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Comments on Initial State Radiation and Multiple Interactions*

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Abstract:
A complex of problems permeates the field of event structure in hadron collisions. In this talk, two aspects are singled out for consideration. Firstly, initial state radiation, where our "backwards evolution" scheme is compared with experimental data and analytical calculations. Secondly, the possibility of multiple parton-parton interactions in a single hadron collision, which may play an important rôle for the proper understanding of "beam jets" and multiplicity distributions.

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Introduction

The event structure in hadronic collisions is very complex, and the task of understanding this structure is correspondingly difficult. In order to make any headway at all, one imposes a separation of the physics into different aspects, each of which may be considered in some approximate scheme. Thus one speaks of the hard central interaction, of final state radiation, of initial state radiation, of structure functions, of beam jets, and of fragmentation. Elsewhere in these proceedings, it is discussed how these pieces come together in the Lund Monte Carlo for hadronic processes (PYTHIA) [1]. This is one of the programs developed to simulate events realistic enough to bear comparison with existing data, and thus be useful to give some indication what might be expected at SSC energies. In this talk, special emphasis will be given to current studies in two of the areas mentioned above, initial state radiation and beam jet structure. This work is largely being carried out in collaboration with Mats Bengtsson and Maria van Zijl.

Initial State Radiation

Some time ago, we developed a model for initial state radiation based on "backwards evolution" [2,3]. Here the hard scattering is selected first, folding the hard scattering matrix element by Q^2-evolved structure functions. The cascades that produced the interacting partons are thereafter reconstructed "backwards" from the hard scattering towards the shower initiators, counter to the naive time direction. This means that the evolution is reconstructed in falling sequence of Q^2 values, from that of the hard scattering towards some low cutoff scale $Q_0^2 \approx 4\ \mathrm{GeV}^2$, with the parton momentum fraction x increasing at each step. Although less well suited for easy understanding of shower development, this description actually offers large advantages in the construction of exclusive events.

The scheme is based on the Altarelli-Parisi equations [4]

$$\frac{df_b}{dt}(x,t) = \frac{\alpha_s(t)}{2\pi} \sum_a \int \frac{dx'}{x'} f_a(x',t) P_{a\to bc}\left(\frac{x}{x'}\right)$$

which express the change in the number of partons b at momentum fraction x as the resolution scale $t = \ln(Q^2/\Lambda^2)$ is increased by splittings $a\to bc$. During a decrease dt, the probability that a parton b originally present becomes "unresolved" is then given by $dP_b = df_b/f_b$. The probability that nothing happens exponentiates to yield the form factor

$$S_b(x,t_{max},t) = \exp\left\{-\int_t^{t_{max}} dt' \frac{\alpha_s(t')}{2\pi} \sum_a \int dz\, P_{a\to bc}(z) \frac{x' f_a(x',t')}{x f_b(x,t')}\right\}$$

giving the probability that a parton b remains at x from t_{max} to $t<t_{max}$. Knowledge of S_b can be used to find the t value, z value and flavour a of a branching $a\to bc$, where $z = x/x' = x_b/x_a$ and $Q^2 =$

$\Lambda^2\exp(t)$ gives the spacelike virtuality of parton b. The branching that produced parton a may now be studied in the same fashion, and so on iteratively until Q_0^2 is reached.

The kinematics of the evolution is not uniquely given. We have chosen to require $\hat{s} = x_1 x_2 s$, not only at the hard scattering Q^2 scale but also for the resolved partons at any lower scale, so that any splitting corresponds to an incrase (moving backwards) of the invariant mass-squared by a factor 1/z. With this choice, it is possible to find the largest timelike virtuality of the parton c at each of the branchings and allow an "associated timelike shower" to develop from this virtuality. Once the actual virtuality has been found the kinematics may be completely reconstructed.

