CERN-TH.4072/84

SQUARK AND GLUINO PRODUCTION AT THE $p\bar{p}$ COLLIDER

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ABSTRACT

We present detailed calculations of the missing energy and jet signatures for squark and gluino pair-production at the CERN $p\bar{p}$ collider, incorporating three improvements over previous analyses. These include smearing in transverse momentum at production, the fragmentation of supersymmetric partons into supersymmetric hadrons, and the hadronization of their decay products as well as of the residual jets. We find that previous estimates of supersymmetric signatures are not substantially altered by the inclusion of smearing and fragmentation. We present rapidity and transverse momentum distributions for both primary and secondary jets, and discuss variations in the trigger conditions which could increase the number of two-jet events and help discriminate between different models for the monojet events. We find that gluinos give fatter jets than do squarks of the same mass, although it is not easy to understand the thinness of the observed monojets if they are due to squarks of mass $O(40)$ GeV. We emphasize that the key feature to clarify the supersymmetric interpretation of missing p_T events is the study of secondary jet production in mono- and dijet events.

CERN-TH.4072/84

December 1984

1. - INTRODUCTION

Much current theoretical effort is directed towards solving one or the other of two fundamental problems: naturalness or flavour. There is little indication of the mass scale at which the flavour problem may be resolved, but naturalness requires some new physics at or below 1 TeV¹⁾. Approaches to this problem fall into two main classes, composite models and supersymmetry. Many representatives of both these scenarios actually suggest the appearance of new particles or other novel phenomena at or below 100 GeV^{2),3)}. The CERN $p\bar{p}$ Collider has recently opened this energy régime up to preliminary exploration. Several categories of puzzling events have already been found, and both composite modellers and supersymmetrists derive encouragement from some of the events observed. We are particularly interested in the missing energy events reported⁴⁾ by UA1, especially the monojets. If these events are to be interpreted in terms of new physics, we favour a supersymmetric interpretation.

It has long been known that supersymmetry had a missing energy signature⁵⁾. Calculations⁶⁾ have suggested that it should be possible to observe pair-production of squarks $\tilde{q}$ or gluinos $\tilde{g}$ at the $p\bar{p}$ Collider, if their masses were below 0(50) GeV. It might also be possible⁷⁾ to observe squarks with mass 0(100) GeV if the gluino mass were only a few GeV. At first sight, one would expect the pair production of squarks or gluinos to yield multi-jet events with missing energy. In fact, we discovered⁸⁾ that under the trigger conditions used by UA1 for collecting their 1983 data⁴⁾, monojets should predominate if $m_{\tilde{q}}$ or $m_{\tilde{g}} \leq 0(40)$ GeV. Even if they were contaminated by backgrounds the relatively small number of missing energy events in the 1983 data could be used to argue that masses for squarks or gluinos below 0(40) GeV could be excluded. We even found⁹⁾ that both squarks and gluons of such a mass were qualitatively consistent with the number and distribution in missing p_T of the observed missing energy events¹⁰⁾. The squark interpretation was possibly slightly preferred^{9),11)} for a combination of two reasons, namely its tendency to give thinner jets with a harder missing p_T spectrum. However, no definite preference could be based on the small statistics of the 1983 data⁴⁾. The 1984 collider run provides a substantial increase in the available number of monojet and other missing energy events, and it is now appropriate to study the signatures of squarks and gluinos in more detail.

Previous calculations^{8),9),12)} were made at the parton level, with no attempt made to incorporate transverse momentum smearing for the two colliding partons, no allowance made for energy reduction when the outgoing spartons fragment into supersymmetric hadrons, and no attempt at hadronization of the final state jets. We expected the previous estimates of jet topologies and missing energy signatures to remain roughly correct even after these effects were taken

into account, but it was difficult to predict the masses or charged multiplicities unless hadronization was considered.

In this paper, we present results from an improved program including initial state smearing in transverse momentum, parton to hadron fragmentation, and hadronization. We have tried two prescriptions for transverse momentum smearing, one in which the produced sparticle pair is given a transverse momentum with a distribution¹³⁾ inspired by results for $W^\pm$ and Z^0 production, and which is subsequently shared equally between the two produced sparticles, and one in which the initial state quarks and gluons are given a transverse momentum distribution computed¹⁴⁾ using QCD. As one might expect, the two prescriptions give similar results, and while they do broaden the missing p_T spectrum, the effect is not large because $\langle p_T^2 \rangle^{1/2} \ll m_{\tilde{q}}, m_{\tilde{g}}$. We have also treated two possible contributions to the supersymmetric parton ($\tilde{q}$ or $\tilde{g}$) to hadron ($\tilde{q}\bar{q}$ or $\tilde{g}g, \tilde{g}q\bar{q}$) fragmentation function, which should be convoluted together. One is the non-perturbative Ansatz of Peterson et al.¹⁵⁾, and the other is perturbative¹⁶⁾ QCD. We find that perturbative QCD corrections are not important for $m_{\tilde{q}} \sim 40$ GeV, but they could be important for light gluinos^{7),17),18)}. While these fragmentation corrections do not greatly soften the missing p_T signature for $\tilde{q}\bar{q}$ production, they can produce small residual jets pointing in the primordial sparton directions. We take these into account in our subsequent analysis, and find that they can give rise to a few three-jet events in $\tilde{q}\bar{q}$ production.

The third new step in our calculations is to hadronize the supersymmetric hadron decay products and also the residual jets mentioned above. This we do by adapting the independent jet fragmentation routines developed and used by the TASSO Collaboration¹⁹⁾ for analyzing data from PETRA. The bulk of these data are very well described by independent jet fragmentation²⁰⁾, which is adequate for our purposes. We discuss the invariant masses and charged multiplicities to be expected for monojets coming from pair-production of squarks with $m_{\tilde{q}} \sim 40$ GeV or of gluinos with $m_{\tilde{g}} \sim 40$ GeV, and also from single production of squarks with $m_{\tilde{q}} \sim 80$ GeV. The apparent thinness of the monojets observed in 1983 is a problem for all three explanations, particularly for $\tilde{g}g$ production and least for single $\tilde{q}$ production. However, a definite conclusion requires a careful analysis of backgrounds as well as more monojet data, and also passing our generated events through the UA1 detector simulation.

