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PROPERTIES OF SUPERSYMMETRIC PARTICLES AND PROCESSES

Author
Barnett, R.M.

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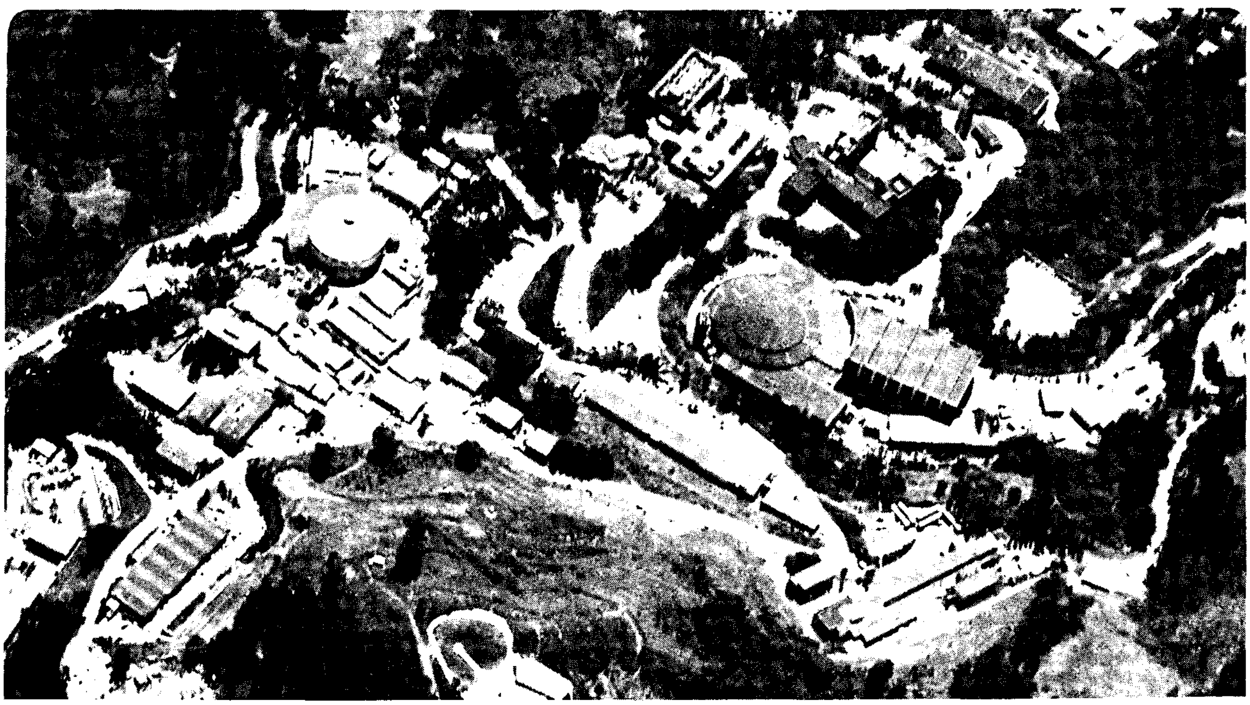

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Properties of Supersymmetric Particles and Processes*

R. Michael Barnett

Lawrence Berkeley Laboratory, University of California
Berkeley, California 94720

Lectures given at the
1985 SLAC Summer Institute on Particle Physics — Supersymmetry
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ABSTRACT

In these lectures I discuss the motivations for experimental searches for supersymmetric particles. The role of R-parity in these searches is described. The production and decay characteristics of each class of supersymmetric particles are investigated in the context of both e^+e^- and hadron machines. There is a detailed presentation of a sample calculation of a supersymmetric process. Emphasis is given to the signatures for detection of supersymmetric particles and processes. The current limits for supersymmetric particles are given.

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1. Introduction and Motivations

Let us assume that supersymmetry exists. We must then ask whether any particles or processes are accessible to experimentalists. This is the subject of my lectures. I will give only limited introductory remarks, since Joe Polchinski has given a general introduction to supersymmetry in his lectures[1] at this Summer Institute. Other good review articles with a phenomenological orientation may be found in Refs. 2–5.

Searches for supersymmetry share the characteristics of other searches for new physics. The experimentalists run the risk of having negative results. While the rewards of positive results may be high, the lifetime of the relevant theory may be short and the theorist's interest may be gone before the experiment begins running. It is, therefore, wise for an experimentalist to ask why he or she should invest significant effort into a search for supersymmetric particles. It is dangerous to ask this question of an enthusiast, so I will make an extra effort to give a balanced answer (but I do remain an enthusiast).

There are good theoretical motivations for supersymmetry such as a possible solution to the gauge hierarchy problem (discussed below). These and other motivations are presented in the lectures by Polchinski[1] and Peskin[6]. But I have heard non-enthusiasts (both theorists and experimentalists) argue that these motivations are obscure and uncertain and that other theories may do as well. The enthusiasts often use the word "elegant" in describing supersymmetry, and certainly the unification of fermions with bosons and the unification of gravity with the other interactions would be good justification for this word. The detractors counter that it would have been elegant if it had unified the known fermions with the known bosons thereby halving the total number of particles. Instead each known particle appears to require an undiscovered superpartner.

The advocates point to the fact that the existence of specific new particles with specific couplings is predicted and that there is good reason to expect masses to be less than 500–1000 GeV. The detractors emphasize that neither relative nor absolute masses are known so that it is impossible to predict the rates. The large production cross-sections for some supersymmetric particles have been reported by the optimists. The pessimists retort that the UA1 and UA2 experiments at CERN may have produced 10^4–10^5 gluinos and yet been unable to definitely detect any of them.

The enthusiasts have noted a number of very dramatic signals which are possible in some cases. But others counter that the dramatic signal found by UA1 (monojets) has been exceedingly difficult to separate from a large number of backgrounds. At a minimum, one would hope that the absence of signals would significantly constrain the theory. However, since we cannot be certain about the identity or mass of the lightest supersymmetric particle, about mixings or about which particle will decay into which, it is difficult to constrain supersymmetry unless one makes certain assumptions whose validity cannot be proven.

Finally, enthusiasts will argue (successfully I believe) that supersymmetry provides a very specific example of *how* to look for *new* physics at high energy machines. One can do real calculations. There are a large variety of types of processes. Ultimately the experimentalist will have to decide whether the theoretical motivations are adequate justification for this effort, but any attempt to find new physics will likely involve signals similar to those described in these lectures.

What is supersymmetry? It is a symmetry relating the fermions and bosons of elementary particle physics. Supersymmetric theories possess a self-conjugate, spin $\frac{1}{2}$ symmetry generator Q:

$$Q\,|\mathrm{Boson}\rangle = |\mathrm{Fermion}\rangle \tag{1.1}$$

and

$$Q\,|\mathrm{Fermion}\rangle = |\mathrm{Boson}\rangle$$

The (fermionic) Q obey anticommutation relations:

$$\left\{Q_\alpha{}^i,\, Q_j{}^{\dagger\beta}\right\} = \left(\gamma^\mu\right)_\alpha^\beta P_\mu \delta_j^i \tag{1.2}$$

where $P_\mu = i\partial_\mu$ is the momentum generator, α (β) is a spinoral index and i=1,2,...,N. N is an internal symmetry index. For phenomenological reasons, we choose "simple" supersymmetry, N=1.

The most frequently quoted theoretical motivations for supersymmetry are that it has fewer divergences than other quantum field theories, that it provides a possible mechanism for unifying gravity with other interactions and that it may offer an understanding of the problem of greatly different scales (called the "gauge hierarchy" problem). This latter problem[7] arises from the observation that the electroweak scale is vastly different from the scales of gravity and of grand unification:

$$\frac{m_W}{m_{Planck}} \sim \mathcal{O}(10^{-17}) \qquad \frac{m_W}{m_{GUT}} \sim \mathcal{O}(10^{-13}) \ . \tag{1.3}$$

One next notes that the calculation of the masses of scalar particles (the Higgs bosons) are quite unstable; the corrections to the bare mass tend to be very large. It takes an "unnatural" fine-tuning to avoid large masses: to keep m(Higgs) at $\mathcal{O}(m_W)$ rather than at $\mathcal{O}(m_{Planck})$. Among the corrections are one-loop radiative corrections to the mass which are quadratically divergent. However, one notes that those corrections involving fermions have the *opposite* sign as those involving bosons (Fermi statistics). Supersymmetry can take advantage of this because it provides a superpartner for each boson and fermion (on a one-to-one basis), and the particle and superpartner have identical couplings. It does require, however, that the mass splitting between particles and their superpartners be limited: less than about 1 TeV. Ultimately one hopes that supersymmetry could solve the gauge hierarchy problem by relating scalar masses to fermion masses which can be kept small via nearly exact chiral symmetries.

In an ideal world each known particle would turn out to be the superpartner of another known particle, in effect halving the total number of particles. Unfortunately, for phenomenological reasons, *no* known particle can be the superpartner of any other particle. No spin-zero elementary particles are known, so the partners of quarks and leptons are unknown. The Higgs boson (when found) will not be the partner of e, μ or τ, since that would lead to e-number or μ-number violation. Nor could the Higgs be a partner of quarks, because that implies Higgs would have color and that would lead to low-energy baryon-number violation. Finally the quarks, e, μ and τ are not partners of gauge bosons because they are not in adjoint representations of SU(3) and SU(2).

Supersymmetry clearly is a *broken* symmetry. Otherwise m(fermion) = m(boson), and this is never observed. The only question is the scale of breaking. As mentioned above, there is reason to expect that supersymmetric masses will be less than 500–1000 GeV.

The set of particles expected in simple supersymmetric theories is shown in Table I. The notation to indicate a supersymmetric particle is to place a tilde over the particle ($\tilde{P}$). Notice that the quarks and charged leptons have both left-handed and right-handed components (e.g., e_L is the SU(2) weak doublet while e_R is the SU(2) weak singlet). As a result there are *two* spin-zero partners of these particles which are labeled with subscripts L and R (e.g., $\tilde{e}_L$ and $\tilde{e}_R$); however, spin 0 particles are not themselves left- or right-handed. The subscripts simply refer to the component of the spin $\frac{1}{2}$ partner.

There is no compelling supersymmetric model which would tell us the *masses* of supersymmetric particles or even the *sequence* of masses. We do get some guidance from supergravity models. Otherwise we need to look to

Table I.

Particle		Spin	Supersymmetric Partner*			Spin
quark	q	$\frac{1}{2}$	scalar-quark	$\tilde{q}_L$,	$\tilde{q}_R$	0
	u	$\frac{1}{2}$		$\tilde{u}_L$,	$\tilde{u}_R$	0
	d	$\frac{1}{2}$		$\tilde{d}_L$,	$\tilde{d}_R$	0
	⋮			⋮		
lepton	ℓ	$\frac{1}{2}$	scalar-lepton	$\tilde{\ell}_L$,	$\tilde{\ell}_R$	0
	e	$\frac{1}{2}$		$\tilde{e}_L$,	$\tilde{e}_R$	0
	μ	$\frac{1}{2}$		$\tilde{\mu}_L$,	$\tilde{\mu}_R$	0
	τ	$\frac{1}{2}$		$\tilde{\tau}_L$,	$\tilde{\tau}_R$	0
neutrino	ν_e	$\frac{1}{2}$	scalar-neutrino	$\tilde{\nu}_e$		0
	ν_μ	$\frac{1}{2}$		$\tilde{\nu}_\mu$		0
	ν_τ	$\frac{1}{2}$		$\tilde{\nu}_\tau$		0
gluon	g	1	gluino	$\tilde{g}$		$\frac{1}{2}$
photon	γ	1	photino	$\tilde{\gamma}$		$\frac{1}{2}$
Z	Z^o	1	Z-ino	$\tilde{Z}^o$		$\frac{1}{2}$
higgs	H_i^o	0	higgsino	$\tilde{H}_i^o$		$\frac{1}{2}$
W	$W^\pm$	1	W-ino	$\tilde{W}^\pm$		$\frac{1}{2}$
higgs	$H_i^\pm$	0	higgsino	$\tilde{H}_i^\pm$		$\frac{1}{2}$

*Mass eigenstates can be mixtures of $\tilde{\gamma}$, $\tilde{Z}^o$, $\tilde{H}_i^o$

or of $\tilde{W}^\pm$, $\tilde{H}_i^\pm$ or of $\tilde{q}_L$, $\tilde{q}_R$ or of $\tilde{\ell}_L$, $\tilde{\ell}_R$

cosmology and experiments to learn more. To make predictions for such experiments, many theorists calculate processes "independent" of supersymmetric masses which is to say as a function of masses. But the sequence of masses is critical: e.g., does the scalar quark decay into the gluino or vice versa.

2. The Lightest Supersymmetric Particle

I will make (in most cases) two important assumptions: a) The photino ($\tilde{\gamma}$) is the *lightest* supersymmetric particle. b) Mixing of the photino with $\tilde{Z}^o$ and $\tilde{H}^o$ is not too important. The first assumption is the most common output of recent models. The results described in this paper would be modified somewhat if this were not a correct assumption; these changes are described in an accompanying paper[8] by Howard Haber in these proceedings.

Other alternatives for the lightest particle are the Goldstino,[9] the scalar neutrino[10] ($\tilde{\nu}$) and the Higgsino[11] ($\tilde{H}$). In the currently popular supergravity models, the Goldstino is absorbed into the gravitino[12] and is, therefore, not a candidate (although the gravitino might be a candidate). The consequence of $\tilde{\nu}$ or $\tilde{H}$ being the lightest will be briefly discussed in the next section.

Charged and strongly-interacting particles are not possible candidates for the lightest supersymmetric particle. As discussed in the next section, the lightest particle will most likely be stable. Charged or strongly-interacting particles produced in the early universe would be dissipated and condensed out on earth and elsewhere. Searches for 1 TeV or less stables are negative at reasonable levels.

3. R-Parity and Its Implications

Most models of supersymmetry have a multiplicatively-conserved quantum number[13] called "R-Parity". Ordinary particles such as u, d or s quarks, e, μ or τ leptons, γ, gluons, or W bosons have R = +1 while their supersymmetric partners ($\tilde{u}$, $\tilde{d}$, $\tilde{s}$, $\tilde{e}$, $\tilde{\mu}$, $\tilde{\tau}$, $\tilde{\gamma}$, $\tilde{g}$, $\tilde{W}$) have R = –1.

This quantum number can be defined by

$$R = (-1)^{2J+3B+L} \tag{3.1}$$

where J = spin, B = baryon number and L = lepton number. To our great misfortune there is also another definition possible:

$$R = (-1)^{E+1} \tag{3.2}$$

where E = 1 for existing particles

E = –1 for non-existent particles.

In general I will assume that R-parity is *not* broken. An accompanying paper[14] by Sally Dawson in these proceedings describes the consequences of R-parity violation.

R-parity has a number of important implications. The first is that the lightest supersymmetric particle has to be stable, since no combination of positive R-parity products can come from a negative R-parity parent. As stated above, I will usually assume the photino is lightest. If $\tilde{\nu}$ or $\tilde{H}$ were lighter, then

$$\tilde{\gamma} \to \nu\,\tilde{\nu} \tag{3.3}$$

or

$$\tilde{\gamma} \to \gamma\,\tilde{H}^{o} \tag{3.4}$$

would be allowed.

The second implication is that supersymmetric particles are produced in *pairs* (since the initial state has R = +1); see the figures below (where the double lines indicate negative R parity):

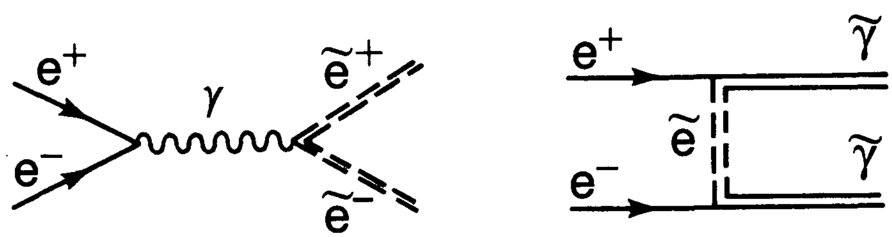

FIG 1

Thirdly, all decays end up with photinos if it is the lightest supersymmetric particle; see for example the figure below:

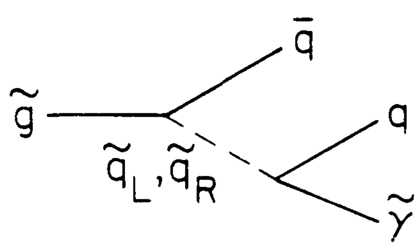

FIG 2

A fourth implication is that the production cross-sections of *light* supersymmetric particles are of order *weak* interaction scale. In this figure

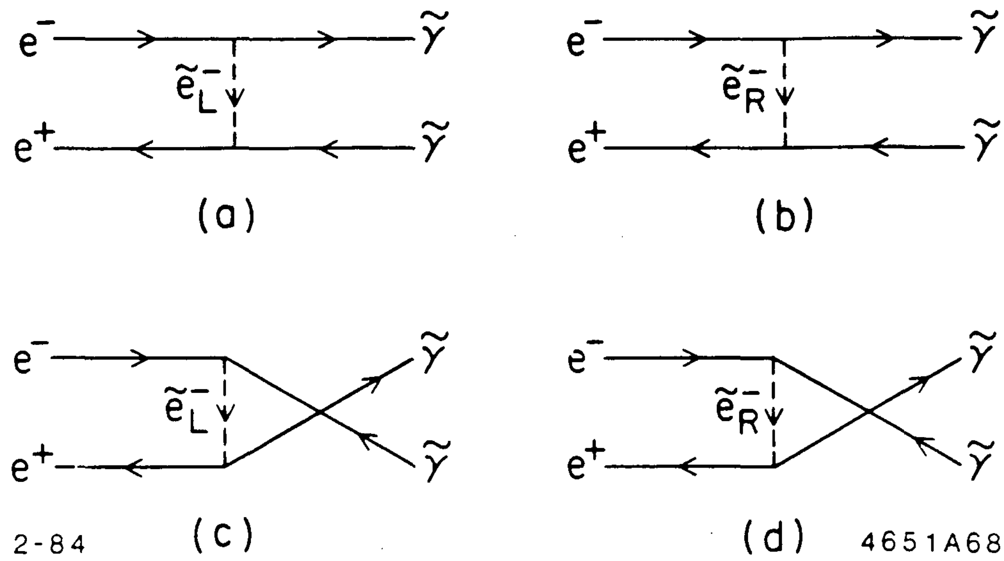

FIG 3

we note (from experiment) that $M(\tilde{e}) \gtrsim \mathcal{O}(m_W)$. A potential exception is the gluino, because it can be produced via the following diagram:

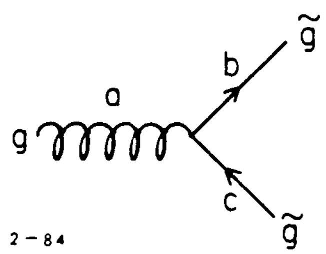

FIG 4

However, we will see later that the gluino may not be light.

The fifth implication of R-parity is that the *photino interaction* cross-section[15] is of order the weak interaction scale. This is because its interaction with matter requires the exchange and/or the production of heavy particles as seen in the diagrams of Fig. 5.

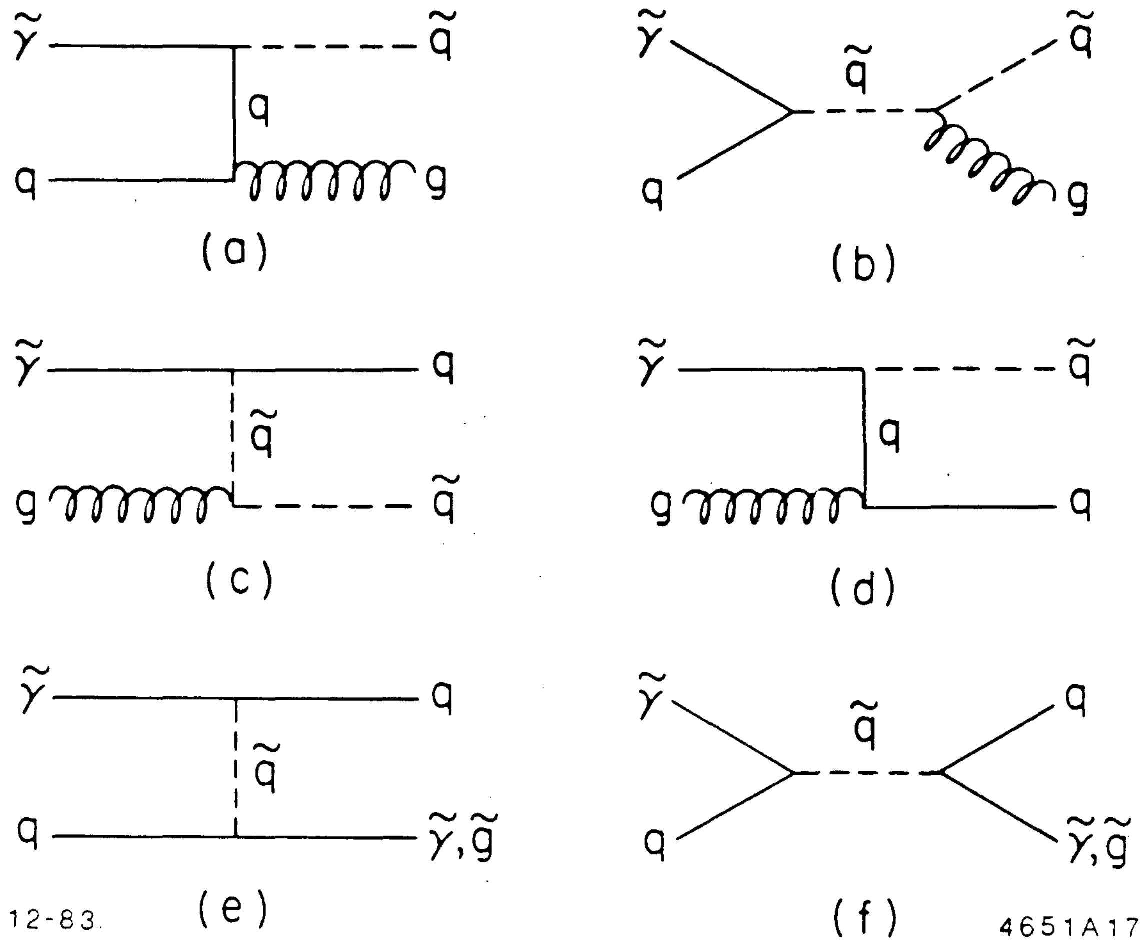

FIG 5

A major consequence of these implications of R-parity is that the *photino* behaves like a *neutrino* (though it is not associated with an e or μ). Like a neutrino, a photino will normally leave detectors without being detected. Since photinos are the eventual output of every process, this means that the key signature for supersymmetric processes is *MISSING ENERGY*.

