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ABSTRACT

Implications of large missing transverse momentum ($\not p_T$) events at the CERN $p\bar{p}$ collider to supersymmetry theories are surveyed. Three popular models where photinos carry missing momentum, A) squark-pair scenario, B) gluino-pair scenario, and C) light-gluino scenario are critically reviewed. Both scenarios A and B predict non-back-to-back dijet plus $\not p_T$ events, whereas scenario C predicts monojet dominance at higher $\not p_T$ ($\gtrsim$ 40 GeV) and back-to-back dijet events at moderate $\not p_T$. Controversial problems in the light gluino scenario are discussed in detail. A model with neutral Higgsinos from the Z boson decay as the carrier of $\not p_T$ is ruled out by e^+e^- experiments at PETRA. Light Goldstinos can also be a source of $\not p_T$ in models where supersymmetry breaks down in the TeV energy region. If a Goldstino is the only light supersymmetric particle, then the large $\not p_T$ event rate at $\sqrt{s}$ = 630 GeV is almost four times larger than that at $\sqrt{s}$ = 540 GeV. If both a Goldstino and a gluino are light (< 1 GeV), then we expect clean monojets with a quite hard $\not p_T$ spectrum.

INTRODUCTION

In most supersymmetry theories[1], there exists an unbroken discrete symmetry called R-parity[2] where all the standard model particles (leptons, quarks, gauge bosons, and Higgs bosons) are assigned an even R-parity while all their supersymmetric partners (sleptons, squarks, gauginos, and Higgsinos) are assigned an odd R-parity. This unbroken symmetry tells us that supersymmetric (R-odd) particles should be pair produced in the collisions of ordinary (R-even) particles and that the lightest R-odd particle should be absolutely stable[3]. Pair produced supersymmetric particles in high energy experiments are both expected to decay into this new stable particle. If it is electrically neutral and uncoloured (e.g. photino, neutral Higgsino, sneutrino or Goldstino) then it interacts feebly with matter and escapes detection like neutrino, giving rise to missing-p_T events in calorimetric experiments. Hence large $\not p_T$ events had been expected and studied rather extensively[2,4-6] as a possible signal of supersymmetry.

It is therefore not surprising that the observation[7] last year of unusually large $\not p_T$ events associated mainly by a single jet at the CERN $p\bar{p}$ collider triggered a number of interpretations[8-13] based on supersymmetry. Early attempts to test various supersymmetry scenarios against the data appeared in Refs. 12,14 and 15.

So far the anomaly observed last year has not been confirmed in the new higher luminosity run of the CERN collider. Preliminary reports[16] indicate the observation of monojet events with a similar

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topology to the ones observed last year and also that of dijet plus large $\not{p}_T$ events where most of the dijets have back-to-back configuration in the transverse momentum plane, while acceptance and background estimates have not been presented. It is clearly most important to determine that the observed signals can in no way come from the standard model sources[17,18].

It would be truly exciting if we are observing the first experimental signal of supersymmetry. Even the slightest experimental information on the masses of superpartners will drastically improve our understanding of the supersymmetry breaking mechanism and will give us a clue to determine if the currently most attractive scenario, the N = 1 supergravity grand unification[19] and its eventual embedding into the superstring theory[20] (the first serious candidate of the theory of quantum gravity) is in a correct direction.

In this talk I will review the implications for supersymmetry of the large $\not{p}_T$ events observed at hadron colliders. First, I discuss the most popular scenario where photinos carry the missing momentum; these are A) squark-pair production ($m_{\tilde{q}} \simeq$ 40-50 GeV) followed by the decay $\tilde{q} \to q\tilde{\gamma}$, B) gluino-pair production ($m_{\tilde{g}} \simeq$ 40-50 GeV) with the $\tilde{g} \to q\bar{q}\tilde{\gamma}$ decays, and C) associate production of a light $\tilde{g}$ ($m_{\tilde{g}} \simeq$ 3-5 GeV) and a heavy $\tilde{q}$ ($m_{\tilde{q}} \simeq$ 100 GeV) followed by the decay $\tilde{q} \to q\tilde{\gamma}$. Only this last scenario explains the observed monojet dominance over multijet plus $\not{p}_T$ events and I will discuss in detail the dynamical problems associated with the possibility of having light gluinos. On the second part, I examine the possibility that neutral Higgsinos carry the missing momentum. Neutral Higgsinos are produced via Z boson decay and we find that such a possibility can be ruled out by present e^+e^- annihilation experiments at PETRA/PEP. Finally, I discuss briefly the possibility that light ($\simeq$ massless) Goldstinos carry the missing momentum.

Before we begin, it is worth noting that a candidate for the $\not{p}_T$ carrier at hadron colliders does not necessarily have to be the lightest R-odd particle. Even if it is unstable, an electrically neutral and uncoloured particle can escape detection in calorimetric experiments either if it lives long enough not to decay in the detectors or if it decays mainly into unobservable modes. For example, an unstable photino can be a candidate if it decays into a photon and say, a Goldstino[21], an axiono[22] or a Higgsino[23] very slowly or if it decays mainly into an invisible mode such as a neutrino and a sneutrino[24]. Such possibility should be kept in mind when we examine further implications of the models.

PHOTINOS AS THE CARRIERS OF MISSING MOMENTUM

In order to have a significant cross section at the CERN $p\bar{p}$ collider, a superparticle-pair should be produced either via $O(\alpha_s^2)$ subprocesses $qq \to \tilde{q}\tilde{q}$, $(q\bar{q},gg) \to (\tilde{q}\bar{\tilde{q}},\tilde{g}\tilde{g})$, and $qg \to \tilde{q}\tilde{g}$, or via $O(\alpha)$ subprocesses, i.e., via W and Z boson decays. All other parton subprocesses, e.g. $O(\alpha_s\alpha)$ processes, such as $q\bar{q} \to (\tilde{g}\tilde{\gamma},\tilde{g}\tilde{\omega},\tilde{g}\tilde{z})$ or $qg \to (\tilde{q}\tilde{\gamma},\tilde{q}\tilde{\omega},\tilde{q}\tilde{z})$, are found to give a negligible contribution to the present large $\not{p}_T$ signal. When squarks and/or gluinos are produced via the $O(\alpha_s^2)$ subprocesses they decay, depending on their relative masses, via either $\tilde{q} \to q\tilde{g}$ ($m_{\tilde{q}} > m_{\tilde{g}}$) or $\tilde{g} \to q\bar{\tilde{q}},\bar{q}\tilde{q}$ ($m_{\tilde{g}} > m_{\tilde{q}}$), followed by either

$\tilde{g} \to q\bar{q}\tilde{\gamma}$ or $\tilde{q} \to q\tilde{\gamma}$ with the photino giving missing p_T. Thus depending on the relative masses of the squark and the gluino, we may have three very different scenarios.

