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Matrix-Element Corrections to Parton Shower Simulation of Deep Inelastic Scattering

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Abstract

We discuss the need to supplement angular-ordered parton showers with a hard matrix element correction, to obtain a complete description of final states in deep inelastic scattering. This has recently been implemented in the HERWIG Monte Carlo event generator, and we show the results.

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1 Introduction

Monte Carlo event generators that include coherent parton showers[1–3] have been extremely succesful in describing the final states of e^+e^- anihilation into hadrons. Observables related to the flow of multiplicity relative to the jet directions are particularly sensitive to coherence phenomena, and models that do not account for this coherence are strongly disfavoured by the data. All parton showers effectively resum the large logarithms that arise in the perturbative expansion from collinear emission, but only if coherence is correctly included are those from soft emission also correctly resummed.

However, since parton showers are based on the soft and collinear limits, getting these right is not sufficient to get uniformly reliable predictions, as many observables are specifically sensitive to non-soft non-collinear emission, for example the thrust distribution, or transverse momentum flows. Thus it is essential that the generated distributions agree with the fixed-order matrix-element cross-section, at least for the hardest emission (ie. the highest transverse momentum). Then, if the hardest emission and all soft or collinear emissions are correct, a good overall description should be achieved. In the colour dipole model[3] this is automatically done, but in the parton cascade picture we must correct the algorithms to ensure the agreement of the hardest emission.

In [4], a method was described to correct the *first* emission of a parton shower to the first-order matrix-element in e^+e^- annihilation. However, because they use a virtuality-ordered parton shower this is not necessarily the hardest emission. Furthermore, as we discuss below, and in more detail in [5], the probabilistic parton shower picture is only justified if the phase-spaces for resolved emission, and unresolved or virtual emission are treated equivalently, which is not the case if corrections are only applied to the first resolvable emission. This manifests itself as an unphysical dependence of the hard correction on the infrared cutoff, since lowering the cutoff reduces the fraction of the time that a given emission is treated first.

When using angular-ordered parton showers as in HERWIG[1], these problems become particularly severe. This is because the usual algorithms[6,7] are only valid when emission is restricted to the forward hemisphere, limiting the choice of Lorentz frames in which the algorithm can be used. This has the result that very hard emission cannot be generated at all by the angular-ordered algorithm. Although this makes the need for a matrix-element correction greater, it also reduces the theoretical difficulties, because one can define a region of phase-space in which the matrix-element is used (followed by showering of the resulting partons), and one in which only the parton shower is used, without any competition between the two. The fraction of events that should be generated in the missing region is sufficiently small that complications from form factors and multiple emission within it can be neglected. We call this the 'hard' matrix-element correction, which was described for e^+e^- annhilation in [8]. In this paper, we describe the equivalent correction for DIS, and show first results from it.

In addition to this hard correction, we would also like to correct the distribution of emissions within the parton shower's phase-space, in analogy to [4]. As we discuss in [5], this can be done while ensuring that the correction enters the form factor, simply by uniformly correcting every emission which is the hardest so far, at the stage at which it is treated. Since this is always true of the first emission, this is treated exactly as in [4], but any later emissions that are harder are also corrected, so the hardest emission in the event always agrees with the first-order matrix-element. We call this the 'soft' correction, and show results of it below. As we shall see, it is not a significant correction for observables studied so far in DIS at HERA, but small differences can be seen in e^+e^- annihilation in measures that separate out the energy and angular dependence of rare regions of phase-space[8].

In the remainder of this paper, we recap the main ingredients of the angular-ordered parton shower algorithms, before describing the hard and soft corrections. We then show the numerical results of implementing these corrections in the HERWIG Monte Carlo event generator and see that, as one would expect, they are important for p_t-weighted quantities such as the 'seagull' plot and E_T flow, but do not affect multiplicity flows or jet rates significantly.

2 Parton Showers

In the presence of an infrared cutoff that defines the observability of partons, the cross-section to emit an additional parton in a hard process can be written

$$d\sigma \propto \frac{\alpha_s}{2\pi}\frac{dq^2}{q^2} dz P(z) e^{-\int_{q^2} \frac{\alpha_s}{2\pi}\frac{dq'^2}{q'^2} dz' P(z')}. \tag{1}$$

The first part is the 'naive' cross-section that would be obtained in neglect of soft and virtual gluons, with $P(z)$ the appropriate DGLAP splitting function. The exponential is the resummed effect of all possible virtual and unresolvable gluon emissions. Because the integrand is identical to the naive cross-section, and its phase-space limits governed by the naive phase-space, this can be interpreted in a time-ordered probabilistic way, as the probability that there was no emission 'earlier' than the one in question. However, this does not specify what the ordering variable q^2 should be: (1) is equally valid for any $q^2 = f(z)m^2$, and this is one of many features that differs between algorithms.