4. Mixing of Neutral Gauge and Higgs Fermions

The states $\widetilde{B}$, $\widetilde{W}_3$, $\widetilde{H}_1{}^o$ and $\widetilde{H}_2{}^o$ are neutral, Majorana spinor, electroweak-interaction eigenstates which mix to give *mass* eigenstates. Crudely speaking, a Majorana fermion is its own antiparticle or more formally a CPT self-conjugate particle. In other words, the helicity ($\uparrow \widetilde{\gamma}$) goes to ($\downarrow \widetilde{\gamma}$) under CPT and then ($\downarrow \widetilde{\gamma}$) returns to ($\uparrow \widetilde{\gamma}$) under a 180° rotation.

The photino is given by

$$\widetilde{\gamma} = \frac{(g' \widetilde{W}_3 + g \widetilde{B})}{\sqrt{g^2 + g'^2}} = \widetilde{W}_3 \sin \theta_W + \widetilde{B} \cos \theta_W \tag{4.1}$$

in analogy to the ordinary photon, where $g = e/\sin \theta_W$ and $g' = e/\cos \theta_W$. The photino may or may not be a mass eigenstate. Discussion of the other gauge and Higgs fermions will be given in Sections 11–12.

5. The Photino

Let us summarize the characteristics of the photino. I shall assume (although it is not necessary) that mixing with $\widetilde{Z}^o$ and $\widetilde{H}^o$ is small, and that the $\widetilde{\gamma}$ is a mass eigenstate. It is taken (here) to be the lightest supersymmetric particle and stable (via R-parity). Its coupling is the same as that of the photon. The photino interacts weakly with matter (heavy masses are involved), and production cross-sections are weak. The only mechanism for "detecting" it is by noting the missing energy its presence implies.

Currently, the only limits on the mass of the photino come from cosmological arguments.[16–22] It is necessary that the mass density of the universe be below the critical density which closes the universe. If the photino is too heavy there is a danger of overclosing the universe. If $M(\tilde{\gamma}) \lesssim 100$ eV, then there is no problem. If $M(\tilde{\gamma})$ is larger, however, one must consider how photinos are created and lost. Pair production of photinos will continue until they are frozen out as the temperature, T, drops below $M(\tilde{\gamma})$. But annihilation processes would continue, and the eventual number of photinos remaining depends on the efficiency of these processes. The rate of these processes is a function of heavy masses (such as $M(\tilde{e})$ and $M(\tilde{q}) \equiv M_{heavy}$) appearing in diagrams such as:

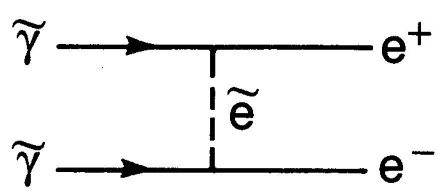

FIG 6

One finds[18]

$$M_{\tilde{\gamma}} \gtrsim \tfrac{1}{2} \text{ GeV} \quad \text{for} \qquad M_{\text{heavy}} = 20 \text{ GeV}$$

$$M_{\tilde{\gamma}} \gtrsim 2 \text{ GeV} \quad \text{for} \qquad M_{\text{heavy}} = 40 \text{ GeV}$$

$$M_{\tilde{\gamma}} \gtrsim 5 \text{ GeV} \quad \text{for} \qquad M_{\text{heavy}} = 100 \text{ GeV}$$

Since we have no evidence that $M_{\text{heavy}} > 100$ GeV, we also have no evidence that $M_{\tilde{\gamma}} > 5$ GeV.

In the next section, I will discuss the computation of photino production in a specific process.

6. A Sample Supersymmetric Calculation

I wish to show that supersymmetric calculations are not that much more difficult than ordinary calculations. They can be tedious, but are straightforward and can be done in a finite time. The process I have chosen is $e^+e^- \to \tilde{\gamma}\,\tilde{\gamma}$ (this cross-section was derived in Ref. 4). I will assume that $\tilde{e}_L$, $\tilde{e}_R$ and $\tilde{\gamma}$ are mass eigenstates. The relevant diagrams are:

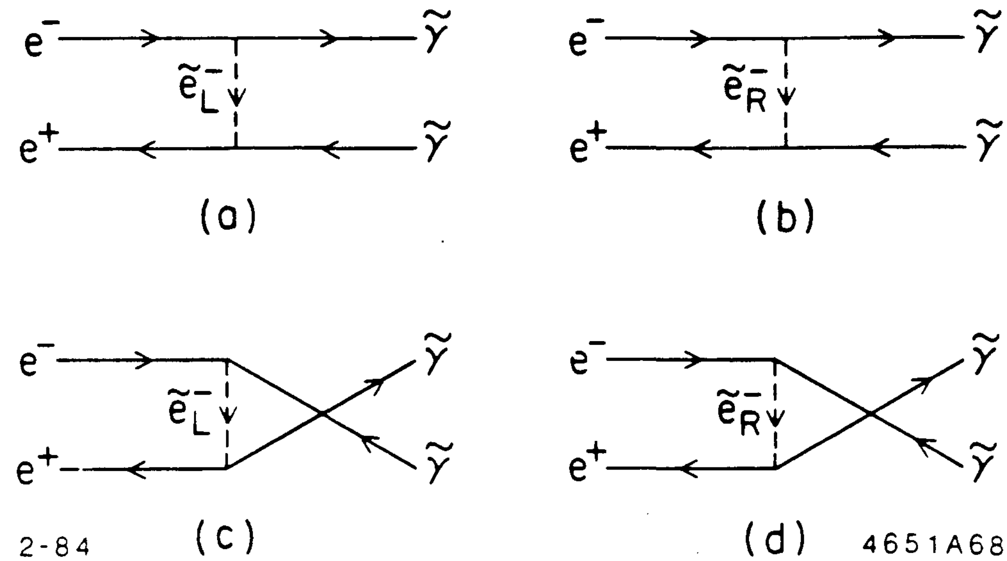

FIG 7

The photino interaction terms needed are (recall that e and $\tilde{\gamma}$ are 4-component spinors)

$$\sqrt{2}\, g \sin\theta_W \left[\bar{\tilde{\gamma}}\, \frac{(1-\gamma_5)}{2}\, e\, \tilde{e}_L{}^* + \bar{e}\, \frac{(1+\gamma_5)}{2}\, \tilde{\gamma}\, \tilde{e}_L - \bar{\tilde{\gamma}}\, \frac{(1+\gamma_5)}{2}\, e\, \tilde{e}_R{}^* - \bar{e}\, \frac{(1-\gamma_5)}{2}\, \tilde{\gamma}\, \tilde{e}_R \right] \qquad (6.1)$$

The Feynman rules needed to do the calculations are shown in Fig. 8 where α and β are spinor indices (watch the direction of arrows).

To obtain the Feynman rules for $\tilde{e}_R$, make the change

$$(1 \pm \gamma_5) \rightarrow (-1 \pm \gamma_5) \ . \qquad (6.2)$$

In general, most Feynman rules may be obtained from Ref. 4.

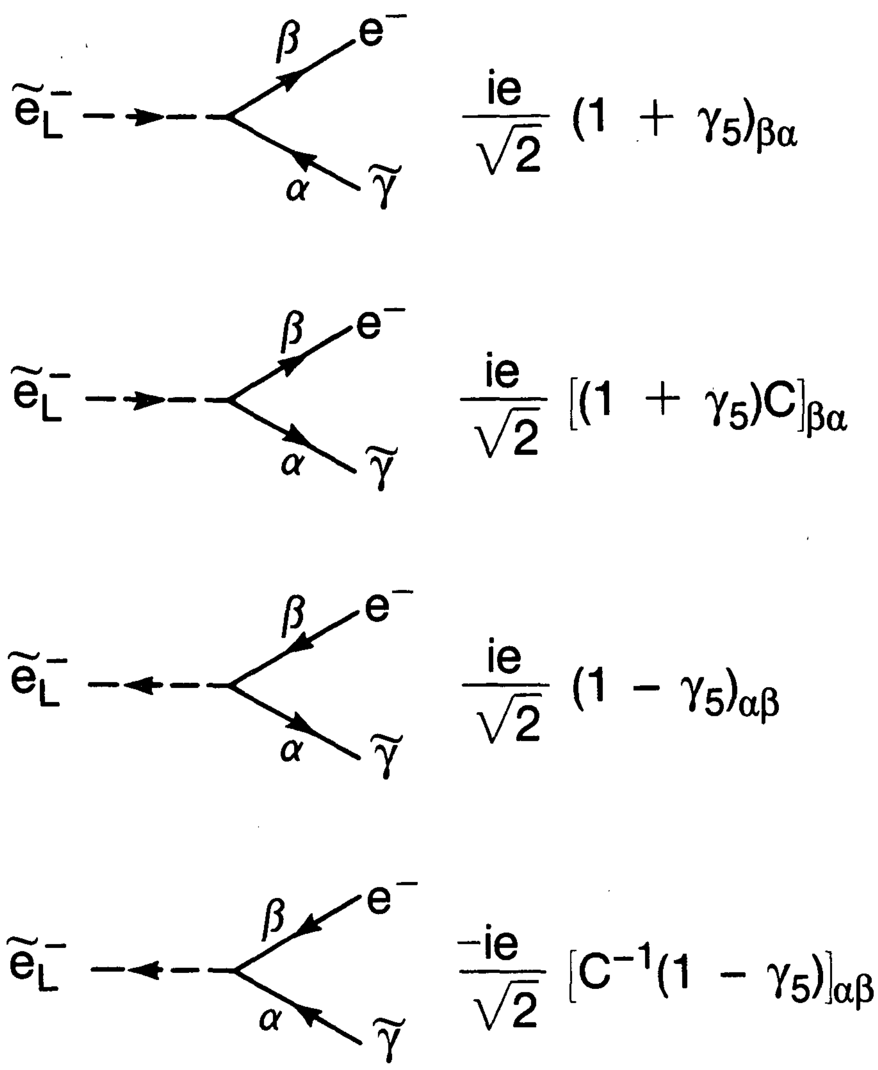

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FIG 8

Let us first calculate the amplitudes for the a and c diagrams where it is $\tilde{e}_L$ which is exchanged. This is equivalent to letting $M(\tilde{e}_R) \rightarrow \infty$. I will come back to the alternative case $M(\tilde{e}_R) = M(\tilde{e}_L)$ later. Let us define momenta by:

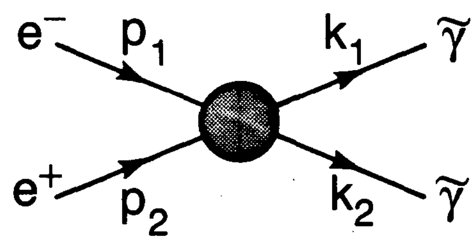


FIG 9

and

$$t = (p_1 - k_1)^2 \tag{6.3}$$

$$u = (p_1 - k_2)^2 \tag{6.4}$$

Then the amplitudes are

$$M_a = \left\{ \frac{e^2}{2\,(M^2_{\tilde{e}_L} - t)} \right\} \bar{u}(k_1)\,(1-\gamma_5)\,u(p_1)\,\bar{v}(p_2)\,(1+\gamma_5)\,v(k_2) \tag{6.5}$$

$$M_c = \left\{ \frac{-e^2}{2\,(M^2_{\tilde{e}_L} - u)} \right\} v(k_2)^T\,C^{-1}\,(1-\gamma_5)\,u(p_1)\,\bar{v}(p_2)\,(1+\gamma_5)\,C\,\bar{u}(k_1)^T \tag{6.6}$$

The u and v spinors satisfy (for spin $s = \pm\,\frac{1}{2}$) the relations:

$$v(k,s) = C\,\bar{u}\,(k,s)^T \tag{6.7}$$

$$u(k,s) = C\,\bar{v}\,(k,s)^T \tag{6.8}$$

Using $C^T = -C$, we can then rewrite Eq. 6.6 as:

$$M_c = \left(\frac{e^2}{2\,(M^2_{\tilde{e}_L} - u)} \right) \bar{u}(k_2)\,(1-\gamma_5)\,u(p_1)\,\bar{v}(p_2)\,(1+\gamma_5)\,v(k_1) \tag{6.9}$$

Now we can proceed by squaring the sum of the two amplitudes and summing over initial and final spins. The amplitude M_c enters the sum with a minus sign relative to M_a because of Pauli statistics. In $e^-e^- \to e^-e^-$ scattering the same sign occurs between t- and u-channel terms. So we wish to find

$$\sum |M_a - M_c|^2 \ .$$

To evaluate it we need to make use of the following relations for the projection operators (where $\not{p} \equiv \gamma \cdot p$ and $\sum\limits_s$ is the sum over spins):

$$\sum_s u(p,s)\,\bar{u}(p,s) = \not{p} + M \tag{6.10a}$$

$$\sum_s v(p,s)\,\bar{v}(p,s) = \not{p} - M \tag{6.10b}$$

$$\sum_s u(p,s)\,v^T(p,s) = (\not{p} + M)C^T \tag{6.10c}$$

$$\sum_s \bar{u}^T(p,s)\bar{v}(p,s) = C^{-1}\,(\not{p} - M) \tag{6.10d}$$

$$\sum_s \bar{v}^T(p,s)\,\bar{u}(p,s) = C^{-1}\,(\not{p} + M) \tag{6.10e}$$

$$\sum_s v(p,s)\,u^T(p,s) = (\not{p} - M)C^T \tag{6.10f}$$

If our calculations did not involve Majorana spinors, we would not have needed Eqs. (6.10c-f). Also useful in evaluating the matrix element squared are:

$$C^\dagger = C^{-1} \qquad\qquad C^T = -C \tag{6.11}$$

$$C^{-1}\,i\gamma_5\,C = (i\gamma_5)^T \qquad\qquad C^{-1}\,\gamma_\mu\,\gamma_5\,C = (\gamma_\mu\,\gamma_5)^T \tag{6.12}$$

$$C^{-1}\,\gamma_\mu\,C = -\gamma_\mu{}^T \qquad\qquad C^{-1}\,\sigma_{\mu\nu}\,C = -\sigma_{\mu\nu}{}^T \tag{6.13}$$

To proceed we write

$$\sum |M_a - M_c|^2 = \sum |M_a|^2 + \sum |M_c|^2 - 2\sum M_a M_c{}^* \tag{6.14}$$

Starting with $\sum |M_a|^2$ and letting

$$A_t \equiv \frac{e^2}{2\,(M^2_{\tilde{e}_L} - t)} \qquad\qquad A_u \equiv \frac{e^2}{2\,(M^2_{\tilde{e}_L} - u)} \tag{6.15}$$

we have

$$\sum_{\text{spins}} |M_a|^2 = A_t{}^2 \sum_{\text{spins}} \bar{u}(k_1)\,(1-\gamma_5)\,u(p_1)\,\bar{u}(p_1)(1+\gamma_5)\,u(k_1) \times \bar{v}(p_2)\,(1+\gamma_5)\,v(k_2)\,\bar{v}(k_2)\,(1-\gamma_5)\,v(p_2) \tag{6.16}$$

Using Eqs. (6.10) we find

$$\sum_{\text{spins}} |M_a|^2 = A_t{}^2\,\mathrm{Tr}\left[(1-\gamma_5)\,(\not{p}_1+m_e)\,(1+\gamma_5)(\not{k}_1+M_{\tilde{\gamma}})\right] \times \mathrm{Tr}\left[(1+\gamma_5)\,(\not{k}_2-M_{\tilde{\gamma}})\,(1-\gamma_5)\,(\not{p}_2-m_e)\right] \tag{6.17}$$

$$= A_t{}^2\,\mathrm{Tr}\left[2\,(1-\gamma_5)\,\not{p}_1(\not{k}_1+M_{\tilde{\gamma}})\right]\,\mathrm{Tr}\left[2\,(1+\gamma_5)\,\not{k}_2(\not{p}_2-m_e)\right] \tag{6.18}$$

where the last line follows from

$$(1 \pm \gamma_5)\gamma_\mu = \gamma_\mu\,(1 \mp \gamma_5)$$

$$(1 + \gamma_5)\,(1 \mp \gamma_5) = \begin{Bmatrix} 0 \\ 2 \end{Bmatrix}(1+\gamma_5) \tag{6.19}$$

Other γ matrix relations allow further simplification:

$$\mathrm{Tr}\,[\text{odd number } \gamma_\mu{}'\text{ s}] = 0$$

$$\mathrm{Tr}\,[\gamma_5\,\not{a}\,\not{b}] = 0 \qquad\qquad \mathrm{Tr}\,[\not{a}\,\not{b}\,] = 4\,a \cdot b \tag{6.20}$$

These give

$$\sum_{\text{spins}} |M_a|^2 = A_t^2\, 4\,(4p_1 \cdot k_1)\,(4\, p_2 \cdot k_2) \tag{6.21}$$

or using

$$t = p_1^2 - 2\, p_1 \cdot k_1 + k_1^2 = -2\, p_1 \cdot k_1 + m_e^2 + M_{\tilde{\gamma}}^2 \tag{6.22}$$

$$\sum_{\text{spins}} |M_a|^2 = \left\{ \frac{e^4}{4\,(M_{\tilde{e}_L}^2 - t)^2} \right\} 16\, \left(t - m_e^2 - M_{\tilde{\gamma}}^2\right)^2 \tag{6.23}$$

Obviously, the $\sum |M_c|^2$ is the same except for the interchange $t \rightarrow u$:

$$\sum_{\text{spins}} |M_c|^2 = \left\{ \frac{e^4}{4\,(M_{\tilde{e}_L}^2 - u)^2} \right\} 16\, \left(u - m_e^2 - M_{\tilde{\gamma}}^2\right)^2 \tag{6.24}$$

Next let us do the Interference Term: $-2 \sum_{\text{spins}} M_a M_c^*$. For t- and u-channel interference terms, it is convenient (for rearranging terms to obtain the trace) to make use of this relation:

$$\begin{aligned} [\overline{v}(p)\,(1-\gamma_5)\,v(k)] &= [\overline{v}(p)\,(1-\gamma_5)\,v(k)]^T \\ &= v^T(k)\,(1-\gamma_5)^T\,\overline{v}^T(p) \end{aligned} \tag{6.25}$$

which follows because this is a scalar quantity. The need for this trick occurs in supersymmetry because of interference between diagrams involving the neutral Majorana fermion $\tilde{\gamma}$ $(\tilde{\gamma} = \overline{\tilde{\gamma}})$. Using this trick we then rewrite the amplitudes:

$$M_a = A_t\, \overline{u}(k_1)\,(1-\gamma_5)\,u(p_1)\,v^T(k_2)\,(1+\gamma_5)^T\,\overline{v}^T(p_2) \tag{6.26}$$

$$M_c^* = A_u\, \overline{u}(p_1)\,(1+\gamma_5)\,u(k_2)\,v^T(p_2)\,(1-\gamma_5)\,\overline{v}^T(k_1) \tag{6.27}$$

We can then write

$$-2 \sum_{\text{spins}} M_a M_c^* = -2 A_t A_u \sum_{\text{spins}} \bar{u}(k_1)(1-\gamma_5) u(p_1) \bar{u}(p_1)(1+\gamma_5) u(k_2)$$

$$\times v^T(k_2)(1+\gamma_5)^T \bar{v}^T(p_2) v^T(p_2)(1-\gamma_5)^T \bar{v}^T(k_1) \quad (6.28)$$

$$= -2 A_t A_u \,\mathrm{Tr}\Big[(1-\gamma_5)(\not{p}_1+m_e)(1+\gamma_5)(\not{k}_2+M_{\tilde{\gamma}}) C^T$$

$$\times (1+\gamma_5)^T (\not{p}_2-m_e)^T (1-\gamma_5)^T C^{-1}(\not{k}_1+M_{\tilde{\gamma}})\Big] \quad (6.29)$$

To reduce this we note that (using Eqs. 6.11–6.13)

$$C^T (1+\gamma_5)^T (\not{p}_2-m_e)^T (1-\gamma_5)^T C^{-1}$$

$$= -C(1+\gamma_5)^T C^{-1} C(\not{p}_2-m_e)^T C^{-1} C(1-\gamma_5)^T C^{-1}$$

$$= -(1+\gamma_5)(-\not{p}_2-m_e)(1-\gamma_5) \quad (6.30)$$

Inserting Eq. (6.30) into Eq. (6.29) we obtain

$$-2 \sum_{\text{spins}} M_a M_c^* = -2 A_t A_u \,\mathrm{Tr}\Big[(1-\gamma_5)(\not{p}_1+m_e)(1+\gamma_5)(\not{k}_2+M_{\tilde{\gamma}})$$

$$\times (1+\gamma_5)(\not{p}_2+m_e)(1-\gamma_5)(\not{k}_1+M_{\tilde{\gamma}})\Big] \quad (6.31)$$

$$= -2 A_t A_u \, 8 M_{\tilde{\gamma}} \,\mathrm{Tr}\Big[(1-\gamma_5)\not{p}_1 \not{p}_2 (\not{k}_1+M_{\tilde{\gamma}})\Big] \quad (6.32)$$

$$= -2 A_t A_u \, 8 M_{\tilde{\gamma}}^2 \, 4 \, p_1 \cdot p_2 \quad (6.33)$$

The last steps involved use of Eqs. (6.19–6.20). By noting that

$$s = 2 p_1 \cdot p_2 + 2 m_e^2 \quad (6.34)$$

we obtain the final form for the interference term

$$-2 \sum_{\text{spins}} M_a M_c^* = -8 e^4 M_{\tilde{\gamma}}^2 \frac{s-2m_e^2}{(M_{\tilde{e}_L}^2 - t)(M_{\tilde{e}_L}^2 - u)} \quad (6.35)$$

Note that the interference term is zero if the photino is massless.