SCENARIO A: $m_{\tilde{q}} \simeq 40$ GeV $< m_{\tilde{g}}$ [12,13]

If the squark is lighter than the gluino, then the produced squark-pair each decay into a quark and a photino giving rise to a $q\bar{q}\tilde{\gamma}\tilde{\gamma}$ final state. Most of the time, the squark pair is produced near the threshold and each squark has small p_T. For squarks of mass 40 GeV, the photinos each have typically $p_T \simeq 20$ GeV. In order for the two photino momenta to add up vectorally to give large $\not{p}_T$ of typically 30-40 GeV, the two momentum vectors should have rather a small opening angle in the transverse momentum plane. This would then lead to a configuration where the two quark momenta are also aligned and the event has a high probability to be identified as a monojet event in the UA1 jet defining algorithm[7,25]. If squark masses are higher, then no particular final state configuration is required to pass the large $\not{p}_T$ experimental cut and the event would typically be registered as a dijet event.

This is clearly seen in Fig. 1 where the sum (a) and the ratio (b) of monojet and dijet cross sections with $\not{p}_T > 40$ GeV at $\sqrt{s} = 630$ GeV are shown as a function of the squark and gluino masses. Here we assumed that the squark masses are degenerate for 5 flavours and 2 chiralities[12,13] in accordance with the low energy constraints from the absence of large flavour changing neutral currents[26] and of anomalous parity violation in nuclei[27]. When a gluino mass is not much larger than a squark mass, the contribution from $\tilde{g}\tilde{q}$ and $\tilde{g}\tilde{g}$

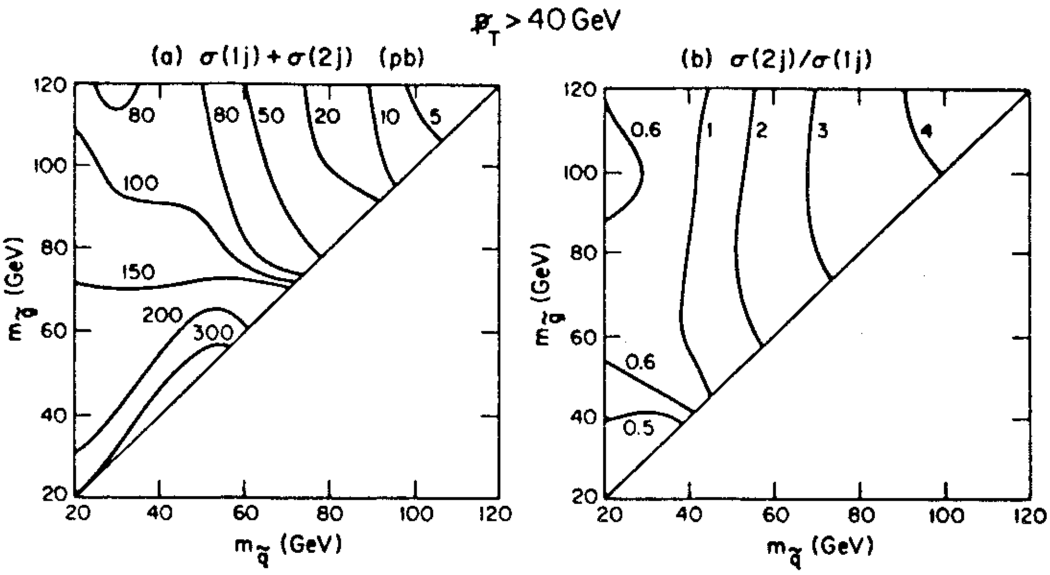

Fig. 1 Predictions for the sum (a) and the ratio (b) of monojet and dijet cross sections with $\not{p}_T > 40$ GeV at $\sqrt{s} = 630$ GeV, taken from Ref. 28.

production processes followed by the cascade decay $\tilde{g} \to q\bar{\tilde{q}}$ or $\bar{q}\tilde{q}$; $\tilde{q} \to q\tilde{\gamma}$ is found to be significant. If we request monojet dominance, then $m_{\tilde{q}} \lesssim 40$ GeV follows from Fig. 1b. Further requirement of $\sigma(\not{p}_T > 40 \text{ GeV}) < 100$ pb would rule out almost all the region in the $m_{\tilde{g}} > m_{\tilde{q}}$ plane except for a tiny window with $m_{\tilde{q}} = 20\text{-}40$ GeV and $m_{\tilde{g}} > 100$ GeV. In the N = 1 supergravity grand unified models[19], there is a severe constraint on the average mass squared of squarks and sleptons

$$\overline{m_{\tilde{q}}^2} - 0.77\, m_{\tilde{g}}^2 = \overline{m_{\tilde{\ell}}^2}$$

which follows[29,30] from the constraint $m_{\tilde{q}} = m_{\tilde{\ell}}$ at the grand unification (GUT) scale via the renormalization group running of the mass parameters[31,30] assuming three generations. This together with the bound $m_{\tilde{\ell}}^2 > (20 \text{ GeV})^2$ from e^+e^- annihilation experiments[32] forbids a gluino mass much larger than a squark mass. The aforementioned tiny window is hence closed for the N = 1 supergravity GUT models. With four generations of quark-lepton flavours, the numerical coefficient 0.77 in the above equation is replaced[30] by 1.65 and the whole region of scenario A disappears.

SCENARIO B: $m_{\tilde{g}} \simeq 40 \text{ GeV} < m_{\tilde{q}}$ [6,8]

If a gluino is lighter than a squark, then it would decay into a photino and a quark-pair. Even though the final state from $\tilde{g}\tilde{g}$ production now consists of four quarks and two photinos, sufficiently light gluinos ($\lesssim$ 40 GeV) are again found to give mainly monojet plus $\not{p}_T$ events due to the same trigger bias as explained for the scenario A. Fig. 2 shows the one- and multi-jet cross sections with $\not{p}_T > 4\sigma$ at $\sqrt{s}$ = 540 GeV, σ being the $\not{p}_T$ resolution[7].

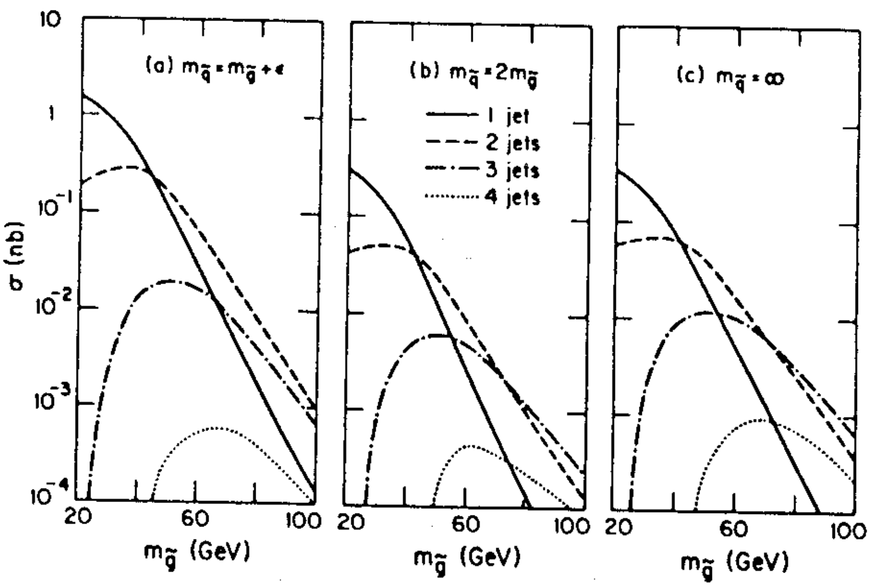

Fig. 2 One- and multi-jet cross sections with $\not{p}_T > 4\sigma$ versus gluino mass at $\sqrt{s}$ = 540 GeV taken from Ref. 14.

