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LIGHT GLUINOS AND HEAVY SQUARKS
CONFRONT THE COLLIDER DATA

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ABSTRACT

We analyze the missing transverse momentum events seen at the CERN $\bar{p}p$ collider within a supersymmetric model with light gluinos, $m(\tilde{g}) \sim \text{few GeV}$, and medium heavy squarks, $m(\tilde{q}) > 80 \text{ GeV}$. Gluino fragmentation, previously ignored, is crucial to the analysis, e.g., the published limit $m(\tilde{g}) > 40 \text{ GeV}$ does not survive. A model with $m(\tilde{g}) \sim 3 \text{ GeV}$, $m(\tilde{q}) \sim 80 \text{ GeV}$, successfully describes all features of the data, including the strikingly small mass of the observed monojets. A model with $m(\tilde{g}) \sim 7 \text{ GeV}$, $0(1 \text{ TeV}) > m(\tilde{q}) > 150 \text{ GeV}$ is also acceptable. If one believes in relations between supermasses derived from models incorporating supergravity, the masses we "determine" lead to a contradiction with cosmological constraints on relic photinos. We choose to doubt these mass relations and we suggest further collider tests of models with light gluinos.

1. - INTRODUCTION AND OUTLOOK

Much ado about "monojet" events, observed by the UA1 collaboration in $p\bar{p}$ collisions at $\sqrt{s} = 540$ GeV, has recently hit the literature. Six events have been recorded¹⁾ consisting of narrow energetic jets (transverse energy > 25 GeV) accompanied by a large compensating "missing transverse momentum" carried away by an unobserved particle (or particles). One hopes these events will not slowly drip down the drain in the form of an increasingly better understood background, such as the production of neutrinos from Z^0 -decay, in association with a "QCD" jet²⁾. Events with missing energy have been known for long to be a possible signature for the production of supersymmetric partners of the known particles³⁾.

A handful of supersymmetric "scenarios" (an adequately theatrical expression) have been mounted on attempts to fit these scarce data⁴⁾⁻⁸⁾. To accept any of these scenarios one must momentarily forget that events containing very hard photons (or photon-like signals) accompanied by hadrons or lepton pairs, have also been observed. These hard photons are inexplicable within existing super-models and would rule them out as "theories of everything". Reluctantly following a tradition of amnesia, we close our eyes to photons, we hope they will not recur in higher statistics data, and we proceed to discuss our favourite super-scenario: the "Madrid" model⁹⁾.

Prior to publication of the monojet data, Herrero et al.⁹⁾ serendipitously (?) announced that if gluinos ($\tilde{g}$) happened to be very light [$m(\tilde{g}) \sim$ a few GeV], squarks ($\tilde{q}$) of masses up to ~ 150 GeV would be copiously produced at CERN's collider in the process $p\bar{p} \rightarrow \tilde{g}\tilde{q} + X$. Squark decays into a quark and a "weakly" interacting photino ($\tilde{q} \rightarrow q\tilde{\gamma}$) would then be obvious candidates for "clean" monojet plus missing p_T events. But, we are told, experiment¹⁾ excludes gluino masses below ~ 40 GeV. This claim is based on the following reasoning. Light gluinos should be very copiously pair produced, $p\bar{p} \rightarrow \tilde{g}\tilde{g} + X$, as in Fig. 1a. The subsequent gluino decays $\tilde{g} \rightarrow q\tilde{q}$ would lead to unreasonably large numbers of events in which the missing transverse momenta of the two photinos add to a large value $|\vec{p}_{T1} + \vec{p}_{T2}|$. This argument is reflected numerically⁵⁾ in the curve labelled "NOFRAG" in Fig. 2, where it is seen that too many missing p_T events would have survived the experimental cuts for $m(\tilde{g}) < 40$ GeV.

In an attempt to avoid the constraint $m(\tilde{g}) > 40$ GeV, Barger et al.⁷⁾ have assumed photinos to be almost as massive as gluinos [$m(\tilde{\gamma}) \sim 0.9 m(\tilde{g})$], effectively lengthening the gluino lifetime to the point where gluino-containing

hadrons have no time to decay while they travel in the "central detector". Barger et al. show that this scenario of light stable gluinos generally describes the data better than competing scenarios with $m(\tilde{g}) \sim m(\tilde{q})$ or $m(\tilde{q}) \ll m(\tilde{g})$. However, an approximate mass degeneracy between photinos and gluinos is at variance with the generally accepted prediction¹⁰⁾ that $m(\tilde{g}) \sim 7 m(\tilde{\gamma}) \sim 0(\alpha_s/\alpha)m(\tilde{\gamma})$. Fortunately, the mass-degeneracy assumption is unnecessary. The point is that the conclusion that light gluinos [$m(\tilde{g}) \sim \text{few GeV}$] are excluded by experiment, may be an incorrect one. The reason, as we shall discuss very explicitly below, is that past calculations of gluino-induced missing p_T signals have ignored fragmentation. A gluino made in a reaction at an invariant energy much bigger than $m(\tilde{g})$ is bound to radiate many "fragmentation" gluons before it picks up a $q\bar{q}$ pair or a "constituent" gluon to become a colourless superhadron $\tilde{G}$, a process depicted in Fig. 1b. This significantly lowers the momentum of the gluino before it decays with the emission of a photino. Fragmentation effectively reduces the momentum of the final photinos, diminishing the number of events with a missing p_T above a fixed minimum. Our results on the effects of fragmentation, to be explicitly discussed in Chapter 2, are shown in Fig. 2. The curve labelled "NOFRAG" is the total cross-section with the experimental cuts on a minimum missing p_T and a minimum jet transverse energy, in a calculation that ignores gluino fragmentation⁵⁾. The curve labelled "FRAG" corresponds to our total results with fragmentation taken into account. The curves labelled "MONOJETS" and "DIJETS" specify the topologies of the events with missing p_T , with the experimentally defined conditions on these "jets". On the basis of the no-fragmentation cross-section plotted in Fig. 2 and the experimental observation of no more than six "missing p_T " events in a run with an integrated luminosity of 113 nb^{-1} , the UA(1) collaboration¹⁾ placed a limit $m(\tilde{g}) \gtrsim 40 \text{ GeV}$. On the same basis, and with gluino fragmentation taken into account, we conclude from Fig. 2 that $m(\tilde{g}) \lesssim 7 \text{ GeV}$ is also a possibility. If we were to ascribe to the theoretical calculation and its mock-up of experimental cuts an uncertainty of half an order of magnitude, a realistic attitude, we would be left with no statement at all to be made on the mass of the gluino. An optimistic and diametrically opposed attitude would be to attribute the observed six events to the pair production of gluinos heavier than 40 or lighter than 7 GeV. But the observed monojets have very narrow width: the invariant mass, m_{CH} , of their charged tracks is $\sim 3 \text{ GeV}$ in one case, less in all other measured ones. This surprisingly narrow widths makes the possibility of attributing the observations to the production of gluino pairs somewhat uncomfortable, even if gluinos are very light. As we shall see in detail in Chapter 2, our predicted average invariant mass for a light gluino jet

is roughly $7 \text{ GeV} + m(\tilde{g})/2$. If $\langle m_{\text{CH}} \rangle \sim \langle W \rangle/2$, the prediction exceeds the experimental indications even for $m(\tilde{g})$ as low as 3 GeV.

Having concluded that pair-production of light gluinos [$m(\tilde{g}) \lesssim 7 \text{ GeV}$] is not excluded by experiment, but is unlikely to be the principal source of high missing p_T events, we turn our attention to the next lightest possible super-channel in the Madrid scenario: $p\bar{p} \rightarrow \tilde{q}\tilde{g} \dots$. We shall conclude that this channel is a better candidate-source of the observed events, provided $m(\tilde{q}) \sim 80 \text{ GeV}$. To understand qualitatively the results we shall present, the following comments may be useful. The dominant production amplitude for $\tilde{q}\tilde{g}$ pairs in the range of masses we analyze is $q\bar{q} \rightarrow \tilde{q}\tilde{g}$, in which a gluino is exchanged. For light gluinos this process is forward-backward peaked. Thus, the decay products of the gluino are unlikely to pass the cuts defining a large missing p_T or a jet with large E_T . The eventual striking signatures mainly come from the decay of the heavy squark. Since the most demanding experimental cuts (on jet energies or missing energy) are transverse, one may think of the squark as being at rest in the lab when computing the effects of the cuts on the cross-sections. The cleanest missing p_T plus monojet signal in squark decay is due to the process $\tilde{q} \rightarrow q\tilde{\gamma}$, whose branching ratio is "electromagnetic" (1.5%, when averaged over $\tilde{u}$ and $\tilde{d}$ charges in a production ratio 2:1). This "clean" channel, depicted in Fig. 1c, is the one analyzed in detail by Barger et al.⁶⁾, who, having assumed light gluinos to be effectively stable, had no other source of missing p_T events in the Madrid scenario. Some of their results are reproduced in Fig. 3, where the positions of the observed events are indicated. The calculations are for $m(\tilde{g}) = 3 \text{ GeV}$, but the results are not very different in the range $m(\tilde{g}) = 3$ to 7 GeV. The predicted total cross-section with experimental cuts, Fig. 3d, indicates that the observed monojets, if they originate in the chain $p\bar{p} \rightarrow \tilde{q}\tilde{g}$, $\tilde{q} \rightarrow q\tilde{\gamma}$, would be compatible with squark masses in the range of 80 to 120 GeV. The small contamination of dijets in Fig. 3d originates from the gluinos primarily produced alongside the squarks ($q\bar{q} \rightarrow \tilde{q}\tilde{g}$), that sometimes are not so collimated along the beam direction as to escape the transverse energy threshold defining a second jet. (This contamination diminishes as one lets the gluino decay: its hadronic debris have a slightly smaller probability to pass the second jet E_T cut. We ignore it for the time being, to discuss it later when we present detailed results.) In a standard scenario wherein gluinos are not effectively stable, the relatively rare $\tilde{q} \rightarrow q\tilde{\gamma}$ decays are not the only source of missing p_T events in $\tilde{q}$ decay. The chain decay $\tilde{q} \rightarrow q\tilde{g}$, $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}$, depicted in Fig. 1c, seems at first sight to be an overwhelming competitor, since both decays are expected to occur with a branching ratio very close to unity. But here, once again, gluino fragmentation plays a very important role. The light gluino is produced in a process at an energy