The organization of this paper is as follows. In Section 2 we introduce and discuss our prescriptions for initial p_T smearing and for sparton to hadron fragmentation. In Section 3 we present results for jet energy distributions in

both monojet and dijet events, as well as missing p_T^2 and rapidity distributions for both primary and secondary jets. We stress that squarks given harder jet p_T and missing p_T spectra than do gluinos of the same mass. Moreover, dijet events pile up against the UA1 cuts in jet p_T and missing p_T , making their relative number very sensitive to the way in which experimental cuts are applied. While a detailed evaluation would require passing our hadronized events through the UA1 detector simulation, we do suggest ways in which the UA1 trigger conditions could be modified in such a way as to enhance the dijet sample if the observed monojets are due to $\tilde{q}\tilde{q}$ or $\tilde{g}\tilde{g}$ production and decay. It is clear that a quantitative understanding of dijet events is essential for discriminating between different models for the monojets. In Section 4 we discuss the hadronization of sparticle decay products and of the residual fragmentation jets. We study the multiplicities and the invariant masses of the charged particles in monojets to be expected in different models, investigating the dependence on different fragmentation prescriptions and how this can be reduced by selecting only charged hadrons above certain p_T cuts. Monojets from $\tilde{g}\tilde{g}$ production are broader than those from $\tilde{q}\tilde{q}$ production and decay, primarily because $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$ decay gives broader jets than $\tilde{q} \rightarrow q\tilde{\gamma}$. However, both are broadened by jet coalescence, and the narrowest jets come from single $\tilde{q}$ production.

Our results could be altered by experimental effects such as the difference between parton jet and hadron energies, errors in hadron energy measurement and beam fragmentation effects. We expect these effects to enhance slightly the missing energy signal, but this could only be confirmed by passing hadronized sparticle production events through a detector simulation. Our program is now at the level of an event generator which can be used for this purpose, and is available to anyone wanting to use it in this way.

2. - TRANSVERSE MOMENTUM SMEARING AND FRAGMENTATION

2.1 Smearing in transverse momentum

In previous calculations^{8)-10),12)} of sparticle production, we and others have neglected the p_T of the initial partons (g or q) colliding to form sparticle pairs ($\tilde{g}\tilde{g}$ or $\tilde{q}\tilde{q}$). We expect the initial partons to have intrinsic non-perturbative p_T convoluted with perturbative QCD corrections due to the emission of quasi-soft gluons at scales up to the subprocess centre-of-mass energy $\sqrt{s}$. As long as the p_T they produce is much less than $\sqrt{s}$, these perturbative QCD corrections factorize and are universal. The tail at $p_T = 0(\sqrt{s})$ is process-dependent,

and has not yet been calculated for sparticle pair-production. The smearing in p_T has been studied intensively for $W^\pm$ and Z^0 production^{13),21)}, where perturbative QCD calculations agree well with the data. A simple parametrization of this p_T smearing is available¹³⁾:

$$\frac{dN}{dp_T^2} \propto \exp\left(\frac{-25 p_T}{\sqrt{s}}\right) \quad (2.1)$$

and has been used in calculations of heavy quark production and decay signatures. The standard prescription²²⁾ has been to use (2.1) for the p_T distribution of the produced pair of heavy objects, then assign $p_{T/2}$ equally to each of them, and finally let them decay independently. This is clearly an incomplete prescription, since the two initial partons combining to give the distribution (2.1) will not in general have equal p_T themselves. One should really generate their p_T distributions independently, and then generate the outgoing heavy objects' momenta using the full subprocess differential cross-section $d\sigma/dt(s,t)$. Moreover, the p_T distributions expected for initial quarks and gluons differ because of their different colour charges. Useful parametrizations are¹⁴⁾:

$$\frac{dN}{dp_T^2} = \frac{A}{(1+p_T^2/B)^{BC}} \begin{cases} \ln A = \left(\frac{.2097}{\ln Q}\right) + 1.347 - 1.588 \ln Q \\ \ln B = \left(\frac{1.207}{\ln Q}\right) + 1.065 + .0143 \ln Q \\ \ln C = \left(\frac{1.290}{\ln Q}\right) - .9189 - .2770 \ln Q \end{cases} \text{ for } q \quad (2.2)$$

$$\begin{cases} \ln A = \left(\frac{-.8901}{\log_{10} Q}\right) + 3.725 - 3.127 \log_{10} Q \\ \ln B = \left(\frac{3.802}{\log_{10} Q}\right) - 4.065 + 1.521 \log_{10} Q \\ \ln C = \left(\frac{-1.103}{\log_{10} Q}\right) + 2.877 - 1.642 \log_{10} Q \end{cases} \text{ for } g$$

We do not expect big errors to be induced by not treating the initial partons' p_T independently, as long as $\langle p_T^2 \rangle^{\frac{1}{2}} \ll m_q, m_g$. We have verified that the two prescriptions (2.1) and (2.2) seem to give similar results for the missing p_T physics of interest here. In what follows, for technical reasons we mainly use the lumped parametrization (2.1).

Figure 1 shows the smeared p_T distribution we obtain for the pair-production of squarks weighing 40 GeV at $\sqrt{s} = 630$ GeV. We find that $\langle p_T^2 \rangle^{\frac{1}{2}} = 8.8$ GeV $\ll m_q$. The effect of this p_T smearing on the missing p_T distribution is shown in Fig. 2. We see that there is not a great deal of difference, although the smeared distribution is of course somewhat wider. We therefore conclude that our previous results were stable against p_T smearing.

2.2 Supersymmetric parton to hadron fragmentation

As mentioned earlier, there are two distinguishable contributions to this fragmentation. The non-perturbative contribution we model using the Ansatz of Petersen et al.¹⁵⁾:

$$D_{NP}(z) = \frac{N}{z(1-\frac{1}{2} - \frac{\epsilon}{1-z})^2} \quad (2.3)$$

where z is the light-like energy-momentum fraction:

$$z = \frac{E_h + p_h''}{E_j + p_j''} \quad (2.4)$$

The other contribution is perturbative QCD gluon bremsstrahlung¹⁶⁾ which gives in the leading logarithmic approximation

$$F(z) = A(1-z)^{A-1} : A = c \ln\left(\frac{\alpha_s(\tilde{m})}{\alpha_s(E_j)}\right) \quad (2.5)$$

where

$$c = \begin{cases} 16/25 & \text{for } \tilde{q} \\ 36/25 & \text{for } \tilde{g} \end{cases} \quad (2.6)$$