Although we can again expect monojet dominance with an appropriate cross section at $m_{\tilde{g}} \simeq 40$ GeV, it is very unlikely that this scenario can explain the narrowness of all the observed monojet events with $\not{p}_T > 35$ GeV[7]. Naively, the monojet in this scenario should be broad since it is a combination of up to four quark jets[12,14]. A detailed study of the expected jet structure in various scenarios was performed by Ellis and Kowalski[33], who used a standard jet fragmentation model in e^+e^- jet studies. Fig. 3 shows the expected mean charged multiplicity and the charged particle invariant mass as a function of the minimum hadron p_T in a jet[33]. The solid, dashed, and dash-dotted lines are the predictions of scenarios A, B and C respectively. Unless we start observing multi-prong monojets as well as dijets in association with large $\not{p}_T$, this scenario should also be discarded.

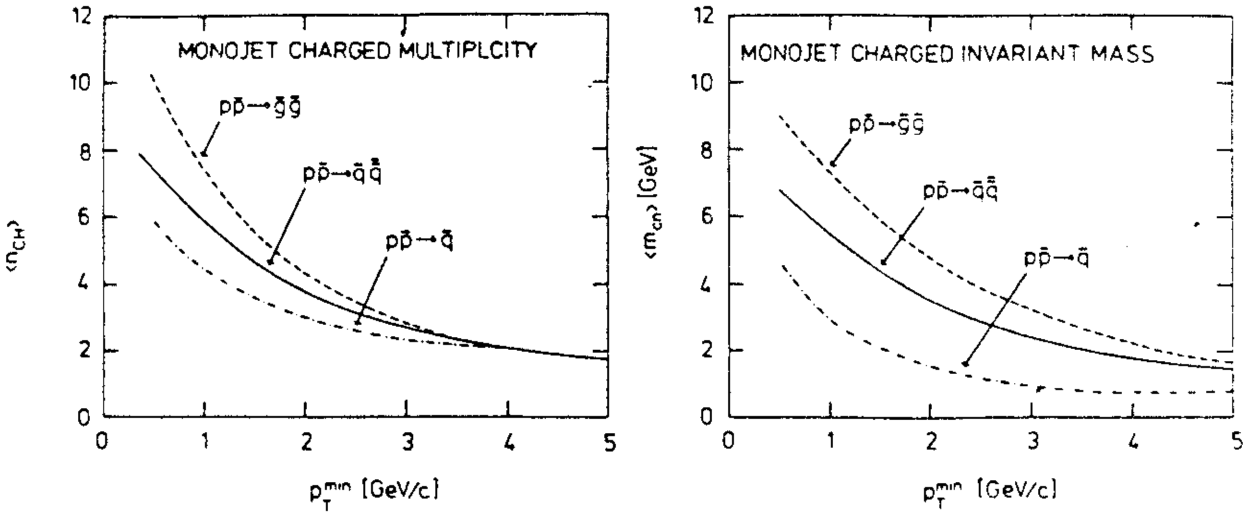

Fig. 3 Mean charged multiplicity and mean charged hadron invariant mass of monojet events for various supersymmetry scenarios shown as a function of the minimal hadron p_T. This figure is taken from Ellis and Kowalski (Ref. 33).

SCENARIO C: $m_{\tilde{g}} \ll m_{\tilde{q}} \simeq 100$ GeV [5,10,11]

If a gluino is sufficiently light, then the $\tilde{g}\tilde{q}$ production cross section which is formally $O(\alpha_s^2)$ can numerically become as large as $O(\alpha_s)$ because of the appearance of a large logarithmic coefficient $\ln(m_{\tilde{q}}^2/m_{\tilde{g}}^2)$. This can be most easily understood if we introduce an effective gluino distribution in a nucleon[34]. Then the heavy squark production cross section via quark-gluino fusion is just $O(\alpha_s)$. The produced squark would mainly decay into a quark and a gluino, but it can also decay into a quark and a photino which gives a clean monojet event. Even after multiplying by the small branching fraction of $O(\alpha/\alpha_s)$, the monojet cross section is still $O(\alpha)$ and could be significant.

This scenario was first proposed by Herrero et al.[5] as a possible clean signal of supersymmetry at hadron colliders. After the observation[7] of monojet events was reported, we[10] were forced to reinvent the same scenario by systematically searching for a supersymmetry scenario which can give rise to a clean monojet. With $m_{\tilde{q}} \simeq 100$ GeV, we can expect a clustering of the events near $\not{p}_T^q \simeq 50$ GeV due to the Jacobian peak. With $m_{\tilde{g}} \simeq 3$ GeV, the monojet cross section is expected to be about 20 pb at the CERN collider[10,14]. Furthermore, we do expect in this scenario that the resulting monojets should be narrower than a typical high p_T jet observed at the collider. There are two reasons for this. First our monojet is always a quark jet whereas a typical high p_T jet at the collider is a statistical mixture of a quark and a gluon jet. Second, the lowest order QCD contribution to the mass of the jet system turns out to be significantly smaller[14,35] than that for the mass of the e^+e^- annihilation jets. Hence we should expect that the monojet in this scenario is narrower than a typical hadron jet at similar p_T observed at the collider or extrapolated from e^+e^- jets. Quantitative significance of such an effect at the hadron level, however, remains unclear.

Implications of this scenario for a light gluino of mass about 3-5 GeV are quite involved. Gluinos, being colour octets, can be quite copiously pair-produced in hadron collisions[36]. Once they decay into photinos, a high p_T gluino-pair should lead to large $\not{p}_T$ at hadron colliders[6], while small p_T gluinos can lead to anomalous signals[36] in the beam dump experiments[37]. Early studies at the collider[7] seemed to rule out the gluino mass all the way up to 40 GeV, whereas the mass values of interest ($m_{\tilde{g}} \simeq 3$ GeV, $m_{\tilde{q}} \simeq 100$ GeV) lie near the boundary of the sensitivity of the present beam dump experiments[37]. It is therefore of great interest to study carefully if a light gluino can still be a viable possibility in the face of the collider data. There are three relevant issues to consider: the gluino distribution in a nucleon[5,11,34,38-40], the gluino fragmentation[35,39-41], and the contribution from the QCD 2 → 3 processes[42].