Any algorithm that correctly implements (1) effectively resums the leading collinear logarithms to all orders, but large logarithms also arise from soft gluon emission and, since many interfering Feynman diagrams contribute, it appears at first sight that the possibility of a probabilistic picture is lost. However, the coherent nature of gluon emission means that a probabilistic picture can be recovered, provided a suitable choice is made for the ordering variable. For coherent final-state parton showers[6], the appropriate choice is

$$q^2 \equiv E^2\xi = E^2(1 - \cos\theta), \tag{2}$$

where E is the energy of the emitter, and θ is the opening angle of the emission. Because energies in the final state decrease away from the hard process, ordering in q^2 is equivalent to the famous ordering of opening angles. The upper limit for q^2 is set by the coherence of radiation from different partons involved in the hard process, and can also be represented as an angular ordering, with the emission from each parton limited to be within an angular cone that extends as far as the parton's colour partner. This can equivalently be represented as

$$q^2 < p \cdot p_{partner}. \tag{3}$$

Although q^2 becomes Lorentz-invariant in the soft and collinear limits, in general it is not, and we must also define the frame in which (1) will be used. In HERWIG, a separate frame is used for each parton in which it is 90° from its colour partner. This means that the initial condition is

$$\xi < 1. \tag{4}$$

In general, z could be defined as any energy-like fraction, but in HERWIG it is defined as the energy fraction in the showering frame. To be specific, when discussing gluon emission from a quark, z refers to the gluon's energy fraction.

For relatively soft emission from initial-state partons, a similar picture to final-state radiation can be derived, with ordering of angles from the fixed incoming hadron direction. One can represent the emission as a 'backward evolution'[9], in which the evolution scale is reduced away from the hard process, tracing the incoming parton back into the original hadron. The probability for each step of the evolution is then

$$d\sigma \propto \frac{\alpha_s}{2\pi}\frac{dq^2}{q^2}dzP(z)\frac{\frac{x}{z}f(\frac{x}{z},q^2)}{xf(x,q^2_{last})}e^{-\int_{q^2}\frac{\alpha_s}{2\pi}\frac{dq'^2}{q'^2}dz'P(z')\frac{\frac{x}{z'}f(\frac{x}{z'},q'^2)}{xf(x,q^2_{last})}}, \tag{5}$$

where x and q^2_{last} are the energy fraction and q^2 of the previous (harder) branching. The algorithm ensures that at each stage of the evolution, the distribution of x and q^2 is identical to that of the input parton distribution function.

However at small x, where the emitted parton can be much harder than the parton entering the hard process, no such simple picture exists. Fortunately, the bulk of the effects can be incorporated in a probabilistic way by replacing angular ordering by transverse momentum ordering. The coherent initial-state parton shower algorithm[7], interpolates the two regions, by defining

$$q^2 \equiv E^2\xi = E^2(1 - \cos\theta), \tag{6}$$

where E is the energy of the parton that split, and θ is the emission angle from the incoming hadron direction. When all emission is soft, all the energies are similar, and q^2 ordering produces angular ordering, and when the emission is hard, E is similar to the emitted parton's energy, so q^2 ordering produces transverse momentum ordering. The initial condition is again

$$q^2 < p \cdot p_{partner}, \tag{7}$$

and the same frame is used, so ξ is this time limited by

$$\xi < z^2, \tag{8}$$

and z refers to the space-like parton's momentum fraction, ie. the quark for gluon emission.

The infrared cutoff is on the transverse momentum in both the initial- and final-state showers, but in the initial-state case a cut must also be applied to q^2, since this is the scale at which the parton distribution functions are evaluated, and most parametrisations are only reliable above 4 or 5 GeV^2. This is an uncomfortably large value, which is responsible for the 'rapidity gap' events that are generated by HERWIG.