Combining Eqs. (6.23), (6.24), and (6.35), the differential cross-section $\frac{d\sigma_1}{d\Omega}(e^+e^- \to \tilde{\gamma}\,\tilde{\gamma})$ for the case where $M_{\tilde{e}_R} \to \infty$ is

$$\frac{d\sigma_1}{d\Omega} = \frac{\alpha^2}{4s}\left(\frac{s-4M_{\tilde{\gamma}}^2}{s-4m_e^2}\right)^{1/2}\left[\frac{(t-M_{\tilde{\gamma}}^2-m_e^2)^2}{(M_{\tilde{e}_L}^2-t)^2} + \frac{2\,M_{\tilde{\gamma}}^2\,(2m_e^2-s)}{(M_{\tilde{e}_L}^2-t)\,(M_{\tilde{e}_L}^2-u)} + \frac{(u-M_{\tilde{\gamma}}^2-m_e^2)^2}{(M_{\tilde{e}_L}^2-u)^2}\right] \tag{6.36}$$

For the case $M_{\tilde{e}_R} = M_{\tilde{e}_L}$, I quote the result for $\frac{d\sigma_2}{d\Omega}(e^+e^- \to \tilde{\gamma}\,\tilde{\gamma})$ derived by Haber and Kane[4] (see also their erratum, to be published):

$$\frac{d\sigma_2}{d\Omega} = \frac{\alpha^2}{2s}\left(\frac{s-4M_{\tilde{\gamma}}^2}{s-4m_e^2}\right)^{1/2}\left[\frac{(t-M_{\tilde{\gamma}}^2-m_e^2)^2+4\,m_e^2\,M_{\tilde{\gamma}}^2}{(M_{\tilde{e}}^2-t)^2} + \frac{8m_e^2M_{\tilde{\gamma}}^2-2s(M_{\tilde{\gamma}}^2+m_e^2)}{(M_{\tilde{e}}^2-t)\,(M_{\tilde{e}}^2-u)} + \frac{(u-M_{\tilde{\gamma}}^2-m_e^2)^2+4\,m_e^2\,M_{\tilde{\gamma}}^2}{(M_{\tilde{e}}^2-u)^2}\right] \tag{6.37}$$

Note that if $m_e = 0$ then

$$\frac{d\sigma_1}{d\Omega} = \frac{1}{2}\,\frac{d\sigma_2}{d\Omega} \tag{6.38}$$

This concludes our calculation of $e^+e^- \to \tilde{\gamma}\,\tilde{\gamma}$.

7. Photino Interaction Cross Section and Beam-Dump Experiments

We have just calculated the production of photinos. Let us now consider their interaction with matter.[15] The interaction cross-section can be written as:

$$\sigma^{\tilde{\gamma}}(\text{interaction}) = \sum_{\text{quarks}} \int dx\, q(x)\, \hat{\sigma}(\hat{s}) \tag{7.1}$$

The quantity q(x) is the quark structure function. The term $\hat{\sigma}(\hat{s})$ describes the hard process $\tilde{\gamma} + q \rightarrow$ anything. The variable $\hat{s}$

$$\hat{s} = sx \tag{7.2}$$

describes the hard process.

One method of searching for light gluinos involves observing photino interactions. In a "beam dump" experiment, one would produce large numbers of gluinos in the dump (via gluon fusion, etc.) if light gluinos exist. These gluinos are expected to decay before interacting in the dump. As always photinos would be among the products. The photinos pass through the shielding and *like neutrinos* can occasionally interact in a neutrino detector. These interactions will appear similar to neutrino neutral-current interactions.

For beam-dump experiments where $\hat{s} << M_{\tilde{g}}^2$ and $M_{\tilde{g}}$ is by hypothesis small, the relevant diagrams are:

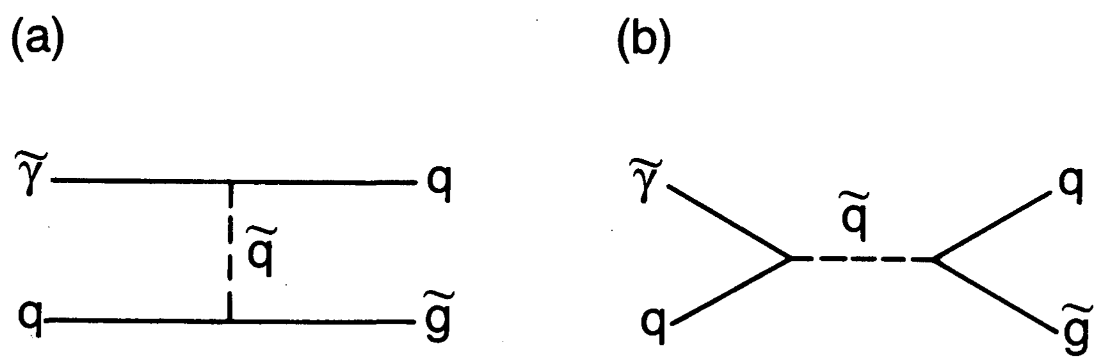


FIG 10

In this case ($\hat{s} << M^2_{\tilde{q}}$) we then have

$$\hat{\sigma} \sim \left(\frac{\alpha \, \alpha_s}{M^4_{\tilde{q}}} \right) \hat{s} \tag{7.3}$$

Written in the form of ν cross-sections, this yields[15,23]

$$\sigma_{Int}^{\tilde{\gamma}} \approx 20 \times 10^{-38} \, E^{lab}_{\tilde{g}} \left(\frac{m_W}{M_{\tilde{q}}} \right)^4 \tilde{F} \, (M^2_{\tilde{g}}, s) \quad cm^2 \tag{7.4}$$

where

$$\tilde{F} = \sum_q \int_{\frac{M^2_{\tilde{g}}}{s}}^{1} dx \; x \; q(x) \left(1 - \frac{M^2_{\tilde{g}}}{xs} \right)^2 \left(1 + \frac{M^2_{\tilde{g}}}{8xs} \right) e_q{}^2 \tag{7.5}$$

Let us evaluate $\tilde{F}$ for the case $M_{\tilde{q}} = 100$ GeV and $\sqrt{s} = 5$ GeV and note that the answer scales with $M_{\tilde{g}}/\sqrt{s}$:

$M_{\tilde{g}}$	0	0.4	1.0	2.0	3.0	(GeV)
$\tilde{F}$	0.15	0.12	0.077	0.021	0.0024	

We can now compare with neutrino cross-sections

$$\sigma(\nu) \approx 0.6 \times 10^{-38} \, E_\nu \, cm^2 \quad , \tag{7.6}$$

and we see that they are roughly comparable. Limits on the mass of the gluino from beam-dump experiments will be discussed later.

8. Charged Scalar Leptons

The subject has been dealt with at length in the lectures of David Burke[24] at this Summer School, so I will only briefly discuss it. At electron-positron colliders (not on the Z^0 resonance), the following processes can be considered:

a) $e^+e^- \to \tilde{e}^+ \tilde{e}^-$

Limited to $M_{\tilde{e}} < \frac{\sqrt{s}}{2}$.

b) $e^+e^- \to e^+\tilde{e}^- \tilde{\gamma} \to e^+e^- \tilde{\gamma}\tilde{\gamma}$ (Ref. 25) See Fig. 11.

Limited to $M_{\tilde{e}} < \sqrt{s}$.

c) $e^+e^- \to e^+e^- \tilde{\gamma}\tilde{\gamma}$ (Ref. 26)

Off-shell $\tilde{e}$.

d) $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$

No signal.

e) $e^+e^- \to \tilde{\gamma}\tilde{\gamma}\gamma$ (Ref. 26,27) See Fig. 12.

Determined via and limited by $\tilde{e}$ propagator.

At Z^o factories we can consider the following processes:

a) $Z^o \to \tilde{e}^+ \tilde{e}^-$

Probably ruled out by current limits on $M_{\tilde{e}}$ unless $M_{\tilde{\gamma}} \geq 13$ GeV.

b) $Z^o \to \tilde{e}^+ e^- \tilde{\gamma}$ (Ref. 28)

Strongly suppressed.

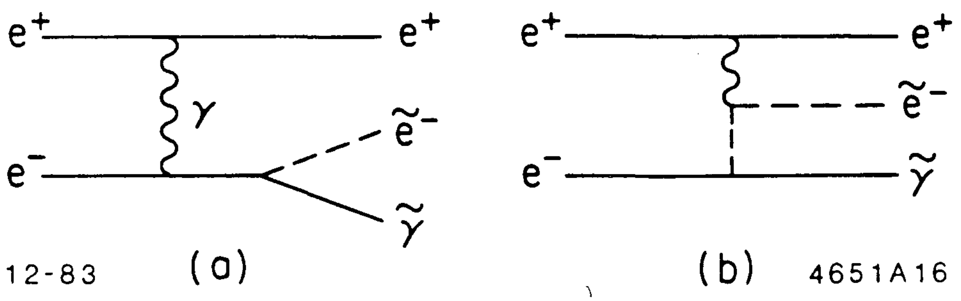


FIG 11

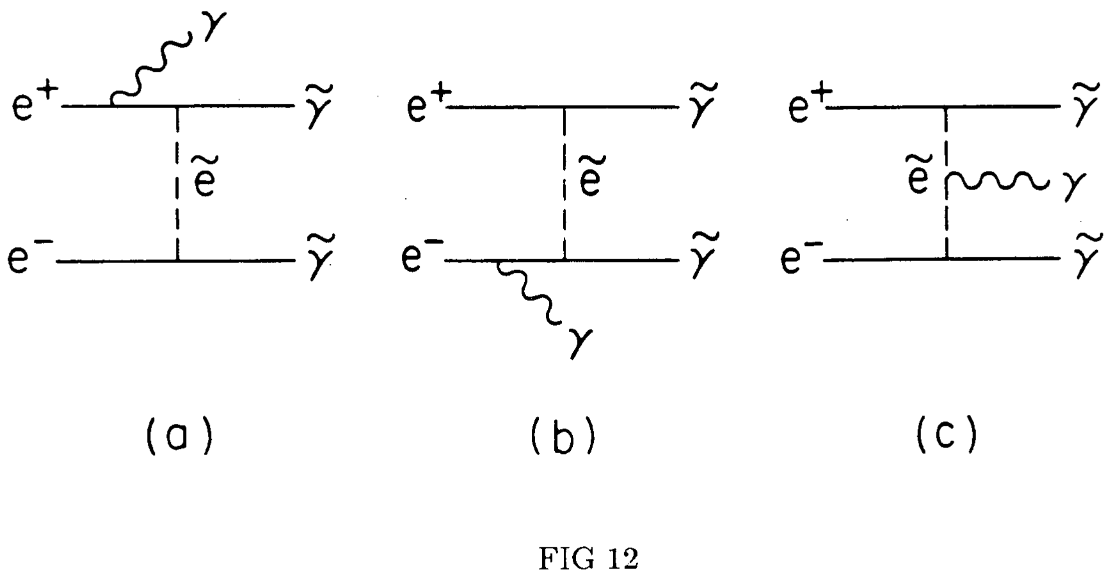


FIG 12

c) $Z^o \rightarrow e^+e^- \tilde{\gamma}\tilde{\gamma}$

Expect strong suppression.

d) $Z^o \rightarrow \tilde{\gamma}\,\tilde{\gamma}$ or $\tilde{\gamma}\tilde{\gamma}\gamma$

Only relevant if there is significant mixing of Higgsino in the "photino".

While it is true that $e^+e^- \rightarrow \tilde{e}^+\,\tilde{e}^-$ via a γ or Z^o has the usual β^3 threshold suppression for spin-zero particles, one should remember that other non-s-channel diagrams are allowed:

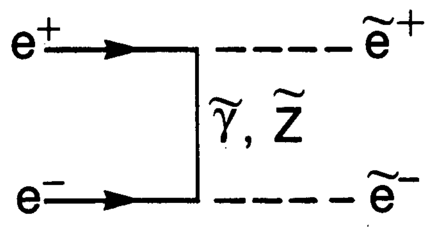


FIG 13

Hadron colliders are not ideal for the detection of scalar electrons or other scalar leptons. One can produce them in $q\bar{q} \rightarrow \gamma \rightarrow \tilde{e}^+\,\tilde{e}^-$ and look for electron pairs and large missing energy. The same process can occur via a physical Z^o boson, but may not be feasible given current limits. In $q\bar{q} \rightarrow W^+ \rightarrow \tilde{e}^+\,\tilde{\nu}$ we would expect for $M_{\tilde{e}} = 50$ GeV, a branching ratio to this mode of about 1%. This would require very large numbers of $W^\pm$ and would probably be obscured by backgrounds.

Limits on the masses of charged scalar leptons can be obtained from e^+e^- colliders, but one must remember that at a minimum they depend on $M_{\tilde{\gamma}}$ and on the ratio $M_{\tilde{e}_L}/M_{\tilde{e}_R}$. For the scalar electron, $\tilde{e}$, the best limits come from the process $e^+e^- \rightarrow \tilde{\gamma}\tilde{\gamma}\gamma$ but they are the most sensitive to $M_{\tilde{\gamma}}$. The limit from the MAC experiment[29] (90% C.L.) is

$$M_{\tilde{e}} > 43.5 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 0, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.1}$$

The limits from the ASP experiment[30] (90% C.L.) are

$$M_{\tilde{e}_L} \gtrsim 42 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 0, \qquad M_{\tilde{e}_R} >> M_{\tilde{e}_L} \tag{8.2}$$

$$M_{\tilde{e}} > 51 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 0, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.3}$$

$$M_{\tilde{e}} \gtrsim 48 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 5 \text{ GeV}, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.4}$$

$$M_{\tilde{e}} \gtrsim 33 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 10 \text{ GeV}, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.5}$$

If $M_{\tilde{\gamma}} > 13$ GeV, then no limit for $M_{\tilde{e}}$ exists from $e^+e^- \rightarrow \tilde{\gamma}\,\tilde{\gamma}\,\gamma$.

From the process $e^+e^- \rightarrow \tilde{e}\,\tilde{\gamma}\,e$, the JADE[31] and CELLO[32] experiments obtain (95% C.L.):

$$M_{\tilde{e}} > 25 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} = 0, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.6}$$

Similar limits are found by the MAC[33] and Mark II[34] collaborations. From $e^+e^- \rightarrow \tilde{e}^+\,\tilde{e}^-$ the JADE[31] and Mark J[35] experiments find (95% C.L.)

$$M_{\tilde{e}} > 22\text{–}23 \text{ GeV} \qquad \text{for } M_{\tilde{\gamma}} \leq 19 \text{ GeV}, \qquad M_{\tilde{e}_R} = M_{\tilde{e}_L} \tag{8.7}$$

Note that this limit holds for $M_{\tilde{\gamma}}$ up to 19 GeV.

The best limits for the scalar muon are for $M_{\tilde{\gamma}} < 15$–16 GeV (95% C.L.):

$$M_{\tilde{\mu}} > 20.9 \text{ GeV} \qquad (8.8)$$

These limits are from the JADE[39,40] and Mark J[35] collaborations although similar limits exist from the MAC,[41] TASSO,[37] and CELLO[36] collaborations.

For the scalar tau ($M_{\tilde{\gamma}} < 13$ GeV and 95% C.L.):

$$M_{\tilde{\tau}} > 18 \text{ GeV} \qquad (8.9)$$

Again these are from JADE[40] and Mark J[35] with similar results from CELLO,[36] TASSO,[42] Mark II[43] and MAC[38] collaborations.

9. Scalar Neutrinos

The characteristics of scalar neutrinos, $\tilde{\nu}$, are:

a) Their interactions are weak since charged scalar leptons have masses greater than 40–50 GeV (assuming $\tilde{\mu}$ and $\tilde{\tau}$ are not lighter than $\tilde{e}$), and since they interact via W or $\tilde{W}$.

b) There appear to be no mass limits from cosmological arguments.[10]

c) The only mass limit currently available is $M_{\tilde{\nu}} > 1.5$ GeV from considerations of τ lepton decay characteristics.[44]

d) Scalar neutrinos can be produced via decays and also directly in e^+e^- annihilation.

e) The most probable decay (see discussion below) is $\tilde{\nu} \to \nu\,\tilde{\gamma}$ which can only occur via triangle diagrams.

f) It is possible (though not found in most models) that $\tilde{\nu}$ is the *lightest* supersymmetric particle, and in that case it would be stable. Note that this would imply $\tilde{\gamma} \rightarrow \nu\,\tilde{\nu}$, so that the photino would also be an unobservable particle in this case (the $\tilde{\nu}$ leaves the detector without interacting).

g) If $M_{\tilde{\nu}} > M_{\tilde{W}}$, then the decays

$$\tilde{\nu} \rightarrow \nu\,\tilde{Z} \quad \text{or} \quad \tilde{\nu} \rightarrow e^-\,\tilde{W}^+ \tag{9.1}$$

are possible.[45]

The decay of $\tilde{\nu}$ is most likely $\tilde{\nu} \rightarrow \nu\,\tilde{\gamma}$ which occurs through the diagrams shown in Fig. 14. To reduce the degrees of freedom in this complex calculation, let us make the simplifying assumptions:

$$M(\tilde{W}) = M(\tilde{H}) = M(H) = M(W) \tag{9.2}$$

$$M(\tilde{Z}) = M(\tilde{H}^{o}) = M(Z) \tag{9.3}$$

Then the result of Barnett, Haber and Lackner[46,47] for the width of $\tilde{\nu}$ is

$$\Gamma(\tilde{\nu} \rightarrow \nu + \tilde{\gamma}) = \left(\frac{M_{\tilde{\nu}}\,\alpha^3}{128\,\pi^2 \sin^4 \theta_W} \right) \left| F\left(\frac{M_{\tilde{\nu}}^2}{m_W{}^2}\,,\,\frac{M_{\tilde{e}}^2}{m_W{}^2} \right) \right|^2 \tag{9.4}$$

or

$$\tau = 1.13 \times 10^{-16} \left(\frac{1\ \text{GeV}}{M_{\tilde{\nu}}} \right) \frac{1}{|F|^2}\ \text{sec} \tag{9.5}$$

where the quantity $|F|^2$ was calculated[47] to be:

$$|F|^2 \approx 0.1 - 10 \tag{9.6}$$

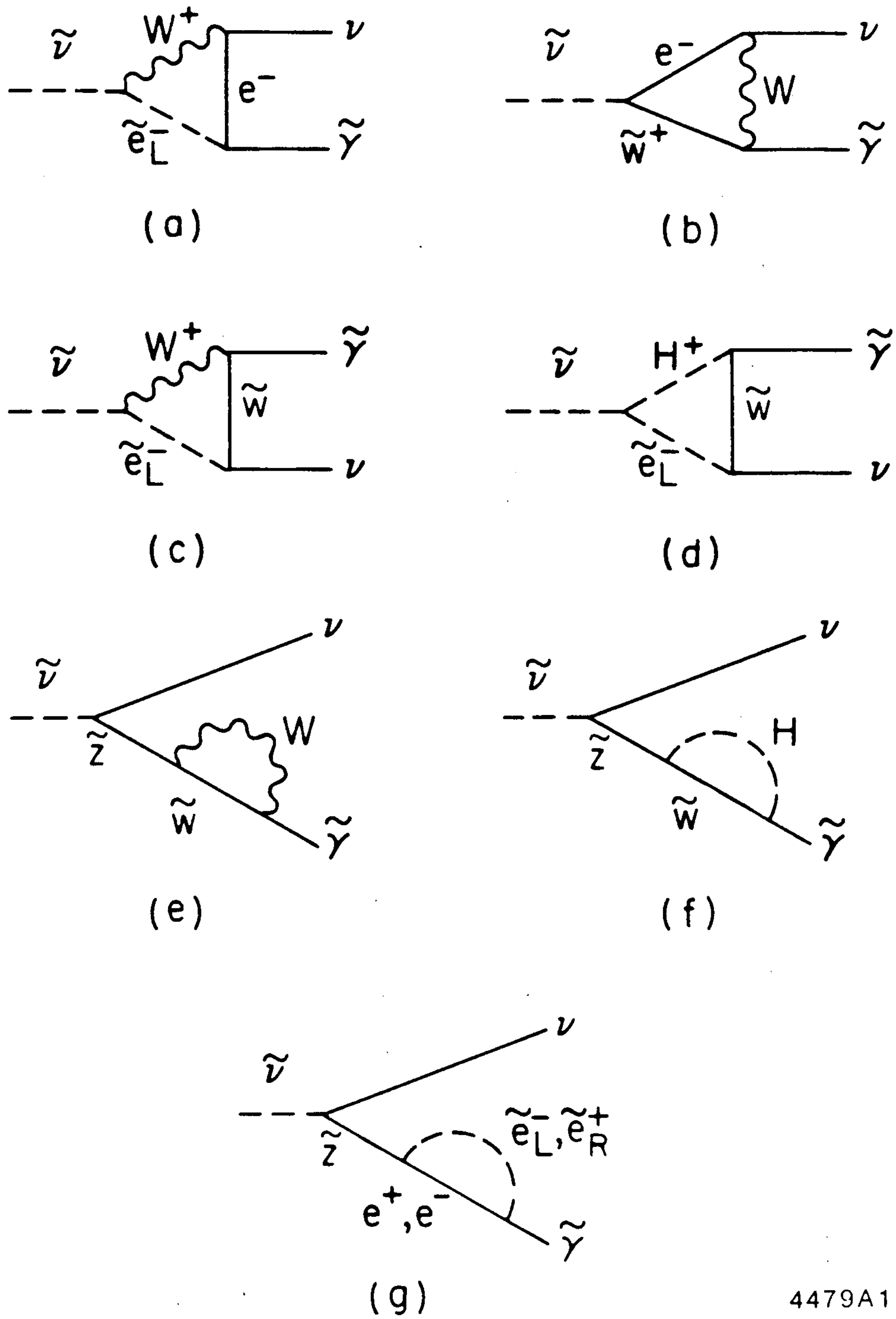


FIG 14

In many models this $\tilde{\nu} \to \nu\tilde{\gamma}$ decay dominates, leaving no visible signal for $\tilde{\nu}$. However, there are 3-body and 4-body decays (e.g., $\tilde{\nu}_e \to \nu_e\, e^+e^-\, \tilde{\gamma}$) which can be non-negligible if certain well-defined conditions occur. These decays are relevant if (and only if) any of the following conditions are met (from Ref. 47):

a) $M_{\tilde{\nu}} \approx M_{\tilde{q}} >> M_{\tilde{g}}$

This implies the decays $\tilde{\nu} \to e^-\, \tilde{g}\, q_1\, \overline{q}_2$ or $\tilde{\nu} \to \nu\, \tilde{g}\, q\, \overline{q}$,

and $\Gamma\,(\tilde{\nu} \to \text{4–body}) \sim \Gamma(\tilde{\nu} \to \nu\tilde{\gamma})$.

b) $M_{\tilde{\nu}} \approx M_{\tilde{e}} << M_{\tilde{q}}$

Here one would find $\tilde{\nu} \to \nu_e\, e^-\, \tilde{\gamma} + \{e^+, \mu^+, \tau^+\}$

with $\Gamma\,(\tilde{\nu} \to \text{4-body}) \sim 10^{-3}\,\Gamma\,(\tilde{\nu} \to \nu\tilde{\gamma})$.

c) $M_{\tilde{\nu}} > M_{\tilde{e}}$ (or $M_{\tilde{q}}$)

Then, of course, 3-body decays would dominate (see Fig. 15).

d) $M_{\tilde{\nu}_\mu} \gtrsim M_{\tilde{\nu}_e} + 0.5$ GeV

In this case one would find $\tilde{\nu}_\mu \to \mu^-\, e^+\, \tilde{\nu}_e$.