It is usually assumed that when one probes a nucleon at much higher momentum scale (Q) than a heavy parton mass (m > 1 GeV) there appears to be an effective parton distribution and that the distribution can be calculated via the Altarelli-Parisi equation[43] with an appropriate decoupling condition[44]. In our case, the probe scale $Q \simeq m_{\tilde{q}} \simeq 100$ GeV seems, at first sight, to be sufficiently larger than the parton mass $m_{\tilde{g}} = 3$-5 GeV to ensure the validity of the approximation where one calculates the squark production cross section via the fusion $q\tilde{g} \to \tilde{q}$ by using an Altarelli-Parisi generated effective gluino distribution in a nucleon. We examined[40] this by generating an effective gluino distribution with the simplest decoupling condition

$$\tilde{g}(x, Q < 2m_{\tilde{g}}) = 0 \ .$$

Surprisingly, this overestimates the squark production cross section by more than a factor of three compared to the lowest-order calculation with the exact $qg \to \tilde{q}\tilde{g}$ kinematics. Because of this discrepancy, a squark production cross section of 2 nb (which is needed to explain the monojet rate) is obtained with $m_{\tilde{g}}$ as large as 20 GeV[11,39] in the $q\tilde{g} \to \tilde{q}$ approximation whereas the same cross section requires $m_{\tilde{g}} < 5$ GeV[10] in the exact lowest order calculation.

To clarify the problem, we examined[40] the $q\tilde{g} \to \tilde{q}$ calculation in the single logarithmic approximation (or equivalent gluino approximation, $E_{\tilde{g}}A$). The total squark production cross section is expressed in the lowest order as

$$\sigma(p\bar{p} \to \tilde{q}X) = \int d\sqrt{\hat{s}} \left(\frac{dL}{d\sqrt{\hat{s}}}\right)_{qg} \hat{\sigma}(qg \to \tilde{q}\tilde{g})$$

where $\hat{\sigma}$ denotes the subprocess cross section and $(dL/d\sqrt{\hat{s}})_{qg}$ is the parton luminosity density[45], $\sqrt{\hat{s}}$ being the invariant mass of the colliding quark-glue system. In the $E_{\tilde{g}}A$, $\hat{\sigma}$ can be expressed as a convolution of the $g \to \tilde{g}$ splitting function[43,34] and the $q\tilde{g} \to \tilde{q}$ fusion cross section. We find that this convolution integral gives a resonable approximation to the exact $\hat{\sigma}$ <u>except</u> in the tiny region near the threshold

$$m_{\tilde{q}} < \sqrt{\hat{s}} \lesssim m_{\tilde{q}} + m_{\tilde{g}}$$

where it gives a non-vanishing contribution whereas the exact cross section vanishes. This small discrepancy in the subprocess cross section $\hat{\sigma}$, however, leads to the major discrepancy in the total cross section σ because the quark-gluon luminosity is rapidly falling with increasing $\sqrt{\hat{s}}$ in the relevant region. We show in Fig. 4a and b the product of the luminosity function and $\hat{\sigma}$ as a function of $\sqrt{\hat{s}}$ with the exact kinematics and with the $E_{\tilde{g}}A$, respectively. The area under the curves show the total cross section. We see that the area under the $m_{\tilde{g}} = 20$ GeV curve (dotted line in Fig. 2b) with the $E_{\tilde{g}}A$ is almost as large as the area under the $m_{\tilde{g}} = 5$ GeV curve (solid line in Fig. 2a) with exact kinematics.

The threshold suppression effect for the heavy quark leptoproduction was carefully studied by Glück et al.[46], whose modified gluon-to-heavy quark kernel was subsequently used in the parametrization of Ref. 45. Shown in Fig. 4c are the $E_{\tilde{g}}A$ results obtained by using their modified kernel. The discrepancy with the exact results is still large. This simply reflects the fact that the kinematics in the leptoproduction (space-like probe) and in hadroproduction (time-like probe) is different. Although further modification of the splitting function might lead to resonable distributions, we find it safe to use the exact lowest order results whenever possible. Similar care should be taken generally when one makes use of the Altarelli-Parisi generated heavy quark distributions.

When the gluino mass is below 10 GeV, the large $\not{p}_T$ (of typically > 25 GeV) comes only from p_T much greater than $m_{\tilde{g}}$, and hence the $\tilde{g} \to \tilde{g}_h$ ($\tilde{g}$-containing-hadron) fragmentation effects should become important. We may use[35,40] the parametrization of Peterson et al.[47] for the heavy quark fragmentation which interpolates well between the charm and bottom fragmentation functions[48]. A notable difference between the gluino and heavy quark fragmentation function is that the former, being colour octet, radiates off more gluons than the latter and receives larger scaling violation effects. Convenient parametrizations of the gluino fragmentation functions are given in Ref. 40.

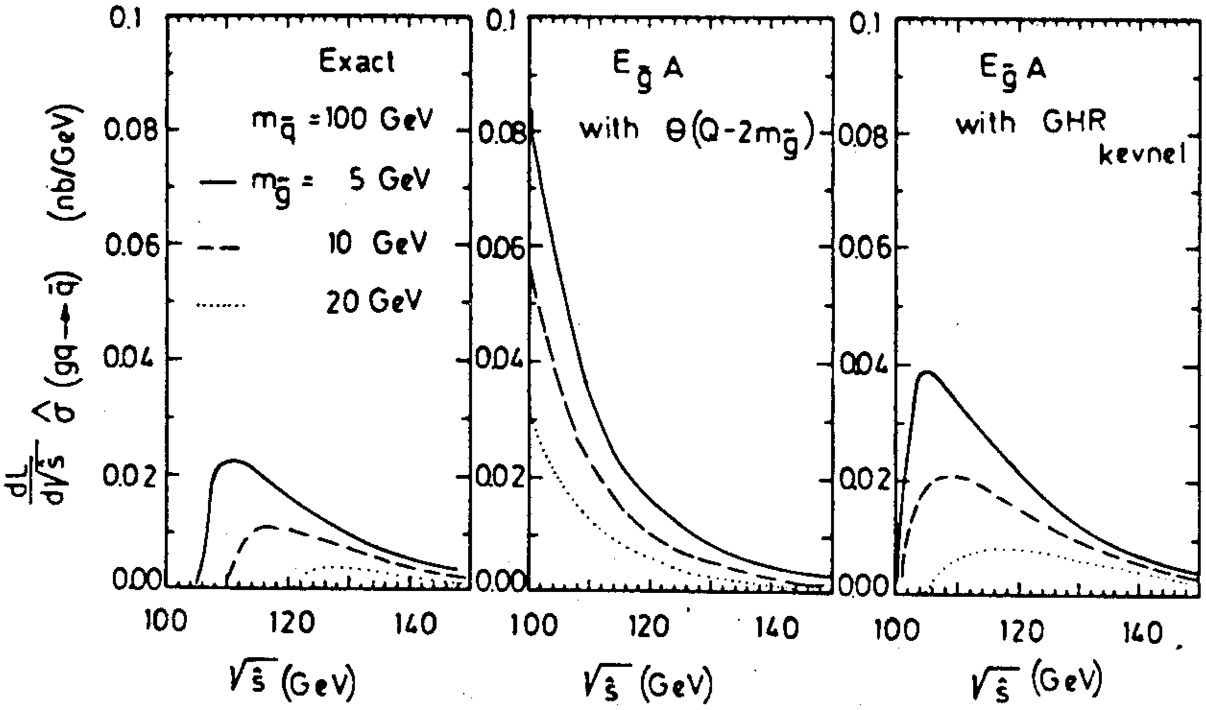

Fig. 4 Product of the subprocess cross section $\hat{\sigma}(ug \to \tilde{u}\tilde{g})$ and the ug luminosity function for $m_{\tilde{u}}$ = 100 GeV plotted against $\sqrt{\hat{s}}$ with (a) exact kinematics, (b) $E_{\tilde{g}}A$ with the theta function threshold, and (c) $E_{\tilde{g}}A$ with the kernel of Glück et al.[46]. Figs. 4a and b are taken from Ref. 40.