scale $m(\tilde{q})$, and has plenty of room to have its momentum degraded before it becomes the constituent of a hadron. By the time the gluino decays, the typical photino momenta are much smaller than the maximum, $p_T^{\max} \sim m(\tilde{q})/2$. Since the experimental missing p_T cut is of the order of 28 GeV, and may be close to $p_T^{\max}$, the fraction of photinos that pass the cut is very small. This is made explicit in Table 1, where we give our results for the number of monojets plus missing p_T events from the sum of the two sources $\tilde{q} \rightarrow q\tilde{g}$ and $q\tilde{g}$, divided by the corresponding number from the "clean" source ($\tilde{q} \rightarrow q\tilde{\gamma}$), as a function of $m(\tilde{q})$ and $m(\tilde{g})$.

$m(\tilde{g})$ $m(\tilde{q})$	3	5	7
80	1.02 (< 0.1)	1.17 (0.1)	1.36 (0.15)
100	1.01 (0.52)	1.10 (1.2)	1.25 (1.7)
120	1.01 (1.4)	1.06 (2.7)	1.15 (3.7)

Table 1: Monojets from squark decay normalized to monojets from the specific decay channel $\tilde{q} \rightarrow q\tilde{\gamma}$, as a function of $m(\tilde{g})$ and $m(\tilde{q})$, in GeV. The numbers in parenthesis are the ratios of dijets to monojets, following the production process $p\bar{p} \rightarrow \tilde{g}\tilde{q}\dots$.

These ratios are close to unity, the "clean" decay $\tilde{q} \rightarrow q\tilde{\gamma}$ is the dominant source of monojets in squark decay in the range of parameters we are considering. Incidentally, the width of the monojets from squark decay is expected to be comfortably narrow. We shall see in Chapter 3 that the predicted mean value of the invariant mass of their charged tracks, in a first order QCD-estimate, is $\langle m_{CH} \rangle \sim 0.018 m(\tilde{q})$, or ~ 1.4 GeV for $m(\tilde{q}) \sim 80$ GeV. This estimate includes the effect of experimental cuts in jet angular spread, that bias the jets towards small invariant masses. The expected jet invariant masses in squark decay are not blatantly contradictory to the data; quite the contrary. Recall that six events have been observed, with a monojet whose transverse energy is an excess of ~ 25 GeV. No dijet events have been reported with a second "small" jet of transverse energy in excess of 12 GeV. The $\tilde{q} \rightarrow q\tilde{g}$ decays depicted in Fig. 1d may

give rise to dijets when the gluino bremsstrahlung and decay products are sufficiently collimated and energetic. The non-observation of dijets places severe limits on the values of $m(\tilde{g})$ and $m(\tilde{q})$, even in the narrow range of possibilities we are presently analyzing. This is shown in Table 1, where we report (in parentheses) our predictions at $\sqrt{s} = 540$ GeV for the dijet to monojet ratio in the $p\bar{p} \rightarrow \tilde{q}\tilde{g}$ channel. If this channel is the dominant source of observed events, squark masses above 100 GeV would be excluded, at least for $m(\tilde{g}) \gtrsim 5$ GeV. For $m(\tilde{q}) \sim 80$ GeV, and gluino masses up to the ~ 7 GeV limit implied by the $\tilde{g}\tilde{g}$ channel, the expected dijet to monojet ratio in squark decay is not in contradiction with the data.

We have argued that light gluinos, $m(\tilde{g}) < 7$ GeV, are not excluded on the basis of the expected signals following $\tilde{g}$ pair production. We have also announced that $\tilde{g}\tilde{q}$ production and the subsequent decays are candidate sources for the narrow monojet plus missing p_T events, provided the supermasses are in the neighbourhood of $m(\tilde{q}) \sim 80$ GeV, $m(\tilde{g}) \sim 3$ GeV. Before we let the reader jump to conclusions (we would not do it ourselves on the basis of a few events), we must discuss the results for the sum of all channels in the Madrid scenario. The relevant channels are $p\bar{p} \rightarrow \tilde{g}\tilde{g}$ and $\tilde{q}\tilde{g}$, the pair production of $\tilde{q}$ being strongly suppressed for $m(\tilde{q}) > 80$ GeV. Adding these channels is not without danger, for they have different theoretical uncertainties. Gluino pairs are produced in $q\bar{q}$ or gg scattering, while $\tilde{q}\tilde{g}$ pairs are produced in qg collisions. The two processes have a different functional dependence on the gluon structure function, which is relatively poorly known. Moreover one may take, as some do⁹⁾, the point of view that light gluinos are not only "made in the collisions" but somehow "pre-exist" in the form of a gluino sea. In our region of parameter space, $m(\tilde{g}) = 3$ to 7 GeV, $m(\tilde{q}) \sim 80$ to 120 GeV, a calculation based on a theoretical estimate of the gluino sea would enhance the $p\bar{p} \rightarrow \tilde{g}\tilde{q}$ cross-section by a factor 2 or 3, relative to a calculation wherein gluinos are made on the spot (the two points of view, it goes without saying, are not orthogonal or incompatible. The "sea" point of view, though somewhat uncertain and beset by threshold problems, is an improvement over the "production" point of view in that it consistently takes into account the interplay between gluon, quark and gluino structure functions). To be specific, we shall give results in the "production" viewpoint, with the quark and gluon structure functions of Duke and Owens¹¹⁾ ($\Lambda_{QCD} \sim 0.3$ GeV) and the supersymmetric cross-sections of Ref. 12). This means that we have not made a premature effort to "fit" the "data".

In Table 2 we give our results, as functions of $m(\tilde{g})$ and $m(\tilde{q})$, and with the experimentally defined cuts, for the total cross-section in nanobarns for the missing p_T plus (at least one) jet events and (in parenthesis) for the ratio of dijets to monojets. A cross-section of 0.1 nb would correspond to ~ 10 events expected in the published UA1 data. Two "corners" of Table 2 seem to be acceptable. The first is $m(\tilde{g}) = 3$ GeV, $m(\tilde{q}) = 80$ GeV where both the cross-section and the dijet/monojet ratio are satisfactory. The second is $m(\tilde{g}) = 5$ to 7 GeV, $m(\tilde{q}) > 150$ GeV, where the dijet/monojet ratio is somewhat uncomfortable. The second corner actually corresponds to a model for which squarks are too heavy to be produced in significant quantities, only the $\tilde{g}\tilde{g}$ channel survives.

$m(\tilde{q}) \backslash m(\tilde{g})$	3	5	7
80	0.13 (0.29)	0.20 (0.35)	0.27 (0.35)
100	0.09 (0.86)	0.16 (0.87)	0.24 (0.86)
120	0.075 (1.2)	0.14 (1.0)	0.20 (0.80)
≥ 150	0.055 (1.2)	0.10 (0.75)	0.15 (0.55)

Table 2: Madrid model predictions for the total production cross-section (in nb) for events with missing p_T above the experimental cuts and one or more jets. Also given (in parentheses) are the expected ratios of dijets to monojets.