However, this condition is not expected to be found in models.

The available 4-body decay modes (depending on masses) are summarized in Fig. 16.

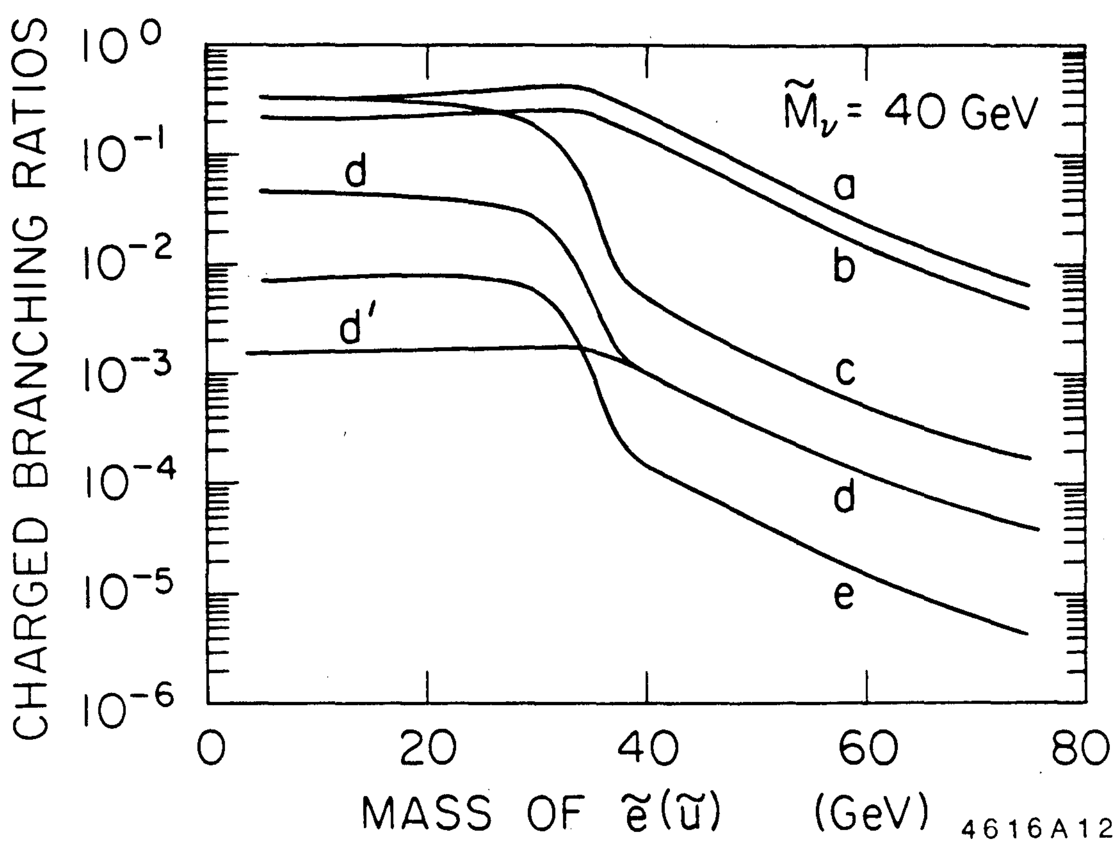
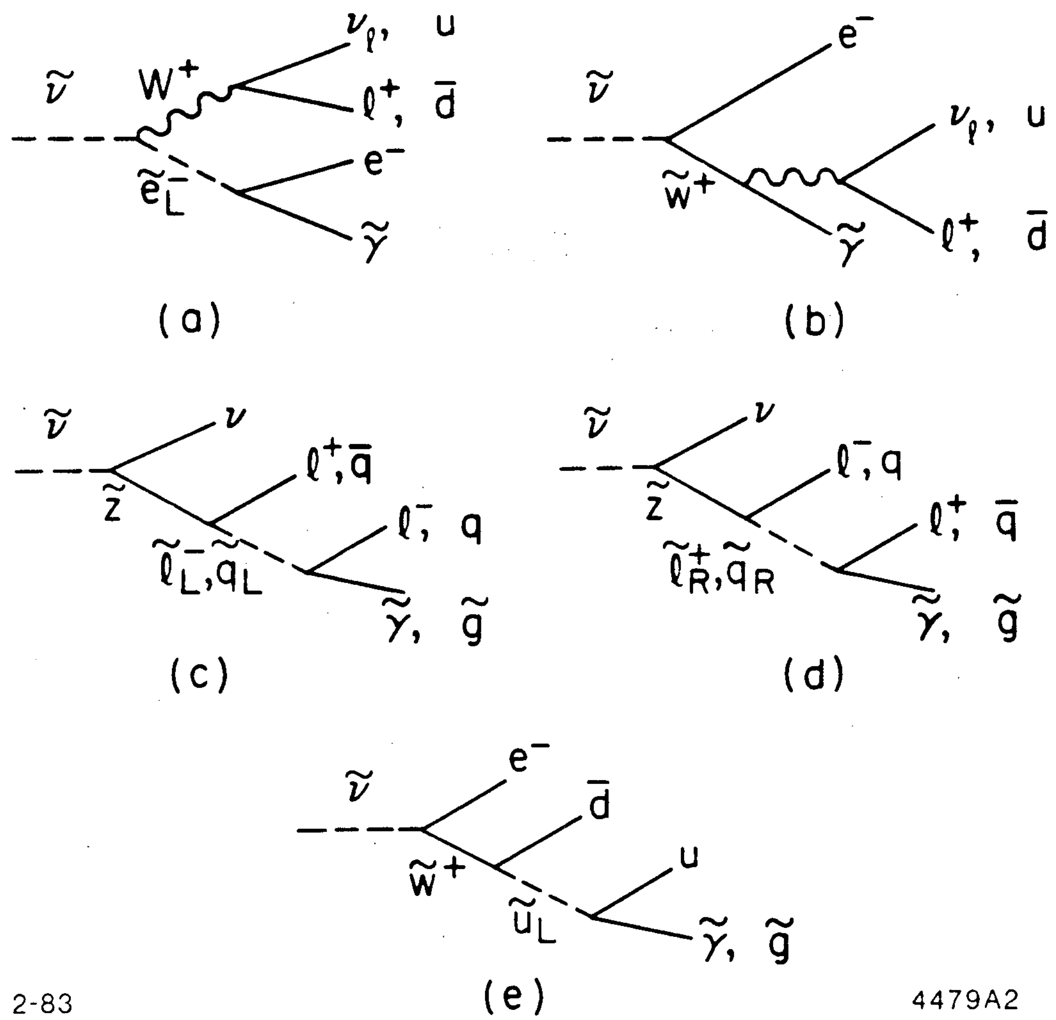

CHARGED BRANCHING RATIOS
10^0
10^{-1}
10^{-2}
10^{-3}
10^{-4}
10^{-5}
10^{-6}
0
20
40
60
80
MASS OF $\tilde{e}(\tilde{u})$ (GeV)
4616A12
$\tilde{M}_\nu$ = 40 GeV
a
b
c
d
d'
e

FIG 15



$\tilde{\nu}$
W^+
ν_ℓ, u
ℓ^+, $\bar{d}$
$\tilde{e}_L^-$
e^-
$\tilde{\gamma}$
(a)
$\tilde{\nu}$
e^-
ν_ℓ, u
$\tilde{w}^+$
$\tilde{\gamma}$
ℓ^+, $\bar{d}$
(b)
$\tilde{\nu}$
ν
ℓ^+,$\bar{q}$
$\tilde{z}$
ℓ^-, q
$\tilde{\ell}_L^-$,$\tilde{q}_L$
$\tilde{\gamma}$, $\tilde{g}$
(c)
$\tilde{\nu}$
ν
ℓ^-, q
$\tilde{z}$
ℓ^+ $\bar{q}$
$\tilde{\ell}_R^+$,$\tilde{q}_R$
$\tilde{\gamma}$, $\tilde{g}$
(d)
$\tilde{\nu}$
e^-
$\bar{d}$
$\tilde{w}^+$
u
$\tilde{u}_L$
$\tilde{\gamma}$, $\tilde{g}$
(e)
2-83
4479A2

FIG 16

The production of $\tilde{\nu}$ can occur via $e^+e^- \rightarrow \tilde{\nu}$ + other. Ideally, the best way to produce $\tilde{\nu}$ would be in Z^o or $W^\pm$ decay.[48–50] For Z^o decay the branching ratio to $\tilde{\nu}\,\tilde{\bar{\nu}}$ is:

$$\frac{\Gamma\,(Z^o \rightarrow \tilde{\nu}\,\tilde{\bar{\nu}})}{\Gamma\,(Z^o \rightarrow \nu\,\bar{\nu})} = \frac{1}{2}\left(1 - 4\,\frac{M_{\tilde{\nu}}^2}{M_Z{}^2}\right)^{3/2} \tag{9.7}$$

For W^+ decay the branching ratio to $\tilde{e}_L{}^+\,\tilde{\nu}_e$ is:

$$\frac{\Gamma\,(W^+ \rightarrow \tilde{e}_L{}^+\,\tilde{\nu}_e)}{\Gamma\,(W^+ \rightarrow e^+\nu_e)} = \frac{1}{2}\left(\frac{(m_W{}^2 - M_{\tilde{\nu}}^2 - M_{\tilde{e}}^2)^2 - 4\,M_{\tilde{\nu}}^2\,M_{\tilde{e}}^2}{m_W{}^4}\right)^{3/2} \tag{9.8}$$

The factors of $\frac{1}{2}$ which occur in Eqs. (9.7)-(9.8) reflect the factor

$$\frac{g_s{}^2}{2\,(g_L{}^2 + g_R{}^2)} \tag{9.9}$$

(where g_s is the coupling of the vector boson to scalars). For weak interactions involving the neutrino, we have $g_S = g_L = g$ and $g_R = 0$ giving the factor $\frac{1}{2}$. Note that for electromagnetic interactions $g_S = g_L = g_R = e$ which gives the famous factor of $\frac{1}{4}$ for scalar particles in e^+e^- annihilation.

Let us now consider the signals for scalar neutrinos from W and Z decay.[48–51]

a) $W^+ \rightarrow \tilde{e}_L{}^+\,\tilde{\nu}$

with $\tilde{e}_L{}^+ \rightarrow e^+\tilde{\gamma}$ and $\tilde{\nu} \rightarrow \nu\,\tilde{\gamma}$.

This signal is similar to that for $W^+ \rightarrow e^+\nu$; however, the final-state electron is softer and has a $\sin^2\theta$ distribution instead of the $(1 \pm \cos\theta)^2$ distribution of $(e^+ \nu)$. It may be difficult to separate from backgrounds[50–54] such as $W^+ \rightarrow \tau\,\nu$ (with $\tau \rightarrow e\,\nu\,\overline{\nu}$) and $b \rightarrow c\,e\,\nu$.

b) $W^+ \rightarrow \tilde{e}_L^{\,+}\,\tilde{\nu}$

with $\tilde{e}_L^{\,+} \rightarrow e^+\tilde{\gamma}$ and $\tilde{\nu} \rightarrow$ 3- or 4-body

Since this latter decay probably implies $M_{\tilde{\nu}} \gtrsim M_{\tilde{e}}$ and since $M_{\tilde{e}} \gtrsim$ 40–50 GeV experimentally, this process is very unlikely to occur (unless $M_{\tilde{\gamma}} \geq 13$ GeV).

c) $Z^o \rightarrow \tilde{\nu}\,\overline{\tilde{\nu}}$

with $\tilde{\nu}$ and $\overline{\tilde{\nu}} \rightarrow \nu\tilde{\gamma}$

This is only observable in "neutrino-counting" experiments[55] where $e^+e^- \rightarrow \gamma$ + nothing visible. One would have to run above the Z^o and search for $\gamma\, Z^o$.

d) $Z^o \rightarrow \tilde{\nu}\,\overline{\tilde{\nu}}$

with $\tilde{\nu} \rightarrow \nu\,\tilde{\gamma}$ and $\overline{\tilde{\nu}} \rightarrow$ 3- or 4-body decay

This decay chain leads to striking one-sided signatures in $e^+e^- \rightarrow Z^o$, as in the simulated event[50] shown in Fig. 17 (where the

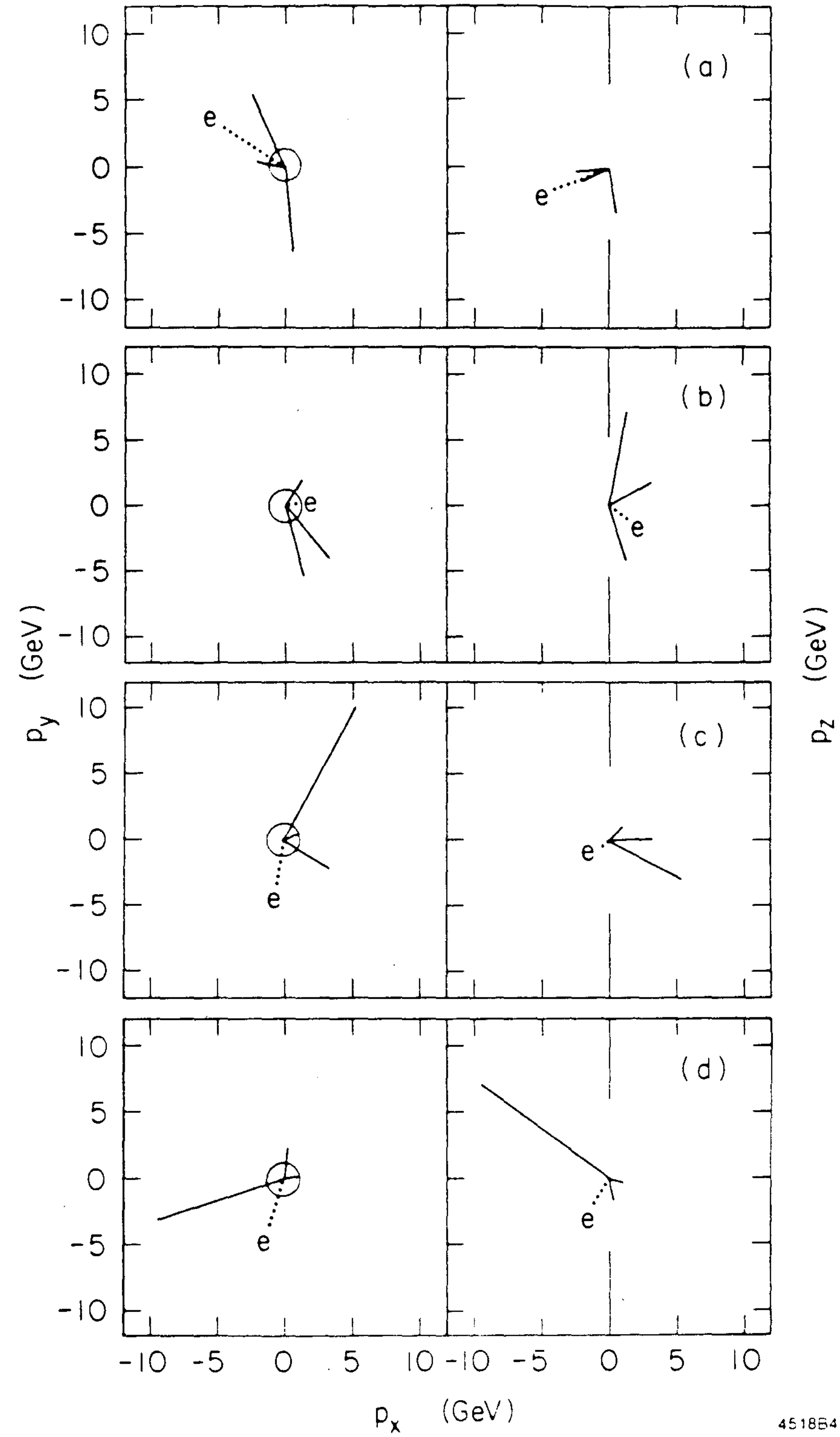


FIG 17

circle represents the beam pipe) (the solid lines are two quarks and a gluino). Here one sees tracks in one hemisphere, and no balancing momenta in the other hemisphere. But again we need $M_{\tilde{\nu}} \gtrsim M_{\tilde{e}} \gtrsim$ 40–50 GeV so that this process may be kinematically forbidden.

e) $Z^o \to \tilde{\nu}_e \bar{\tilde{\nu}}_e$

with $\tilde{\nu}_e \to \nu\tilde{\gamma}$ and $\bar{\tilde{\nu}}_e \to \nu_\mu e^- \mu^+ \tilde{\gamma}$

This is my personal favorite since it yields a spectacular signature of an e^- and a μ^+ on one side, and nothing at all on the opposite side:

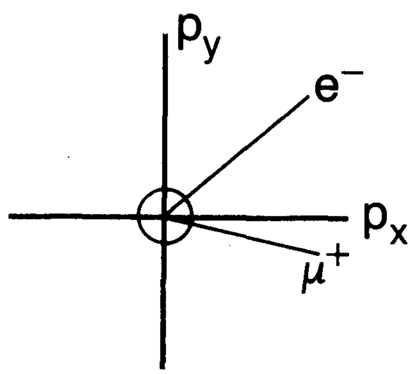


FIG 18

In any circumstances it is an extremely rare event; for $M_{\tilde{\nu}}$ = 30–50 GeV and $M_{\tilde{e}}$ = 50 GeV it occurs for 10^{-3}–10^{-4} of all $\tilde{\nu}$ decays. As the limit on $M_{\tilde{e}}$ increases, this will become impractical from the Z^o boson.

These processes can also occur in continuum e^+e^- as in these diagrams:

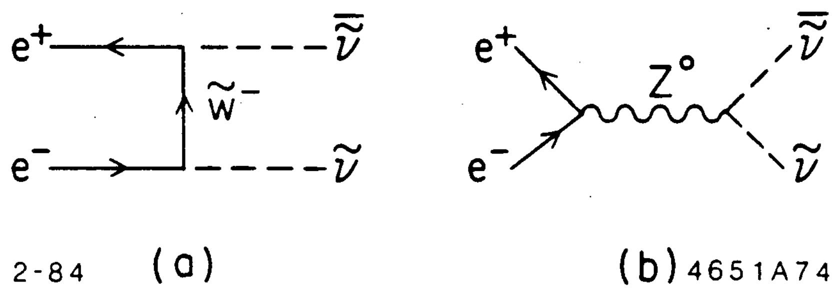


FIG 19

To get a 3–4-body decay (a visible decay) requires $M_{\tilde{\nu}} \gtrsim (M_{\tilde{e}} - 10\ \text{GeV}) \gtrsim$ 20–30 GeV (see Fig. 15). Therefore PEP and PETRA may have insufficient energy. However in a machine with $s = 2\ m_W^2$, we have calculated (Ref. 50) assuming $M_{\tilde{\nu}}^2 << s$ and $M_{\tilde{W}} \approx m_W$ that

$$R_{\tilde{\nu}\tilde{\bar{\nu}}} \equiv \frac{\sigma(\tilde{\nu}\,\tilde{\bar{\nu}})}{\sigma_{\mu\mu}\ (\text{point})} \approx 1.5 \tag{9.10}$$

where $R_{total} \approx 10$. The ratio $R_{\tilde{\nu}\tilde{\bar{\nu}}}$ may be enhanced in some models, and one may observe some of these interesting one-sided events.

The process $e^+e^- \to \gamma\, \tilde{\nu}\, \tilde{\bar{\nu}}$ through the diagrams below

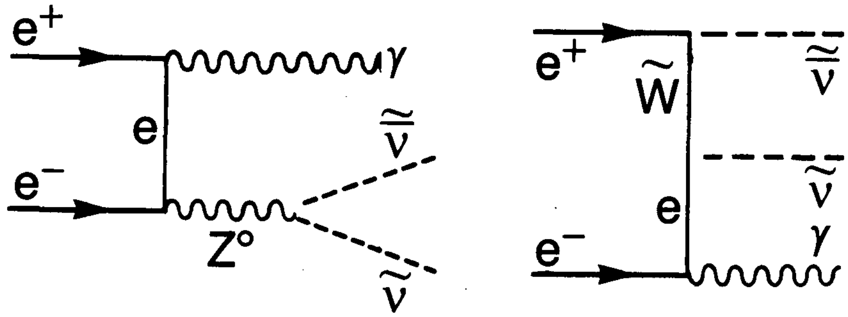


FIG 20

has been calculated by Ware and Machacek[56] and is discussed in David Burke's lectures[24] at this Summer Institute.

The only mass limits for $\tilde{\nu}$ come from the paper of Kane and Rolnick.[44] They consider the decay $\tau^+ \to \tilde{\nu}_\tau \ell^+ \tilde{\nu}_\ell$ (see below).

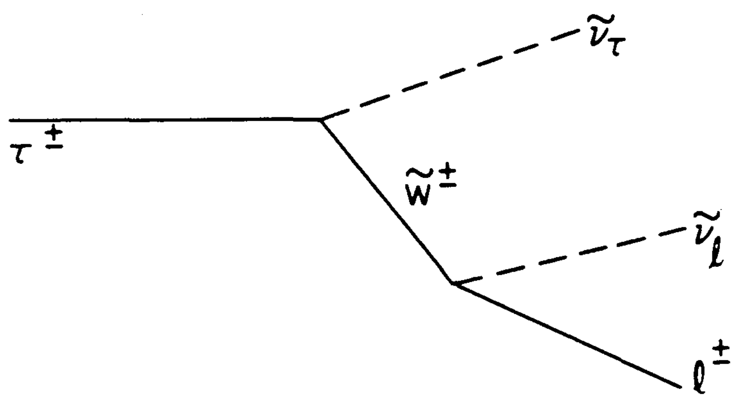


FIG 21

Depending on the value of $M_{\widetilde{W}}$ they conclude that $M_{\tilde{\nu}} \gtrsim m_\tau$; otherwise the decay properties of τ would be modified.