Recently, Herzog and Kunszt[42] pointed out that the QCD 2 → 3 processes, $gg \to g\tilde{g}\tilde{g}$ and $qg \to q\tilde{g}\tilde{g}$, give rise to large $\not{p}_T$ events with the rate larger than the 2 → 2 ($gg \to \tilde{g}\tilde{g}$ and $q\bar{q} \to \tilde{g}\tilde{g}$) contributions for smaller gluino masses ($m_{\tilde{g}}$ = 10 to 20 GeV). Fig. 5 shows schematically the five typical momentum configurations for the process $gg \to g\tilde{g}\tilde{g}$ where the matrix element becomes large. Curly lines and solid lines denote gluon and gluino three-momentum vectors, respectively, in the colliding gluon c.m. frame. The configurations may be labelled as (a) $g \to \tilde{g}$ splitting, (b) $\tilde{g}$-excitation, (c) gluon emission colinear to a gluino, (d) gluon emission colinear to initial gluons, and (e) soft gluon emission. Among these, the latter three configurations give similar final states to the leading order one, i.e. a back-to-back gluino pair. Hence the loop correction in the same order is required to know the actual magnitude of the correction. In particular, the leading logarithmic terms in the configuration (c) and (d) are already taken into account by the scaling violations of the $\tilde{g}$ fragmentation function and the g distribution in a nucleon, respectively. On the other hand, the configurations (a) and (b) appear only in the α_s^3 or higher orders. The 2 → 3 processes hence give the leading contributions to these two configurations where two

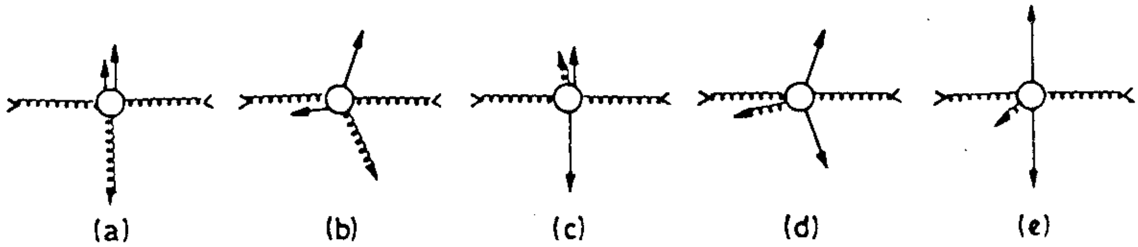

Fig. 5 Schematical view of the five typical three momentum configurations of the subprocess $gg \to g\tilde{g}\tilde{g}$ where the matrix element is large. Curly lines denote gluons and solid lines gluinos.

high p_T gluinos are almost collinear (a) or only one gluino has high p_T (b). In fact, Herzog and Kunszt observed that the collinear gluino configuration (a) gives the dominant source of large $\not{p}_T$ for lighter gluinos because the magnitudes of two photino transverse momenta add up to give large $\not{p}_T$ in this collinear configuration whereas they tend to cancel out in the $2 \to 2$ contributions where only the difference of the two photino transverse momenta gives $\not{p}_T$ in the back-to-back configuration.

To examine this $g \to \tilde{g}$ splitting effect quantitatively for the relevant mass range, $m_{\tilde{g}}$ = 3-5 GeV, we introduce a scalar source for two gluons[49] and a spinor source for a quark and gluon system. By attaching a gluino-pair to a gluon leg emitted from these sources (see Fig. 6), we obtain very simple amplitudes with the correct leading $g \to \tilde{g}$ splitting behaviour in the collinear configuration with exact $2 \to 3$ kinematics. Normalizing the magnitude and aligning the jet axis to those in the dominant $2 \to 2$ subprocesses ($gg \to gg$ and $qg \to qg$), we obtain a simple $2 \to 3$ cross section which simulates the exact $2 \to 3$ cross section in the collinear gluino-pair configuration (Fig. 4a) but gives negligible contributions in all the other configurations shown in Fig. 4.

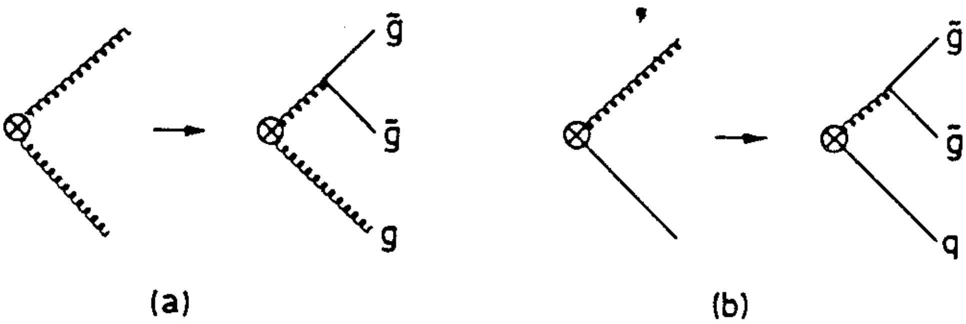

Fig. 6 Feynman diagrams for a scalar source to gg and $g\tilde{g}\tilde{g}$ (a) and for a spinor source to qg and $q\tilde{g}\tilde{g}$ (b).