In Table 3 we give results for the ratio of monojets from the $\tilde{q}\tilde{g}$ channel, to those from the $\tilde{g}\tilde{g}$ channel. The numbers in parentheses are our estimates for the average invariant mass of the monojets, $\langle W \rangle$, with all channels in the Madrid scenario taken into account. To the extent that one may expect $\langle W \rangle \sim 2\langle m_{CH} \rangle$ only the corner $m(\tilde{g}) < 5$ GeV, $m(\tilde{q}) < 100$ GeV seems to be acceptable. In particular the possibility of dominance of the $\tilde{g}\tilde{g}$ channel, $m(\tilde{q}) > 150$ GeV, seems to be excluded for all gluino masses in the range we study. It must be remembered, however, that the definition of a jet invariant mass is a fairly tricky enterprise, both theoretically and experimentally. In particular, the association of an extra soft track to a jet may very significantly affect its measured mass.

Unlike in other supersymmetric scenarios, our results for jet invariant masses are not in violent disagreement with observation, even in the case where the $\tilde{g}\tilde{g}$ channel dominates. The exclusion of this possibility should therefore be taken cum grano salis.

$m(\tilde{g}) \backslash m(\tilde{q})$	3	5	7
80	3.2 (3.9)	1.4 (5.0)	1.0 (6)
100	1.0 (6.0)	0.48 (7.6)	0.30 (8.9)
120	0.32 (8.5)	0.13 (10)	0.09 (11.3)
>150	0 (9.9)	0 (10.9)	0 (11.9)

Table 3: Ratio of the number of monojets from the $(\tilde{q}\tilde{g})$ channel to the corresponding number from the $(\tilde{g}\tilde{g})$ channel. Also shown (in parentheses) are the total (charged plus neutral) expected average widths of the monojets, in GeV.

To summarize, two possibilities survive the comparison of the Madrid scenario with the experimental indications. One corresponds to masses $3 \text{ GeV} < m(\tilde{g}) < 5 \text{ GeV}$, $80 \text{ GeV} < m(\tilde{q}) < 100 \text{ GeV}$, wherein all properties of the data are nicely reproduced. For the other allowed domain $5 \text{ GeV} < m(\tilde{g}) < 9 \text{ GeV}$, $m(\tilde{q}) \gtrsim 140 \text{ GeV}$, the invariant mass of the observed monojets is narrower than the theoretical expectation by a factor $\sim 0(3)$.

In the next two Chapters we develop in unabashed detail the arguments presented in this Introduction. Our considerations are based on analytic approximations, rather than black box Monte-Carlo's. This means we have not tried to pronounce the last word on the Madrid scenario, but rather many comprehensible ones. The last word is to the experimentalists to utter in unison.

2. - GLUINO PAIR PRODUCTION

Gluinos are colour-octet massive spinors. This characterization suffices to estimate the shape of their fragmentation function, with guidance from our experimental knowledge and theoretical understanding of the fragmentation functions of gluons and massive quarks. Let z be the ratio of longitudinal momenta between the final photino and the original gluino, as in Fig. 1b. In a collinear approximation adequate to the description of light gluinos produced at high energy [$\langle s \rangle \gg m^2(\tilde{g})$], the z distribution function $D(z, s)$ is the convolution of three functions: $D = D_1 * D_2 * D_3$. The first of them is a perturbative fragmentation function, D_1 , describing gluon emission by the originally produced gluino. The second is a non-perturbative "dressing" function, D_2 , that reflects the conversion of the gluino into a gluino-containing hadron, labelled $\tilde{G}$ in Fig. 1b. The third function, D_3 , describes $\tilde{g} \rightarrow \tilde{\gamma}$ decay. For D_1 we take the form

$$D_1(x, \hat{s}) = A (1-x)^{A-1} \quad (2.1a)$$

$$A = \frac{36}{25} \ln \left\{ \ln(\hat{s}/\Lambda^2) / \ln(4m^2(\tilde{g})/\Lambda^2) \right\} \quad (2.1b)$$

which is a simple but fairly precise description (in the sense of the first few x -moments) of the "leading-log" fragmentation of a colour-octet particle. Our results will be given for $\Lambda = 0.3$ GeV. For D_2 we have three the ansätze, that lead to very similar results on the photino energy distribution

$$D_2(x) = \mathcal{N} x [x-1 - \epsilon x/(1-x)]^{-2} \quad (2.2a)$$

$$D_2(x) = \delta(x - x_0) \quad (2.2b)$$

$$D_2(x) = (2x_0 - 1)^{-1} \theta(x + 1 - 2x_0) \quad (2.2c)$$

$$\epsilon \equiv [0.6 \text{ GeV}/m(\tilde{g})]^2; \quad x_0 \equiv m(\tilde{g})/[m(\tilde{g}) + 2m(q)] \quad (2.2d,e)$$

Here $m(q) = 0.3$ GeV is a light quark "constituent mass" and the first expression for D_2 is that advocated by Peterson et al.¹³⁾. Our form of a gluino fragmentation function, $D_1 * D_2$, is presumably realistic since it closely paraphrases, but for a rescaling of colour charges, a similar calculation of charmed and bottom quark fragmentation that correctly reproduces the data¹⁴⁾. Finally, the $\tilde{g} \rightarrow \tilde{\gamma}$ decay function, D_3 , can be explicitly computed. In the approximation $m^2(\tilde{q}) \gg m^2(\tilde{g})$, $m^2(q) \ll m^2(\tilde{g})$, the result is

$$D_3(x) = \frac{1}{3} (5 - 9x^2 + 5x^3) \quad (2.3)$$

The fractional energy distribution of photinos is given by the triple convolution

$$D(\bar{z}, \hat{s}) = D_1 * D_2 * D_3 = \int_0^1 dx_1 dx_2 dx_3 \delta(\bar{z} - x_1 x_2 x_3) D_1(x_1) D_2(x_2) D_3(x_3) \quad (2.4)$$

Let z_1 and z_2 , as in Fig. 1a, be the momentum fractions of the two photinos in the process $p\bar{p} \rightarrow \tilde{g}\tilde{g} \dots \rightarrow \tilde{\gamma}\tilde{\gamma} \dots$. The elementary processes that contribute to gluino pair production are $q\bar{q} \rightarrow \tilde{g}\tilde{g}$ and $g\bar{g} \rightarrow \tilde{g}\tilde{g}$. Let $d\sigma_i/d\cos\theta$ ($i = q, g$) be the corresponding computed c.m.s. cross-sections at an invariant squared energy $\hat{s} = sx_1 x_2$, with x , the momentum fraction of the original partons. The missing transverse momentum distribution is simply

$$\frac{d\sigma}{dp_T} = \frac{2x_T}{\sqrt{s}} \int_{x_T}^{2-x_T} d\bar{z}_+ \int_{-\cos\bar{\theta}}^{\cos\bar{\theta}} \frac{d\cos\theta}{\sin^2\theta} \int_H^B d\bar{z}_- \bar{z}_-^2 D(\bar{z}_1, \hat{s}) D(\bar{z}_2, \hat{s}) \frac{\sum_i d\sigma_{ii}(\hat{s}=s\tau)}{d\cos\theta} \quad (2.5a)$$

$$* \int_{\tau}^1 dx_2 x_2^{-1} F_i(x_2, \hat{s}) F_i\left(\frac{\tau}{x_2}, \hat{s}\right)$$

$$\left[p_T = (\sqrt{\hat{s}}/2) \sin\theta (z_1 - z_2) \right] \quad (2.5b)$$

$$x_T = 2p_T/\sqrt{s} \quad (2.5c)$$

$$\bar{z}_{+,-} = \bar{z}_1 \pm \bar{z}_2 \quad (2.5d)$$

$$\tau = (x_T/\bar{z}_- \sin\theta)^2 \quad (2.5e)$$

$$H = x_T/\sin\theta \quad (2.5f)$$

$$B = \min \{ 2 - \bar{z}_+, \bar{z}_+, x_T\sqrt{s}/(2m(\tilde{g})\sin\theta) \} \quad (2.5g)$$

$$\cos\bar{\theta} = (1 - x_T^2/[\min(\bar{z}_+, 2 - \bar{z}_+)]^2)^{1/2} \quad (2.5h)$$

where F_i are the q and g proton structure functions¹¹⁾. To ascertain the effect of gluino fragmentation on the "missing p_T signal" induced by the $\tilde{g}\tilde{g}$ channel we have proceeded as follows. First, with the benediction of our UA(1) colleagues¹⁵⁾, we have substituted their elaborate missing p_T cut by an effective cut $p_T > 28$ GeV, that adequately reflects the more sophisticated analysis. Second, we have computed the ratio R between the numbers of events above the cut with fragmentation taken into account, and the corresponding number with no fragmentation, $D_1 * D_2 = \delta(1-x)$. Third, we have multiplied the total missing p_T signal for this channel as computed in Ref. 5) (in a no-fragmentation model with allowance for experimental cuts) by R, to obtain an estimate for the signal as correctly computed with gluino fragmentation. The results are shown in Fig. 2 and labelled "NO FRAGMENTATION"⁵⁾ and "FRAGMENTATION". Fragmentation reduces the missing p_T signal by as much as an order of magnitude at low gluino mass. Finally, we have extracted from the total missing p_T signal those events that would count as monojets and dijets. "Monojets" are events where the hadronic debris of one of the gluinos have $E_T > E_T^0 = 25$ GeV [$\sqrt{s} \sin\theta(1-z)/2 > E_T^0$], "dijets" are events where the hadronic debris of the second gluino have $E_T > E_T' = 12$ GeV. These are not the complete set of constraints used by UA1 in their definitions of jets, since they also demand that the particles in a jet be located within a certain "cone". This cone of approximately 30° to 45° aperture in geometrical opening angle is wide enough for the energetic gluino debris to fall within it in more than 70% of the cases, for $m(\tilde{g}) < 10$ GeV. We have neglected these small corrections in our results as plotted in Fig. 2. As we announced in the introduction, the same numerology that implies a bound $m(\tilde{g}) \gtrsim 40$ GeV when fragmentation is forgotten, now allows for a second "window" at $m(\tilde{g}) \gtrsim 7$ GeV.