Once the parton shower has terminated, a model for how the remaining hadron remnant breaks up is required. In HERWIG, if the backward evolution has not resulted in a valence quark, further emission resulting in a valence quark is forced. In the current version, 5.7, this forced emission is at the cutoff scale, which can generate a considerable amount of transverse momentum because the scale is so large. In general, one would expect that if the parton's evolution could be traced all the way down to the infrared cutoff scale, it would be evolved into a valence quark, so an improved treatment would be to generate a $\log q^2$ value uniformly between the q^2 cutoff and the infrared cutoff. As we shall see, this makes a sizeable difference to the predictions. Finally, the valence quark is given an intrinsic transverse momentum, with adjustable Gaussian width, which defaults to zero.

After generating initial- and final-state showers, the original partons are off mass-shell, so their energy and momentum cannot be individually conserved. This is solved by applying a separate boost to each jet along its direction, such that energy and momentum are globally conserved, although the jet momenta are no longer equal to the parton momenta. In relating the showering variables z and ξ to the kinematic variables, it is useful to note that the virtuality of the jet, and transverse momenta of emissions relative to the jet, are conserved by this reconstruction procedure.

3 Hard Corrections

As discussed in [10], the corrections to any quark scattering process,

$$\mathrm{q}(q_1) + X \quad\rightarrow\quad \mathrm{q}'(q_2) + X' \tag{9}$$

to first ordered, at tree level,

$$\begin{aligned} \mathrm{QCDC:}\quad \mathrm{q}(p_1) + X &\rightarrow \mathrm{q}'(p_2) + \mathrm{g}(p_3) + X', \\ \mathrm{BGF:}\quad \mathrm{g}(p_3) + X &\rightarrow \mathrm{q}'(p_2) + \bar{\mathrm{q}}(p_1) + X', \end{aligned} \tag{10}$$

can be written in terms of the lowest order matrix element, $\mathcal{M}$, as

$$\begin{aligned} \frac{d\sigma_{\mathrm{QCDC}}}{d\sigma_0} &= \frac{C_{\mathrm{F}}\alpha_s}{2\pi}\frac{-x_{\mathrm{B}}x_1 f_q(-x_{\mathrm{B}}x_1,Q^2)}{x_{\mathrm{B}}f_q(x_{\mathrm{B}},Q^2)}\frac{dx_1dx_2}{-x_1^3(-1-x_1)(1-x_2)}\frac{d\phi}{2\pi} \\ &\quad \left\{x_1^2\,\frac{|\mathcal{M}(r_1,q+r_1)|^2}{|\mathcal{M}(q_1,q_2)|^2} \;+\; x_2^2\,\frac{|\mathcal{M}(r_2-q,r_2)|^2}{|\mathcal{M}(q_1,q_2)|^2}\right\}, \\ \frac{d\sigma_{\mathrm{BGF}}}{d\sigma_0} &= \frac{T_{\mathrm{R}}\alpha_s}{2\pi}\frac{-x_{\mathrm{B}}x_3 f_g(-x_{\mathrm{B}}x_3,Q^2)}{x_{\mathrm{B}}f_q(x_{\mathrm{B}},Q^2)}\frac{dx_1dx_2}{x_3^2(1-x_2)}\frac{d\phi}{2\pi} \\ &\quad \left\{x_1^2\,\frac{|\mathcal{M}(r_1,q+r_1)|^2}{|\mathcal{M}(q_1,q_2)|^2} \;+\; x_2^2\,\frac{|\mathcal{M}(r_2-q,r_2)|^2}{|\mathcal{M}(q_1,q_2)|^2}\right\}, \end{aligned} \tag{11}$$

where $x_i \equiv \frac{2p_i\cdot q}{q\cdot q}$ are the longitudinal momentum fractions in the Breit frame, with $-x_1+x_2+x_3=2$ for QCDC, and $x_1+x_2-x_3=2$ for BGF. The massless vectors $r_1=-p_1/x_1$ and $r_2=p_2/x_2$ are parallel to p_1 and p_2 but with unit longitudinal momentum fraction.