Shown in Figs. 7a and b are, respectively, the 2 → 2 and the 2 → 3 subprocess contributions to the $\not{p}_T$ distributions[50] for $m_{\tilde{g}}$ = 3 GeV at $\sqrt{s}$ = 630 GeV. The importance of the $\tilde{g} \to \tilde{g}_h$ fragmentation effect is clearly seen in these figures. Also shown is the significance of the experimental $\not{p}_T$ resolution (4σ) cut; dotted, dashed and solid lines denote the results obtained with no 4σ cut, by imposing the 4σ cut with spectator E_T = 25 GeV, and with the 4σ cut

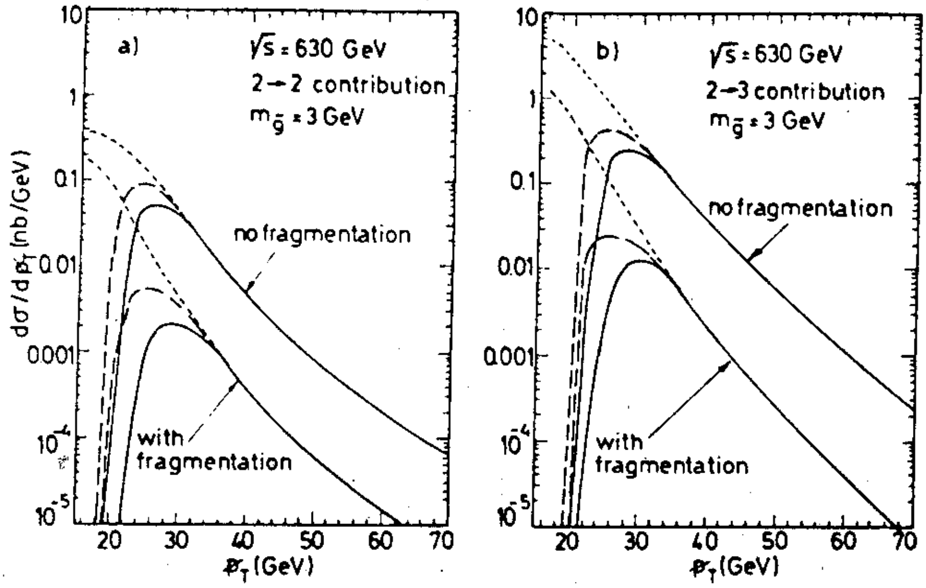

Fig. 7 Missing p_T distributions from the 2 → 2 (a) and the 2 → 3 (b) processes for $m_{\tilde{g}}$ = 3 GeV in $p\bar{p}$ collisions at $\sqrt{s}$ = 630 GeV taken from Ref. 50. See text for more details.

with spectator E_T = 40 GeV, respectively. The significance of the possible large spectator E_T effect on the experimental $\not{p}_T$ resolution cut was emphasized by Haber[51]. Hence the solid curves with fragmentation and large spectator E_T give the most conservative estimates. It is clear that the 2 → 3 processes give by far the dominant contribution to the large $\not{p}_T$ events. We find for $m_{\tilde{g}}$ = 3 GeV (5 GeV) the integrated rates 25 pb (60 pb) from the 2 → 2 processes whereas 120 pb (210 pb) arise from the 2 → 3 processes. Roughly half of the events are classified as monojet and the remaining half as dijet events according to the UA1 jet selection criteria[7,25]. These contributions are peaked around $\not{p}_T \simeq$ 30 GeV and have almost always back-to-back jet activities, that is, $\not{p}_T$ vector is not isolated in the transverse plane. Almost no events survive above $\not{p}_T \simeq$ 40 GeV. Hence a typical event may look just like a jet-jet fluctuation of an ordinary dijet event where one of the jet p_T is overestimated and the other underestimated. The large expected rates (50-100 events with the accumulated integrated luminosity of 0.4 pb^{-1}), however, means that either this signal should be observed or a light gluino should be ruled out in the near future. The estimates for the rate are very conservative since (i) we use smaller values for the QCD coupling, (ii) we have not taken into account the 'excitation' type configuration (Fig. 4b), and (iii) we expect a significant enhancement[52] from higher order corrections.

Needless to say the above arguments do not apply if gluinos are rather stable[10,53]. However, for gluinos of mass around 3-5 GeV, we should requite a near degeneracy of gluino and photino masses[10]. A theoretically more attractive possibility is that all the neutral gauginos receive a common Majorana mass at the GUT scale, in which case we expect[31] $m_{\tilde{\gamma}}/m_{\tilde{g}} \simeq 8\alpha/3\alpha_s(m_{\tilde{g}}) \simeq 1/6$. Hence we expect a light ($\sim$ 0.5 GeV) photino for $m_{\tilde{g}} \lesssim$ 3 GeV, which would be a cosmological embarrassment[54] if the photino is stable because it would contribute too much to the mass density of the universe. The simplest way to avoid this problem would be to expect photinos to decay[21-23] either very slowly or into an invisible decay mode[24] as discussed in the introduction.

HIGGSINOS AS THE CARRIERS OF MISSING MOMENTUM

Order α processes, that is, the W and Z decays into supersymmetric particle pairs can in principle give rise to monojet events. Here the colourless supersymmetric particles can be produced significantly. In order to produce monojets, one of the produced pair should be neutral and long-lived to escape detection and the other should decay into hadrons. Narrow high p_T monojets can be expected only when both supersymmetric particles are rather light (< 10 GeV). This rules out[32] charged superparticles as candidates and leaves the neutral decay modes of the Z boson. Since the Z boson does not couple to photino and zino at the tree level, the unique possibility is the Z decay into Higgsinos.

The decay mode $Z \to \tilde{h}_1\tilde{h}_2$ was, to my knowledge, first studied by Ellis et al.[4] and its implication to the collider monojets studied in Ref. 55. We assume $\tilde{h}_1$ to be light (< a few GeV) and stable or long-lived to escape detection. The heavier Higgsino $\tilde{h}_2$ would then decay into $\tilde{h}_1$ and a quark-pair via a virtual Z exchange. We show in Fig. 8 the Feynman diagram of the subprocess. The monojet production rate at the CERN collider has been estimated[55] to be about 20 pb for $\not{p}_T$ > 35 GeV with the maximum $Z\tilde{h}_1\tilde{h}_2$ coupling[56].

Once such light Higgsinos exist, they can be produced in the present e^+e^- annihilation experiments at PEP/PETRA via a virtual Z exchange. Depending on the mass difference between $\tilde{h}_1$ and $\tilde{h}_2$, a typical event would look like a monojet plus $\not{p}_T$ or a dijet plus $\not{p}_T$ in lower energy e^+e^- experiments. Also shown in Fig. 8 is the experimental bound for the mass and coupling of the Higgsinos obtained by the JADE collaboration[57] from the nonobservation of such events. This bound is so stringent that virtually no monojet event can be expected from this source at the CERN collider.

Perhaps the most important lesson learned from this example is that it exemplifies the fruitful interactions between the hadron collider experiments at higher energies and the e^+e^- collider experiments at lower energies but in a cleaner environment. The aforementioned JADE bound was obtained in the region where the Higgsino pair production cross section should be more than two order of magnitude smaller than that at the