In the rest of this Chapter, we discuss the question of the width of jets in the $\tilde{g}\tilde{g}$ channel. In four out of the six missing p_T plus monojet events reported in Ref. 1), the invariant mass m_{CH} of the tracks in the jet is measured. One of these jets contains a single charged track, and for the others $m_{CH} = 0.79 \pm 0.12$, 3.14 ± 0.38 , 0.517 ± 0.06 GeV. This pattern of fairly low invariant masses, if significant, disagrees with the expectations for gluino pair production, even for very light gluinos. A complete QCD estimate, with experimental cuts, of the jet invariant masses in gluino pair production is a difficult task, involving the calculation of gluon emission by the original, intermediary and final partons. But we know that in QCD perturbation theory, the dominant contributions to the invariant mass distribution of a parton jet (gluino or otherwise) originate in

the emission of gluons by this very parton, a process favoured by singular propagator effects. This dominance is particularly overwhelming when a cut in jet angular spread is applied: gluons emitted by other partons in the process are not correlated with the parton-jet direction. This implies that a good estimate of jet properties can be obtained by analyzing a simpler process than $\tilde{g}\tilde{g}$ production in $p\bar{p}$ collisions, for instance the production of gluino pairs by a colourless vector source. Let such a vector source, at rest, produce gluino pairs at an invariant mass Q , comparable to or bigger than the transverse mass of the UA1 monojet events. We shall estimate the invariant mass of the jets produced by this source, as a function of $m(\tilde{g})$, with an "experimental cut" $\theta < \theta_m$ in the opening angle between the parton jets in the process. The experimental UA1 jet opening cut is not just an angular cut, it varies with pseudorapidity, it is not applied in the centre-of-mass system of the colliding partons, etc. None of this is too worrisome, since our conclusions shall not dramatically depend on θ_m in the generous range $30^\circ < \theta_m < 60^\circ$. To zeroth order in α_s , our gluino source decays into two back-to-back gluinos, as in Fig. 4a. Define $z = 2\tilde{\gamma} \cdot \tilde{g} / m^2(\tilde{g})$, a measure of the photino momentum in gluino decay, with particle four-momenta denoted by the corresponding particle's name. The invariant mass of the $\bar{q}q$ pair in $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}(z)$ decay is $W = m(\tilde{g})(1-z)^{1/2}$. But, in our parton language, only quarks whose relative angle is less than θ_m contribute to a measured $W > 0$. Let $F(z)$ be the fraction of quark pairs that materialize within an angle $\theta < \theta_m$:

$$F(z=0) ; \beta^2(q\bar{q}) \equiv 1 - \frac{4m^2(\tilde{g})(1-z)}{Q^2(1-z/2)} < \frac{1+\cos\theta_m}{2} \quad (2.6a)$$

$$F(z) \simeq \left[1 - \frac{[1 - \beta^2(q\bar{q})][1+\cos\theta_m]^2}{\beta^2(q\bar{q})\sin^2\theta_m} \right]^{1/2} ; \beta^2(q\bar{q}) > \frac{1+\cos\theta_m}{2} \quad (2.6b)$$

where $\beta(q\bar{q})$ is the velocity of the $\bar{q}q$ system at its average energy, with the pair of quarks considered as a single entity of mass W . For $Q \sim 100$ GeV and $m(\tilde{g}) < 7$ GeV, the condition on $\beta(q\bar{q})$ on the second of the above equations is amply satisfied. The fraction $F(0)$ is essentially unity: it ranges from 0.97 to 0.85 at $\theta_m = 30^\circ$, and from 0.99 to 0.97 at $\theta_m = 60^\circ$, as $m(\tilde{g})$ varies from 3 to 7 GeV at $Q = 100$ GeV. The function $F(z)$ for $z > 0$ is even closer to unity than $F(0)$. In the very good approximation $F(z) = 1$, the jet invariant mass distribution to zeroth order in α_s is

$$\frac{d\Gamma^{(0)}(z)}{\Gamma dW^2} \simeq \delta(W^2 - m^2(\tilde{g})(1-z)) \quad (2.7)$$

This equation means that the gluino hadronic decay products fully contribute to the jet invariant mass, not a particularly surprising result for light gluinos produced at high energy.

We proceed to discuss the $O(\alpha_s)$ corrections to Eq. (2.7). To first order in α_s the relevant process, depicted in Fig. 4b, is $Q \rightarrow \tilde{g}g$, where Q stands for our colourless vector source of mass Q , followed by gluino-decay $\tilde{g} \rightarrow q\bar{q}\tilde{\gamma}(z)$. The jet invariant mass, W , is now that of the $q\bar{q}g$ ensemble. The matrix element singularly favours fast gluinos, meaning that once again the $q\bar{q}$ gluino decay products will be well collimated within an angle $< 30^\circ$. For a jet opening angle $\theta_m > 30^\circ$ it is a good approximation to take the $q\bar{q}$ pair to be collinear with their parent gluino in combining their four-momentum with that of the gluon to compute W . Consider first the case $z = 0$, in which the photino carries away no momentum and the $(q\bar{q})$ pair carries all of the parent gluino four-momentum. The $q\bar{q}g$ invariant mass is that of the $\tilde{g}g$ pair. The invariant mass distribution is identical to the corresponding distribution in $q\bar{q}g$ production in e^+e^- annihilation, except for a rescaling of triplet to octet colour charge. In this way we obtain

$$f(\omega) \equiv \frac{Q^2 d\Gamma^{(1)}(z=0)}{\Gamma dW^2} = \frac{3\alpha_s}{2\pi\omega} [I(x_+) - I(x_-)] \theta(\omega_m - \omega) \left[1 + O\left(\frac{m(\tilde{g})}{Q}\right) \right] \quad (2.8a)$$

$$\omega \equiv W^2/Q^2 \quad (2.8b)$$

$$0 < \omega < \omega_m = \frac{3 + \cos\theta_m - 2\sqrt{2}(1 + \cos\theta_m)}{1 - \cos\theta_m} \quad (2.8c)$$

$$I(x) = -[1+(1-\omega)^2] \ln(1-x) - x - x^2/2 \quad (2.8d)$$

$$x_{+,-} = \frac{1}{2} \left[1+\omega \pm \sqrt{(1+\omega)^2 - 8\omega/(1+\cos\theta_m)} \right] \quad (2.8e)$$

Consider next the case $z \neq 0$. The $(\bar{q}q)$ pair now carries on average a fraction $1 - \frac{1}{2}z$ of the momentum of their fast parent gluino. The square of the invariant mass of the $gq\bar{q}$ pair is reduced by the same fraction, relative to what it would have been at $z = 0$. We can therefore obtain the distribution of invariant masses at $z \neq 0$ from a simple change of variables, $W^2 \rightarrow W^2(1-z/2)$, in the $z = 0$ results of Eqs. (2.8), to wit

$$\frac{Q^2 d\Gamma^{(1)}(z)}{\Gamma dW^2} = \frac{2}{2-z} \rho \left[\frac{2}{(2-z)} \frac{W^2}{m^2(\tilde{q})} \right] \quad (2.9)$$

From Eqs. (2.7) and (2.9), we compute the average invariant mass of a gluino jet. In the range $30^\circ < \theta_m < 60^\circ$, the result can be well fit by the simple expression

$$\langle W(z) \rangle = m(\tilde{q}) \sqrt{1-z} + \frac{\alpha_s}{t_c} Q \sqrt{1-\frac{z}{2}} F(\theta_m) \quad (2.10c)$$

$$F(\theta_m) = 0.885 + 6.15 (1 - \cos\theta_m) ; \quad 30^\circ < \theta_m < 60^\circ \quad (2.10b)$$

The last step is to average over the momentum distribution of photinos in $\tilde{g} \rightarrow q\bar{q}\gamma$ decay

$$\frac{1}{\Gamma} \frac{d\Gamma}{dz} = 2z^2(3-2z) \quad (2.11)$$

We obtain^{*)}

$$\langle W \rangle \sim \frac{32}{63} m(\tilde{g}) + 0.804 \frac{\alpha_s}{\pi} Q F(\theta_m) \quad (2.12a)$$

$$\langle W \rangle \sim \frac{m(\tilde{g})}{2} + Q * \begin{cases} 0.07 & , \theta_m = 30^\circ \\ 0.11 & , \theta_m = 45^\circ \\ 0.16 & , \theta_m = 60^\circ \end{cases} \quad (2.12b)$$