For arbitrary lepton-quark scattering, the ratios of matrix-elements have identical small-y_{B} limits, and in general these are already reasonable approximations for $y_{\mathrm{B}} \lesssim 0.5$. We use this limit throughout, since it allows us to formulate the corrections completely independently of the process that scattered the quark. Switching to the more convenient variables $x_p = Q^2/(\hat{s}+Q^2)$ and $z_p = p_2\cdot P/q\cdot P$, where P is any massless vector parallel to the incoming hadron, we obtain

$$\begin{aligned} \frac{d\sigma_{\mathrm{QCDC}}}{d\sigma_0} &= \frac{C_{\mathrm{F}}\alpha_s}{2\pi}\frac{\frac{x_{\mathrm{B}}}{x_p}f_q(\frac{x_{\mathrm{B}}}{x_p},Q^2)}{x_{\mathrm{B}}f_q(x_{\mathrm{B}},Q^2)}\frac{dx_pdz_p}{(1-x_p)(1-z_p)}\frac{d\phi}{2\pi} \\ &\quad \left\{1+x_p^2(x_2^2+\tfrac{3}{2}x_\perp^2)\frac{(1-\sin\theta_2\cos\phi)^2}{1+\frac{1}{2}\sin^2\theta_2}\right\}, \\ \frac{d\sigma_{\mathrm{BGF}}}{d\sigma_0} &= \frac{T_{\mathrm{R}}\alpha_s}{2\pi}\frac{\frac{x_{\mathrm{B}}}{x_p}f_g(\frac{x_{\mathrm{B}}}{x_p},Q^2)}{x_{\mathrm{B}}f_q(x_{\mathrm{B}},Q^2)}\frac{dx_pdz_p}{1-z_p}\frac{d\phi}{2\pi} \\ &\quad \left\{x_p^2(x_1^2+\tfrac{3}{2}x_\perp^2)\frac{(1+\sin\theta_1\cos\phi)^2}{1+\frac{1}{2}\sin^2\theta_1}+x_p^2(x_2^2+\tfrac{3}{2}x_\perp^2)\frac{(1-\sin\theta_2\cos\phi)^2}{1+\frac{1}{2}\sin^2\theta_2}\right\}. \end{aligned} \tag{12}$$

where $\sin\theta_i$ is the Breit frame angle between parton i and the current direction,

$$\sin\theta_i = \frac{x_\perp}{\sqrt{x_i^2 + x_\perp^2}}, \tag{13}$$

$x_\perp$ is the scaled transverse momentum,

$$x_\perp^2 = \frac{4p_t^2}{Q^2} = 4(1-x_p)(1-z_p)z_p/x_p, \tag{14}$$

and ϕ is the azimuth between the hadronic and leptonic planes. First order events can then be generated by first generating lowest-order events, and then generating x_p, z_p according to (12) within a given phase-space region.

The important point about a hard matrix-element correction to a parton shower is that their phase-spaces be complementary: that the matrix element only generates events that the parton shower cannot, and vice versa, with no configurations that cannot be produced by either. To this end, we must relate the parton shower variables, z, ξ, to the matrix-element variables, x_p, z_p, using the fact that the virtuality of the jet, and momenta transverse to it are conserved during the kinematic reconstruction. For final-state emission we have

$$\begin{aligned} \tfrac{1}{2}Q^2\frac{z^2(1-z)^2}{1-2z(1-z)\xi}\xi(2-\xi) &= p_t^2 = Q^2(1-x_p)(1-z_p)z_p/x_p, && (15)\\ Q^2 z(1-z)\xi &= m^2 = Q^2(1-x_p)/x_p, && (16) \end{aligned}$$

which lead to

$$\begin{aligned} x_p &= \frac{1}{1+z(1-z)\xi}, && (17)\\ z_p &= \tfrac{1}{2}\left(1-\frac{2z-1}{\sqrt{1-2z(1-z)\xi}}\right), && (18) \end{aligned}$$

or

$$\begin{aligned} z &= \tfrac{1}{2}\left(1-(2z_p-1)\sqrt{(3x_p-2)/x_p}\right), && (19)\\ \xi &= \frac{2(1-x_p)}{1-x_p+2(3x_p-2)z_p(1-z_p)}. && (20) \end{aligned}$$

The limit $\xi < 1$ translates to

$$x_p > \frac{1+4z_p(1-z_p)}{1+6z_p(1-z_p)}. \tag{21}$$

For initial-state emission we have

$$\begin{aligned} \tfrac{1}{2}Q^2\frac{(1-z)^2}{z^2}\xi(2-\xi) &= p_t^2 = Q^2(1-x_p)(1-z_p)z_p/x_p, && (22)\\ -Q^2\frac{1-z}{z^2}\xi &= m^2 = -Q^2(1-z_p)/x_p, && (23) \end{aligned}$$