Convoluting the two contributions:

$$D(z) = D_{NP}(z) \otimes F(z) \quad (2.7)$$

is a delicate exercise for several reasons. One obvious remark is that since we will be interested in squark and gluino jets with energies not much larger than the squark and gluino masses, the leading logarithmic approximation (2.5) is not really appropriate, and one would be better advised to use a complete lowest order QCD perturbation theory calculation of $\tilde{q}\tilde{q}g$ and $\tilde{g}g\tilde{g}$ final states. Such a calculation is not available, and we can only hope that (2.5) is, because of finite mass effects, an upper bound on perturbative QCD corrections. Another difficulty is that the perturbative QCD contribution (2.5) requires a renormalization scale μ_0 for its specification. This should presumably be taken as the scale at which the non-perturbative Ansatz (2.3) is the full fragmentation function. The parametrization (2.3) has been developed for c and b quarks with a PEP beam energy of 14.5 GeV in mind¹⁵⁾, and has been used successfully at a range of comparable PETRA beam energies. Good fits to these data are given by

$$\epsilon = 0.15 \left(\frac{m_c}{m_h} \right)^2 \quad (2.8)$$

where m_h is the heavy parton/hadron mass. However, when we seek to apply the non-perturbative Ansatz (2.3) to heavier squarks and sleptons, there is no unambiguous prescription for the jet energy scale where it is directly applicable, or in other words what value of μ_0 to use in evaluating the perturbative correction (2.5).

We take $\mu_0 = 15$ GeV for light gluinos: $m_{\tilde{g}} \sim 5$ to 10 GeV, and

$$\mu_0 = m_{\tilde{q}, \tilde{g}} + 10 \text{ GeV} \quad (2.9)$$

for squarks and gluinos weighing 40 GeV or more. Figure 3 shows the non-perturbative (2.3) and (2.8) and perturbative (2.5) and (2.6) contributions to 40 GeV $\tilde{q}$ fragmentation at $\sqrt{s} = 630$ GeV. We see that the non-perturbative contribution is considerably more important. Therefore, the uncertainties in the treatment of perturbative QCD are not likely to be very important in this case.

There is another difficulty in using non-trivial fragmentation functions of the type (2.3) or (2.5), namely energy-momentum conservation. To the extent that we do not distinguish between the heavy parton and hadron masses, either E or $p_{\parallel}$ or both must be violated whenever $z < 1$. This problem is inherent in the treatment of jet fragmentation at PETRA, but there it is not very important, because m_c or $m_b \ll E_{\text{beam}}$. The resulting violations of energy-momentum conservation are relatively small and easily compensated by adjustments in the routines used for hadronizing the jets¹⁹⁾. By the same token, there would be no difficulty here for $m_{\tilde{g}} \lesssim 10$ GeV. However, the unwary could encounter substantial violations of energy-momentum conservation when using (2.3) and (2.6) for $m_{\tilde{q}, \tilde{g}} \sim 40$ GeV. We evade this difficulty by brutalizing the fragmentation functions in two ways. First we observe that the non-perturbative Ansatz (2.3) has only been checked for z in the region of the peak in the distribution, and that there is no experimental evidence for the precise form of its tail at low z . therefore, we truncate fragmentation to $z > z_0$:

$$z_0 = 1 - 10\epsilon \quad (2.10)$$

which is well below the peak in Fig. 3. Now the violations of energy-momentum conservation are relatively small, and almost suitable for handling by the hadronization routine discussed in Section 4. However, in that section we hadronize

the pair of residual jets left over after the two squark (or gluino) jets have fragmented. To do this we need the four-momentum of the system of two residual jets to be time-like. Therefore, we reject events in which the two choices of $E + p_{\parallel}$ fractions z leave behind a space-like residual jet-jet system. This technical modification does not appear to bias significantly the individual fragmentation functions.

The introduction of non-trivial sparton to shadron fragmentation (Fig. 3) does not seem to dilute significantly the missing p_T signature for $m_{\tilde{q}}$ or $m_{\tilde{g}}$ above ~ 10 GeV, see Table 1. However, it is potentially important if $m_{\tilde{g}} \sim 5$ to 10 GeV. In this case one would expect missing energy events from gg , $q\bar{q} \rightarrow \tilde{g}\tilde{g}$ production, as well as from $\tilde{g}q \rightarrow \tilde{q}$ production. Due to computer time limitations, it is difficult to get reliable numerical results in this case, since only a very small fraction ($O(10^{-4})$) of the light $\tilde{g}\tilde{g}$ pairs produced survives all the UA1 cuts. There can only be a narrow window of acceptable light gluino masses (if any at all) because beam dump experiments²³⁾ suggest $m_{\tilde{g}} \gtrsim 6$ GeV for $m_{\tilde{q}} \sim 100$ GeV. Moreover, conventional models with equal gaugino masses in the limit of grand unification would require $m_{\tilde{Y}} \sim 1/10 m_{\tilde{g}} \lesssim 1$ GeV, while the mass range $m_{\tilde{Y}} \lesssim 5$ GeV is excluded²⁴⁾ by cosmological limits on the relic photino mass density if $m_{\tilde{q}} \sim 100$ GeV. Therefore, we have not studied in detail the missing energy signatures of light $\tilde{g}\tilde{g}$ pair-production.

3. - MISSING TRANSVERSE MOMENTUM SIGNATURES

3.1 Introduction

In the rest of this section we discuss the topological cross-sections, missing transverse momentum signatures, jet energy and rapidity distributions for monojet and dijet events due to the pair-production of $\tilde{q}$ and $\tilde{g}$ with masses in the range 40 to 50 GeV. We will finish with a suggestion for triggering the UA1 detector in a manner which should net more dijet events. First, however, in this subsection we describe the technical assumptions we have made to model the present UA1 and UA2 trigger conditions.