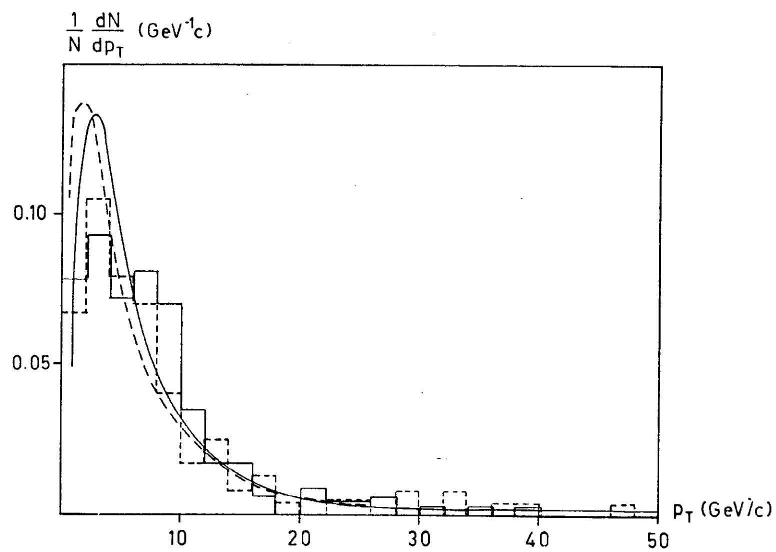

Fig. 1. W transverse momentum spectra. Histograms: UA1 [5] full and UA2 [6] dashed. Curves: model results, without smearing dashed and with full.

Recently, results on W and associated jet production have been presented by UA1 [5] and UA2 [6], based on an increased statistics. We have compared with those data and also attempted to include effects introduced by calorimetric fluctuations, electron isolation criteria etc. à la UA1. Results for the W transverse momentum spectrum are shown in Fig. 1, and for the fraction of events containing at least one jet in Fig. 2. More detailed comparisons with jet production properties have also been made [3]. At present level of accuracy, there is thus every indication that our model provides a sensible description of initial state radiation.

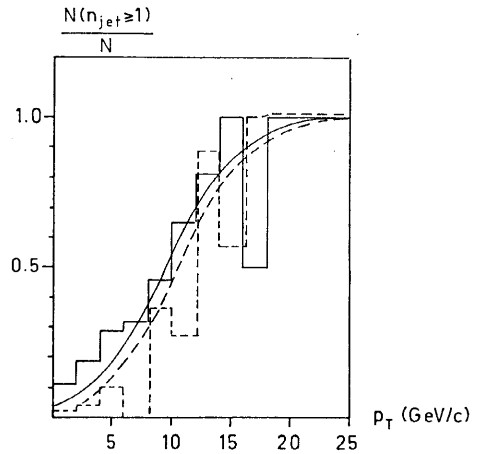

Fig. 2. Fraction of W events with at least one jet as a function of the W transverse momentum. Notation as in Fig. 1.

The sensitivity to the different assumptions made has been explored by running the basic program with various alternative choices at some ten different points. Taking the mean transverse momentum of the W $\langle p_{TW}\rangle$ as one measure, it turns out that the sensitivity to each point separately is typically 5-10%, giving $\langle p_{TW}\rangle$ shifted upwards or downwards about equally often. Not surprisingly, the most sensitive choice is to associate the maximum allowed spacelike virtuality with M_W^2, rather than some arbitrary multiple of this. There is also some dependence e.g. on the choice of structure functions and hence Λ, and on whether associated timelike showers are allowed to develop or not, whereas results are fairly insensitive to the lower cutoff of the spacelike evolution or to the inclusion of a primordial k_T contribution.

Of special interest is the choice of scale in α_S. As discussed elsewhere [7], there are indications that the naive leading log scale Q^2 becomes modified to $k_T^2 \approx (1-z)Q^2$ when loop corrections are included. Our standard choice has been to use $Q^2/4$ as scale, which

does give a somewhat higher $\langle p_{TW}\rangle$ than with Q^2 as scale. As it turns out, results with k_T^2 come very close to those with $Q^2/4$: on the one hand, the structure of the Altarelli-Parisi kernels favours z close to 1, on the other, only branchings with z well separated from 1 contribute significantly to the total p_{TW}. Still unresolved is the question whether the evolution should be carried out entirely in terms of k_T^2 rather than Q^2, i.e. whether one should use an evolution parameter $t = \ln(k_T^2/\Lambda^2)$. In practice, the results obtained are very similar, particularly when compared with the case above when k_T^2 was already used as argument in α_S. The stability discussed here means that our limited knowledge is no hindrance in constructing a fairly reliable model for initial state radiation, but also that it will be a tough task to settle any of the open theoretical issues by using experimental data as arbiter.