For $\theta = 30^\circ$, $Q = 100$ GeV, the average invariant mass of a gluino jet is 7 GeV plus half the gluino mass. If $\langle m_{CH} \rangle \sim \frac{1}{2} \langle W \rangle$ is a measure of the expected "charged" mass of a jet, this result disagrees with the experimental indications, even for $m(\tilde{g})$ as low as 3 GeV, for which we expect $\langle m_{CH} \rangle \sim 4.25$ GeV. There is a more model independent way to state this result. The second term in Eq. (2.12),

^{*)} A result such as Eq. (2.12) for an average invariant mass with a jet angular cut is "infra-red safe" in the usual QCD sense. Yet, in deriving it, we have ascribed to the decaying gluinos a certain fraction of the energy of the source Q . In the leading $O(\alpha_s)$ term, for instance, the gluino energy was taken to be $E(\tilde{g}) = Q/2$ in the derivation of Eqs. (2.6). The simultaneous fixing of an invariant mass to the hadron debris and a fraction of energy to the gluino does not lead to an infra-red safe expression for the corresponding differential cross-section: the singularities must be resummed (i.e., in the leading log approximation) before they begin to make sense. What the above abstruse statements mean in a more physical language is the following. Only gluinos that have emitted fairly energetic and non-collinear gluons contribute to the $O(\alpha_s Q)$ term in Eq. (2.12). But in a complete calculation the momentum fraction carried by the gluino in the first term of Eq. (2.12) would also be affected by the emission of soft and collinear gluons, the ones that enter into the calculation of the gluino fragmentation function and lower its average momentum well below the maximum possible one: the real gluino containing hadron in a real jet scarcely ever carries all the available momentum, even if the jet mass is close to the gluino mass. This means that we ought to have reduced the energy of the gluino in the calculation of Eqs. (2.6) by x ($Q/2 \rightarrow xQ/2$) and averaged over the gluino fragmentation function $D(x)$. In principle this reduces the fraction F of gluino decay products within θ_m , in practice the reduction is negligible given that z distribution of Eq. (2.6) is so peaked near $z = 1$. We stand by Eq. (2.12).

divided by $\langle 1 - z/2 \rangle = 0.804$, is also the theoretically expected $\langle W \rangle$ of a gluon jet. In unbiased events, gluon jets dominate at $p_T \sim 0(30 \text{ GeV})$. Light gluino jets are therefore expected to have an average mass which is 80% that of a conventional jet, plus $m(\tilde{g})/2$, the contribution of the gluino decay products.

3. - GLUINO-SQUARK PRODUCTION

If gluinos are light, the first threshold for the strong production of superhadrons is $\tilde{g}\tilde{g}$, the next is $\tilde{g}\tilde{q}$. Squarks have not been observed in e^+e^- annihilation experiments, and should have masses $m(\tilde{q}) \gtrsim 20 \text{ GeV}$. A $\tilde{q}\tilde{g}$ pair decays into states containing photinos and is a potential source of missing p_T events:

$$p\bar{p} \rightarrow \tilde{q}\tilde{q} \quad \left\{ \begin{array}{l} \tilde{q} \rightarrow q\bar{q}\tilde{\gamma} \\ \tilde{q} \rightarrow \left\{ \begin{array}{l} q\tilde{\gamma} \\ q\tilde{q} \end{array} \right. , \tilde{g} \rightarrow q\bar{q}\tilde{\gamma} \end{array} \right.$$

Of all these branches, $\tilde{q} \rightarrow q\tilde{\gamma}$ is, for relatively heavy squarks ($M \sim 100 \text{ GeV}$), the cleanest potential source of monojet plus missing p_T events. A variety of predictions⁷⁾ for this process is shown in Fig. 3, with the gluino mass fixed at 3 GeV. The predictions are based in a calculation that paraphrases, at the parton level, the experimental cuts. The assumed q and g structure functions are those of Duke and Owens¹¹⁾ and the cross-sections those of Ref. 12). Since $m(\tilde{g}) \ll m(\tilde{q})$, a doubling of the gluino mass affects the predictions of Fig. 3 by factors smaller than their theoretical uncertainty (at least a factor of two, originating in the neglect of higher order QCD corrections and in our relative ignorance of the gluon structure function). A look at Fig. 3 indicates that, if $\tilde{q}\tilde{g}$ production followed by $\tilde{q} \rightarrow q\tilde{\gamma}$ decay were the only relevant process, it would provide a fair description of the data for $m(\tilde{q}) = 80$ to 120 GeV , $m(\tilde{g}) \lesssim 7 \text{ GeV}$ (the range allowed by our previous analysis of the $\tilde{g}\tilde{g}$ channel). Our gluinos, unlike those of Barger et al., rapidly decay into "missing" photinos: before we reach any conclusions, we must do some more work. On our way to the study of the dominant $\tilde{q} \rightarrow q\tilde{\gamma}$ branch, we pause for two extra paragraphs concerning the $\tilde{q} \rightarrow q\tilde{\gamma}$ channel.

It is easy to give a simple but accurate estimate of the effect of experimental cuts on the total cross-section for missing p_T plus monojet events in the $\tilde{q} \rightarrow q\tilde{\gamma}$ channel. The effective missing p_T cut is $p_T > p_T^0 \sim 28 \text{ GeV}$, the monojet

cut is $E_T > E_T^0 \sim 25 \text{ GeV} < p_T^0$. The very transverse and energetic jet in this channel is a quark jet, very unlikely to fail the generous experimentally defined cut on jet opening angle. For the masses we are considering, the $\tilde{q}\tilde{g}$ pairs are produced with relatively small transverse momentum (smaller than 12 GeV in $\sim 80\%$ of the cases). Squarks are scalars. The cuts are transverse. All of the above implies that the ratio of the cross-section with cuts to the theoretical uncut cross-section is simply

$$R = \frac{\sigma(\tilde{q} \rightarrow \tilde{g}, \text{CUT})}{\sigma(\tilde{q} \rightarrow \tilde{g}, \text{UNCUT})} \sim \int_0^1 d\cos\theta \theta \left[\sin\theta - \frac{2\max(p_T^0, E_T^0)}{m(\tilde{q})} \right] = \left[1 - \left(\frac{2p_T^0}{m(\tilde{q})} \right)^2 \right]^{1/2} \quad (3.1)$$

while the p_T distribution shows the characteristic Jacobian peak of Fig. 3a.

We expect the squarks produced in the $gq \rightarrow \tilde{g}\tilde{q}$ process to be mainly the roughly degenerate pair $\tilde{u}$ and $\tilde{d}$ (or their antiparticles). In $\bar{p}p$ collisions the ups should be about twice as copious as the downs. The average branching ratio for squark decay into photinos is thus

$$B = \frac{\langle \tilde{q} \rightarrow q\tilde{\gamma} \rangle}{(\tilde{q} \rightarrow q\tilde{g})} \sim \frac{1}{3} \left[2\left(\frac{2}{3}\right)^2 + \left(\frac{1}{3}\right)^2 \right] \frac{3\alpha}{4\alpha_s} \sim \frac{1}{67} \quad (3.2)$$

The decay $\tilde{q} \rightarrow \tilde{g}q$ is some 67 times more frequent than the decay $\tilde{q} \rightarrow q\tilde{\gamma}$. The chain $\tilde{q} \rightarrow q\tilde{g}$, $\tilde{g} \rightarrow \tilde{\gamma}q\bar{q}$, depicted in Fig. 1d, is a potential source of missing p_T plus multijet events. Fragmentation of the gluino prior to its decay slows down the photino, and more than compensates (above the experimental missing p_T cut) for the favourable branching ratio of squark decay into a gluino. Let z be the ratio of the final photino energy to the original gluino energy $z \approx 2E(\tilde{\gamma})/m(\tilde{q})$. In this definition, we have ignored the fact that the squark may not be at rest, but moving forward or backward close to the beam axis. Not to worry: longitudinal boosts are innocuous, since all cuts to be considered are transverse (the jet opening angle cuts are practically ineffective for the light gluinos under

consideration). Let $D(z)$ be the photino energy distribution, with D as in Eq. (2.4). Let θ be the decay angle, relative to the beam axis, of the squark (at rest) into a gluino and a quark jet. The missing p_T cut is

$$p_T(\tilde{\gamma}) = \frac{m(\tilde{q})}{2} z \sin \theta > p_T^0 \quad (3.3)$$

The (quark) monojet trigger condition is

$$p_T(q) = \frac{m(\tilde{q})}{2} \sin \theta > E_T^0 \quad (3.4)$$

which is always less restrictive than the missing p_T cut for $p_T^0 > E_T^0$ (a monojet should be seen if the photino passes its test). The condition for the gluino hadronic debris to be ascertained as a second jet is

$$p_T(q\bar{q} + g's) = \frac{m(\tilde{q})}{2} (1-z) \sin \theta > E_T^1 \quad (3.5)$$

The total fraction, T , of $\tilde{q} \rightarrow q\tilde{g}$ decays that pass the missing p_T plus mono or dijet cuts, and the fraction, D , of events that appears as dijets, can now be simply estimated. They are:

$$T = \int_a^1 D(z, m^2(\tilde{q})) \left(1 - \frac{a^2}{z^2}\right)^{1/2} dz \quad (3.6a)$$

$$D = \int_a^{z_1} dz D(z, m^2(\tilde{q})) \left(1 - \frac{a^2}{z^2}\right)^{1/2} + \int_{z_1}^{1-b} dz D(z, m^2(\tilde{q})) \left(1 - \frac{b^2}{(1-z)^2}\right)^{1/2} \quad (3.6b)$$

$$a \equiv \frac{2p_T^0}{m(\tilde{q})} ; \quad b \equiv \frac{2E_T^1}{m(\tilde{q})} ; \quad z_1 \equiv \frac{a}{a+b} \quad (3.6c,d,e)$$

where we have exploited the fact that $p_T^0 > E_T^0$.