The UA1 trigger transverse momentum resolution is²⁵⁾

$$\sigma = 0.7 \sqrt{E_T} \quad (3.1)$$

where E_T is the total transverse energy. Analogously to our previous work^{8),9)}, we take E_T to be the sum of the transverse energies of the partons produced in sparticle decays, plus the residual jet energies, plus 20 GeV as an allowance for the minimum bias background. Following UA1⁴⁾, we only retain events with missing p_T

$$p_T > \max(15 \text{ GeV}, 4\sigma) \quad (3.2)$$

As an improvement on our previous analysis, we now make some attempt to simulate the lack⁴⁾ of calorimetry close to the vertical azimuthal directions in the UA1 detector. We throw away all events in which the p_T vector points within $\pm 20^\circ$ of the vertical azimuthal directions. We also discard all events in which one of the final state partons of transverse energy above 10 GeV points within $\pm 5^\circ$ of the vertical azimuthal directions. In this way we attempt to mimic the rejection of events in which the missing transverse momentum vector points somewhere near a "hole", or in which there is substantial energy flow into a "hole". As before, we mimic the operation of the UA1 jet algorithm²⁵⁾ as follows. We take the most energetic parton, coalesce with it the most energetic parton within $\Delta\phi^2 + \Delta\eta^2 < 1$, then add on the most energetic parton within $\Delta\phi^2 + \Delta\eta^2 < 1$ of the coalesced jet, and iterate as long as possible. The resulting coalesced jet becomes jet 1. We then repeat the procedure, starting with the most energetic remaining parton and building up jet 2. And so on for more jets until all decay partons and residual jets are included in jets. Because we now have residual jets from fragmentation as well as decay partons, $\tilde{q}\tilde{q}$ pair-production, which formerly could only yield mono- or dijet configurations, can now also give three or more jets. In principle, a crucial distinction between $\tilde{q}\tilde{q}$ or $\tilde{q}\tilde{g}$ or $\tilde{g}\tilde{g}$ production on the one hand, and single $\tilde{g}q \rightarrow \tilde{q}$ production on the other hand, is that the former want to provide multijet configurations with missing energy. Naively, single $\tilde{g}q \rightarrow \tilde{q}$ production followed by $\tilde{q} \rightarrow q\tilde{\gamma}$ decay can only give monojets, but gluon bremsstrahlung can complicate this simple picture. Moreover, there could also be dijet events from light $\tilde{g}\tilde{g}$ production or from $\tilde{q} \rightarrow q\tilde{g}$ decay. Clearly, the existence of multijet events with missing transverse momentum is a crucial one for supersymmetric models. However, most UA1 data has been taken^{4),25)} with a single jet trigger set at a threshold which is nominally 25 GeV in the 1984 run. We have found^{8),9)} that this induces a strong trigger bias in favour of monojet

configurations, and disfavours configurations with two jets. Indeed, monojets dominate for $m_{q,g} \sim 40$ GeV. The dividing line between mono- and dijet events is here set where the second jet has transverse energy 12 GeV. This agrees with the UA1 publication⁴⁾, whereas previously^{8),9)} we had taken this boundary to be 10 GeV. Results are in the next subsection.

The UA2 detector is isotropic in azimuth, but has a polar angle acceptance limited to $|\eta| < 2.0$. We evaluate the supersymmetric cross-sections for events where within $|\eta| < 0.85$, the fiducial region of the UA2 central calorimeter, there is a trigger jet of energy $E_T^1 > 30$ GeV and either no secondary jet with transverse energy above 10 GeV (monojets) or a second jet with $E_T^2 > 10(20)$ GeV at an azimuthal angle separation $\Delta\phi < 130^\circ$, and with net missing $p_T > 20$ GeV (dijets). Results are in the next subsection.

3.2 Topological cross-sections

We calculate $\tilde{q}\tilde{q}$ and $\tilde{g}\tilde{g}$ production with the same theoretical assumptions about α_s , the choice of renormalization scale, and the number of contributing squark flavours as previously. Here we just comment briefly on our choice of five flavours each of left- and right-handed squarks. Fashionable supergravity models^{2),11)} give squark mass spectra which are largely degenerate:

$$m_q^2 = m_0^2 (1 + C_i \xi^2) + O(m_q m_0) + m_q^2 + O(m_{q'}^2) \quad (3.3)$$

where $\xi = m_{1/2}/m_0$ with $m_{1/2,0}$ denoting primordial supersymmetry breaking masses for spin- $\frac{1}{2}$ gauginos and spin-0 squarks respectively. The coefficients C_i are different for left-handed squarks and for right-handed charge $+2/3$ or $-1/3$ squarks²⁶⁾. However, they are universal for different generations, and the differences are in any case small because the C_i are dominated by gluino contributions which are the same for all colour triplets. The $O(m_q)$ and m_q^2 corrections in (3.3) only involve the masses of quarks with the same electric charge as the squarks, and hence are not very important for the five lightest flavours, though they could be important²⁷⁾ for $\tilde{t}$ squarks. The $O(m_q^2)$ corrections²⁸⁾ in (3.3) involve Kobayashi-Maskawa mixing angles and the masses of quarks with different electric charge, e.g., $\propto m_t^2$ in the square of the $\tilde{b}$ squark mass, which is potentially important. However, the coefficient of this term is numerically small if one fixes m_0 and ξ by demanding consistency with the $m_{q,g} \sim 40$ GeV interpretation of the monojet data¹¹⁾. Therefore, the inclusion in our cross-section calculations of the $\tilde{b}$ squarks and the exclusion of $\tilde{t}$ squarks is at least self-consistent. One or more of the observed monojets⁴⁾ may actually contain a charged

lepton from $\tilde{c}$ or $\tilde{b} \rightarrow c$ or b decay. We do not exclude the possibility that $\tilde{t}$ squarks may have interesting signatures²⁹⁾ outside the monojet event sample.

Table 1a lists the topological n -jet cross-sections at $\sqrt{s} = 630$ GeV for missing p_T events due to $\tilde{g}\tilde{g}$ and $\tilde{q}\tilde{q}$ pair production, incorporating all the previously discussed improvements on our work and applying the cuts inspired by UA1 mentioned in the previous subsection. We observe that fragmentation has negligible effect on the cross-sections for higher masses. For lower masses it diminishes the cross-section up to a factor 2, but since cross-sections without fragmentation are very high, the fragmentation effects do not modify any conclusion from our previous work. Table 1b lists the corresponding topological cross-sections obtained applying the UA2-inspired cuts mentioned in the previous subsection. We see that the results are quite comparable with those in Table 1a. The distributions plotted subsequently in this paper are with the UA1 cuts applied to obtain Table 1a: distributions obtained with UA2 cuts are similar.

The lack of perfect calorimetry affects more strongly the dijet than the monojet cross-section, e.g. for squarks with $m_{\tilde{q}} = 40$ GeV we expect a cross-section with perfect calorimetry of 150 pb for monojets and 120 pb for dijets. After the effects of the UA1 holes were taken into account, the cross-sections for monojets decreased by 20% but for dijets by ~40%.