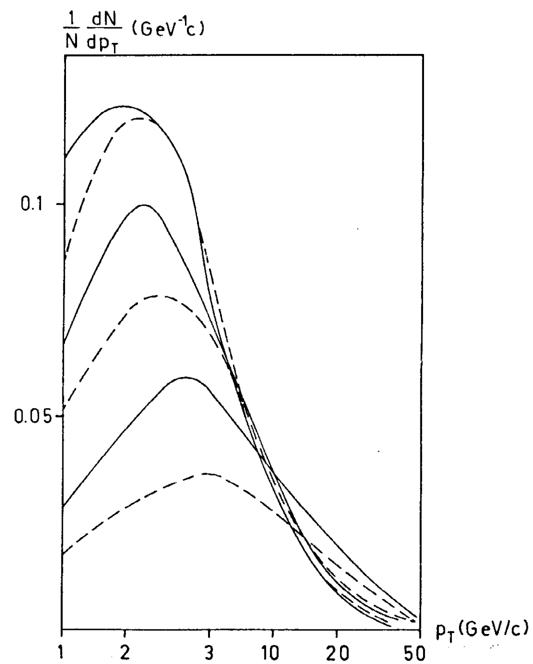

Fig. 3. Transverse momentum distribution of centrally produced W. Full model results, dashed Altarelli et al. [8], at CM energies 0.63, 1.6 and 10 TeV, from top to bottom.

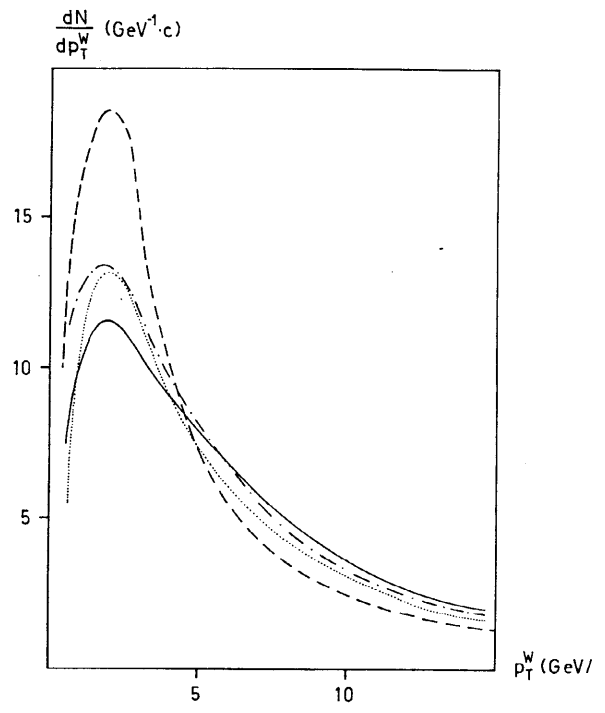

Fig. 4. Transverse momentum distribution of W at 540 GeV. Dashed and dotted the results of Davies et al. [9] with Duke-Owens set 1 and set 2 structure functions, dash-dotted and full model results for same cases.

Several analytic calculations of the W p_T spectrum have been presented. In Fig. 3 we compare our results with the Altarelli et al. [8] calculations, and in Fig. 4 the comparison is with the calculation of Davies et al. [9]. First it should be noted that our method does not give information about K-factors; for comparison with the Davies et al. curves an arbitrary factor K=2 has therefore been introduced. At around 600 GeV, our Monte Carlo results end up somewhere in between these two analytical calculations, as follows. The peak of the Altarelli et al. dN/dp_{TW} curve is almost 1 GeV above that of our results, so that their curve is somewhat less sharply peaked. When comparing with Davies et al., the position of the dN/dp_{TW} peak agrees

well, but their curve is considerably more strongly peaked. Further, we do not obtain anywhere near as strong a dependence on the structure function choice as Davies et al.. At a CM energy of 10 TeV, the shape of our curve agrees well with Altarelli et al. out to a p_{TW} of 50 GeV, but the normalization is wrong, indicating that we underestimate the probability of having p_{TW}>50 GeV. This is traceable to the assumption that the scale of the hard interaction is set solely by M_W^2, which is not true for the tail where $p_{TW} \geqslant M_W/2$; here one could rather use the matrix elements for $qg \to Wq'$ and $q\bar{q}' \to Wg$ to give the basic hard process, which may then become modified by initial state radiation effects in its turn.

Multiple parton-parton interactions

The differential cross section for a hard parton-parton interaction is given by perturbative OCD, as a convolution of hard scattering matrix elements and the structure functions of the incoming hadrons. The integrated cross section of all interactions with $p_T > p_{Tmin}$, $\sigma_{hard}(p_{Tmin})$, is divergent for $p_{Tmin} \to 0$. At present Collider energies, $\sigma_{hard}(p_{Tmin})$ becomes comparable with the total cross section for $p_{Tmin} \approx 1.5$ GeV. This need not lead to contradictions: $\sigma_{hard}(p_{Tmin})$ does not give the hadron-hadron cross section but the parton-parton one. Each of the two incoming hadrons may be viewed as a beam of partons, with the possibility of several parton-parton interactions when the hadrons pass through each other.