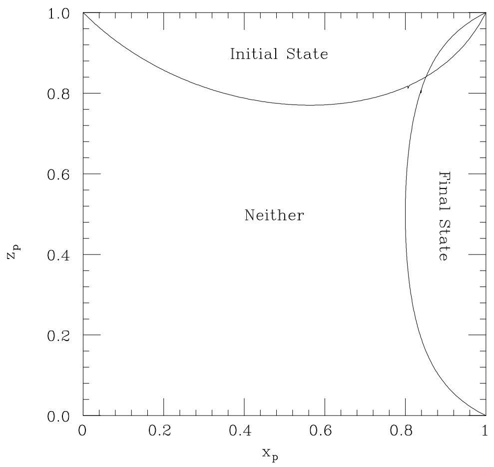

Figure 1: Phase-space limits of HERWIG's parton showers in DIS.

which lead to

$$x_p = \frac{z^2 + (1-z)\xi - \sqrt{(z^2-(1-z)\xi)^2 + 2z^2(1-z)^2\xi(2-\xi)}}{2(1-z)\xi}, \quad (24)$$

$$z_p = \frac{z^2 - (1-z)\xi + \sqrt{(z^2-(1-z)\xi)^2 + 2z^2(1-z)^2\xi(2-\xi)}}{2z^2}, \quad (25)$$

or

$$z = \frac{-\sqrt{x_p}\left(\sqrt{x_p} - \sqrt{x_p + 2(1-z_p)(1-z_p+x_pz_p)}\right)}{1-z_p}, \quad (26)$$

$$\xi = \frac{\left(\sqrt{x_p} - \sqrt{x_p + 2(1-z_p)(1-z_p+x_pz_p)}\right)^2}{1 + x_p - z_p - \sqrt{x_p(x_p + 2(1-z_p)(1-z_p+x_pz_p))}}. \quad (27)$$

The limit $\xi < z^2$ translates to

$$z_p > 1 - \tfrac{2}{3}x_p\left(1 + \Re\left[a^{1/3}\left(\tfrac{1}{2} + \tfrac{\sqrt{3}}{2}i\right)\right]\right), \quad (28)$$

$$a = 10 - 45x_p + 18x_p^2 + 3i\sqrt{3(9 + 66x_p - 93x_p^2 + 12x_p^3 - 8x_p^4 + 24x_p^5 - 8x_p^6)}. \quad (29)$$

For the evolution of quarks back in to gluons, corresponding to BGF in the matrix element, we obtain exactly the same relations, but with no final-state radiation component of course. Note that these limits are independent of all parameters, even x_{B} and Q^2. They are shown in Fig. 1, which is similar to that found for e^+e^- annihilation[11].

A simple Monte Carlo procedure can be obtained by first integrating the parameter-independent part of the matrix-element (ie. everything except the α_s factor and ratio of distribution functions) over the region of phase-space not populated by the parton shower. Then in each event, an x_p, z_p, ϕ point is generated within that region. The parameter-dependent part is calculated and multiplied by the integral to give the probability of keeping the generated point. The fraction of generated events kept varies with $x_{\rm B}$ and Q^2, but is typically around 10% for $Q^2 \sim 1000$ GeV^2 rising to around 25% for $Q^2 \sim 10$ GeV^2 at HERA. It is dominated by the region $p_t \sim Q/4$. This fraction is much higher than in e$^+$e$^-$ annihilation, where it is 2.5%, hardly varying with Q^2.

4 Soft Corrections

As well as filling phase-space regions not populated by the parton shower, we would also like to correct the distribution of emissions within its phase-space. However, theoretical considerations become more important for corrections within the parton shower phase-space. This is because the whole approach of using a probabilistic parton shower picture is only valid if we can guarantee that the exponentiated factor from unresolved real and virtual emission is identical to the probability for non-emission within the resolved phase-space region. This means that we cannot use concepts such as the 'first' emission, since this depends on the choice of gauge, ordering variable and on the infrared cutoff. However, as we prove in [5], the general idea of [4] can be used, with a simple modification.

The first stage of the correction is to relate the cross-section produced by HERWIG, to the matrix-element cross-section. For final-state emission, we obtain

$$dz d\xi = \frac{\sqrt{(3x_p - 2)/x_p}}{z(1-z)x_p^2} dx_p dz_p, \tag{30}$$

and for initial-state emission,

$$dz d\xi = \frac{z^3(1 - x_p - z_p + 2x_p z_p)}{x_p^2(1-z)(z + \xi(1-z))} dx_p dz_p. \tag{31}$$