In making these calculations we have assumed 100% $\tilde{q} \rightarrow q\tilde{\gamma}$ or $\tilde{g} \rightarrow q\tilde{q}\tilde{\gamma}$ decay respectively. It may be that $m_{\tilde{W}^\pm, \tilde{Z}^0} < m_{\tilde{q}}$, in which case $\tilde{q} \rightarrow q(\tilde{W}^\pm \text{ or } \tilde{Z}^0)$ decays can also occur. They could reduce the quoted cross-sections while adding other signatures containing final state leptons³⁰⁾. If one looks for a consistent choice of sparticle masses when $m_{\tilde{q}} \sim 40$ GeV one finds¹¹⁾ that $m_{\tilde{g}} \approx m_{\tilde{q}}$. If $m_{\tilde{g}}$ is only slightly less than $m_{\tilde{q}}$, then $\tilde{q} \rightarrow q\tilde{\gamma}$ decay will still dominate. If $m_{\tilde{g}}$ is even slightly greater than $m_{\tilde{q}}$, then $\tilde{g} \rightarrow \tilde{q}\tilde{q}$ decay followed by $\tilde{q} \rightarrow q\tilde{\gamma}$ will dominate. If $m_{\tilde{g}} - m_{\tilde{q}} \lesssim \text{few GeV}$, the $\tilde{q}$ in $\tilde{g} \rightarrow \tilde{q}\tilde{q}$ decay will be soft, $\tilde{g}\tilde{g}$ pair-production will give mono-, di- and trijet events in similar ratios to $\tilde{q}\tilde{q}$ pair-production, and associated $\tilde{g}\tilde{q}$ production would also be similar. These extra production mechanisms could give an overall increase in the different topological cross-sections in the case where $m_{\tilde{q}}$ and $m_{\tilde{g}} \sim 40$ GeV.

3.3 Jet energy, missing p_T and rapidity distributions

We now turn to differential distributions within the different event categories which may serve to discriminate between different interpretations of the monojet events. Shown in Fig. 4a are the distributions in monojet energy for $\tilde{q}\tilde{q}$ pair-production with $m_{\tilde{q}} = 40, 50$ GeV. We see that the events are peaked against the jet trigger cut if $m_{\tilde{q}} = 40$ GeV, but peak above 40 GeV if $m_{\tilde{q}} = 50$ GeV. Shown in Fig. 4b are the corresponding distributions in monojet energy for $\tilde{g}\tilde{g}$ pair-production with $m_{\tilde{g}} = 40, 50$ GeV. Here the monojet transverse energies are peaked against the trigger cut for both 40 and 50 GeV gluino masses, and in both cases the distributions are peaked at smaller p_T than was the case for 40 GeV squarks. The monojets observed⁴⁾ in 1983 include some which are peaked against the cuts and are consistent with background estimates, one event interpretable as $W \rightarrow \tau\nu$ decay, and 5 "interesting" events with $p_T > 40$ GeV.

Shown in Fig. 5a are the distributions in p_T of the more energetic jets in dijet events due to $m_{\tilde{q}} = 40, 50$ GeV, while Fig. 5b shows the corresponding p_T distributions for the primary jets in dijet events due to $\tilde{g}\tilde{g}$ production. Although the p_T distributions are harder for heavier sparticle masses, there are no big differences between 40 and 50 GeV, or between squarks and gluinos. The few dijet events seen in 1983 are close to the cuts and compatible with background. As seen in Table 1, qualitatively we expect a larger abundance of dijets relative to monojets, especially for larger sparticle masses. However, one must be careful not to overstate this potential difficulty, since the dijet events pile up against the cuts and are especially sensitive to the way in which they are applied. In particular, we have been calculating jet energies at the parton level whereas they are measured experimentally at the hadron level. If a secondary jet energy is mismeasured, an event can be shifted from the dijet category to the monojet category, or vice versa. And relative measurement errors are probably larger for the less energetic jets.

The most clear difference between squarks and gluinos is probably shown by the p_T distribution of the second energetic jet of the missing p_T events, see Figs. 6a and 6b. In the gluino case, this distribution is peaked around 10(15) GeV for $m_{\tilde{g}} = 40(50)$ GeV. However, in the squark case this distribution has a clear peak towards 0 GeV. The difference between the two distributions seems to be large enough to survive the smearing effects of the experimental measurements.

Figure 7a shows the distribution in missing p_T^2 for monojets from $\tilde{q}\tilde{q}$. Clearly they mirror to a large extent the primary jet p_T distributions of Fig. 4a. Squarks of 40 GeV peak against the cuts, while squarks of 50 GeV stand away from them, perhaps as desired by the data. The missing p_T^2 distributions for monojets from gluinos shown in Fig. 7b are peaked much closer to the cuts, perhaps at lower values of missing p_T^2 than wanted by the data. The missing p_T^2 distributions for dijets shown in Fig. 8a,b are considerably softer than for the monojets shown in Fig. 7. Particularly soft are the missing p_T^2 distributions for gluinos shown in Fig. 8b. Since they are closest to the p_T missing cut (3.2), these events are the most difficult to distinguish from background sources.

Now we turn to the rapidity distributions shown in Fig. 9a,b. We see from Fig. 9a that the primary jets in squark monojet events with $m_{\tilde{q}} = 40$ GeV have a pseudorapidity distribution which is sharply peaked around $\eta = 0$:

$$\langle \eta_1^2 \rangle^{\frac{1}{2}} = 0.65 \quad (3.3)$$

while the secondary minijets with p_T between 5 and 12 GeV have a somewhat broader distribution:

$$\langle \eta_2^2 \rangle^{\frac{1}{2}} = 1.2 \quad (3.4)$$

The primary and secondary jets in gluino monojet events have slightly more similar rapidity distributions:

$$\langle \eta_1^2 \rangle^{\frac{1}{2}} = 0.64 \quad (3.5)$$

$$\langle \eta_2^2 \rangle^{\frac{1}{2}} = 1.0 \quad (3.6)$$

As can be seen in Fig. 9b, the corresponding rapidity distributions for primary and secondary jets in $\tilde{q}$ dijet events are essentially indistinguishable:

$$\langle \eta_1^2 \rangle^{\frac{1}{2}} = 0.67, \quad \langle \eta_2^2 \rangle^{\frac{1}{2}} = 0.81 \quad (3.7)$$

and the same is true in $\tilde{g}$ dijet events:

$$\langle \eta_1^2 \rangle^{\frac{1}{2}} = 0.65, \quad \langle \eta_2^2 \rangle^{\frac{1}{2}} = 0.73 \quad (3.8)$$

We also obtained similar results for $m_{\tilde{q}}, m_{\tilde{g}} = 50$ GeV respectively. The few 1983 monojet events seen⁴⁾ have a fairly narrow rapidity distribution with $\langle \eta^2 \rangle^{\frac{1}{2}}$ comparable to (3.3) or (3.5). The secondary minijets in the 1983 monojet events have a wider rapidity distribution, perhaps favouring squarks (3.4) over gluinos (3.6). This might be an important trend to pursue in the 1984 data. Note that higher order QCD would give secondary jets with a much broader rapidity distributions, peaked forward and backward instead of centrally.