In [10] we argue that Collider data indicate a significant probability for multiple interactions at 540 GeV. This conclusion is based on the assumption of jet universality, i.e. that the fragmentation mechanism in hadron collisions is no different from that in e^+e^- annihilation. In the latter process, the Lund string fragmentation model [11] provides an accurate description of most phenomenology known to date, whereas it fails to describe the Collider data. Specifically, the predicted multiplicity distribution is far too narrow, and forward-backward multiplicity correlations are almost absent.

If different parton interactions above p_{Tmin} are assumed to take place (essentially) independently of each other, one obtains a Poissonian multiplicity distribution in the number of interactions, with mean given by $\sigma_{hard}(p_{Tmin})/\sigma_{tot}$, where σ_{tot} is the total inelastic and nondiffractive cross section. With a varying number of interactions, the multiplicity fluctuations are increased, and strong forward-backward multiplicity correlations are introduced. Results are sensitive to the choice of p_{Tmin} value, see Fig. 5, with a reasonable description obtained for p_{Tmin} = 1.6 GeV. Phenomena like "hot spots" and the dependence of mean transverse momentum on the charged multiplicity are also explained in this framework.

Multiple interactions offer an appealing picture of hadron collisions, with a logical continuity from minimum bias to hard interaction events, but many points remain to be studied further. Particularly

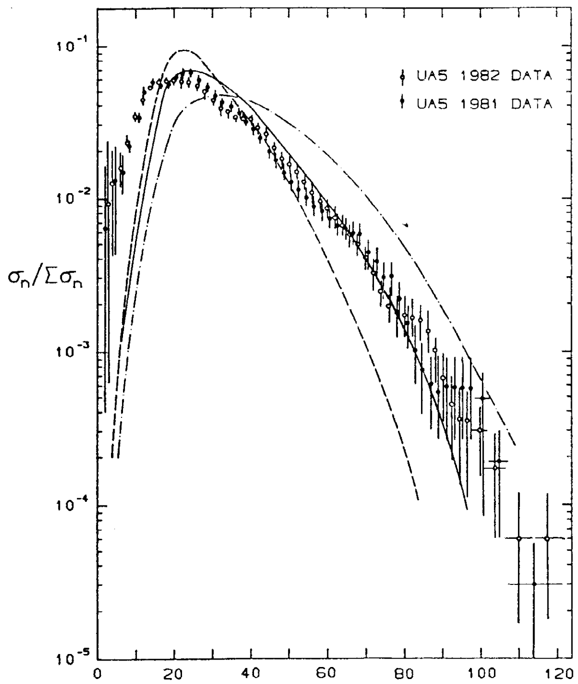

Fig. 5. Charged multiplicity distribution at 540 GeV, UA5 results [12] compared with multiple interaction model: dashed minimum transverse momentum 2.0 GeV, full 1.6 GeV and dash-dotted 1.2 GeV.

nontrivial is the question of string drawing: with a number of partons being scattered into different directions, leaving behind hadron beam remnants, how are all these different objects connected with strings (colour flux tubes or vortex lines)? We have tried some different alternatives without observing much difference, but a number of possibilities remain unexplored.

The parton-parton scattering cross sections given by perturbative QCD are divergent for $p_T \to 0$. An exchanged gluon with small p_T will not be able to resolve the individual colour charges in a colour neutral hadron, however. More appealing than the present sharp cutoff at p_{Tmin} would then be to introduce a scale p_{T0}, such that factors of $\alpha_s(p_T^2)$ should be replaced by $\alpha_s(p_{T0}^2+p_T^2)$ and matrix elements should be

multiplied by $p_T^4/(p_{T0}^2+p_T^2)^2$. A finite cross section is then obtained for $p_T \to 0$.

We have so far assumed that all hadron collisions are equivalent, whereas in fact each collision is also characterized by a varying impact parameter b. A small b value corresponds to a large overlap between the two hadrons, and hence an enhanced probability for multiple interactions. A large b, on the other hand, corresponds to a grazing collision, with a large probability that no parton interactions at all take place. This factor will tend to broaden the minimum bias multiplicity distribution at higher energies. At present energies these effects are not large, since the mean number of interactions anyhow is small. It may explain the "pedestal effect", however: events containing a hard interaction are biased towards small impact parameters, and hence a larger than average multiple interaction probability.