10. Gluinos and Scalar Quarks

10.A. General Characteristics

Let us first summarize the general characteristics of gluinos ($\tilde{g}$) and scalar quarks ($\tilde{q}$ or squarks).

a) They are strongly interacting particles.

b) Strong lower limits for $\tilde{q}$ and $\tilde{g}$ masses have recently been derived;[57] see discussion below.

c) I will generally assume $M_{\tilde{q}_L} = M_{\tilde{q}_R}$. One expects that the masses of the scalar quarks $\tilde{u}$, $\tilde{d}$, $\tilde{s}$, $\tilde{c}$, and $\tilde{b}$, will be approximately equal. Supersymmetric models[58] predict this, and this approximate equality is necessary to suppress flavor-changing neutral currents.[59] The masses $M_{\tilde{t}_L}$ and $M_{\tilde{t}_R}$ might be different from the others.[60,61]

d) They are easy to produce. This is especially true for gluinos which can be produced via the diagram.

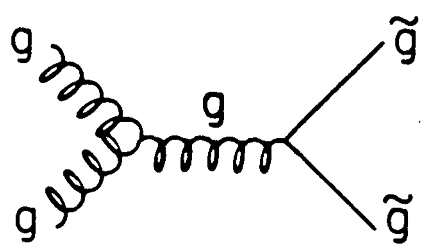


FIG 22

which has no heavy intermediate particle and large color factors. It is suppressed only by phase space.

e) The best means of detection is via missing energy (in decays).

f) The gluino is a self-conjugate, Majorana fermion with C = –1.

g) The decays of both $\tilde{q}$ and $\tilde{g}$ depend on which one is heavier.

h) If $M_{\tilde{q}} > M_{\tilde{W}}$ or $M_{\tilde{Z}}$, then scalar quarks can decay into $\tilde{W}$ or $\tilde{Z}$. I refer the reader to the paper of Baer et al.[62] for a discussion of this case.

10.B. Decays of Gluinos and Scalar Quarks

Let us continue with the subject of decays. If the gluino is heavier than the scalar quark ($M_{\tilde{g}} > M_{\tilde{q}}$) and if the scalar quark is lighter than $\tilde{W}$ and $\tilde{Z}$, then the scalar quark can only decay into a quark and photino:

$$\tilde{q} \rightarrow q\,\tilde{\gamma} \tag{10.1}$$

The gluino can decay into the scalar quark:

$$\tilde{g} \rightarrow \tilde{q}\,q \tag{10.2}$$

In the opposite case where $M_{\tilde{g}} < M_{\tilde{q}}$, the gluino has to decay via a virtual scalar quark:

$$\tilde{g} \rightarrow q\,\overline{q}\,\tilde{\gamma} \tag{10.3}$$

Here the scalar quark can decay into the gluino; however, the suppressed decay into a photino carries off much more missing energy and can be the important decay for detection purposes. These decays are:

$$\tilde{q} \rightarrow q\, \tilde{g} \qquad (10.4a)$$

$$\tilde{q} \rightarrow q\, \tilde{\gamma} \qquad (10.4b)$$

Their relative decay widths have been calculated by Haber and Kane:[4]

$$\Gamma\, (\tilde{q} \rightarrow q\, \tilde{g}) \;=\; \frac{4}{3}\, \alpha_s\, \frac{(M^2_{\tilde{q}} - M^2_{\tilde{g}} - m_q{}^2)}{M^2_{\tilde{q}}}\; \frac{\lambda^{1/2}\, (M^2_{\tilde{q}}, M^2_{\tilde{g}}, m_q{}^2)}{2\, M_{\tilde{q}}} \qquad (10.5a)$$

$$\Gamma\, (\tilde{q} \rightarrow q\, \tilde{\gamma}) \;=\; \alpha\, e_q{}^2\, \frac{(M^2_{\tilde{q}} - M^2_{\tilde{\gamma}} - m_q{}^2)}{M^2_{\tilde{q}}}\; \frac{\lambda^{1/2}\, (M^2_{\tilde{q}}, M^2_{\tilde{\gamma}}, m_q{}^2)}{2\, M_{\tilde{q}}} \qquad (10.5b)$$

where $\lambda(x,y,z) \equiv (x^2 - y^2 - z^2)^2 - 4y^2z^2$.

If we took $M_{\tilde{g}} \approx M_{\tilde{\gamma}}$, or $M_{\tilde{g}}$ and $M_{\tilde{\gamma}} << M_{\tilde{q}}$, then their ratio is

$$\frac{3\, \alpha\, e_q{}^2}{4\, \alpha_s} \qquad (10.6)$$

If $M_{\tilde{\gamma}} << M_{\tilde{g}} << M_{\tilde{q}}$, then Haber and Kane[4] found for the gluino width

$$\Gamma\, (\tilde{g} \rightarrow q\, \bar{q}\, \tilde{\gamma}) = \frac{\alpha\, \alpha_s\, e_q{}^2}{96\, \pi}\, M^5_{\tilde{g}} \left(\frac{1}{M^4_{\tilde{q}_L}} + \frac{1}{M^4_{\tilde{q}_R}} \right) \qquad (10.7)$$

Other decays are possible if $M_{\tilde{W}}$ or $M_{\tilde{Z}} < M_{\tilde{q}}$ or $M_{\tilde{g}}$.

10.C. Production in e^+e^- Annihilation

Turning to the *production* of scalar quarks in e^+e^- annihilation (where gluino production is not practical), the production cross section due to one-photon exchange is given by

$$\frac{d\sigma}{d\Omega} = \left(\frac{3\, \alpha^2}{8\, s} \right) e_q{}^2\, \beta^3 \sin^2 \theta \qquad (10.8)$$

with $\beta^2 = (1 - 4M^2_{\tilde{q}}/s)$. The $\beta^3 \sin^2 \theta$ factor is the usual factor for scalars. If the decay $\tilde{q} \rightarrow q\tilde{\gamma}$ dominates, one might see events such as shown below (where the circle represents the beam pipe):

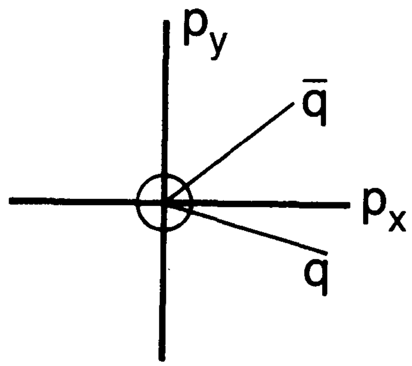


FIG 23

The hadronized final-state may have two acoplanar jets with missing energy. Recently reported preliminary results from the JADE collaboration[63] searching for such events (and measuring R_{hadron}) set the limits (95% C.L.):

$$M_{\tilde{q}} > 21.4 \text{ GeV} \qquad \text{for} \quad M_{\tilde{q}_1} = M_{\tilde{q}_2} \tag{10.9}$$

$$M_{\tilde{q}_1} < 3.2 \quad \text{or} \quad M_{\tilde{q}_1} > 21.0 \text{ GeV} \qquad \text{for} \quad M_{\tilde{q}_1} << M_{\tilde{q}_2} \tag{10.10}$$

where they take $M_{\tilde{\gamma}} < 10$ GeV.

If however the decay $\tilde{q} \to \tilde{g}\, q$ dominates, then the resulting events are more spherical. For $M_{\tilde{g}} = 3$–10 GeV and $M_{\tilde{\gamma}} = (1/6)M_{\tilde{g}}$, JADE finds[63]

$$M_{\tilde{q}} > 19.2 \text{ GeV} \qquad \text{for} \quad M_{\tilde{q}_1} = M_{\tilde{q}_2} \tag{10.11}$$

$$M_{\tilde{q}_1} < 11.3 \quad \text{or} \quad M_{\tilde{q}_1} > 17.8 \text{ GeV} \qquad \text{for} \quad M_{\tilde{q}_1} << M_{\tilde{q}_2} \tag{10.12}$$

In finding such limits it is probably safe to assume that the first five flavors of quarks appear at the same threshold (within resolution) thereby enhancing the rate. One is, of course, subject to the usual factor of $\frac{1}{4}$ which appears for scalars.

10.D. Beam-Dump Experiments

If gluinos are very light, beam-dump experiments may be the best means to search for them. Beam-dump experiments were described in detail in Sec.7. Here, it is the results which concern us. There are new results from the BEBC experiment[64] which set the limits (90% C.L.):

$$M_{\tilde{g}} > 3 \text{ GeV} \quad \text{for} \quad M_{\tilde{q}} \leq 150 \text{ GeV} \tag{10.13}$$

$$M_{\tilde{g}} > 4 \text{ GeV} \quad \text{for} \quad M_{\tilde{q}} \leq 65 \text{ GeV} \tag{10.14}$$

It should be noted that these experiments cannot rule out extremely small $M_{\tilde{g}}$ (approximately 0.5 GeV or less) or large $M_{\tilde{q}}$, because the lifetime of the gluino would increase until it can interact in the dump before it decays ($\tau \gtrsim 5 \times 10^{-11}$ sec) and eventually until it can leave the BEBC experiment before it decays ($\tau \gtrsim 10^{-9}$ sec).

10.E. Production at Hadron Colliders and Mass Limits

Let us turn now to a subject we will consider in detail: the production of gluinos and scalar quarks at high-energy hadron colliders. The reader may have noted an enormous literature[65–78] on this subject. The cause of this excitement was the observation in the 1983 and 1984 runs of the UA1 experiment[79,80] at CERN of events which are labeled "monojets". These are events in $p\bar{p}$ collisions in which a cluster of hadrons (i.e., a "jet") emerges on one side of the beam, but no balancing momenta appear on the opposite side.

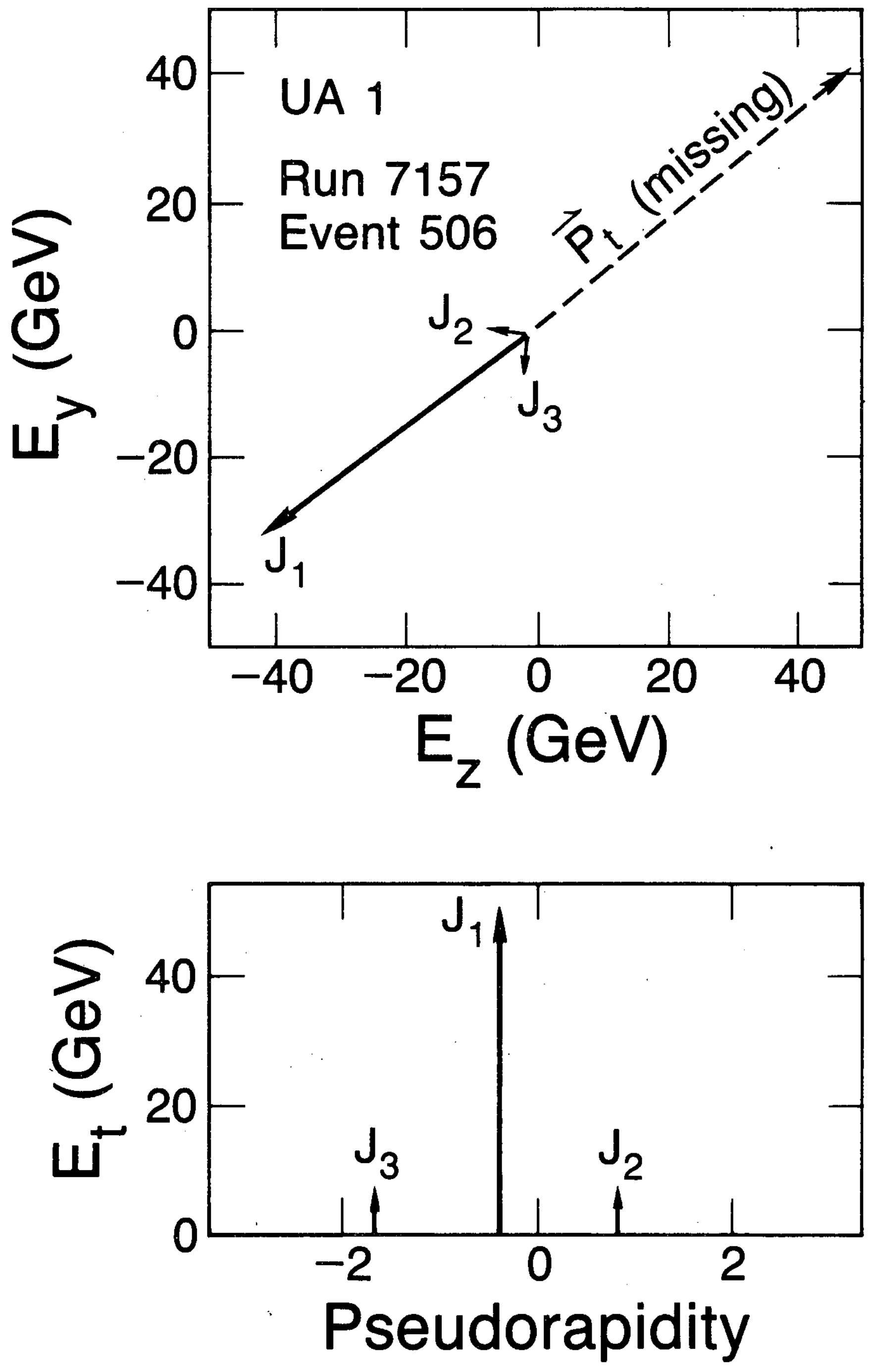
This leaves the impression of missing energy which is, as stated earlier, a key signature of supersymmetry. But while missing energy is a necessary sign for supersymmetry, it is certainly not sufficient evidence to prove supersymmetry. Detailed calculations both of supersymmetry and of backgrounds were needed to sort out this possible signal.

In $p\bar{p}$ collisions the missing energy can be measured accurately only in the plane *transverse* to the beam. This is because the beam and the longitudinal spray from the collision require small holes in the detector (along the beam). Although these holes account for only 0.2° in the UA1 detector,[79] half the total energy of each event tends to escape through the holes. Since this longitudinal energy is lost, only transverse missing energy is a useful quantity.

What the UA1 Collaboration observes[79,80] are events with large missing transverse energy ($E_T^{miss} = 15$–50 GeV) and a hadronic jet with $E_T^{jet} > 15$–25 GeV depending on the trigger. If no other hadronic cluster has

$$E_T(\text{jet } 2) \geq 12 \text{ GeV} \quad , \tag{10.15}$$

the event is labeled a "monojet". If a *secondary* jet does satisfy this criterion, it is labeled a "dijet" (with large missing transverse energy). Events with more than one *secondary* jet, I will call "multijets". An example of a UA1 monojet is shown in Fig. 24. This event shows that the distinction between a monojet and a dijet can be a fine one for theorists. Most events (like this


UA 1
Run 7157
Event 506
$\vec{P}_t$ (missing)
J_1
J_2
J_3
40
20
0
−20
−40
E_y (GeV)
−40
−20
0
20
40
E_z (GeV)
J_1
J_2
J_3
40
20
0
E_t (GeV)
−2
0
2
Pseudorapidity

one) have a secondary cluster (or jet); a number of factors can easily add or subtract from this jet's energy thereby affecting whether or not it satisfies condition (10.15).

The question arises as to whether the events with missing energy result from the production of gluinos and/or scalar quarks. Some sample cross sections from Dawson, Eichten, Quigg[5] for gluinos for various masses and interaction energies are shown in Figs. 25 and 26.

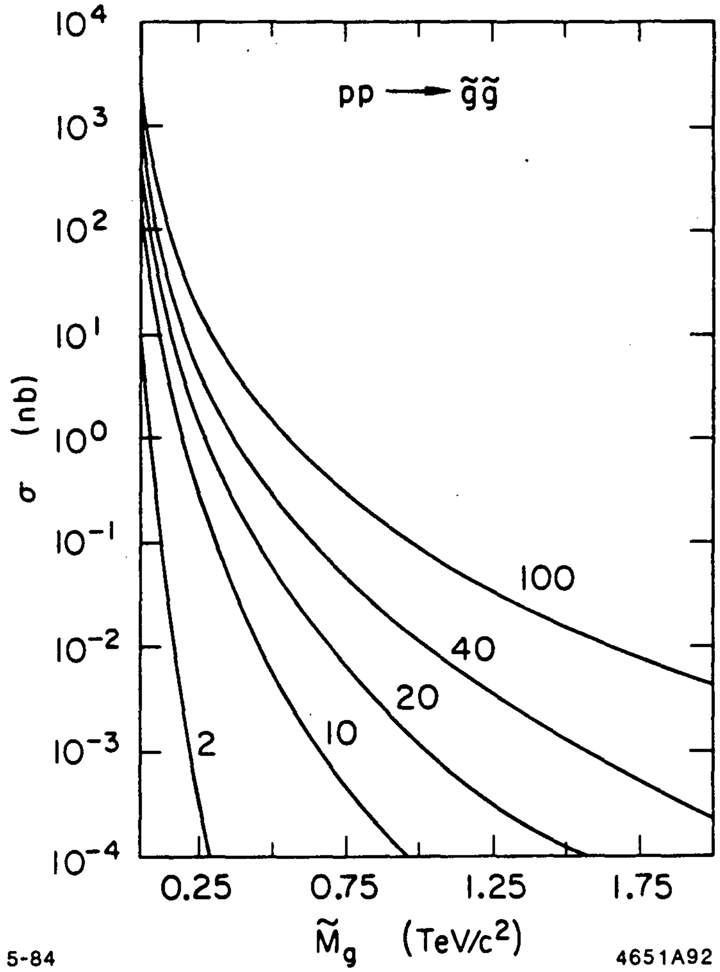
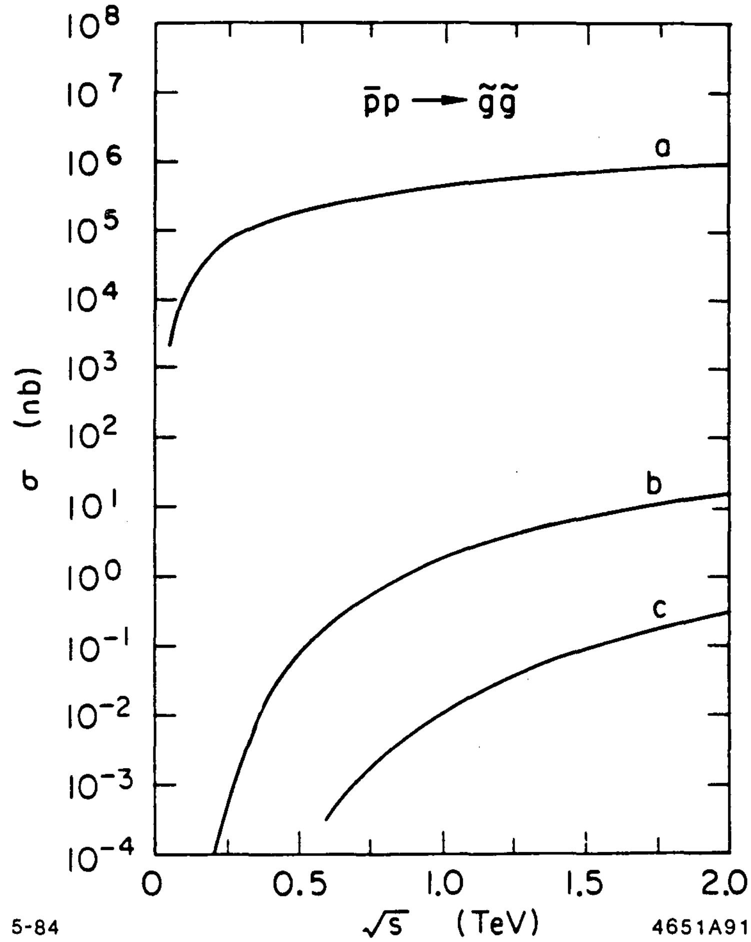


FIG 25



$\bar{p}p \longrightarrow \tilde{g}\tilde{g}$
a
b
c
10^{8}
10^{7}
10^{6}
10^{5}
10^{4}
10^{3}
10^{2}
10^{1}
10^{0}
10^{-1}
10^{-2}
10^{-3}
10^{-4}
σ (nb)
0
0.5
1.0
1.5
2.0
$\sqrt{s}$ (TeV)
5-84
4651A91

FIG 26

The next step is to consider which processes lead to the production of gluinos and/or scalar particles. These are:

(a) gluino pair production (see Fig. 27a-b)

$$gg \rightarrow \tilde{g}\tilde{g} \tag{10.16a}$$

$$q\bar{q} \rightarrow \tilde{g}\tilde{g} \tag{10.16b}$$

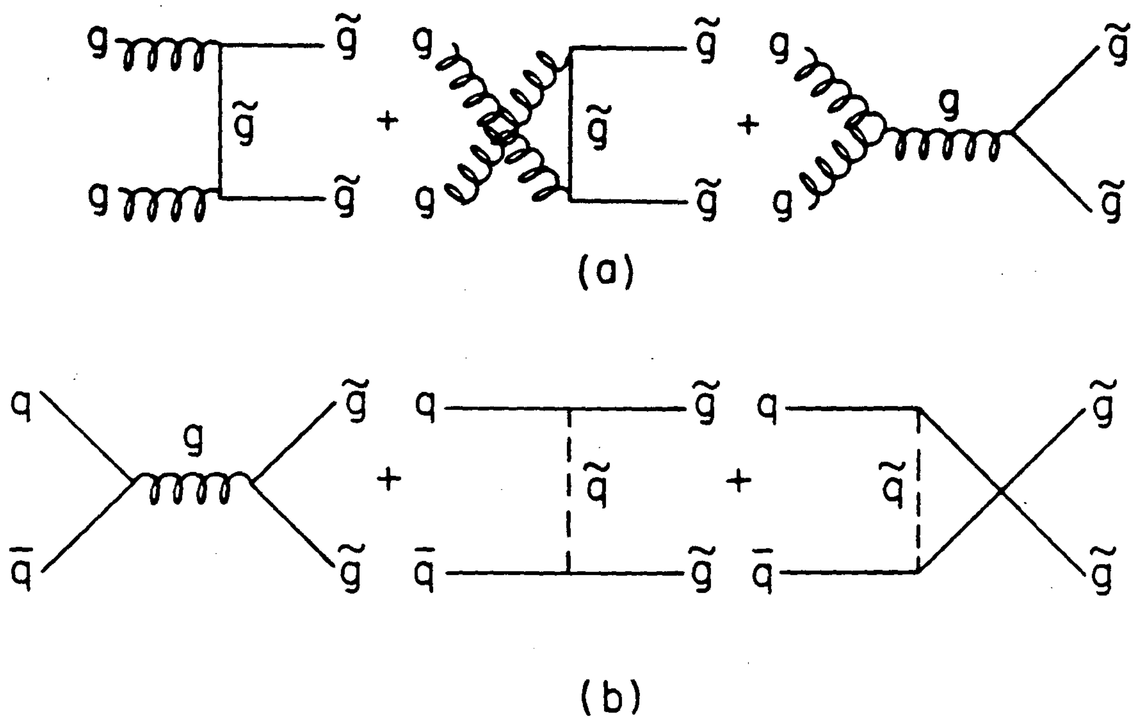


FIG 27a-b

(b) scalar-quark pair production (see Fig. 27c-f)

$$gg \rightarrow \tilde{q}\bar{\tilde{q}} \tag{10.17a}$$

$$q\bar{q} \rightarrow \tilde{q}\bar{\tilde{q}} \tag{10.17b}$$

$$qq \rightarrow \tilde{q}\tilde{q} \tag{10.17c}$$

$$q\bar{q} \rightarrow W \rightarrow \tilde{q}\bar{\tilde{q}} \tag{10.17d}$$

$$q\bar{q} \rightarrow Z \rightarrow \tilde{q}\bar{\tilde{q}} \tag{10.17e}$$

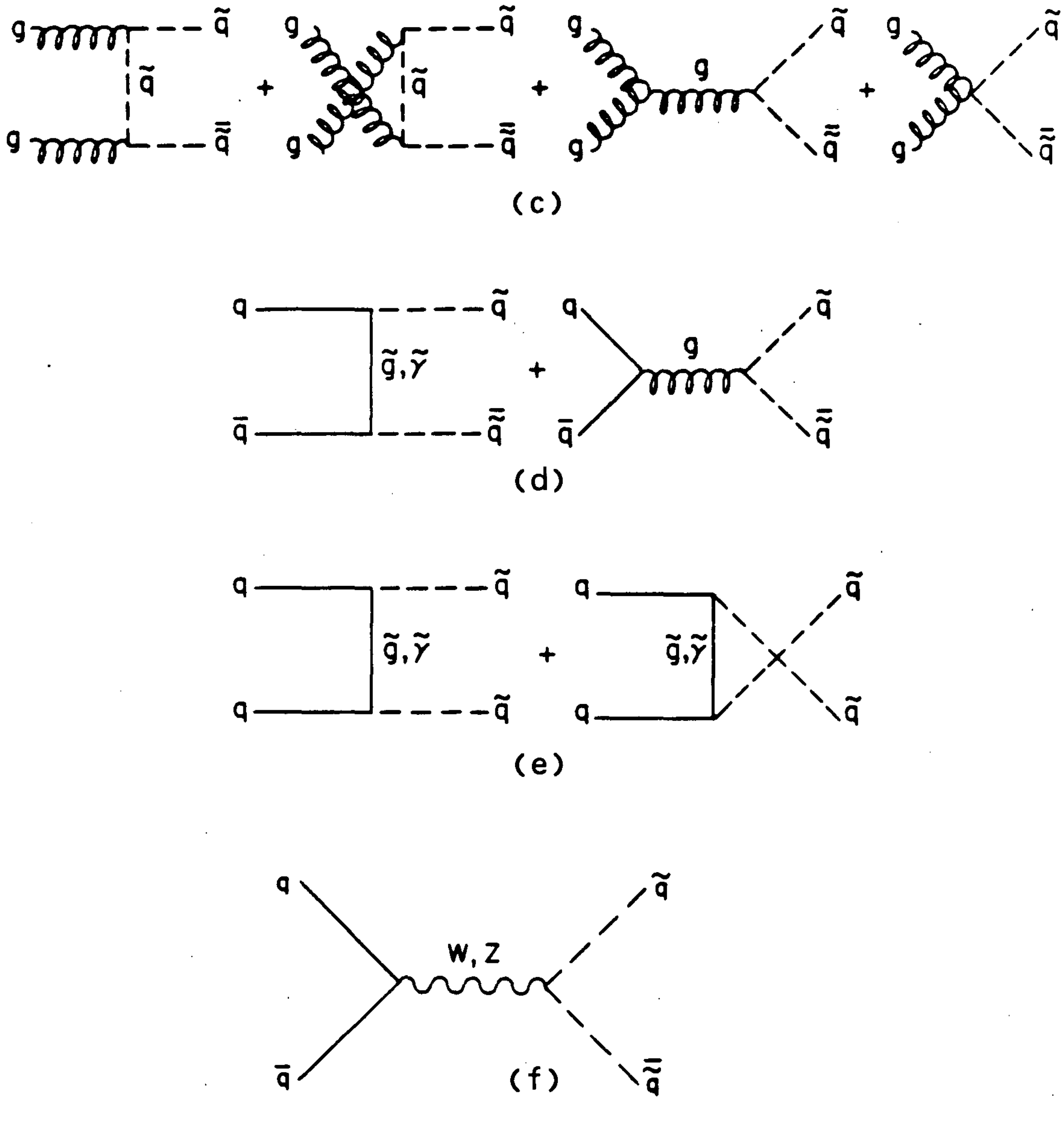


FIG 27c-f

(c) Associated production of gluinos and scalar-quarks (see Fig. 28a)

$$q g \rightarrow \tilde{q}\tilde{g} \tag{10.18}$$

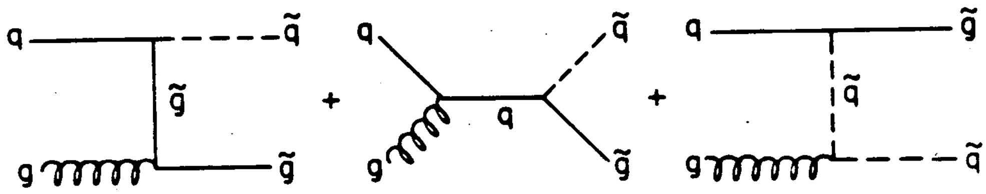


FIG 28a

(d) Associated production of gluinos and photinos (see Fig. 28b)

$$q\bar{q} \rightarrow \tilde{g}\tilde{\gamma} \tag{10.19}$$

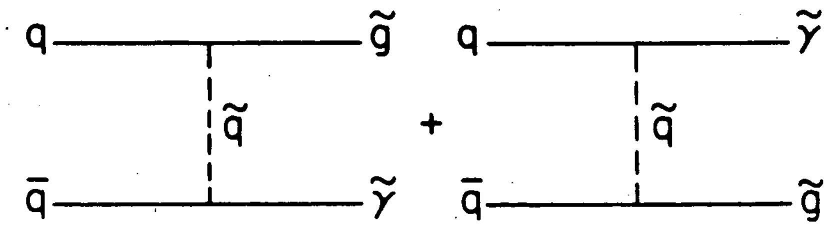


FIG 28b

(e) Associated production of scalar-quarks and photinos (see Fig. 28c)

$$q g \rightarrow \tilde{q}\tilde{\gamma} \tag{10.20}$$

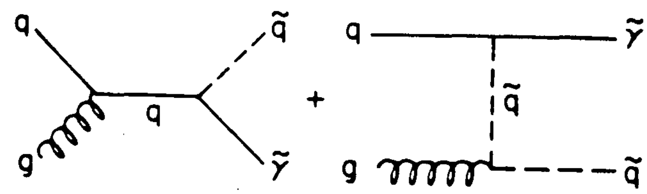


FIG 28c

Processes (10.19) and (10.20) produce a photino directly rather than in a decay. They are suppressed by a power of α, but they are more likely to pass experimental cuts and triggers requiring large missing transverse energy.

These diagrams all need to be convoluted with the decay diagrams for gluinos and scalar quarks (which depend on which one is heavier), see processes (10.1)-(10.4) and Fig. 29. It is the full process (parton + parton to multi-quark and multi-photino final state) which theorists must calculate.

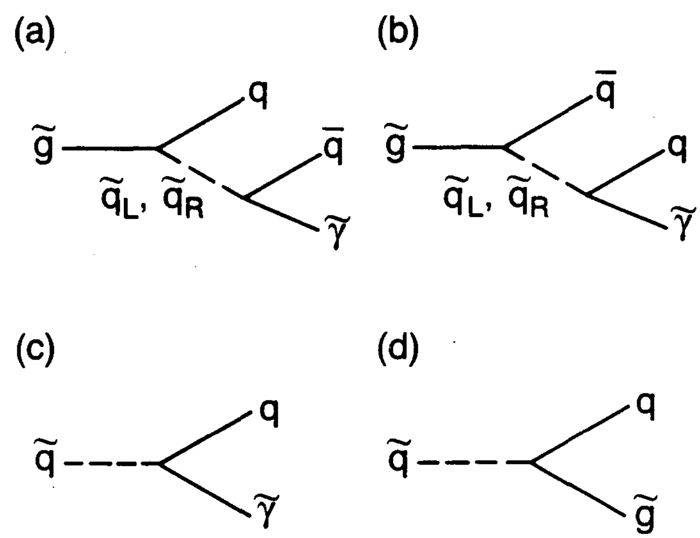


FIG 29

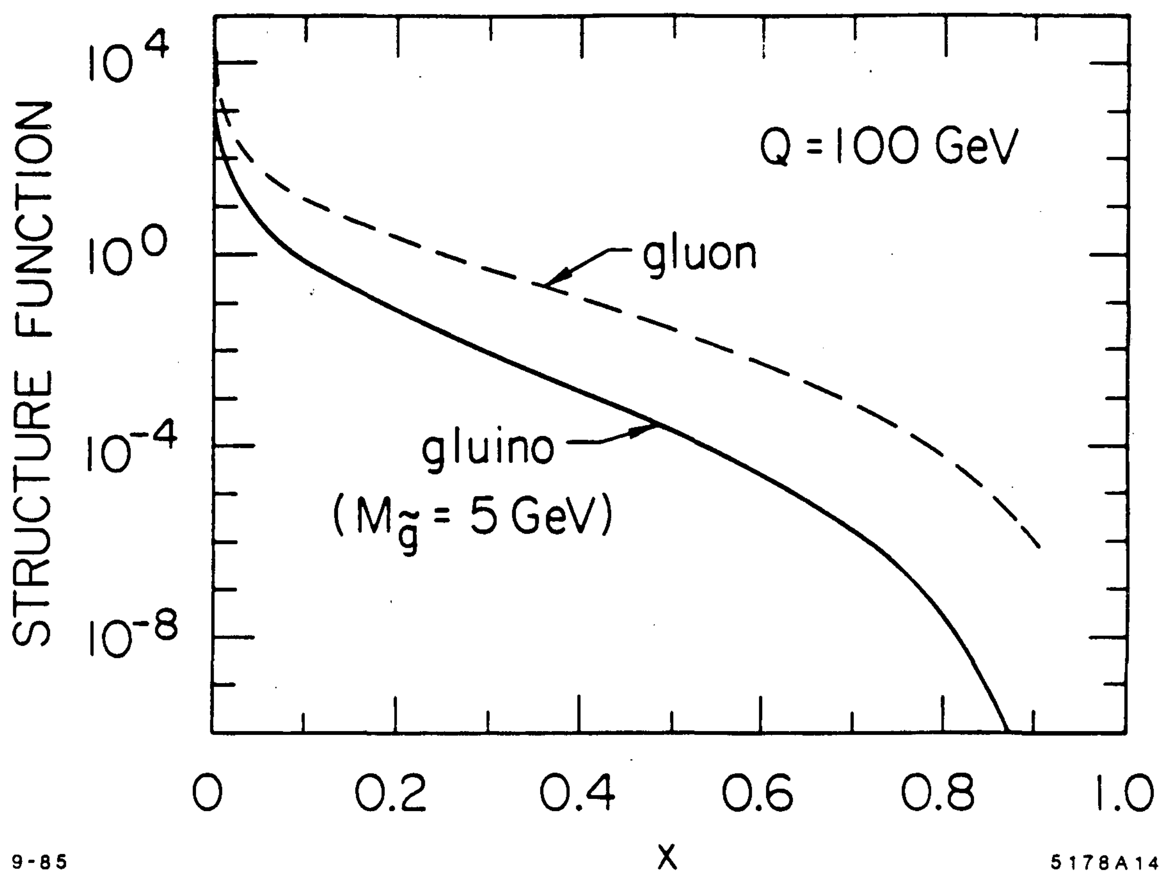
There are higher-order processes such as $gg \to \tilde{g}\tilde{g}g$ which in some cases turn out to be quite important. This is discussed below and I refer the reader to Refs. 81 and 75 for more information.

Another type of process which could be very important[76] if the mass of the gluino is small (3–15 GeV) originates with a perturbatively generated gluino component in the proton.[82] This gluino component allows processes such as

$$\tilde{g} + q \to \tilde{q} \tag{10.21}$$

where the incoming $\tilde{g}$ is treated as an initial-state parton. This process can be very efficient for the production of events with large E_T^{miss}. Because the gluino is a color octet (compared with the color triplet b-quark), the gluino component[83] in the proton is a factor of six times larger than that for an equal-mass b-quark. The resulting gluino distribution function can be almost 1% of that for gluons (if $M_{\tilde{g}} \approx 5$ GeV), see Fig. 30.

An important issue regarding the gluino distribution function has been raised recently by Barger et al.[74] The use of the Altarelli-Parisi equations to find the gluino distribution function may not justified at these Q^2. Those equations sum, to leading log, the emission of multiple gluons thereby incorporating processes beyond the simple $2 \to 2$ processes. However, non-leading-log effects may be important for these Q^2. There is reason to believe that the presently calculated gluino distribution function will give results which are too large; until better techniques are available, the $2 \to 2$ processes



Q = 100 GeV
gluon
gluino
($M_{\tilde{g}}$ = 5 GeV)
STRUCTURE FUNCTION
10^4
10^0
10^-4
10^-8
0
0.2
0.4
0.6
0.8
1.0
X
9-85
5178A14

FIG 30

$(g + g \rightarrow \tilde{g} + \tilde{q})$ probably give a reasonable estimate (or slight underestimate) of the signal. As it turns out, other processes are probably sufficient to rule out light gluinos anyway.

Why would the supersymmetric production and decay processes described above (Eqs. 10.16 –10.21 and 10.1–10.4) lead to monojet-type events? We have previously shown that one expects missing energy, but it is not immediately obvious why events should contain only one jet when a typical process has 2–4 outgoing jets.

Let us consider what is in fact one of the most difficult sources of monojets: the production (and decay) of light gluinos. A very small fraction of these events (1 in 10^4) have sufficient missing energy to pass the experimental requirement $E_T^{miss} > 15$ GeV (plus another cut based on resolution). However, given this requirement of large missing transverse energy, certain implications follow.

The decay of a very light gluino at rest cannot produce significant missing energy. Therefore, in the events which lead to large E_T^{miss}, the gluinos must emerge with significant momenta. Ordinarily the two gluinos will be roughly back-to-back. The two photinos (one from the decay of each gluino) also tend to be back-to-back, and their momenta cancel. Therefore, to obtain large missing transverse energy, one photino (say $\tilde{\gamma}_1$), must carry a large fraction of its gluino's ($\tilde{g}_1$) momentum, and the other photino ($\tilde{\gamma}_2$) must carry very little of its gluino's ($\tilde{g}_2$) momentum. This in turn implies that the jets from the decay of of $\tilde{g}_1$ are weak (and fail to pass the experimental requirement that $E_T^{jet} \geq 12$ GeV). Meanwhile, the jets from the decay of $\tilde{g}_2$ coalesce and are energetic. This results in the appearance of a single jet with large missing transverse energy which we call a "monojet," see Fig. 31.

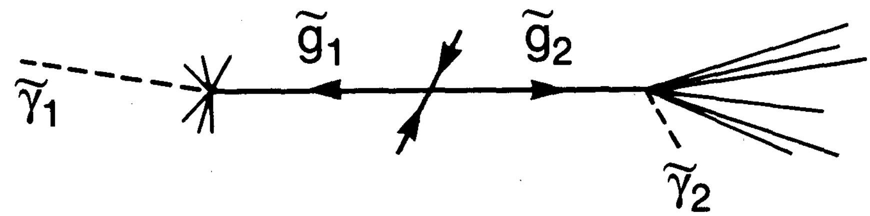


FIG 31

It was Herczog and Kunszt[81] who observed that in the higher-order process gg → $\tilde{g}\tilde{g}$g with a hard outgoing gluon, that this gluon may be back-to-back with the two gluinos. If the two gluinos are going in the same direction, then their resulting photinos' momenta can add rather than cancel. This is why the higher-order processes can, for light gluinos, be more effective at passing experimental cuts.

The actual theoretical analysis requires a complex Monte Carlo calculation. Among the problem theorists have found themselves forced to deal with are:[83]

a) Technical details of experimental cuts and triggers (which can be very important when one is dealing with events which are on the tails of distributions).

b) Resolution and efficiency effects (also very important on tails of distributions).

c) Fragmentation and gluon Bremsstrahlung (relevant to light gluinos).

d) The narrow distinction between monojets and dijets.

An example from problem (a) is the UA1 cut $E_T^{miss} > 4\ \sigma$ where σ is the calorimeter resolution and is approximately $\sigma = 0.7\ \sqrt{E_T\,(\text{total})}$. This cut is designed to eliminate "fake" missing-energy events which are due to non-uniform calorimetry and other mismeasurements. The problem is that the total scalar transverse energy, E_T (total), is not all the same as the total transverse energy in the jets. Each event typically has a substantial amount of accompanying transverse energy not in the jets. Although this energy effects the missing-energy resolution, theorists cannot calculate E_T (total) and must make use of other data for guidance (for example one can study the E_T (total) in ordinary QCD events with hard jets[84,85]). Extrapolation of CERN Sp$\bar{p}$S results to the Tevatron and SSC will be difficult as long as we lack knowledge on how to extrapolate E_T (total) to these higher energies.

Let us elaborate on problem (c). Although in processes (10.16–10.21) gluinos and scalar quarks are shown as final-state particles, we know that they cannot emerge as free particles since they carry color (they must hadronize). This is entirely analogous to c and b quark decays where we know that the resulting hadron (containing *c* or *b*) has less momentum than the *c* or *b* quark started with. Similarly the hadron containing the gluino will not have the full momentum of the gluino. This implies that the photino in gluino decay will have less energy, that there will be less missing energy and that fewer events will pass the E_T^{miss} cuts. It is the lighter particles that have the greatest fragmentation effects.

The primary sources of photino energy loss are:

a) gluon bremsstrahlung:

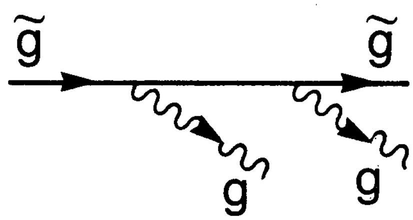


FIG 32

The gluino has large gluon bremsstrahlung because of the large gluon-gluino coupling (color-octet couplings). This effect has been parameterized by De Rújula and Petronzio[77] as:

$$D(z_1) = A\,(1 - z_1)^{A-1} \tag{10.22}$$

where

$$A = \frac{36}{25}\,\log[\log(\hat{s}/\Lambda^2)/\log(4M_{\tilde{g}}^2/\Lambda^2)] \tag{10.23}$$

b) fragmentation into hadrons:

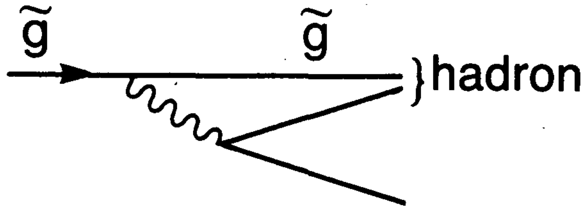


FIG 33

This has been parameterized by Peterson et al.[86] as:

$$D\,(z_2) = \frac{C(1-z_2)^2}{\left[(1-z_2)^2 + \left(\frac{0.6}{M_{\tilde{g}}}\right)^2 z_2\right]^2} \tag{10.24}$$

These two functions are convoluted to obtain the final fragmentation function.

A natural question to ask at this point is: Where does this "lost" energy go? The answer is that it can either add to the hadronic energy of the jet, or it can add to the total scalar energy E_t (total). The former seems more likely for light, fast-moving gluinos.

The question of the distinction between monojets and dijets (problem (d)) leads me to advocate combining the data for monojets, dijets and multijets (thereby considering all "missing-energy" events). These events appear to have large quantities of excess energy. Some energy may have left the jet cones, and some of the excess energy may overlap into the jet cones. UA1 observes[79,80] that most of these events have secondary jets of 6–12 GeV.

Nonperturbative QCD effects can change monojets to dijets and vice versa.

Clearly it is preferable to use all the data and increase the statistical significance while decreasing the sensitivity to theoretical and experimental limitations. While dijets are most subject to backgrounds, most of these backgrounds can be eliminated by making cuts against "back-to-back" dijets.

Many analyses[65–78] of the implications of the UA1 data for supersymmetry have been done by theorists. These generally apply to the 1983 UA1 run[79] of luminosity 115 nb^{-1}. The 1984 UA1 run[80] (with luminosity of 270 nb^{-1}) had modified cuts, triggers and experimental conditions. So far, the only analysis of the 1984 run is that of Haber, Kane and myself[75] (BHK). We performed a comprehensive analysis of *every* supersymmetric process (Eqs. 10.16–10.21 convoluted with Eqs. 10.1–10.4). Since the values of $M_{\tilde{q}}$ and $M_{\tilde{g}}$ are not predicted, our calculations were done for every combination of $M_{\tilde{g}}$ and $M_{\tilde{q}}$. In each case BHK calculated the monojet, dijet and multijet rates.

Figures 34–37 (from BHK) show as contour plots the event rates for the *sums* of all supersymmetric processes subject to the 1984 UA1 experimental conditions. Figs. 34 and 36 show the monojet rate while the other two plots show the sum of monojet, dijet plus multijet. One can make a more severe cut on E_T^{miss} in order to eliminate more backgrounds. The result of the cut $E_T^{miss} > 40$ GeV is shown in Figs. 36 and 37.