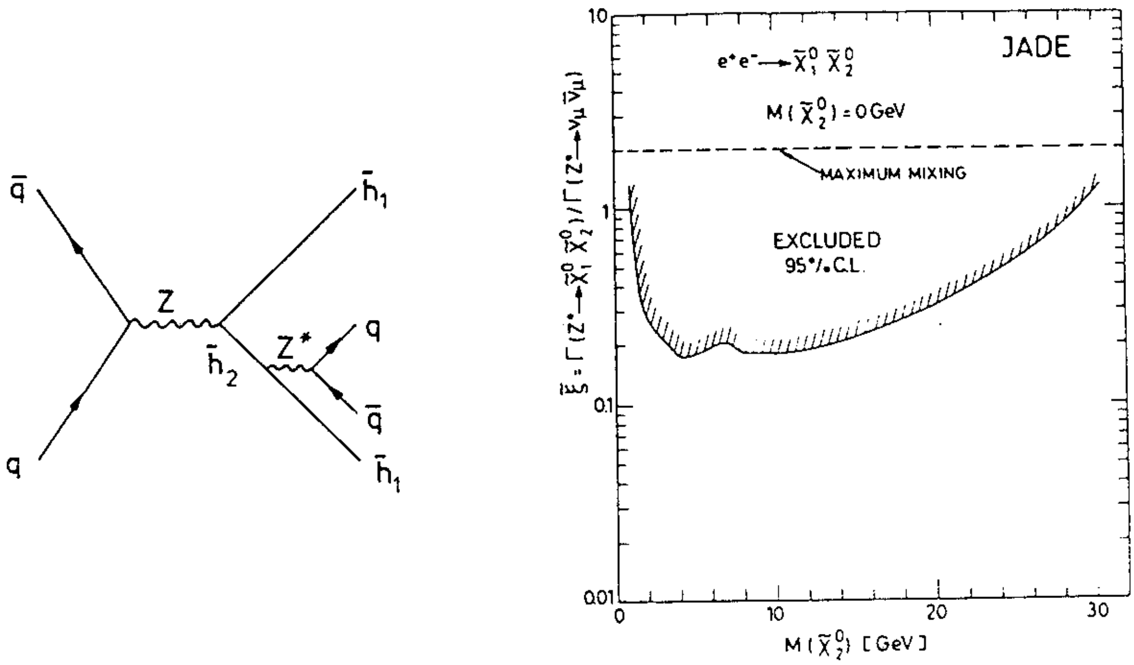

Fig. 8 Feynman diagram and the experimental bound from Ref. 57 for the Higgsino-pair production and decay. $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$ of Ref. 57 correspond, respectively, to $\tilde{h}_1$ and $\tilde{h}_2$.

CERN collider. In e^+e^- collider experiments it would also have been easy[55] to distinguish between monojets originating from a fermion (Higgsino) pair or from a scalar (Higgs boson) pair[58] by using the monojet angular distribution whereas such a determination would be extremely hard at hadron colliders. The slightest hint from hadron collider experiments could be examined in detail in e^+e^- collider experiments.

GOLDSTINOS AS THE CARRIERS OF MISSING MOMENTUM

In a locally supersymmetric theory (supergravity theory)[19], a massless Goldstino is absorbed into the helicity ±1/2 component of a massive gravitino (denoted by $\tilde{G}$ like the Goldstino for reasons given below) due to the super-Higgs effect[59]. A general relation[59] between the gravitino mass ($m_{\tilde{G}}$) and the supersymmetry breaking scale (Λ_{ss}) reads

$$m_{\tilde{G}} \sim \Lambda_{ss}^2/M_P$$

where $M_P \sim 10^{19}$ GeV is the Planck mass. In the most attractive class of supergravity models[19] $m_{\tilde{G}} \sim m_W$ and we do not expect the gravitino to play a role in high energy physics apart from the possible cosmological problem[60,61] due to its small gravitational coupling ($\sim 1/M_P$).

A very light gravitino ($m_{\tilde{G}} < 1$ keV) is at least a logical possibility, as a consequence of a naive first guess $\Lambda_{ss} \sim 1$ TeV, and is cosmologically safe[60]. Although we do not yet have a realistic model with $\Lambda_{ss} \sim 1$ TeV, the possibility of an almost massless gravitino should not be ignored experimentally. One may ask how can ordinary high energy experiments detect the effects of a gravitino which interacts only with gravitational coupling. A crucial observation was made by Fayet[62]: when the gravitino mass is much smaller than the physical mass scale of the process its helicity ±1/2 components couple strongly. The reason is the wave function factor proportional to $1/m_{\tilde{G}}^2 \sim M_P^2/\Lambda_{ss}^2$ cancels the gravitational coupling squared $(1/M_P)^2$ to give a rate proportional to $1/\Lambda_{ss}^2$. The dominant piece of the cross section is then calculated by using only its helicity ±1/2 components, namely the Goldstino. Hence the consequences of a very light gravitino in high energy experiments follow from massless Goldstino dynamics[62,63]. This is analogous to the well known fact[64] that the high energy interaction of massive gauge bosons in spontaneously broken gauge theories is dominated by their helicity zero components and the dominant piece is determined by the Higgs sector without the gauge interaction. Hence we examine the possibility that massless Goldstinos ($\tilde{G}$) carry away the missing momentum.

The simplest case, at least as far as the number of participating new particles is concerned, is when all the superpartners, except the Goldstino, are heavy at the energy scale probed by the CERN collider. In this case supersymmetry is realized nonlinearly[65] simply because no supermultiplets appear at our energy and high dimensional interactions including two Goldstinos appear. One of such lowest-dimension operators gives the four fermion coupling between a quark-pair and a Goldstino-pair as depicted in Fig. 9. One gluon is attached to the vertex because the process $q\bar{q} \to \tilde{G}\tilde{G}$ does not lead to any $\not{p}_T$; the process $q\bar{q} \to g\tilde{G}\tilde{G}$ leads to an observable monojet plus $\not{p}_T$ events. Consequences of this type of interaction at hadron colliders were studied by Nachtman et al.[66]. Fig. 9 shows only a partial contribution from the $q\bar{q} \to g\tilde{G}\tilde{G}$ subprocess to the $\not{p}_T$ spectrum at two collider c.m. energies, 540 GeV and 630 GeV, for illustration. First of all the $\not{p}_T$ spectrum is soft, which is essentially determined by the one gluon emission from the initial parton legs. Secondly, there is a striking energy dependence which gives an almost 4 times larger rate for the events with $\not{p}_T > 40$ GeV at $\sqrt{s} = 630$ GeV as compared to the one at $\sqrt{s} = 540$ GeV. This striking energy dependence is a consequence of the high dimensionality of the four-fermion operator; since each Goldstino couples derivatively according to the low energy theorem[67], the $\bar{q}q\tilde{G}\tilde{G}$ operator has dimensionality eight and the coupling is proportional to $1/\Lambda_{ss}^4$ as compared to the Fermi coupling which is proportional to $1/m_W^2$. Because of this, the subprocess cross sections scale as

$$d\hat{\sigma}(q\bar{q} \to g\tilde{G}\tilde{G}) \propto \hat{s}^3/\Lambda_{ss}^8$$

which gives rise to the strong energy dependence. The qualitative behaviour of the monojet cross section shown in Fig. 9 should not change on including all the leading order ($\sim \alpha_s/\Lambda_{ss}^8$) contributions. Since there is no hint[16] of an increase in the monojet cross section