Using these expressions, the matrix-element cross-sections (12), and HERWIG's cross-sections,

$$\begin{aligned}
\frac{d\sigma_{\rm final-state}}{d\sigma_0} &= \frac{C_{\rm F}\alpha_s}{2\pi}\frac{d\xi}{\xi}dz\frac{1+(1-z)^2}{z}, \\
\frac{d\sigma_{\rm initial-state}}{d\sigma_0} &= \frac{C_{\rm F}\alpha_s}{2\pi}\frac{\frac{x_{\rm B}}{z}f_q(\frac{x_{\rm B}}{z},Q^2)}{x_{\rm B}f_q(x_{\rm B},Q^2)}\frac{d\xi}{\xi}dz\frac{1+z^2}{1-z}, \\
\frac{d\sigma_{\rm gluon-splitting}}{d\sigma_0} &= \frac{T_{\rm R}\alpha_s}{2\pi}\frac{\frac{x_{\rm B}}{z}f_g(\frac{x_{\rm B}}{z},Q^2)}{x_{\rm B}f_q(x_{\rm B},Q^2)}\frac{d\xi}{\xi}dz(z^2+(1-z)^2),
\end{aligned} \tag{32}$$

we obtain the ratio of cross-sections, $d\sigma_{\rm HW}/d\sigma_{\rm ME}$. Unfortunately, unlike the algorithm of [4], this is not always larger than 1, so our algorithm must be able to enhance emission as well as reduce it.

It is straightforward to show that one can uniformly enhance emission by an integer factor n, simply by choosing n values of q^2 according to the usual algorithm, and then keeping the largest. This multiplies the emission cross-

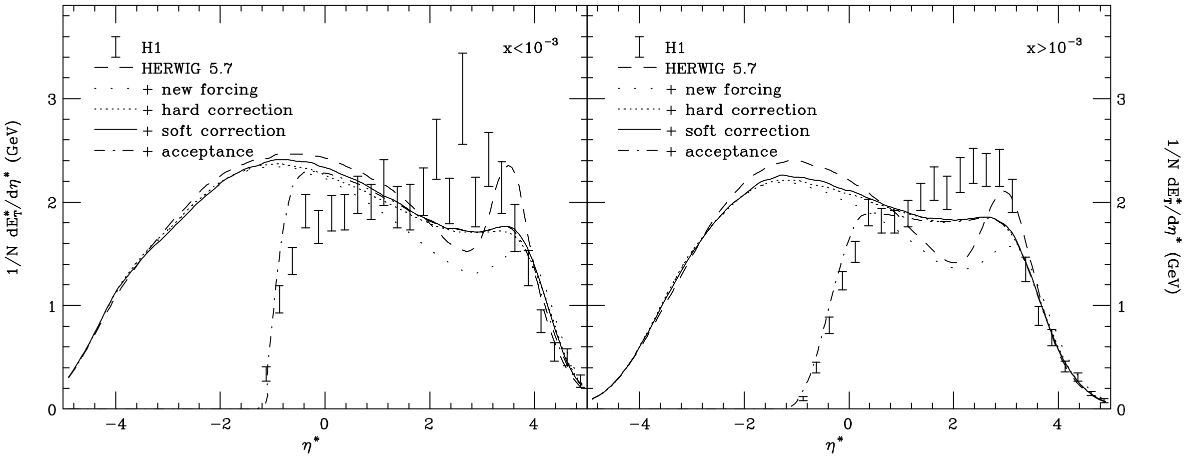

Figure 2: Transverse energy flows in the hadronic cms[12], compared with the predictions of the previous version of HERWIG (dashed), after adding the new procedure for forcing non-valence quark splitting (dotted), the hard matrix-element correction (fine-dotted), the soft matrix-element correction (solid), and the experimental acceptance (dot-dashed). The cuts are as in [12], and the proton direction is to the left.

section by n in both the resolved phase-space, and the form factor[5]. In our case, the cross-section ratio is always larger than 1/2, so we set $n = 2$.

When comparing cross-sections from matrix elements and parton showers, it is important to note that the cross-section describes the inclusive production rate. So when comparing with the first-order matrix element, all emission in which the emitting current is the same as the lowest-order current (ie. in which recoil from other emissions is negligible) should be included. This can be done, while ensuring that the correction enters into the form factor, by uniformly correcting every emission that is the hardest so far[5]. This is always true of the first emission, which is therefore treated exactly like the algorithm of [4]. However, in virtuality-ordered parton showers it sometimes happens that a later emission is harder than the first, and in angular-ordered parton showers it often happens. In our algorithm, the hardest emission in the event is always corrected to the first-order matrix element, while in the algorithm of [4], this is often but not always the case.