With all of the above artillery, we may scale our predictions for monojets and dijets in the sequence $\bar{p}p \rightarrow \tilde{g}\tilde{q}$, $\tilde{q} \rightarrow q\tilde{g}$, $\tilde{g} \rightarrow \tilde{\gamma}q\bar{q}$, to the more elaborate calculations of Barger et al. for the $\tilde{q} \rightarrow q\tilde{\gamma}$ channel. Let "Monojet ($\tilde{q} \rightarrow q\tilde{\gamma}$)" refer to their predicted total number of events above a fixed (or variable) missing transverse energy cut, p_T^0 . The number of monojets from the $\tilde{q} \rightarrow q\tilde{g}$ channel is then

$$\text{Monojets}(\tilde{q} \rightarrow q\tilde{g}) = \text{Monojets}(\tilde{q} \rightarrow q\tilde{\gamma}) \frac{(T-D)}{BR} \quad (3.7)$$

with R, B, T and D as in Eqs. (3.1), (3.2) and (3.6), respectively. Similarly, the number of dijets is

$$\text{Dijets}(\tilde{q} \rightarrow q\tilde{g}) = \text{Monojets}(\tilde{q} \rightarrow q\tilde{\gamma}) \frac{D}{BR} \quad (3.8)$$

Equations (3.7) and (3.8) are the ones we have used in producing the numbers of Table 1 of Section 1.

If and when the missing p_T signals become more abundant, one may be interested in the detailed shapes of their p_T distributions, not only in the integrated quantities we have just discussed. For the $\tilde{q} \rightarrow q\tilde{\gamma}$ channel this distribution is dominantly a monojet signal with the shape

$$\frac{1}{\sigma_{TOT}} \frac{d\sigma(\tilde{q} \rightarrow q\tilde{\gamma})}{dp_T} \simeq \frac{2\Gamma B p_T}{\pi} \int \frac{dW \theta(p_T - E_T^0)}{(m(\tilde{q}) - W)^2 + \Gamma^2/4} \cdot \frac{1}{W \sqrt{W^2 - 4p_T^2}} \quad (3.9)$$

where σ_{TOT} is the total cross-section for the production of $\tilde{q}\tilde{g}$ pairs, $\Gamma \simeq (2/3)\alpha_s m(\tilde{q})$ is the squark total width, and all other quantities have been previously defined. The dijet (D) signal from the $\tilde{q} \rightarrow q\tilde{g}$ channel has the shape

$$\frac{1}{\sigma_{TOT}} \frac{d\sigma^D(\tilde{q} \rightarrow q\tilde{g})}{dp_T} \approx \frac{2\Gamma p_T}{\pi} \int \frac{dW \theta(1-\alpha-\beta)}{(m(\tilde{q})-W)^2 + \Gamma^2/4} \cdot \frac{1}{W^2} \int_{\alpha}^{\delta} \frac{dz D(z)}{z \sqrt{z^2 - \alpha^2}} \quad (3.10)$$

$$\alpha \equiv \frac{2p_T}{W}; \quad \beta = \frac{2E_T^1}{W}; \quad \gamma \equiv \frac{2E_T^0}{W}; \quad \delta = \min \left\{ \frac{\alpha}{\gamma}, \frac{\alpha}{\alpha+\beta} \right\} \quad (3.10b,c,d,e)$$

Here $D(z)$ is the photino distribution of Eq. (2.4). To obtain the total monojet plus dijet signal in the $\tilde{q} \rightarrow q\tilde{g}$ channel, substitute in Eq. (3.10) δ by $\min[1, \alpha/\gamma]$ and the step function by $\theta(1-\alpha)$. In Eqs. (3.9)-(3.10), we have taken into account the p_T -smearing induced by the width of the squark and neglected the smaller effect of the p_T distribution of the produced squarks. The mono and dijet signals from the $\tilde{q} \rightarrow q\tilde{g}$ channel fall steeply with p_T , since $D(z)$ is a very rapidly falling function of its argument. For sufficiently large p_T a source of dijets that we have so far neglected takes over. This is the production of a $\tilde{q}\tilde{g}$ pair, followed by $\tilde{q} \rightarrow q\tilde{\gamma}$ decay, with the original gluino produced at a sufficient transverse momentum to have its decay products pass the "second jet" condition $E_T > E_T^1 \sim 12$ GeV. The $\tilde{q}\tilde{g}$ pairs are typically produced at an invariant energy not too close to its threshold value. This means that the effects of gluino fragmentation are important and the hadron debris of the gluino carry most of the gluino energy. The effective second jet condition is approximately given by $q_T(\tilde{g}) > E_T^1$. The photino in $\tilde{g} \rightarrow \tilde{\gamma}$ decay carries a very small fraction of the gluino energy; its contribution to the missing p_T is negligible, compared to that of the other photino $\tilde{q} \rightarrow q\tilde{\gamma}$. With the above information one may easily estimate the dijet signal accompanying $\tilde{q} \rightarrow q\tilde{\gamma}$ decay:

$$\frac{1}{\sigma_{TOT}} \frac{d\sigma^D}{dp_T} \approx \frac{B\Gamma p_T}{2\pi^2 \sigma_{TOT}} \int \frac{dW}{(m(\tilde{q})-W)^2 + \Gamma^2/4} \left\{ \int_0^{\cos\theta_1} d\cos\theta \int_{\alpha_0}^{|W\sin\theta|} dq_T + \int_{\cos\theta_1}^{\cos\theta_2} d\cos\theta \int_{\alpha}^{|W\sin\theta|} dq_T \right\} \frac{d\sigma}{\sqrt{H} dq_T} \quad (3.11a)$$

$$H \equiv \left[\frac{W q_T \sin\theta}{2} \right]^2 - \left[p_T^2 - \left(\frac{q_T}{2} \right)^2 - \left(\frac{W \sin\theta}{2} \right)^2 \right]^2 \quad (3.11b)$$

$$\alpha_0 = \max\{|2p_T - |W\sin\theta||, E_T^1\} ; \alpha = \max\{\alpha_0, \sqrt{2(E_T^0)^2 + 2p_T^2 - |W\sin\theta|}\} \quad (3.11c,d)$$

$$\cos\theta_1 = \sqrt{1 - \frac{2(E_T^0)^2 + 2p_T^2}{M^2}} ; \cos\theta_2 = \sqrt{1 - \frac{E_T^2 + p_T^2}{M^2}} \quad (3.11e,f)$$

In the above expression $d\sigma/dq_T$ is the $\bar{p}p \rightarrow g\tilde{q}$ cross-section at a fixed transverse momentum of the outgoing partons, expressible in terms of the elementary $gg \rightarrow g\tilde{q}$ differential cross-section by the standard methods. To obtain Eq. (3.11) it is simplest to think in the reference system wherein the squark has no longitudinal momentum. The total monojet plus dijet signal is simply obtained by setting E_T^1 to zero in Eq. (3.11b). The resulting expression is simply a version of Eq. (3.9), corrected for q_T -smearing. We shall discuss all of the above p_T distributions in our concluding chapter.