Finally, shown in Fig. 10 are the distributions of the azimuthal angle difference between the first and the second jet in dijet events for the gluino and squark cases. We see that, especially in the squark case, the two jet directions and missing p_T vector tend to form Mercedes star configurations. There is less probability for the back-to back configuration. Note that, in contrast, the dijet events from light $\tilde{g}q \rightarrow \tilde{q} + q\tilde{g}$ or from light $\tilde{g}\tilde{g}$ pair-production are expected to be predominantly coplanar giving back-to-back configurations peaked at $\Delta\phi = 180^\circ$.

To summarize, we have seen that the study of distributions in jet E_T and missing p_T , rapidity and relative azimuthal angle should distinguish between different supersymmetric hypotheses and are probably characteristic enough to prove or disprove supersymmetric interpretations of missing p_T events. Of particular importance is the study of the secondary jets down to the lowest possible p_T .

3.4 Triggering for dijets

The basic reason why most of the missing energy events were expected to be monojets if $m_{\tilde{q}}$ or $m_{\tilde{g}} \lesssim 40$ GeV was the single jet trigger of 0(25) GeV adopted by UA1 in 1983. This biases one towards configurations where one $\tilde{\gamma}$ is soft while the other must be hard to give large missing p_T . This trigger also favours events where partons are clustered into more energetic jets. To overcome this trigger bias one should look at events with two moderately large deposits of energy, rather than just one large deposit of energy. We have studied this problem quantitatively in the case of squarks, as seen in Table 2 and Fig. 11. The first row shows cross-sections for monojets and dijets from $\tilde{q}\tilde{q}$ production in the absence of any jet trigger. The second row shows the size of cross-sections remaining when a single jet trigger is applied: most of the dijet events are lost. The third row shows what cross-sections are available if one adds a two-jet trigger: $E_{T_1}, E_{T_2} > 15$ GeV. We see that most of the dijet cross-section is recovered. The same point is made in Fig. 11, where the numbers represent the

observable cross-sections for different ranges of jet transverse jet energies, in units of the total monojet cross-section. We see that most of the monojets have no clearly observable minijets ($p_{T_2} < 5$ GeV) while the number of dijet events ($p_{T_2} > 12$ GeV) is about half the total number of monojets. The number of extra dijet events accessible with the two-cluster $E_{T_1}, E_{T_2} > 15$ GeV trigger increases the dijet cross-section above the monojet cross-section. Taking data with such a two-cluster trigger could provide useful extra information about missing energy events.

4. - HADRONIZATION

4.1 Prescriptions for hadronization

The third step in our improvement programme is to incorporate hadronization of the shadron decay products and the residual jets. In principle, all of the calculations presented in the previous section should be carried out with hadronized jets instead of the parton jets we have used. In practice, to do this properly would require putting the hadrons we generate through the UA1 detector simulation. Accordingly, here we just describe the prescriptions we have adopted for hadronization, and describe the results they give for the hadronization of monojets. In particular, we calculate the charged multiplicities and charged particle invariant masses of monojets in the $gg, q\bar{q} \rightarrow \tilde{q}\tilde{q}$ or $\tilde{g}\tilde{g}$ and $\tilde{g}q \rightarrow \tilde{q}$ scenarios. This enables us to address quantitatively the problem of the apparent thinness of the 1983 monojets.

The gluino presumably forms a gluino hadron ($\tilde{g}g$) or ($\tilde{g}q\bar{q}$), and then decays: $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$ becomes $(\tilde{g}g) \rightarrow (\tilde{g}q\bar{q}) + \tilde{\gamma}$ or $(\tilde{g}q\bar{q}) \rightarrow (q\bar{q}q\bar{q}) + \tilde{\gamma}$. The hadronic final state in the $\tilde{g}$ rest frame presumably consists of two jets with axes parallel to the momenta of the q and $\bar{q}$ decay products, and a mass given approximately by the zeroth-order $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$ decay matrix element. We have therefore hadronized the gluino hadron analogously to a $q\bar{q}$ two-jet system in e^+e^- annihilation. Indeed, to do this we have used the two-jet independent quark fragmentation routine developed by the TASSO Collaboration¹⁹⁾ to describe e^+e^- annihilation into hadrons at PETRA.

The squark presumably forms a squark hadron, ($\tilde{q}\bar{q}$) or ($\tilde{q}qq$), and then decays: $\tilde{q} \rightarrow q\tilde{\gamma}$ becomes $(\tilde{q}\bar{q})$ or $(\tilde{q}qq) \rightarrow (q\bar{q})$ or $(qqq) + \tilde{\gamma}$. The hadronic final state in the $\tilde{q}$ rest frame presumably consists of one jet oriented along the momentum axis of the q decay product. The invariant mass of this decay jet determined by the

zeroth-order decay matrix element would just be m_q . We try two Ansätze for attributing a realistic mass to the squark decay jet. One is just to hadronize it in the style of the independent jet fragmentation model, generating at each stage in the cascade a new antiquark with a reduced fraction Z of the $E + P_{\parallel}$ of the jet, and $p_T = O(300)\text{MeV}$. The last $\bar{q}$ in the cascade is taken as the spectator in the $(\tilde{q}\bar{q})$ hadron. This procedure¹⁹⁾ does not exactly conserve four-momentum, nor assign a unique mass to the squark hadron, but it is surely adequate for our needs. Another Ansatz, which we think is valuable for giving a perspective on the important question of the monojet mass, although we do not prefer it over the previous non-perturbative Ansatz, is to use perturbative QCD to estimate the invariant mass of the squark decay jet. Gluon bremsstrahlung in $\tilde{q}$ decay has been calculated^{11),12)}, and one knows the resulting distribution in $q\bar{q}$ invariant masses:

$$\frac{1}{\Gamma_{\text{tot}}} \frac{d\Gamma}{dy} = \frac{\alpha_s(y m_{\tilde{q}}^2)}{6\pi} \left(6 + y - 4 \ln y - \frac{7}{y} - \frac{4 \ln y}{y} \right) \quad (4.1)$$

where $y \equiv m_{q\bar{g}}^2 / m_{\tilde{q}}^2$. We generate two-jet final states with invariant masses distributed according to Eq. (4.1) using $\alpha_s = \alpha_s(m_{q\bar{g}})$, and cutting off the distribution below y_0 :

$$\frac{1}{\Gamma_{\text{tot}}} \int_{y_0}^1 dy \frac{d\Gamma}{dy} = 1 \quad (4.2)$$

In this procedure the four momentum of the resulting two-jet system automatically gives the input squark mass when combined with the recoiling photino.