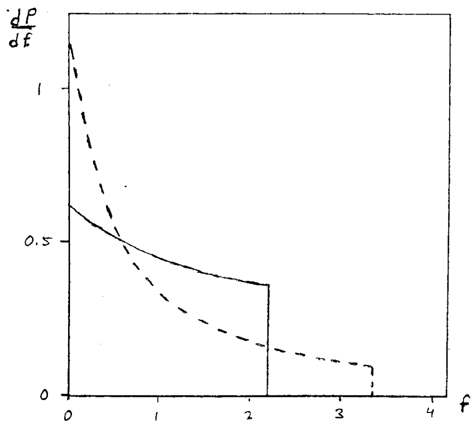

Fig. 6. Distribution of events in relative multiple interaction probability f due to varying impact parameters, full at 540 GeV and dashed at 40 TeV.

As a toy model, one may assume the hadronic matter distribution to be given by a three-dimensional Gaussian, and the parton interaction probability in a given point to be proportional to the product of the matter densities of the two colliding hadrons. Integrating over time and the dimension along the collision axis, and with a suitable choice of scales, the total overlap becomes $\tilde{O}(\tilde{b}) \sim \exp(-b^2)$, i.e. a Gaussian in impact parameter. Further, assume that for a given b the number of parton interactions is distributed according to a Poissonian with mean

$$\langle n_{int}(b)\rangle = k\, e^{-b^2} \sim \tilde{O}(b)$$

where k is an energy-dependent factor. The probability to have at least one interaction when two hadrons pass each other by, and hence to have an observable event, is

$$P_{acc}(\bar{b}) = 1 - P(n_{int}(b)=0) = 1-e^{-ke^{-b^2}}$$

This gives the distribution of events in b values. Combining the two equations above, one obtains the distribution in multiple interaction probability, normalized to the average one, Fig. 6. As the CM energy is increased, so is k. Hence more peripherical collisions are allowed, and the spread in multiple interaction probability is increased.

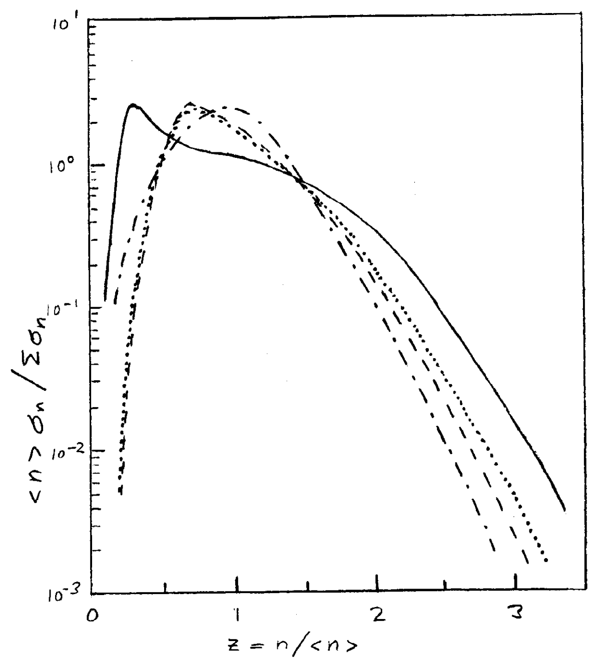

Fig. 7. Scaled multiplicity distribution. Dashed and dotted 540 GeV results with fixed and with varying impact parameter, dash-dotted and full the same for 40 TeV.

In Fig. 7 the KNO distribution is shown, comparing a fixed versus a varying multiple interaction probability. In particular, note the spike at low z for the varying probability case at 40 TeV, which is related to the spike of peripheral events with small interaction probability in Fig. 6; it remains to be seen how sensitive that is to the choice of Gaussian matter distributions.

In this approach, variations in multiplicity do not come from one single source, but from a mixture of different effects: in the basic string fragmentation, in the sharing of energy within a coloured beam remnant, in the flavour and kinematic variables of hard interactions, in associated initial and final state radiation, in the number of scatterings per event for fixed impact parameter, in the distribution of impact parameters and (yet to be studied) in the admixture of diffractive events. This may seem a discouragingly complex scenario, but probably does reflect the multitude of issues that must be understood for accurate predictions of SSC physics!

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