We now turn to the question of the width to be expected for the monojets in $\tilde{q}$ decay. Table 1 indicates that these monojets almost entirely arise, in our domain of parameters, from the "clean" decay $\tilde{q} \rightarrow q\tilde{\gamma}$. At their fairly large mass, squarks are dominantly produced close to their mass-shell or, in other words, fragmentation is not a crucial issue to their fate. This means that the invariant mass distribution of the monojets in the decay $\tilde{q} \rightarrow q\tilde{\gamma}$ is approximately the one expected for the decay of an on-shell squark. We proceed to estimate this mass distribution to first order in QCD. The relevant process is one-gluon emission, $\tilde{q} \rightarrow q\tilde{\gamma}g$, depicted in Fig. 5, for which the differential decay rate is¹⁶⁾

$$\frac{d\Gamma(\tilde{q} \rightarrow qg\tilde{\gamma})}{\Gamma(\tilde{q} \rightarrow q\tilde{\gamma}) dx dy} = \frac{2\alpha_s}{3\pi} \frac{1-x}{1-y} \frac{1+(1-z)^2}{z^2} \quad (3.12a)$$

$$x \equiv \frac{2\tilde{q} \cdot q}{m^2(\tilde{q})} ; y \equiv \frac{2\tilde{q} \cdot \tilde{\gamma}}{m^2(\tilde{q})} ; z \equiv \frac{2\tilde{q} \cdot g}{m^2(\tilde{q})} \quad (3.12b)$$

where four momenta have been denoted by the corresponding particle masses. The invariant mass of the hadronic debris is $W^2 = m^2(\tilde{q})(1-y)$. The distribution $d\Gamma/dW$ is affected by experimental cuts, it would be misleading to compute it directly from Eq. (3.12) without cuts. First of all only fast photinos (y not close to

zero) pass the missing p_T cut. This biases the events towards large y : small W^2 . This bias turns out to be unimportant given the existence of a second stronger bias towards small W^2 : the one induced by the maximum angular spread in the definition of a jet. Let θ_m be maximum angle of two particles (in our approximation, a quark and a gluon) that may coalesce into a single jet. The invariant mass distribution in the decay of a squark at rest with the condition $\theta(gq) < \theta_{\max}$, can be computed from Eq. (3.12) to be

$$\frac{d\Gamma}{\Gamma dW^2} = \left(1 + O\left(\frac{\alpha_s}{\pi}\right)\right) \delta(W^2) + \frac{2\alpha_s}{3\pi W^2} \left[I(x_+) - I(x_-) \right] \theta(\omega_m - \omega) \quad (3.13a)$$

$$\omega \equiv W^2 / m^2(\tilde{q}) \quad (3.13b)$$

$$I(x) = \frac{2\omega}{x} + 2(1+\omega) \ln x - (2+\omega)x + \frac{x^2}{2} \quad (3.13c)$$

where ω_m and $x_{+,-}$ are defined as in Eqs. (2.8c,e). The average invariant mass corresponding to Eqs. (3.12) can be fit, in the range $25^\circ < \theta_m < 55^\circ$ by the simple expression

$$\langle W \rangle \approx \frac{\alpha_s}{\pi} m(\tilde{q}) \left[1.492 + 1.528 \cos \theta_m - 2.611 \cos^2 \theta_m \right] \quad (3.14)$$

For $\theta_m = 30^\circ$, $m(\tilde{q}) = 100$ GeV, $\langle W \rangle = 3.6$ GeV. This means that we expect monojets in squark decay to have an average charged invariant mass of order $m_{CH} \sim 1.8$ GeV, in satisfactory agreement with the experimental indications. To obtain a more model-independent statement, we can compare this QCD estimate with the corresponding one for unbiased (gluon) jets in the same range of p_T (or Q) and with the same cut in angular spread, see Eq. (2.12) and the comments that follow. We expect monojets in squark decay to be thinner than ordinary jets by a considerable factor. In the sense of average invariant masses, and for $\theta_m = 30^\circ$, $Q \sim 100$ GeV, this factor is ~ 2.4 .

4. - A BEAM DUMP SQUEEZE AND A COSMIC CONUNDRUM

We have argued that a supermodel with $m(\tilde{g}) \sim 3$ GeV, $m(\tilde{q}) \sim 80$ GeV, is in agreement with the UA1 indications, while a model with $m(\tilde{g}) \sim 7$ GeV, $m(\tilde{q}) \gtrsim 150$ GeV is not in serious disagreement with them. Beam dump experiments have searched for a photino contamination in a neutrino beam¹⁷⁾. In the standard view wherein $m^2(\tilde{\gamma}) \ll m^2(\tilde{g})$, their negative results imply (model-dependent) limits on the possible masses of squarks and gluinos. These limits are shown in Fig. 6 along with the domains of parameters "determined" by our analysis of our monojets. The allowed region corresponding to the smaller masses may soon be put in difficulty by forthcoming data from a bubble chamber experiment¹⁸⁾. The region corresponding to bigger masses is temporarily safe, but may be accessible to experimental verification at future hadron colliders¹⁹⁾. In Fig. 6 we have not let the squarks heavier than 150 GeV of our high mass "solution" reach arbitrarily large masses, otherwise the gluino lifetime would become long enough for gluino decay vertices to be easily observable in the monojet events, as we discuss in more detail in the next Chapter. We do not believe that the striking signature of a flying gluino would have escaped attention.

We have, throughout this paper, accepted the standard view that for light gluinos $m(\tilde{g}) \sim 7m(\tilde{\gamma})$. Having "determined" certain narrow domains of gluino masses, we have an idea of what mass photinos should have: $m(\tilde{\gamma})$ lies somewhere between 0.4 and 1.2 GeV. Cosmological constraints on photino masses are in contradiction with this photino mass range, as we proceed to discuss. Photinos may significantly contribute to the "observed" universal average energy density $\rho \lesssim 2 \cdot 10^{-29} \text{g/cm}^3$, if their annihilation cross-section is insufficient to have depleted their present abundance²⁰⁾. Photino pairs annihilate into quark or lepton pairs via the exchange of squarks or sleptons, that in models with light gluinos, are expected¹⁰⁾ to have similar masses, m_{sf} . The "weak" photino annihilation cross-section $\sigma_{\tilde{\gamma}}$ scales as $(m_{sf})^{-4}$. Grosso modo the present photino universal energy density $\rho(\tilde{\gamma})$ scales as $[\sigma_{\tilde{\gamma}}]^{-1} \sim m_{sf}^4$. Ellis et al.²⁰⁾ estimate that, if $m_{sf} \sim 20$ GeV, the limit $\rho(\tilde{\gamma}) < \rho$ implies $m(\tilde{\gamma}) > 0.5$ GeV. This means that in the case $m(\tilde{g}) \sim 3$ GeV, $m(\tilde{q}) \sim m_{sf} \sim 80$ GeV, the cosmological bound would be violated by two orders of magnitude (4^4), a fairly unacceptable situation. The case for $m(\tilde{g}) \sim 7$ GeV, $m(\tilde{q}) \gtrsim 150$ GeV, $m(\tilde{\gamma}) \sim 1$ GeV is similarly uncomfortable. The arguments leading to the cosmological bound are generally accepted, though the details of the calculations may be more subject to criticism. In particular, a quark model calculation of an annihilation cross section of very slow particles at a practically fixed value of $\sqrt{s}$ somewhere between 1 and 2 GeV,

a region plagued by resonances, may be no more than a rough estimate. Yet, it is hard to believe that the cross-section estimates fail by two orders of magnitude, or that $\sigma_{\tilde{\gamma}}$ at $\sqrt{s} \approx 2m(\tilde{\gamma})$ happens to be resonant. Compromises can be found by pushing the parameters here and there, but the cosmological bound cannot be fully and graciously avoided.

We have argued that the only known supersymmetric explanation of the UAl events that naturally accommodates very narrow monojets disagrees with cosmology. But, in imposing theoretical constraints on our phenomenological analysis, are we not getting slightly carried away? We are dealing with a supermodel none of whose predicted particles have ever been seen. We are trusting mass relations between gluinos and photinos, squarks and sleptons, that follow from supergravitational models that cannot be said to be fully understood or irresistibly compelling. To diminish the stricture of the cosmological bound to an acceptable level it would suffice, for instance, to reduce the slepton masses from the assumed value $\sim m(\tilde{q})$, down to their present experimental lower limit. Yes, we were getting slightly carried away.

5. - CONCLUSIONS AND FURTHER TESTS

We have seen that light gluinos, and squarks perhaps suspiciously close in mass to intermediate vector bosons, suggest themselves as a supersymmetric explanation of the high missing p_T events of UAl. Two narrow islands in the domain of mass parameters are acceptable, one in the neighbourhood of $m(\tilde{g}) \sim 3$ GeV, $m(\tilde{q}) \sim 80$ GeV, the other along the isthmus $m(\tilde{g}) \sim 7$ GeV, $m(\tilde{q}) > 150$ GeV. The first "solution" very naturally accommodates very narrow monojets, and a dominance of monojets over dijets. The second solution is simple: it only requires one new particle, the gluino, to do the job²¹⁾. It does not so naturally accommodate very narrow monojets.

Guided by absolute faith in UAl, and in the mass relations imposed by supersymmetric models incorporating supergravity, we have determined the masses of all sfermions to be either ~ 80 GeV or > 150 GeV, and the mass of the photino to lie somewhere between 0.4 and 1.2 GeV. In these domains of parameters we encounter a serious cosmological problem: too much mass should be present in the Universe in the form of photinos having survived since the Big Bang. At a more down-to-earth level, we may reduce our faith in theoretical mass relations and ponder what

other predictions of the Madrid scenario could be subject to imminent experimental test. The next and last four paragraphs deal with these predictions.