To fragment the residual jets left over after sparton to shadron fragmentation, we just fragment the two residual jets as a linked two-jet system. This procedure does not reflect correctly the colour structure of the event. However, we do not need the minimum bias background when we are discussing monojet properties, so we are not interested in the hadronization of the residual jets in the p and $\bar{p}$ directions. The procedure we adopt is the simplest way of ignoring them: any error we make is presumably restricted to low p_T hadrons of little interest. Another concern is energy-momentum conservation in the fragmentation of heavy partons, a problem we discussed in Section 2. There we described how we have butchered the original Ansatz (2.3) and (2.5) for sparton fragmentation. Because of the liberties we have taken with the residual jets, we will be on the watch for monojet properties which are insensitive to the inclusion of the residual jets.

4.2 Results

In Table 3 we have listed the average charged multiplicities and invariant masses for monojets due to squarks with mass 40 GeV. The columns and rows correspond to different hadronization prescriptions and different p_T thresholds respectively. The first column is for our reference prescription: conventional non-perturbative fragmentation of the $\tilde{q} \rightarrow q$ jet, and the residual jets included. In the second column the residual jets have been left out, while they are back in the third column, but now we adopt the perturbative $\tilde{q} \rightarrow (qg)$ jet hadronization prescription. In all three cases we construct hadronized jets by determining the coalesced jet axes as describe previously at the parton level, and then gathering in all the hadrons emerging within $\Delta\phi^2 + \Delta\eta^2 < 1$ of this coalesced parton jet axis, regardless of which hadronized parton they were originally emitted from. We find that jet transverse energies are not greatly altered by this hadronization procedure.

We see that above $p_T = 0.5$ GeV the first two columns in Table 3 are very similar, encouraging us to believe that our results are independent of our butchery of the residual jets. The (qg) prescription in the third column gives jets with slightly fewer charged particles with smaller invariant masses at lower p_T , but higher mean multiplicites and masses when a higher p_T cut-off is used. We guess that the model-dependence may be minimized if one takes a p_T cut-off around 1 GeV.

In Fig. 12 we show the mean monojets charged multiplicities for $q\bar{q}$ pair production, $g\bar{g}$ pair production and $gq \rightarrow \tilde{q}$ fusion as a function of the p_T cut-off, using the first prescription of Table 3 for $\tilde{q}$, including residual jets in the $\tilde{g}$ monojets, and using naked $\tilde{q} \rightarrow q$ one-jet hadronization without residual jets or overlap for the $gq \rightarrow \tilde{q}$ fusion model. Fig. 13 shows the corresponding charged particle invariant masses. The $g\bar{g}$ monojets are fattest, followed by the $q\bar{q}$ monojets, with the single $\tilde{q}$ jets thinnest. The $g\bar{g}$ monojets are broad before coalescence, while both $g\bar{g}$ and $q\bar{q}$ monojets are broadened by coalescence.

The 1983 monojets⁴⁾ would appear to be thinner than any of these models would prefer, though one should be cautious in making this deduction. The total number of events is not large, and two of the monojets have two or more unresolved collimated charged particles which may have high energies. However, this is clearly a point to watch. It is interesting to note that our hadronization prescriptions do indicate a non-negligible probability for shadron monojets to contain just one charged particle with p_T above 1 GeV. This is relevant for the

"photon" event H of Ref. 4), whose final state also contains a charged particle of $p_T = 1.2$ GeV emitted within $\Delta\phi^2 + \Delta\eta^2 < 1$ of the electromagnetic cluster. We think it possible that this event was just an example of unusual monojet hadronization.

5. - PROSPECTS

Although further background studies are clearly necessary and desirable, the 1983 monojet events⁴⁾ of UA1 could be signatures of supersymmetry^{7)-12),17),18)}.

We have argued^{9),11)} in favour of an interpretation with $m_{\tilde{q}} \sim 40$ GeV, but other interpretations cannot be excluded using this first event sample. We have tried not to introduce too much theoretical bias into the analysis of this paper, and our results can be used to infer conservative lower bounds on either squark or gluino masses. It turns out¹¹⁾ that in many supergravity models with $m_{\tilde{q}} \sim 40$ to 50 GeV, $m_{\tilde{g}}$ is not much greater than $m_{\tilde{q}}$. In this case $\tilde{g}$ exchange increases the cross-section for $q\bar{q}$ production, and gq and $g\bar{q}$ production add significantly to the missing p_T signature. We find that qualitative features of the monojet and dijet signatures discussed here remain unchanged in such a scenario, and that in this particular theoretical context a more stringent lower bound on $m_{\tilde{q}}$ can be derived from existing monojet data. However, any such bound would be quite model-dependent, and we have preferred to keep this paper relatively model-independent and concentrate on introducing technical innovations into the phenomenological analysis of missing energy events. In this paper we have incorporated three improvements over our previous calculations, namely initial p_T smearing, supersymmetric parton to hadron fragmentation, and jet hadronization. We have shown that our previous results were not strongly affected by the first two of these modifications, have stressed that dijet events with missing transverse momentum may be the key to establishing the supersymmetric interpretation of missing p_T events, and have suggested how to trigger to get more of them. We have incorporated hadronization in our calculations, and estimated the thickness of monojets. The 1983 monojet events seem to be thinner than we would expect, though the significance of this tendency is not clear.

We hope that the issues discussed here can be resolved by the 1984 data. A crucial next step is to incorporate hadronization into the calculations of missing p_T event signatures. To do this effectively requires a detailed detector simulation. The program used to obtain the results presented here can be used to generate events which can be put directly into such a simulation.

ACKNOWLEDGEMENTS

We thank many members of the UA1 and UA2 Collaborations, in particular J. Rohlf and P. Jenni, for many interesting and valuable conversations.