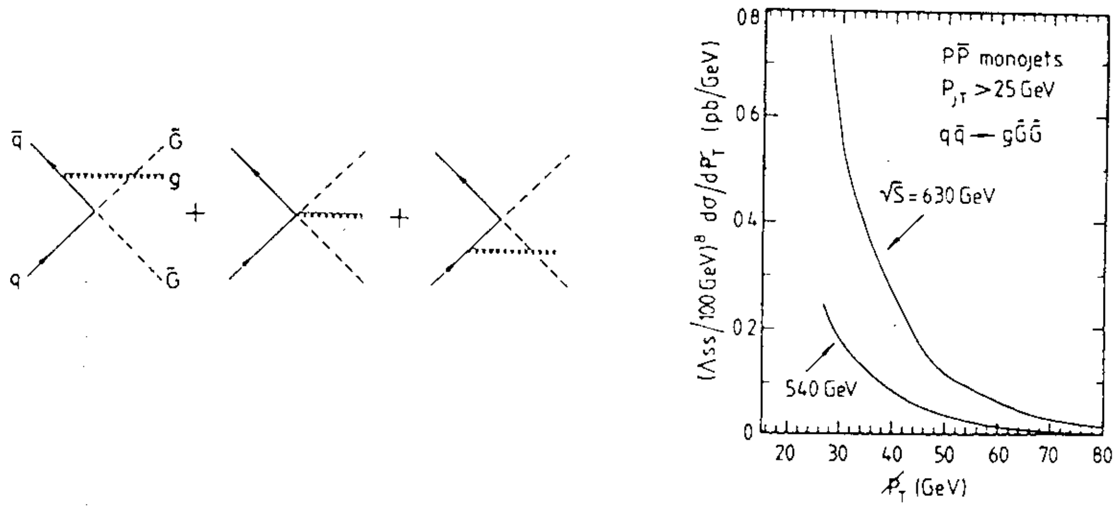

Fig. 9 Feynman diagrams of the process $q\bar{q} \to g\tilde{G}\tilde{G}$ and its contribution to the monojet p_T distribution at $\sqrt{s}$ = 540 GeV and 630 GeV.

at $\sqrt{s}$ = 630 GeV, this interesting possibility cannot be regarded as an explanation of the observed[7] monojet events.

Finally, there is a possibility of clean monojets if in addition to the massless Goldstino a very light ($\lesssim$ 100 MeV) gluino exists. In this case the Goldstino couples to the derivative of the gluon-gluino supercurrent and its low energy interaction can be expressed by the effective Lagrangian[68],

$$L_{eff} = - \frac{1}{2\Lambda_{ss}^{2}} \overline{\psi_{\tilde{G}}}\, \sigma_{\mu\nu}\, \not{\partial}\, (F^{a\mu\nu} \psi^{a}_{\tilde{g}})$$

where $F^{a\mu\nu}$ denotes the usual gauge covariant gluon field strength. It is then straightforward to calculate the cross sections for the processes $q\bar{q} \to \tilde{g}\tilde{G}$ and $gg \to \tilde{g}\tilde{G}$ which scale as $\hat{s}/\Lambda_{ss}^{4}$. The produced gluino would then hadronize into a colour singlet bound state[69] $(\tilde{g}g)$ or $(\tilde{g}q\bar{q})$ and the $\tilde{g}$-jets would just look like ordinary jets. We show in Fig. 10 the Feynman diagram and the partial contribution[70] to the monojet cross section at $\sqrt{s}$ = 540 GeV and 630 GeV from the subprocess $q\bar{q} \to \tilde{g}\tilde{G}$. The $\not{p}_T$ spectrum is very hard and the energy dependence is moderate. The supersymmetry breaking scale Λ_{ss} of about a few hundred GeV would then lead to a desirable monojet rate. In this scenario, it would be natural to expect the photino to also be light and the analogous process $q\bar{q} \to \tilde{\gamma}\tilde{G}$ and $e^+e^- \to \tilde{\gamma}\tilde{G}$ would occur with the relative rate α/α_s. The subsequent $\tilde{\gamma} \to \gamma\tilde{G}$ decay[21] leads to a single photon plus $\not{p}_T$ event. Probably the best place to test this scenario is again in e^+e^- annihilation processes at high energies.

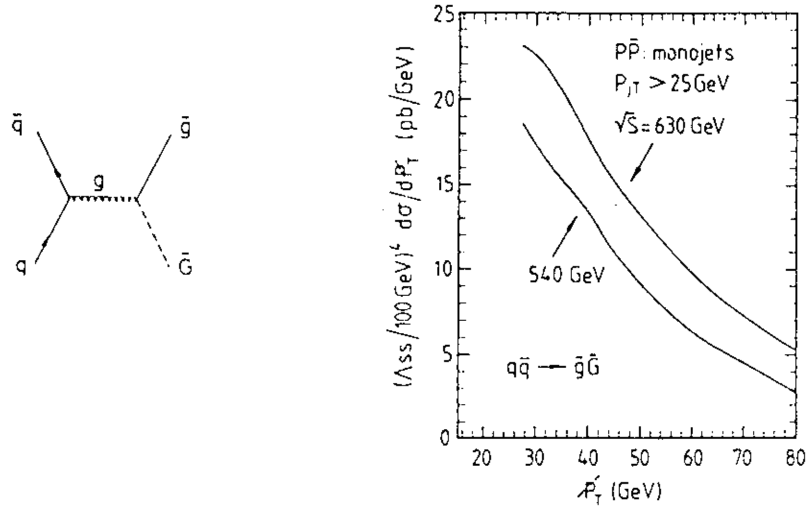

Fig. 10 Feynman diagram of the process $q\bar{q} \to \tilde{g}\tilde{G}$ and its contribution to the monojet p_T distribution at $\sqrt{s}$ = 540 GeV and 630 GeV.

CONCLUSIONS

I have reviewed various attempts to understand the observed monojet events at the CERN collider[7,16] as the first experimental signal of supersymmetry. The monojet dominance and the narrowness of the jets seem to rule out the two most attractive scenarios where either squarks[12,13] or gluinos[6,8] of mass around 40-50 GeV are pair produced and decay into photinos. A scenario where the Z boson decays into two Higgsinos[4,55] is ruled out by e^+e^- experiments at PETRA. An interesting scenario[66] where the Goldstino is the only light superparticle fails to account for the apparently non growing rate with energy. Among the possibilities I examined, only two scenarios do not immediately contradict the observations. These are the scenario[5,10] where a heavy ($\sim$ 100 GeV) squark and a relatively light ($\sim$ 5 GeV) gluino are produced and the scenario[70] where a massless Goldstino and a very light ($\lesssim$ 100 MeV) gluino are produced. These scenarios should be able to be tested in the near future by examining carefully the $\not{p}_T$ events with back-to-back dijet configuration for the former one and by searching a single photon and $\not{p}_T$ events at e^+e^- annihilation for the latter.

I concentrated only on the $\not{p}_T$ plus hadronic jet signal of various supersymmetry scenarios because it is the only observed anomaly which hints at new physics. Recently a number of authors studied the implications of supersymmetry at hadron colliders by examining signals with lepton plus jet[71] and also the purely leptonic[72] signals. Detailed study of these pure- and semi-leptonic signals of supersymmetry is worth pursuing since the relative cleanness of such events could beat the lower rate and also because we can expect a substantial improvement in luminosity in the near future.

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