with the available data. As repeatedly stressed, only the contribution to the E_T distribution from partonic sub-processes is taken into account in the approach discussed so far in this paper. In particular the transverse energy carried by the proton and antiproton fragments cannot be computed in QCD and was omitted. Fortunately a reliable estimate of this effect can be obtained from the observed E_T distribution of minimum bias events⁽¹⁹⁾. In fact the whole, or almost so, transverse energy measured in minimum bias events is precisely due to

the proton and antiproton fragmentation products near the forward direction. Thus the minimum bias transverse energy distribution ($1/\sigma_{MB} d\sigma_{MB}/dE_T$) can be extracted from the data and convoluted with the partonic $d\sigma/dE_T$ computed in the previous section:

$$\frac{d\sigma_{TOT}}{dE_T} = \int \frac{1}{\sigma_{MB}} \frac{d\sigma_{MB}}{dE}(E_1) \frac{d\sigma}{dE}(E_2) \delta(E_T - E_1 - E_2) dE_1 dE_2 \quad (25)$$

Once this convolution is made the only leftover effect is the difference between the E_T of the partons radiated in the subprocess and the E_T of the subset of hadrons originating from their hadronization. We tentatively neglect this effect and directly compare with the data $d\sigma_{TOT}/dE$ as obtained from eq.25.

In fig. 5 the data on the E_T distribution of minimum bias events at $\sqrt{s}=630$ GeV, obtained⁽¹⁹⁾ by the UA1 collaboration are shown together with a fit^(f2), which is then inserted in the convolution of eq.25. The tail at large E_T is deliberately disregarded in the fitting procedure. It is in fact established that this tail is due to the small fraction of events with jets. Since events with jets are produced by hard processes, it would be inappropriate to include them in a modeling of the E_T spread of the $p\bar{p}$ debris.

In fig. 6 the convolution of the minimum bias distribution with the QCD partonic E_T distribution for $\Lambda_{\overline{MS}} = 100-300$ MeV is compared with the preliminary data obtained by the UA1 collaboration in the 1984 run⁽²⁰⁾. The normalized

(f2) The explicit form of this fit is given by

$$\frac{1}{\sigma_{MB}} \frac{d\sigma_{MB}}{dE} = \frac{4E^2}{c^3\sqrt{\pi}} \exp(-E^2/c^2)$$

with $c=24$ GeV. The corresponding average value of E is

$$\langle E \rangle = (2/\sqrt{\pi}) c \approx 27 \text{ GeV.}$$

fit to the minimum bias data is also shown for comparison. The substantial displacement of the W production data with respect to the minimum bias distribution ($\langle E_T^W \rangle \approx 46$ GeV; $\langle E_T^{MB} \rangle \approx 32$ GeV) is reasonably well reproduced by the resulting curves. Within the limitations of our crude approach a value of $\Lambda_{\overline{MS}} \approx 200$ MeV is preferred by the data.

We cannot directly compare our results with the E_T distributions measured by the UA2 collaboration. In fact, the UA2 detector is not sensitive to hadrons emerging near the beam direction with $\theta \leq \theta_{\text{cut}}$ ($\theta_{\text{cut}} \approx 40^\circ$). One can use the model of ref.11 to approximately describe the dependence of the soft gluon exponential on θ_{cut} . We have checked that in the UA2 experimental conditions most of the soft gluon contribution is lost. As a result in that case a rapidly decreasing E_T distribution is obtained which is similar to the perturbative result.

- Conclusion

The successful prediction of the q_T distributions of W's and Z's is an important confirmation of the QCD improved parton model. In fact the relatively large observed value of the average q_T make the intrinsic and non perturbative components quite negligible. Thus the measured distribution is really a test of the perturbative dynamics. Both the soft gluon resummation and the perturbative component are strongly supported by the data. Encouraged by this success we have attempted here a more difficult task, i.e. the prediction of the E_T distribution for the hadrons associated with W/Z⁰ production. Contrary to the case of q_T the E_T distribution cannot be completely calculated in perturbation theory. In fact, only the E_T contribution from the partonic subprocesses can be reliably predicted. The most important left-over effect is the E_T contribution from the p- $\bar{p}$ debris. We have

made no attempt to calculate this term. Rather we extracted it from the data on minimum bias events. On the other hand the E_T distribution from the partonic subprocesses can be computed reliably. The level of theoretical uncertainty is somewhat larger in this case than for the q_T case. The additional uncertainty is mainly from the imaginary part of the Sudakov exponent $S(b,y)$. In spite of this ambiguity we find it quite remarkable that the large difference between the E_T distributions in minimum bias events and in W/Z^0 production is quite well reproduced by our calculation for a value of $\Lambda_{\overline{MS}} \simeq 200$ MeV. This value is quite consistent with the commonly accepted value of $\Lambda_{\overline{MS}}$.

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- Figure Captions

1. Comparison of the E_T and q_T distributions for $\Lambda_{\overline{MS}}=200$ MeV. Their common first order perturbative limit is also shown. This figure, as well as the other ones, refers to the cross sections at $y=0$ for W^+W^- (y is the W rapidity), without inclusion of the $W \rightarrow e\nu$ branching ratio. The first non leading corrections to the soft gluon exponent S are included ($\sqrt{S}=630$ GeV).
2. Energy dependence of the E_T distribution. The curves correspond to $\Lambda_{\overline{MS}} \simeq 200$ MeV, and $\sqrt{S}=0.54, 0.63, 1.6$ TeV.
3. The solid curves refer to the E_T distribution at $\sqrt{S}=630$ GeV for $\Lambda_{\overline{MS}}=200$ MeV with (curve B) and without (curve A) the next-to-leading corrections to the soft gluon exponent S (terms D and F in eqs. 21, 22). The curves A and B are computed by using in the exponent α_s at one and two loop accuracy respectively. The dashed curve is the uncorrected curve, analogous to A, but with $\Lambda_{\overline{MS}}=300$ MeV.
4. The E_T distributions (including the D and F terms in eqs.21,22) for $\Lambda_{\overline{MS}}=(100-300)$ MeV.
5. The UA1 preliminary data on the E_T distribution of minimum bias events at $\sqrt{S}=630$ GeV together with a simple fit of the points at small E_T . The large E_T tail was deliberately disregarded. In fact, it mainly corresponds to events with jets (a hard process), while we are interested in a parametrization of the soft component. The analytic form of the fit is given in footnote 2.
6. The preliminary UA1 data on the hadronic transverse energy produced in association with $W^\pm$ at $\sqrt{S}=630$ GeV are compared with our resulting theoretical predictions for $\Lambda_{\overline{MS}}=(100-300)$ MeV. The curves are obtained by convoluting according to eq.25 a fit of the E_T distribution of minimum bias events (also shown in the figure) with the computed partonic E_T distribution.

Fig. 1

Fig. 2

Fig. 3

Fig. 4

Fig. 5

Fig. 6