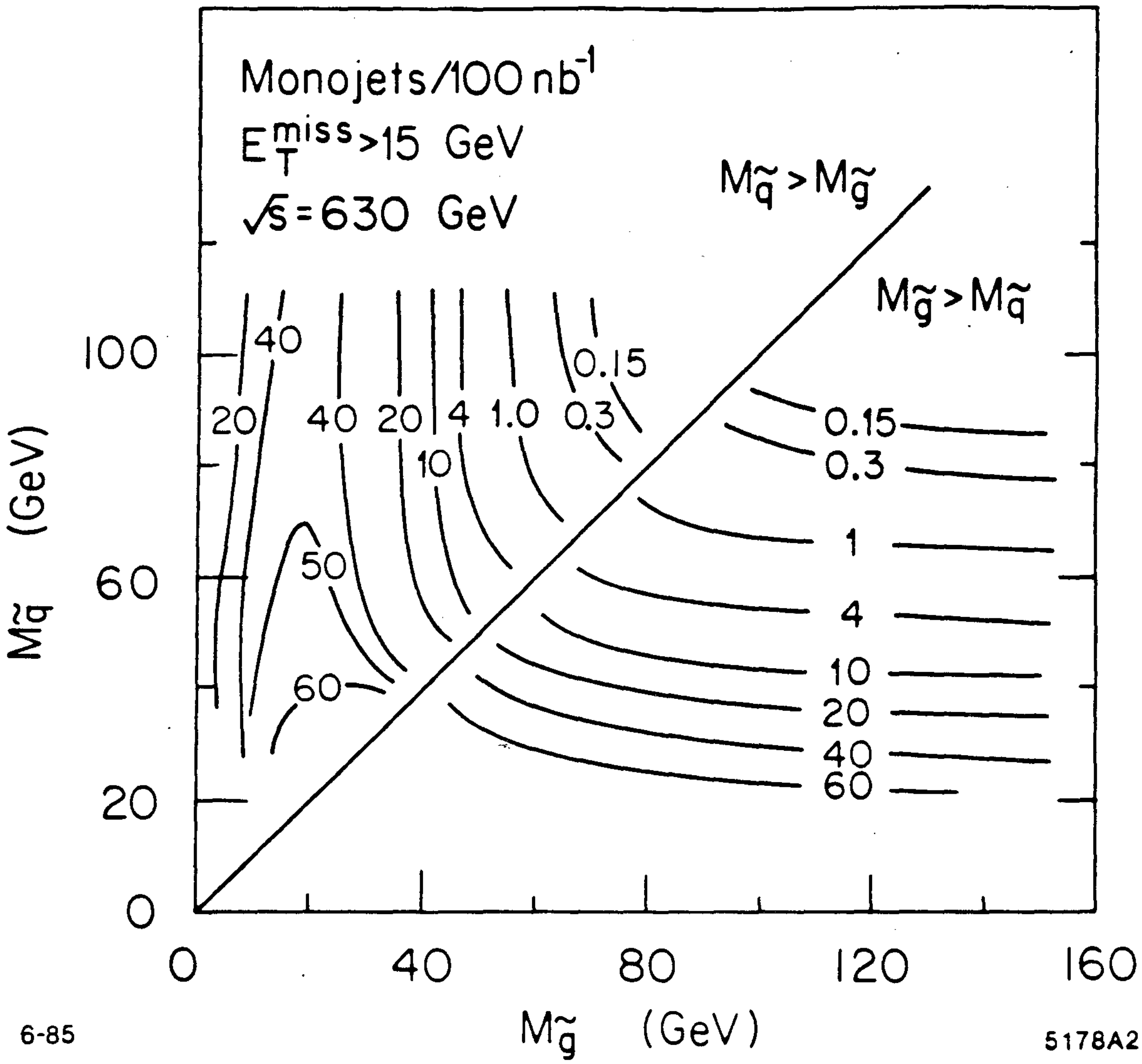
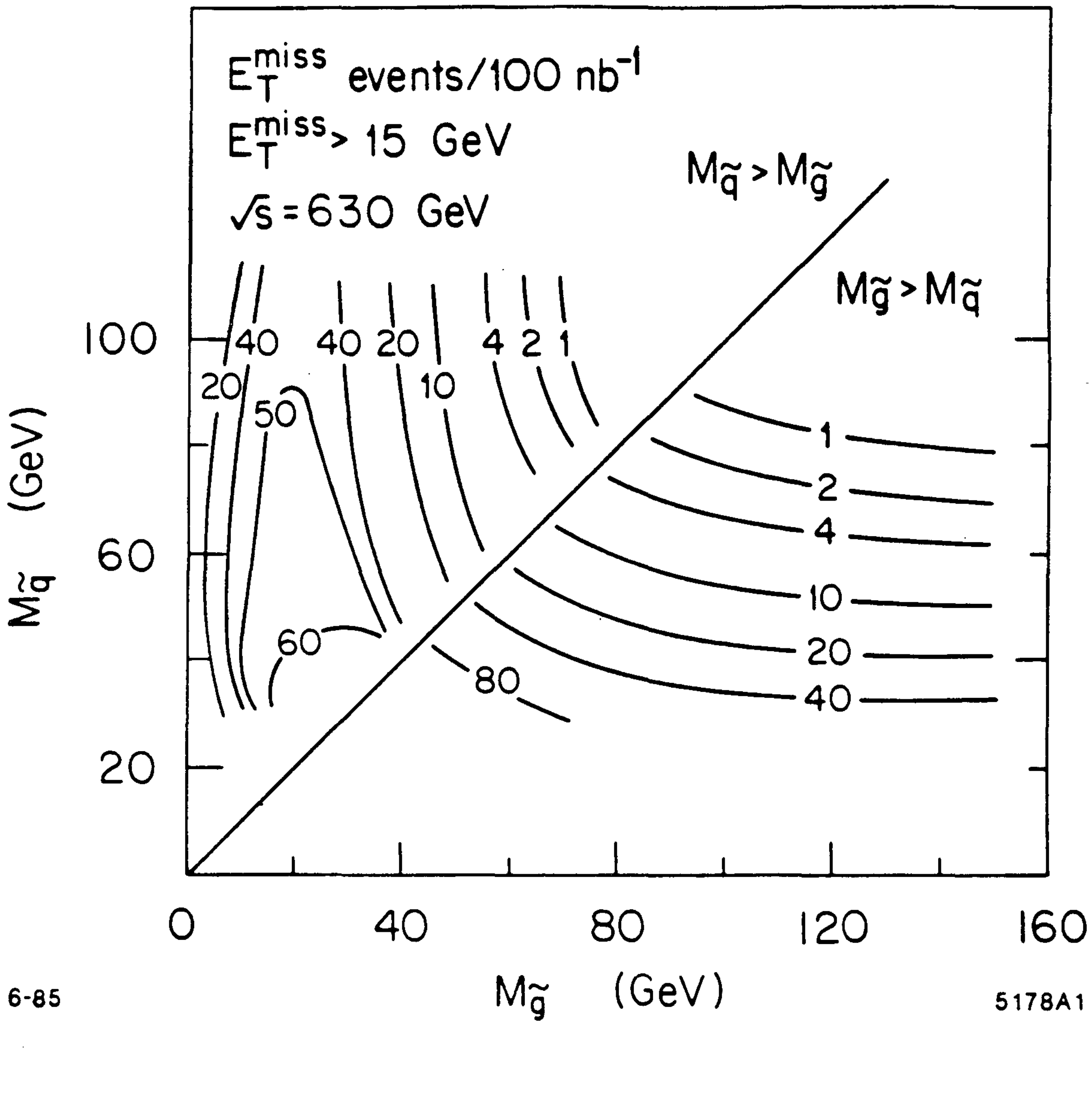

Monojets/100 nb^{-1}
E_T^{miss}>15 GeV
$\sqrt{s}$=630 GeV
$M_{\tilde{q}} > M_{\tilde{g}}$
$M_{\tilde{g}} > M_{\tilde{q}}$
$M_{\tilde{q}}$ (GeV)
$M_{\tilde{g}}$ (GeV)
0
20
60
100
40
80
120
160
40
20
40
20
4
1.0
0.3
0.15
10
50
60
0.15
0.3
1
4
10
20
40
60
6-85
5178A2

FIG 34



E_T^{miss} events/100 nb^{-1}
E_T^{miss} > 15 GeV
$\sqrt{s}$ = 630 GeV
$M_{\tilde{q}} > M_{\tilde{g}}$
$M_{\tilde{g}} > M_{\tilde{q}}$
20
40
50
60
40
20
10
4
2
1
80
1
2
4
10
20
40
100
60
20
$M_{\tilde{q}}$ (GeV)
0
40
80
120
160
$M_{\tilde{g}}$ (GeV)
6-85
5178A1

FIG 35

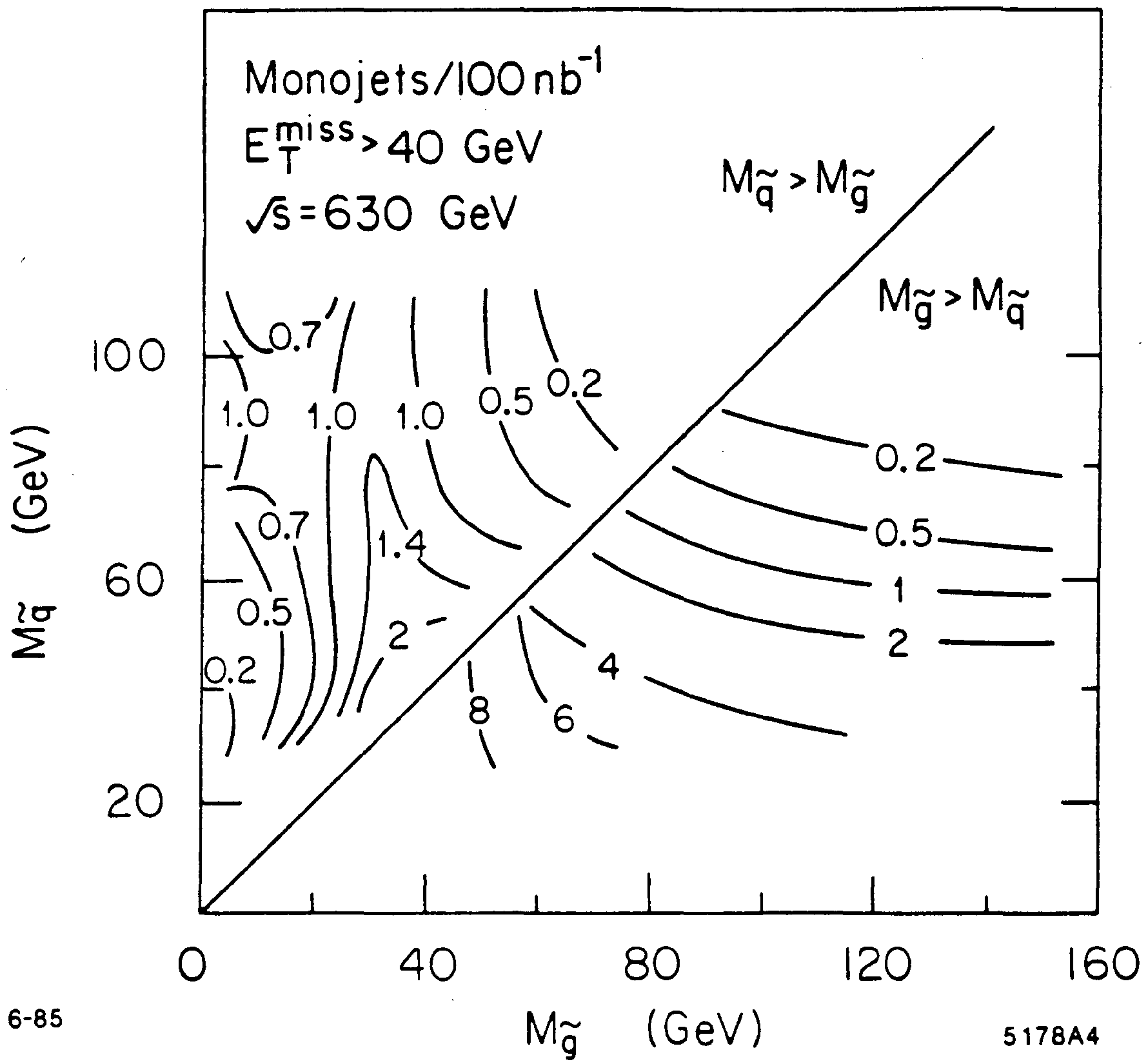

Monojets/100 nb⁻¹
E_T^miss > 40 GeV
√s = 630 GeV
M_q̃ > M_g̃
M_g̃ > M_q̃
0.7
1.0
1.0
1.0
0.5
0.2
0.2
0.5
1
2
0.7
1.4
0.5
2
0.2
8
6
4
100
60
20
M_q̃ (GeV)
0
40
80
120
160
M_g̃ (GeV)
6-85
5178A4

FIG 36

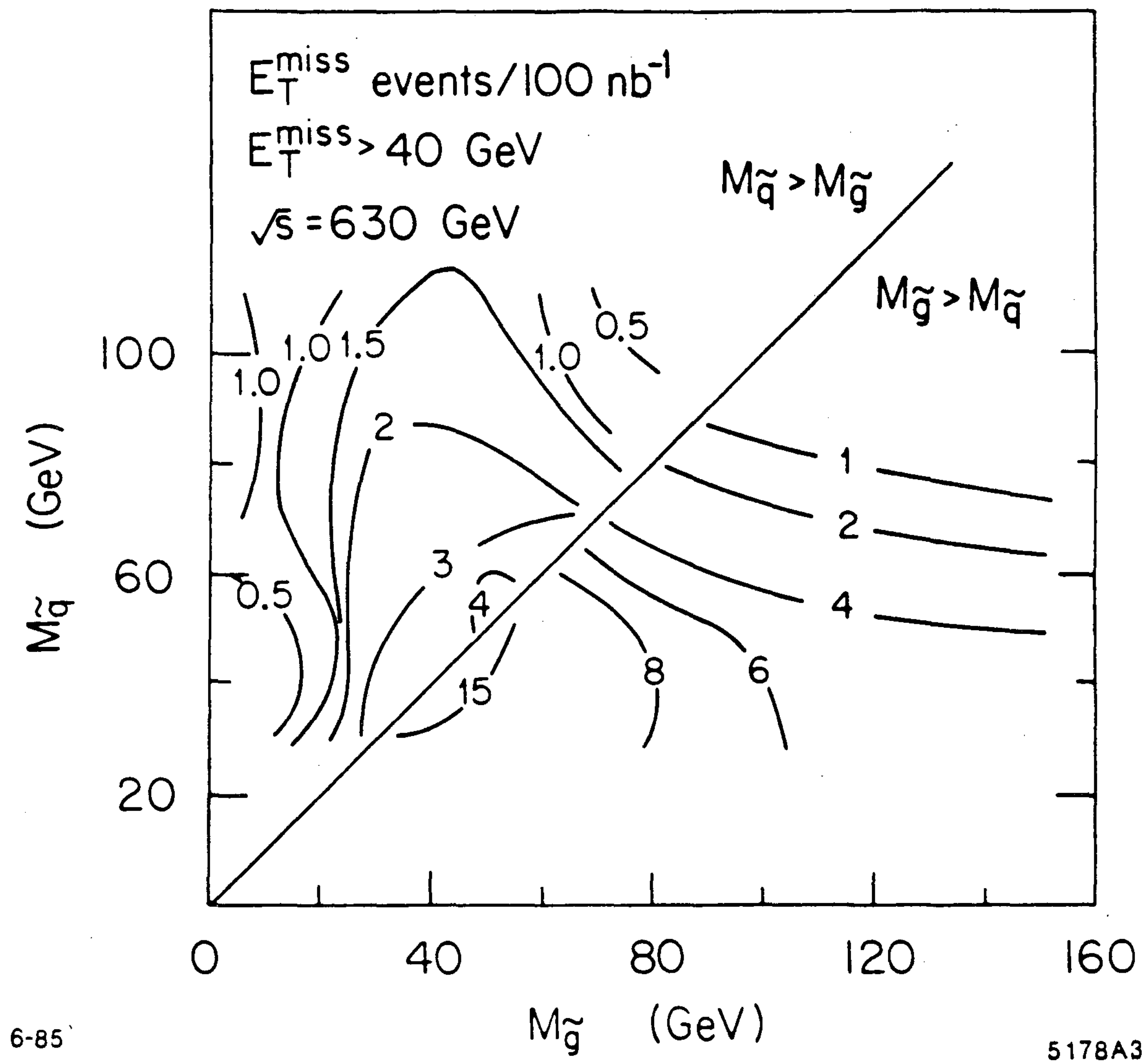

E_T^{miss} events/100 nb^{-1}
E_T^{miss} > 40 GeV
$\sqrt{s}$ = 630 GeV
$M_{\tilde{q}} > M_{\tilde{g}}$
$M_{\tilde{g}} > M_{\tilde{q}}$
1.0
1.0
1.5
1.0
0.5
2
1
2
4
3
4
0.5
15
8
6
100
60
20
0
40
80
120
160
$M_{\tilde{q}}$ (GeV)
$M_{\tilde{g}}$ (GeV)
6-85
5178A3

FIG 37

Next we must ask what are the data. The UA1 Collaboration reports from their 1984 run[80] with a luminosity of about 270 nb^{-1} a total of 23 mono-jets events. These are classified as:

9 identified as $W \to \tau\nu$ (with $\tau \to \nu$ + jet)

6–8 estimated from various backgrounds

6–8 possible excess

They also found 2 dijet events meeting the missing-energy and back-to-back criteria which they classify as:

2 estimated from various backgrounds.

0 possible excess

Since the question of whether there actually is any excess is subject to higher statistics and further experimental and theoretical analysis, BHK chose[75] to use the data to set 90% confidence level limits. These limits are (for $E_T^{miss} > 15$ GeV):

$$\text{Monojets: } 4.8 \text{ events}/100 \text{ nb}^{-1} \tag{10.25}$$

$$\text{Dijets: } 2 \text{ events}/100 \text{ nb}^{-1} \tag{10.26}$$

These are fairly conservative limits.

These limits can (using Figs. 34 and 35) be translated into limits on the masses of gluinos and scalar quarks. The combined monojet, dijet plus multijet rate gives significantly better limits than does the monojet rate. This occurs because as $M_{\tilde{q}}$ and $M_{\tilde{g}}$ increase, the events from supersymmetric processes begin to be dominated by dijets rather than monojets. In fact the best limits can be obtained at the largest masses simply by examining the dijet rate:[75]

$$M_{\tilde{q}} > \begin{cases} 65 & M_{\tilde{g}} \approx 150 \text{ GeV} \\ 75 & M_{\tilde{g}} \approx 80 \text{ GeV} \end{cases} \tag{10.27}$$

$$M_{\tilde{g}} > \begin{cases} 60 & M_{\tilde{q}} \approx 100 \text{ GeV} \\ 70 & M_{\tilde{q}} \approx 80 \text{ GeV} \end{cases} \tag{10.28}$$

A number of people have suggested[76,75,77,78,68] based on the 1983 UA1 data,[79] that there might be a "window" at $M_{\tilde{g}} =$ 3–5 GeV (with $M_{\tilde{q}} \approx 100$ GeV) where fragmentation effects are so great that we could not rule out these light gluinos. Recall that fragmentation effects can greatly reduce the amount of missing energy. However, because of a new trigger in the 1984 UA1 run[80] (which was extremely sensitive to some supersymmetric processes) and because of inclusion of the higher-order process $gg \to \tilde{g}\tilde{g}g$, BHK[75] now pretty well rule out this region. Certainly the "window" is reduced to a "peephole" at most.

The next issue is whether supersymmetry could explain any excess of monojets above calculated backgrounds, if such excess actually exists. The answer is: No! a) As just mentioned, very light gluinos can be ruled out. b) In the scalar-quark and/or gluino mass region around or below 40 GeV, BHK predict far too many monojets compared with what is observed.[80] c) In the mass region around 60–70 GeV it is possible to get correct monojet rates, but the dijet to monojet rate[75] is 2–1 to 3–1, unlike the data.[80] Also the predicted distributions are harder than those observed. d) Of course if $M_{\tilde{q}}$, $M_{\tilde{g}} \gtrsim$ 80 GeV no monojets are predicted. Therefore, the production of gluinos and scalar quarks is in no way associated with the observed monojets (in the standard scenario described here).

The two observed dijet events have a surprising amount of E_T^{miss}. If we assume that all monojets are due to backgrounds, it remains possible[75] that these two events originate from production of 70–90 GeV scalar quarks. But a large increase in statistics is required before any further speculation is warranted.

We should not finish this discussion without mentioning that there are other possible supersymmetric origins[87–94] for missing-energy events such as the production (and decay) of $\widetilde{W}$ and/or $\widetilde{Z}$. Furthermore it is important to emphasize that there are "standard model" backgrounds[95–98] for monojets which may account for all observed events. These include:

a) QCD dijet events can appear as monojets due to detector limitations.[79,80]

b) $p\bar{p} \to W^{\pm}$ + anything, with $W^{\pm} \to \tau\nu$ and $\tau \to \nu$ + jet. Because $W \to e\nu$ is measured one can accurately calculate this background. The experimentalists[80] can also identify some monojets as from this source due to their characteristics.

c) $p\bar{p} \rightarrow Z^o + g(\text{or } q) + \text{anything}$, with $Z^o \rightarrow \nu\bar{\nu}$.

QCD calculations have been done, but when higher statistics are available, $Z^o \rightarrow e^+e^-$ will provide an accurate check.

d) $p\bar{p} \rightarrow W^{\pm} + g(\text{or } q) + \text{anything}$, with $W^{\pm} \rightarrow e\nu$ where the electron is missed because it is slow or inside the jet.

e) The production of $c\bar{c}$, $b\bar{b}$ and $t\bar{t}$ pairs followed by semileptonic decays with the neutrino taking most energy and the lepton either taking very little energy or being missed.

f) Other backgrounds.

Even if it is established that there are missing-transverse-energy events which cannot be described by any conventional physics, we will not have evidence for supersymmetry until we can show that the events are not caused by *explainatons:*

FIG 38

These insidious particles are *designed* to explain whatever data are reported; and if the data change, then they too change.

10.F. Other Techniques to Find Gluinos and Scalar Quarks

Another technique for finding gluinos is to look for $\tilde{g}\tilde{g}$ bound states.[99–101] Consider a color singlet state $\tilde{G}$. Since the parity of the Majorana $\tilde{g}$ is imaginary, we find for $\tilde{G}$:

$$P = (-1)^{\ell+1} \qquad C = (-1)^{\ell+s} \tag{10.29}$$

but $C(\tilde{g}) = 1$ implies

$$C(\tilde{G}) = +1 \tag{10.30}$$

which implies ℓ+s is even. Therefore, no ψ-like state ($^3S_1$) exists, but $\eta_{\tilde{g}}$ does exist.

Goldman and Haber[99] have calculated the width for this state:

$$\Gamma\,(\eta_{\tilde{g}} \to gg) = \frac{36\ \alpha_s^{\,2}}{M^2_{\eta_{\tilde{g}}}}\ \frac{1}{2}\,|\tilde{R}\,(0)|^2 \tag{10.31}$$

For large $M_{\tilde{g}}$ (the Coulombic approximation), $\frac{1}{2}M_{\eta_{\tilde{g}}} = M_{\tilde{g}} \gtrsim 50$ GeV, they find

$$\Gamma\,(\eta_{\tilde{g}} \to gg) \approx \frac{(3\ \alpha_s)^5}{8}\,M_{\eta_{\tilde{g}}} \approx 4 \times 10^{-4}\,M_{\eta_{\tilde{g}}} \gtrsim 40\ \text{MeV} \tag{10.32}$$

This is rather wide compared to quarkonium, due to extra color factors. Goldman and Haber[99] advocate a search for resonant 2-gluon scattering at large p_T. This may be possible at hadron colliders.