5 Numerical Results

The most striking need for a hard matrix-element correction can be seen in comparisons published by H1 on transverse momentum distributions[12]. Not only is HERWIG in some disagreement with the data, but its strange shape clearly shows that some regions of phase-space are missing.

In Fig. 2, we show the predictions of the previous version of HERWIG for the E_T flow in the hadronic cms, together with versions that include the new treat-

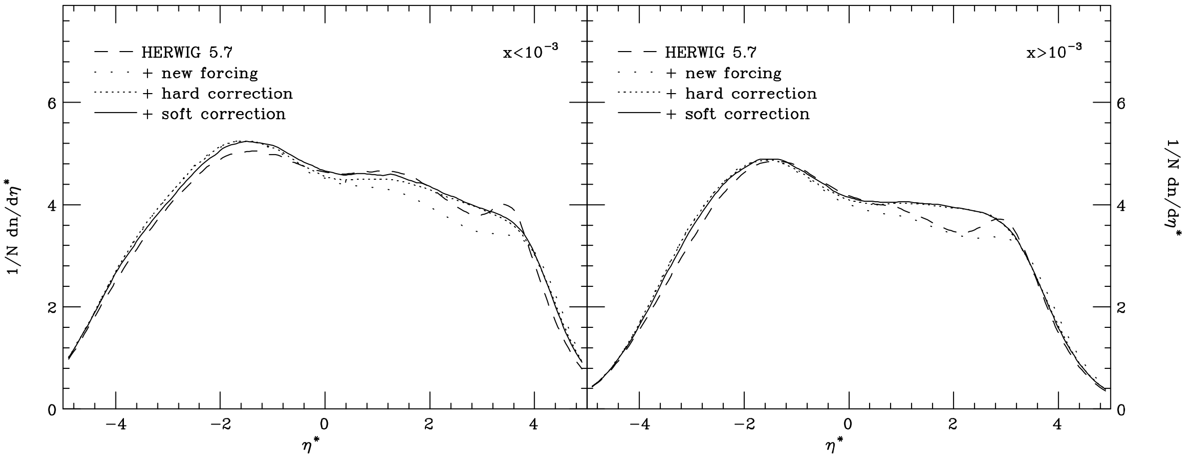

Figure 3: Multiplicity flows in the hadronic cms. Curves are as in Fig. 2.

ment of forced sea quark splitting, the hard correction and the soft correction, as well as the data from [12]. It is clear that both the forced splitting and hard correction are important, while the soft correction makes little difference. Going to higher Q^2 values, one would expect the importance of the forced splitting to decrease rapidly, and that of the hard correction to decrease slowly. For ease of comparison, we set all parameters at the values used in [12], but we have found that returning them to their default values (which are tuned to e^+e^- data[13]) improves the overall agreement with H1's data.

In Fig. 3, we show the equivalent plots for the multiplicity flows, where the corrections are not as important, as one would expect.

In Fig. 4, we show the multiplicity distribution in x_F, which is almost unchanged, and the'seagull' plot, where the shape has become more reasonable, and the agreement with data improved.