In Figs. 7 and 8 we display the expected missing p_T distributions of mono- and dijets for the $\tilde{g}\tilde{g}$ channel and for the $\tilde{q}\tilde{q}$ channel respectively. These are simply obtained from Eqs. (2.5), (3.10) and (3.11). The p_T distributions from squark decay have Jacobian peaks that may make the signals proudly tower over a dying background, as emphasized by Barger et al.⁷⁾. Observation of such a peaked structure would be a convincing argument against the interpretation of the missing p_T events as the tail of some background, e.g., the mere "QCD" production of a jet- Z^0 pair, followed by $Z^0 \rightarrow \nu\nu$ decay²⁾, or even the prosaic decay $W \rightarrow \tau\nu$ ²²⁾. We have refrained from adding the missing p_T distributions of Figs. 7 and 8, since this sum would be even more dependent on the assumed details of gluon (or eventual gluino-sea) structure functions than the sum of total cross-sections above a fixed missing p_T , discussed and criticized at the end of Chapter 1. In Fig. 8 we have separately plotted "decay" dijets and "spectator" dijets. The first are given by Eq. (3.10) and correspond to the chain $\tilde{q} \rightarrow q\tilde{g}$, $\tilde{g} \rightarrow q\tilde{q}\tilde{\gamma}$. Spectator dijets, given by Eq. (3.11), are those for which the squark decays electromagnetically, $\tilde{q} \rightarrow q\tilde{\gamma}$, and the second jet consists of the debris of the gluino originally made with the squark. Decay dijets should be coplanar with the beam axis, with the big jet in the hemisphere opposite that of the small jet and the missing p_T , as in Fig. 9. Spectator dijets need not show similar correlations, but as their missing p_T increases they have a tendency for the jets to be in the hemisphere opposite the photino, and to be collinear in the plane transverse to the beam axis. In our approximations, the normalized monojet and decay-dijet signals of Fig. 8 are insensitive to moderate variations in the collider energy. This is less so for the spectator dijet signal, that has been plotted in Fig. 8 for $\sqrt{s} = 630$ GeV. Notice the dominance of dijets at very large p_T .

There is a hope to determine directly the mass of the squark from eventual coplanar "decay" dijet events. The simple trick is explained in Fig. 9. Attribute to the escaping photino a longitudinal momentum such that its total momentum lies on top of the smaller jet, as in the figure. Reconstruct the total mass of the jet-jet missing $\vec{p}$ system. It should peak at a fixed value, $W = m(\tilde{q})$.

Suppose photinos are indeed what is missing in the missing p_T events. One may then liken their discovery to that of the neutrino [by Pauli²³⁾, not by Reines and Cowen²⁴⁾]. The difference is that nuclei were seen to decay, but hadrons have not been "seen". The lifetime of the squark is supposed to be too

short to see a squark live and die. For the gluino the situation is more interesting. The gluino lifetime²⁵⁾ and its typical track length at collider energies are:

$$\tau(\tilde{g} \rightarrow q \bar{q} \tilde{\chi}) \simeq \frac{48\pi}{\alpha \alpha_s} \left(\frac{m(\tilde{q})}{m(\tilde{g})} \right)^4 \frac{1}{m(\tilde{g}) \sum_q e_q^2} \quad (5.1a)$$

$$\simeq (1.7 \cdot 10^{-14} \text{ s}) \left(\frac{0.2}{\alpha_s} \right) \left(\frac{m(\tilde{q})}{80 \text{ GeV}} \right)^4 \left(\frac{3 \text{ GeV}}{m(\tilde{g})} \right)^5 \quad (5.1b)$$

$$c\gamma\tau \simeq (10^{-2} \text{ cm}) \left(\frac{0.2}{\alpha_s} \right) \left(\frac{m(\tilde{q})}{80 \text{ GeV}} \right)^4 \left(\frac{3 \text{ GeV}}{m(\tilde{g})} \right)^6 \left(\frac{E(\tilde{g})}{50 \text{ GeV}} \right) \quad (5.1c)$$

where the numbers are for three quark flavours (u,d,s) and we have normalized Eq. (5.1c) to a fairly hefty gluino boost. In the scenario with $m(\tilde{q}) \sim 80 \text{ GeV}$ there is little hope to see the decay of the gluino-containing hadron. But in the scenario with $m(\tilde{q}) > 150 \text{ GeV}$ anything may happen. For $m(\tilde{g}) = 7 \text{ GeV}$, $c\gamma\tau = 1.3 \text{ cm}$ at $m(\tilde{q}) = 1 \text{ TeV}$, 0.8 mm at $m(\tilde{q}) = 0.5 \text{ TeV}$. Such a lengthy mismatch between the interaction point and the gluino decay vertex would presumably have been noticed. This implies an upper limit $m(\tilde{q}) < 0.5\text{--}1 \text{ TeV}$. It need not be even said that the vertex of missing p_T events should be scrutinized surgically.

If real gluinos are as light as the ones we have considered, sleptons and squarks are said to have similar masses^{10),26)}. Sleptons of mass 80 GeV , or heavier than 150 GeV , are not efficiently made at CERN's collider. Sleptons should not be found. We consider this to be one of the safest predictions of the Madrid scenario.

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FIGURE CAPTIONS

Fig. 1 : Different processes contributing to missing p_T events in a model with light gluinos and medium-heavy squarks. a) The $\tilde{g}\tilde{g}$ channel. b) A detailed view of gluino fragmentation. c) $\tilde{q}\tilde{g}$ production followed by $\tilde{q} \rightarrow \tilde{\gamma}$ decay. d) $\tilde{q}\tilde{g}$ production followed by $\tilde{q} \rightarrow \tilde{g}$ decay.

Fig. 2 : Cross-sections for missing p_T events from the $\bar{p}p \rightarrow \tilde{g}\tilde{g}$ channel, at $\sqrt{s} = 540$ GeV with "UAl cuts", as functions of the gluino mass. The curve⁵⁾ labelled "NO FRAG" does not take into account gluino fragmentation, the curve labelled "FRAG" does. These curves refer to total cross-sections for missing p_T + jets. The curves labelled "MONOJETS" and "DIJETS" are cross-sections for the corresponding experimentally defined events, with inclusion of the effects of fragmentation. The one-event and ten-event levels, with an integrated luminosity $\sim 100 \text{ nb}^{-1}$, are indicated.

Fig. 3: Calculations by Barger et al.⁷⁾ of differential and total cross-sections with "UAl cuts", in a model with effectively stable gluinos of mass 3 GeV. The signals originate in the process $\bar{p}p \rightarrow \tilde{g}\tilde{q}$, followed by $\tilde{q} \rightarrow q\tilde{\gamma}$. Only the monojet signal is shown in b). Elsewhere solid lines refer to monojets, dashed lines to dijets.

Fig. 4 : Decays of a colourless source of mass Q into a) $\tilde{g}\tilde{g}$ and b) $\tilde{g}\tilde{g}\tilde{g}$. The "jet invariant mass" is that of the hadronic debris within an opening angle $\theta < \theta_m$.

Fig. 5 : Diagrams for the decay $\tilde{q} \rightarrow q\tilde{g}\tilde{\gamma}$.

Fig. 6 : Limits on $m(\tilde{g})$ and $m(\tilde{q})$ (in GeV) from beam dump experiments¹⁷⁾. Also shown are the allowed domains of parameters "determined" by our analysis of UAl data. The curves labelled SSC and LHC correspond to the regions explorable by photino searches at future colliders¹⁹⁾. The line labelled " 10^{-8}s " corresponds to searches for "stable" particles¹⁶⁾. The other gluino lifetime cutoffs stem from possible reinteractions of the gluino-containing hadron in the dump.

Fig. 7 : Missing p_T distribution for dijets and monojets originating from the $\bar{p}p \rightarrow \tilde{g}\tilde{g}$ channel, for the three indicated values of $m(\tilde{g})$.

Fig. 8 : Missing p_T distributions for dijets and monojets originating from the $\bar{p}p \rightarrow \tilde{q}\tilde{g}$ channel, for the indicated values of $m(\tilde{g})$ and $m(\tilde{q})$. We have distinguished "decay" dijets (dashed lines) from "spectator" ones (dotted lines).

Fig. 9 : A possible way to hunt for a value of the squark mass in dijet events produced in the chain $\bar{p}p \rightarrow \tilde{q}\tilde{g}$, $\tilde{q} \rightarrow q\tilde{g}$, $\tilde{g} \rightarrow \tilde{\gamma}$.

Fig. 1

Fig. 2

Fig. 3

(a)

(b)

Fig. 4

Fig. 5

Fig. 6

Fig. 7

Fig. 8

Fig. 9