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Table 1a

Topological cross-sections for missing p_T events in pb evaluated with cuts inspired by UA1 (see text). The values in brackets were calculated without fragmentation effects.

# Jets m [GeV]		1	2	3
10	$\tilde{g}$	450 ± 140 (800 ± 100)	150 ± 30 (290 ± 80)	0
	$\tilde{q}$	—	—	—
20	$\tilde{g}$	380 ± 70 (560 ± 70)	190 ± 40	80 ± 80
	$\tilde{q}$	—	—	—
30	$\tilde{g}$	220 ± 20 (235 ± 16)	100 ± 10 (107 ± 7)	4 ± 1 (5 ± 1)
	$\tilde{q}$	290 ± 10	85 ± 2	0.2 ± 0.02
40	$\tilde{g}$	93 ± 6 (95 ± 6)	86 ± 5 (82 ± 4)	11 ± 1.5 (9 ± 1.5)
	$\tilde{q}$	120 ± 8	75 ± 5	1.5 ± 0.2
50	$\tilde{g}$	23 ± 1.5 (27 ± 1.7)	48 ± 2 (43 ± 1)	13 ± 1 14 ± 1
	$\tilde{q}$	38 ± 1	46 ± 1	0.1 ± 0.01
60	$\tilde{g}$	4.5 ± 0.5	15 ± 1	8.5 ± 0.5
	$\tilde{q}$	11 ± 0.2	22 ± 0.3	0.0

Table 1b

Topological cross-sections for missing p_T events in pb evaluated with cuts inspired by UA2 (see text).

<div style="display: inline-block; transform: rotate(-45deg);"> # jets m (GeV) </div>	1	2		
		$E_T^2 > 10 \text{ GeV} : E_T^2 > 20 \text{ GeV}$		
30	$\tilde{g}$	154	50	22
	$\tilde{q}$	215	35	11
40	$\tilde{g}$	80	42	11
	$\tilde{q}$	112	49	12
50	$\tilde{g}$	19	23	13
	$\tilde{q}$	33	27	15

Table 2

Effect of E_T jet trigger cut

$\sqrt{s} = 630 \text{ GeV}$, $m_{\tilde{q}} = 40 \text{ GeV}$

Jet trigger	σ (1 jet)	σ (2 jet)	σ (3 jet)
none	111 ± 23	165 ± 19	9 ± 2
$E_{T_1} > 25\text{GeV}$	121 ± 16	65 ± 16	7 ± 2
$E_{T_1}, E_{T_2} > 15\text{GeV}$	121 ± 16	131 ± 15	9 ± 2

Table 3
Comparison of different hadronization schemes

$\min$ p_T	Standard	no residual jets	qg invariant mass frag'n.
$\langle n_{ch} \rangle$ 0.5	7.5	7.3	5.6
$\langle m_{ch} \rangle$	6.7	6.5	6.3
$\langle n_{ch} \rangle$ 1	5.7	5.6	4.9
$\langle m_{ch} \rangle$	5.4	5.5	5.9
$\langle n_{ch} \rangle$ 2	3.7	3.7	3.8
$\langle m_{ch} \rangle$	3.6	3.9	4.9
$\langle n_{ch} \rangle$ 3	2.7	2.8	3
$\langle m_{ch} \rangle$	2.5	2.8	3.8
$\langle n_{ch} \rangle$ 5	1.8	1.8	2
$\langle m_{ch} \rangle$	1.4	1.6	2.2

FIGURE CAPTIONS

- Fig. 1 : Initial transverse momentum smearing generated by Monte Carlo using the phenomenological parametrization(2.1). The QCD parametrizations (2.2) give similar results within errors.
- Fig. 2 : Effect of initial transverse momentum smearing on the missing p_T spectrum for monojets.
- Fig. 3 : Fragmentation functions for 40 GeV squarks pair-produced at $\sqrt{s} = 630$ GeV.
- Fig. 4 : Transverse momentum distributions for monojets due to (a) $\tilde{q}\tilde{q}$ and (b) $\tilde{g}\tilde{g}$ production. This and subsequent distributions are for $\sqrt{s} = 630$ GeV.
- Fig. 5 : Transverse momentum distributions for the most energetic jets in dijet events due to (a) $\tilde{q}\tilde{q}$ and (b) $\tilde{g}\tilde{g}$ production.
- Fig. 6 : Transverse momentum distributions for the secondary jets in "monojet" ($p_{T_2}^{\text{Jet}} < 12$ GeV) and "dijet" ($p_{T_2}^{\text{Jet}} > 12$ GeV) events for (a) $m_{\tilde{q}}, m_{\tilde{g}} = 40$ GeV and (b) $m_{\tilde{q}}, m_{\tilde{g}} = 50$ GeV.
- Fig. 7 : Missing p_T squared distributions for monojets due to (a) $\tilde{q}\tilde{q}$ and (b) $\tilde{g}\tilde{g}$ production.
- Fig. 8 : Missing p_T squared distributions for dijets due to (a) $\tilde{q}\tilde{q}$ and (b) $\tilde{g}\tilde{g}$ production.
- Fig. 9 : Jet rapidity distributions for (a) monojets and accompanying mini-jets with $12 \text{ GeV} > p_{T_2}^{\text{Jet}} > 5$ GeV and (b) dijets due to $\tilde{q}\tilde{q}$ production.
- Fig. 10 : Azimuthal angles between jets in dijet events due to $\tilde{q}\tilde{q}$ and $\tilde{g}\tilde{g}$ production.

- Fig. 11 : Relative event rates in different domains of the (p_{T_1}, p_{T_2}) plane, normalized to 1 for the total number of monojets with $p_{T_1} > 25$ GeV, $p_{T_2} < 12$ GeV.
- Fig. 12 : Monojet mean multiplicities for charged particles above $p_T^{\min}$, using the "standard" hadronization prescription in Table 3.
- Fig. 13 : Monojet mean invariant masses for charged particles above $p_T^{\min}$, using the "standard" hadronization prescription in Table 3.

Fig. 1

Fig. 2

Fig. 3

Fig. 4a

Fig. 4b

Fig. 5a

Fig. 5b

Fig. 6a

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Fig. 6b

Fig. 7a

Fig. 7b

Fig. 8a

Fig. 8b

Fig. 9a

Fig. 9b

Fig. 10

Fig. 11

Fig. 12

Fig. 13