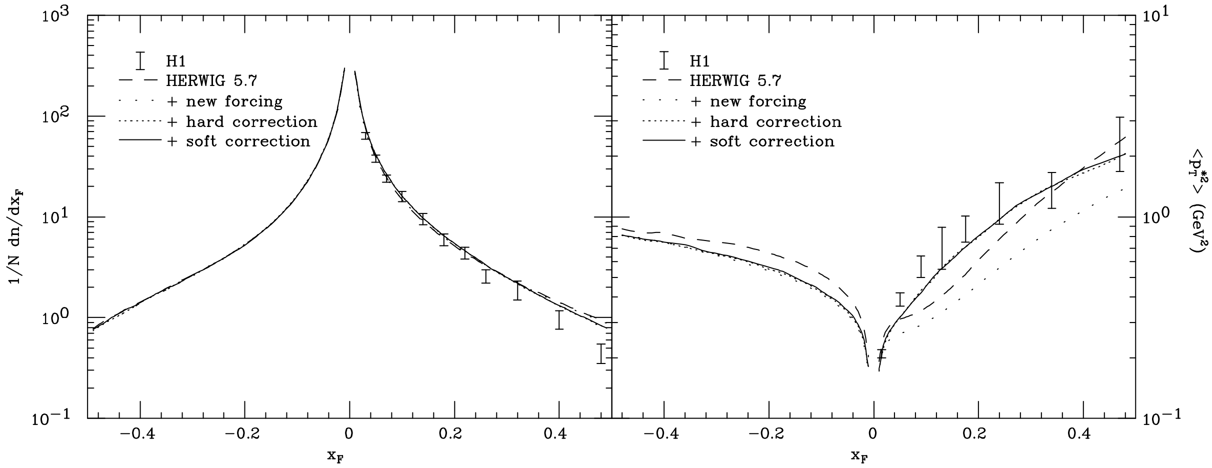

Figure 4: The longitudinal momentum spectrum of charged particles in the hadronic cms (left) and their average transverse momentum (right). Curves are as in Fig. 2, data and cuts are from [12].

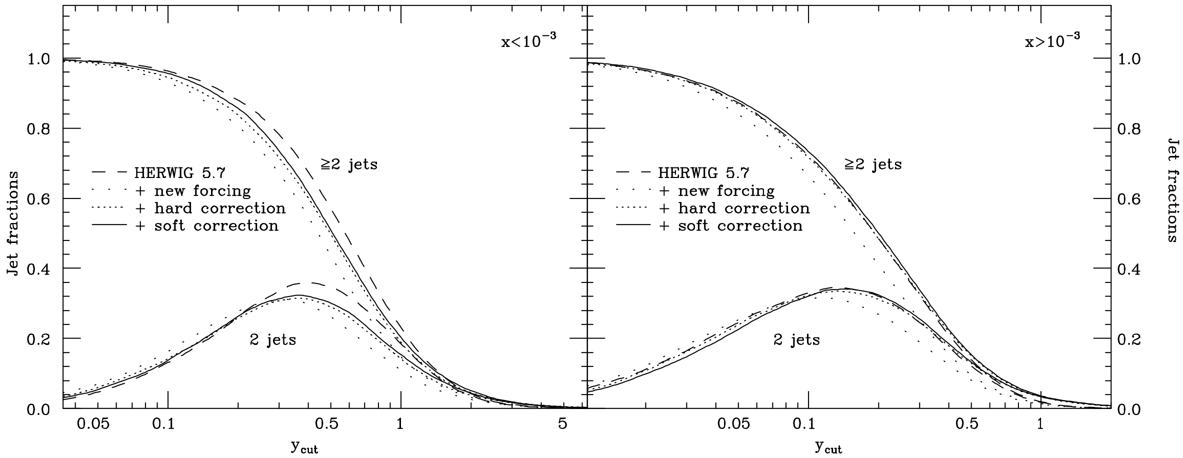

Figure 5: The fraction of events containing two or more jets (in addition to the beam remnant) according to the factorisable $k_\perp$ jet algorithm of [14]. Curves are as in Fig. 2.

Although the E_T flow and seagull plot show that the distribution of high-p_t emissions has changed considerably, we see in Fig. 5 that the jet rate is hardly affected, but is slightly softened by the new treatment of forced splitting, and hardened again by the matrix-element correction, leading to a similar overall rate in the large $x_{\rm B}$ sample, and somewhat softened at small $x_{\rm B}$.

6 Summary

Despite the success of angular-ordered parton showers in describing multiplicity distributions, they must be supplemented by matrix element corrections in order to describe final-state features that are sensitive to hard emission. We have described how these corrections can be implemented without spoiling the theoretical consistency of the model. For the hard correction, corresponding to phase-space regions that cannot be filled by the parton shower, this is achieved by using complementary kinematic cuts in the matrix element. For the soft correction, which corrects the distribution of emissions within the parton shower's phase-space, it is achieved by uniformly correcting each emission that is harder than any previous emissions.

The hard matrix-element correction makes a significant difference to experimental observables that are p_t-weighted, such as the E_T flow and 'seagull' plot, but hardly affects multiplicity distributions or jet rates. The soft correction makes very little difference to any of the distributions we have studied.

We have also seen that a modification to the way in which sea quark splitting is forced has a significant impact on the predictions. This is an unfortunate effect of the fact that most parton distribution sets are only reliable above a few GeV^2. Its significance should rapidly decrease as higher-Q^2 data becomes available.

The combined effect of the three corrections has been to improve the predictions for p_t-weighted quantities, both in their overall shape, and in their agreement with data. The corrections, together with a soft correction for e^+e^- annihilation, should soon be available in the next version of HERWIG, provisionally called 6.0.

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