Another technique for gluino searches is in the decays of the $^3S_1$ state of the $t\bar{t}$. Some authors[102–104] have considered

$t\bar{t} \rightarrow \tilde{g}\tilde{g}$ (10.33)

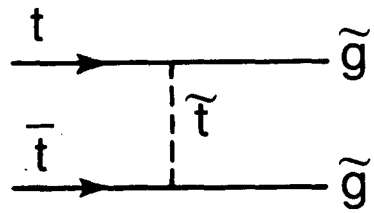


FIG 39

Since $^3S_1$ has $C = -1$ and $\tilde{g}$ has $C = -1$, the C-violation occurs via the mass splitting

$$\Delta M^2_{\tilde{q}} \equiv M^2_{\tilde{q}_L} - M^2_{\tilde{q}_R} \tag{10.34}$$

These authors[102] have found:

$$\frac{\Gamma\,(^3S_1 \rightarrow \tilde{g}\tilde{g})}{\Gamma\,(^3S_1 \rightarrow \mu\mu)} = \frac{4}{9\,e_q^{\,2}} \left(\frac{\alpha_s}{\alpha}\right)^2 \left(1 - \frac{M^2_{\tilde{g}}}{m_q^{\,2}}\right)^{3/2} F^2 \tag{10.35}$$

where

$$F = \frac{\Delta M^2_{\tilde{q}}\, m_q^{\,2} \cos 2\,\theta}{(m_q^{\,2} - M^2_{\tilde{g}} + M^2_{\tilde{q}_1})\,(m_q^{\,2} - M^2_{\tilde{g}} + M^2_{\tilde{q}_2})} \tag{10.36}$$

and θ = mixing angle for $(\tilde{q}_L\,,\,\tilde{q}_R) \longleftrightarrow (\tilde{q}_1\,,\,\tilde{q}_2)$. At best this ratio (10.35) is 10%. One would have to look for a missing energy signal. It may now be that $M_{\tilde{g}} > M_t$ so this would be ruled out.

Keung[105] and Haber[4] have suggested the decays $^3S_1\,(t\bar{t}) \rightarrow \gamma\tilde{g}\tilde{g}$, $g\tilde{g}\tilde{g}$ and $g\tilde{g}\tilde{\gamma}$:

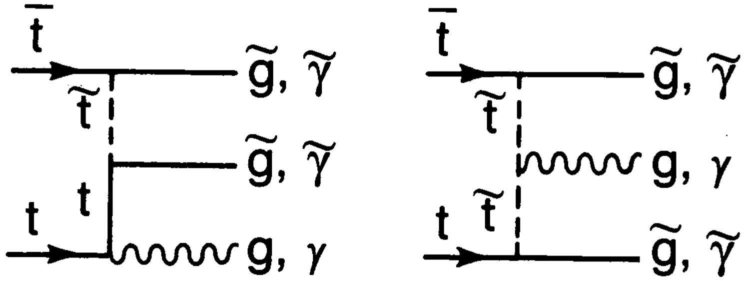


FIG 40

Taking $M_{\tilde{t}_L} = M_{\tilde{t}_R} >> \sqrt{s}$ and $M_{\tilde{g}}$ small gives

$$\Gamma\,(^3S_1 \rightarrow g\tilde{g}\tilde{g}) = \frac{5}{24}\,\frac{\alpha_s}{\alpha\, e_t^2}\,\Gamma(^3S_1 \rightarrow g\tilde{g}\tilde{\gamma})$$

$$= \frac{1}{6(\pi^2-9)}\left(\frac{2\,m_t}{M_{\tilde{t}}}\right)^4 \Gamma\left(^3S_1 \rightarrow ggg\right) \qquad (10.37)$$

This can be relatively large if $M_{\tilde{t}} \sim m_t$. The process $^3S_1 \rightarrow g\tilde{g}\tilde{\gamma}$ leads to a superior signal and may give less phase space trouble.

Others have considered processes[102–103,106–107] such as

$$^3S_1\,(t\bar{t}) \rightarrow gg\tilde{g}\tilde{g}$$

$$\tilde{\gamma}\tilde{\gamma}\,, \quad \text{or} \quad \gamma\tilde{\gamma}\tilde{\gamma}$$

$$\tilde{q}\tilde{\bar{q}}$$

This latter process may be ruled out by $M_t < M_{\tilde{q}}$.

A process[99–101] which may be very useful in ruling out extremely light gluinos ($1 < M_{\eta_{\tilde{g}}} < 8$ GeV) is:

$$^3S_1 \rightarrow \gamma + \eta_{\tilde{g}} \tag{10.38}$$

The branching ratio from a $t\bar{t}$ state would be very small, but from Υ it is

$$B\,(\Upsilon \rightarrow \gamma + \eta_{\tilde{g}}) \approx (1\text{–}3) \times 10^{-4} \tag{10.39}$$

It may therefore be possible eventually to rule out $M_{\tilde{g}} < 4$ GeV.

11. Fermionic Partners of Charged Gauge and Higgs Bosons

Mass eigenstates will in general be mixtures of $\tilde{W}^{\pm}$, $\tilde{H}_1^{+}$, and $\tilde{H}_2^{-}$. Some authors[4] refer to the mass eigenstates as "charginos" or $\tilde{\chi}^{\pm}$.

Some possible decays of $\tilde{W}$ (or W-ino) are shown in Fig. 41 (this figure can be generalized by replacing $\tilde{W}^{+}$ with $\tilde{\chi}^{+}$ and $\tilde{\gamma}$ with $\tilde{\chi}^{o}$; note that the $\tilde{\nu}$ in Fig. 41c is assumed to be an on-shell particle). The present mass limits for

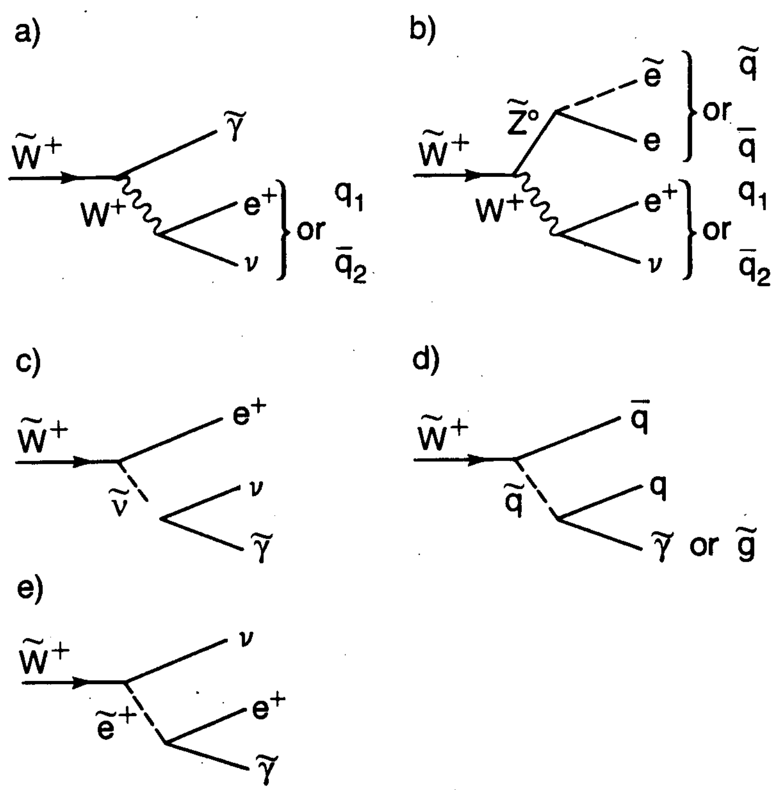


FIG 41

$\tilde{q}$, $\tilde{g}$ and $\tilde{e}$ imply that many of these modes could be comparable, so it is very difficult to make predictions. We remain completely ignorant also about chargino mixings. As a result any mass limits for W-inos are very model-dependent. It is essential to state carefully the assumptions which go into any quoted mass limits.

Charginos can be produced in e^+e^- machines[90,92,108–110] via the diagrams:

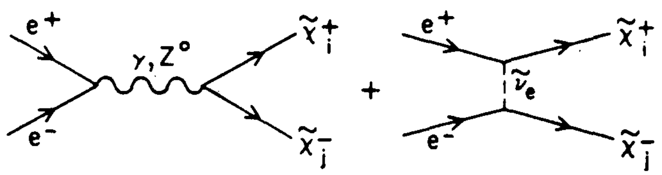


FIG 42

The γ, Z-exchange diagram has only a β (not β^3) threshold suppression (unlike the scalar leptons and quarks). The $\tilde{\nu}_e$-exchange diagram can be large if $M_{\tilde{\gamma}}$ is small.

The analysis of these processes is complicated by the different angular distributions expected from the two diagrams (whose relative size is model dependent). If the Z^o contribution is significant, then a large forward-backward asymmetry is possible.

At hadron colliders,[51,90–92] the following diagrams may be contribute (where $\tilde{\chi}^{\pm}$ may be $\tilde{W}^{\pm}$ or $\tilde{H}_i^{\pm}$, and $\tilde{\chi}^o$ may be $\tilde{\gamma}$, $\tilde{Z}^o$ or $\tilde{H}_i^o$):

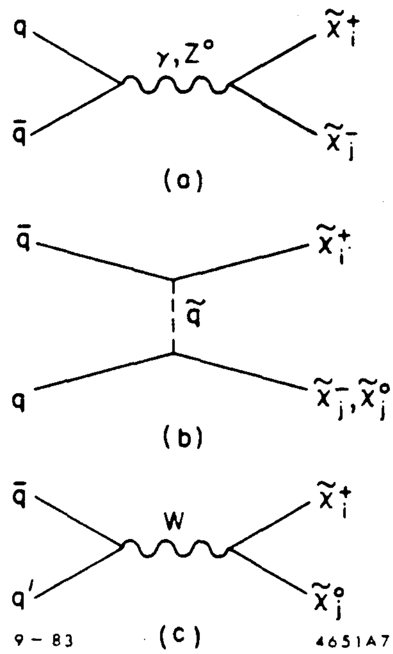


FIG 43

The $\widetilde{W}^+ \widetilde{W}^-$ processes are unlikely to have a good signature. However, $\widetilde{W}^+ \tilde{\gamma}$ may lead to identifiable characteristics such as

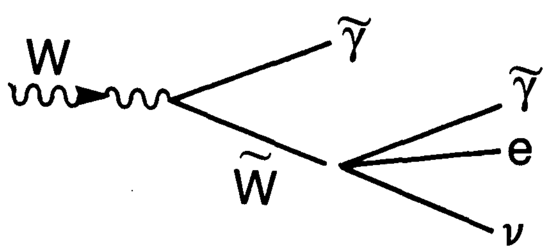


FIG 44

Another technique for producing charginos (and neutralinos) is by producing W and Z bosons which may (for appropriate $M_{\tilde{\chi}^{\pm}}$ and $M_{\tilde{\chi}^o}$) decay into $\tilde{\chi}^o$ and $\tilde{\chi}^{\pm}$. I will not discuss this technique here but refer the reader to the extensive literature (Refs. 87–89, 111–115).

Recently, the JADE,[116] Mark J[117] and CELLO[40] collaborations have reported new mass limits. I will quote those from JADE. I repeat that all mass limits are assumption dependent. If $M_{\tilde{\nu}} < 18$ GeV and if all charginos decay into lepton plus $\tilde{\nu}$, then (95% C.L.)

$$M_{\tilde{\chi}^{\pm}} (\ell\, \tilde{\nu}) \gtrsim 22 \text{ GeV} \tag{11.1}$$

When hadronic decays occur, then we find (95% C.L.) from the expected increase in total hadronic cross-section:

$$M_{\tilde{\chi}^{\pm}} (\text{R (hadronic)}) > 16.5 \text{ GeV} \tag{11.2}$$

For other decay modes, the limits are more dependent on $M_{\tilde{\gamma}}$ and on the branching ratios to $\ell\, \tilde{\nu}\tilde{\gamma}$ and $q_1\bar{q}_2\tilde{\gamma}$. In general, though for large branching ratios for a given mode, they can exclude a region from 1–13 GeV (depending on $M_{\tilde{\gamma}}$) up through 23 GeV. I refer the reader to Ref. 116 for figures which show these results more clearly.

12. Fermionic Partners of Neutral Gauge and Higgs Bosons

The weak interaction eigenstates

$$\tilde{\gamma}, \quad \tilde{Z}^o, \quad \tilde{H}_1{}^o, \quad \tilde{H}_2{}^o$$

mix to give mass eigenstates[109] which are sometimes referred to a "neutralinos" or $\tilde{\chi}^o$. As for the charginos, the unknown masses and mixings imply significant uncertainty in couplings and cause difficulties in calculations.

The neutralino decays[87–89,109] shown below are possible (where $\chi^{\pm}$ may be $W^{\pm}$ or $\widetilde{H}_1^{\pm}$, $\widetilde{\chi}^o$ may be $\widetilde{\gamma}$, $\widetilde{Z}^o$ or $\widetilde{H}_1$ and f refers to quarks or leptons):

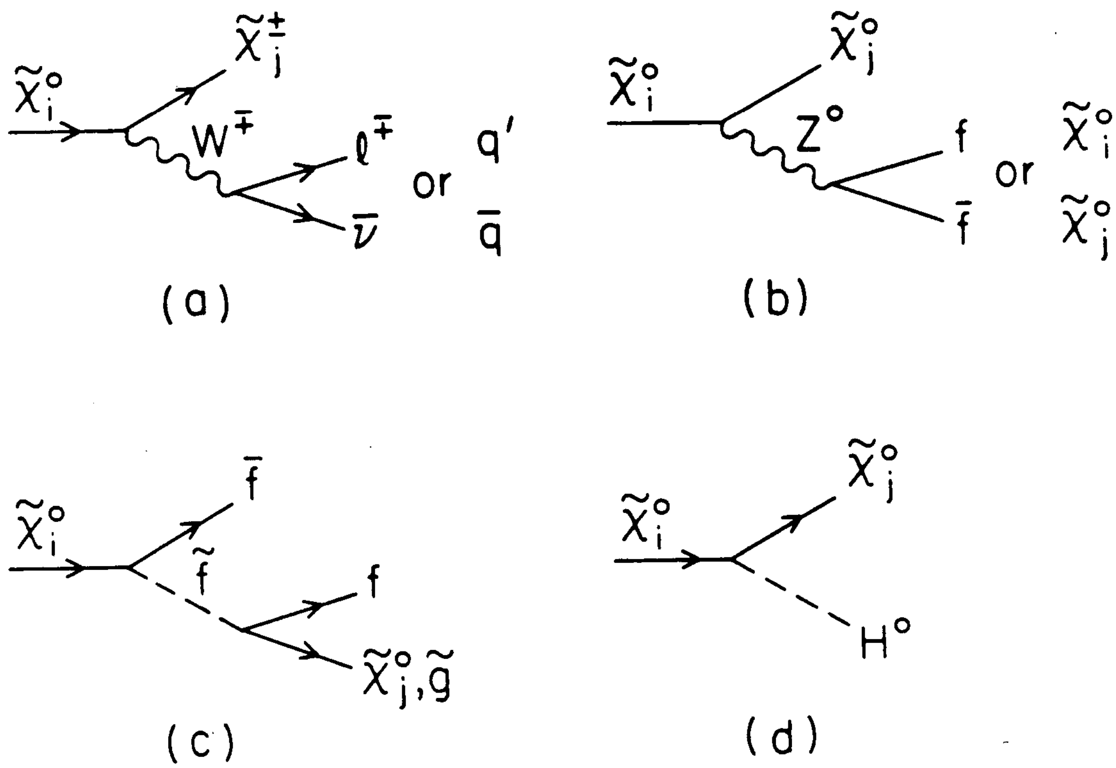


FIG 45

The $\widetilde{e}$ and $\widetilde{q}$ exchange diagrams (c) are comparable if $M_{\widetilde{e}} \approx M_{\widetilde{q}}$.

Neutralinos (such as $\widetilde{Z}^o$) can be produced in e^+e^- annihilation[88,90–91,108,118–119] via the diagrams

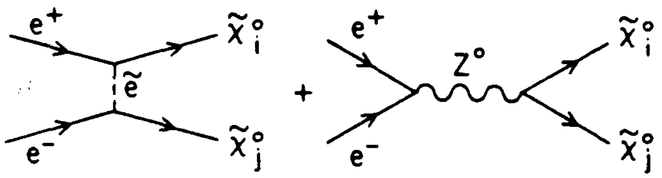


FIG 46

The process $e^+e^- \rightarrow \tilde{\gamma}\,\tilde{Z}^o$ can lead to dramatic one-sided events[91,108,118] as shown (where the circle represents the beam pipe):

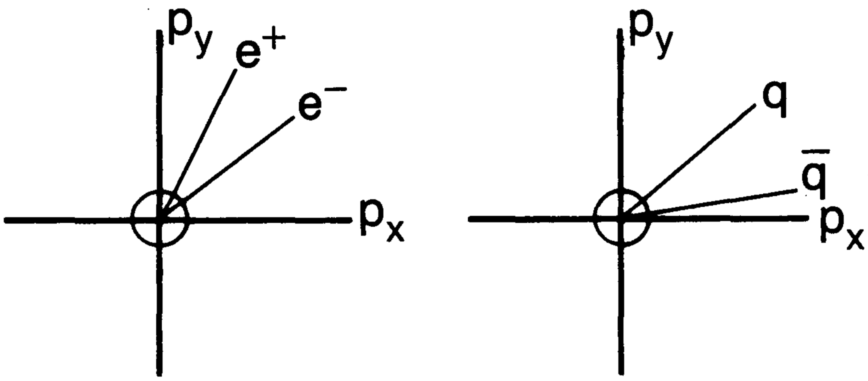


FIG 47

The first type events will typically have $\frac{2}{3}$ of the total energy missing. The backgrounds include $\tau^+\tau^- \rightarrow e^+e^-$; however these background events would be more back-to-back, and one would also see $e^+\mu^-$ events from this background. Another background is 2γ events, but here the missing momentum ordinarily points in the beam direction. The $\tilde{Z}^o\,\tilde{Z}^o$ events could give $e^+e^-\mu^+\mu^-$ events with $\frac{1}{3}$ missing energy (again the circle represents the beam pipe):

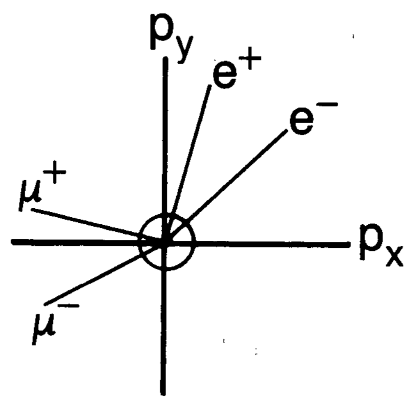


FIG 48

At hadron colliders[120–121] possible processes include:

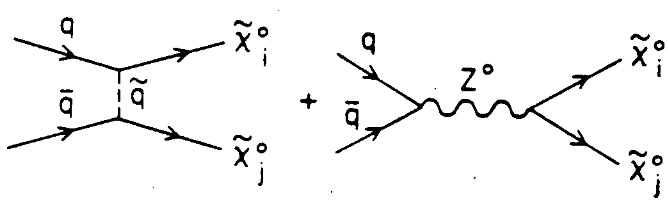


FIG 49

The Z^o exchange may be either on-shell or off-shell. The cross-section for $\tilde{\gamma}\,\tilde{Z}^o$ production is

$$\sigma\,(p\bar{p} \rightarrow \tilde{\gamma}\,\tilde{Z}^o + \text{anything}) \sim 0.1 - 1.\ \text{pb} \tag{12.1}$$

The signatures are similar to those described for e^+e^- annihilation.

Another mechanism for producing $\tilde{\chi}^o$ is via the production of W and Z bosons which (for appropriate $M_{\tilde{\chi}^o}$) in turn decay into neutralinos. This mechanism is discussed in Refs. 88–89,91,111–115,120–121.

Recent mass limits from the JADE[122] and Mark J[117] Collaborations have been reported. I repeat that all mass limits are assumption dependent. If one assumes $M_{\tilde{\gamma}} = 2\text{–}5$ GeV, $M_{\tilde{e}} \lesssim 50$ GeV, $B(\tilde{Z}^o \to ee\,\tilde{\gamma}) \approx 5\%$ and $B(\tilde{Z}^o \to \mu\mu\tilde{\gamma}) \approx 5\%$, then Mark J excludes $M_{\tilde{Z}}$ between about 6 and 33 GeV. The JADE Collaboration finds comparable limits for the decay modes $e^+e^-\,\tilde{\gamma}$, $\mu^+\mu^-\tilde{\gamma}$, $q\bar{q}\,\tilde{\gamma}$ and $q\bar{q}\,\tilde{g}$. The limits depend on $M_{\tilde{g}}$, $M_{\tilde{e}}$ and on the branching ratios to each mode, see Ref. 122. Limits can also be obtained for $\tilde{\chi}^o$ from monojet searches in e^+e^- annihilation; see Refs 123–125.

A technique for finding $\tilde{H}^o$ has been suggested by Campbell, Scott and Sundaresan[103] who assume $M_{\tilde{\gamma}} = O$ and no mixing. They calculated the width for a heavy quark state $^3S_1$ to decay radiatively to $\tilde{H}^o$:

$$\frac{\Gamma\,(^3S_1 \to \tilde{\gamma}\,\tilde{H}^o)}{\Gamma\,(^3S_1 \to ggg)} = \frac{81\,\pi^2\,\alpha\,e_Q{}^2\,G_F\,(m_V{}^2 - M_{\tilde{H}}^2)^2\,(m_V{}^2 + 2\,M_{\tilde{H}}^2)}{20\,\sqrt{2}\,\alpha_2{}^3\,(\pi^2 - 9)\,M_{\tilde{Q}}^4} \tag{12.2}$$

where $m_V \equiv m(^3S_1)$. This could yield a dramatic signal.

13. Flavor-Changing Neutral Currents

As discussed in Sec. 10, the phenomenological reasons for needing small splittings for the scalar quarks ($\tilde{u}$, $\tilde{d}$, $\tilde{s}$, $\tilde{c}$ and $\tilde{b}$) originate in the small magnitude of flavor-changing neutral currents.[58,59] The relevant diagrams include:

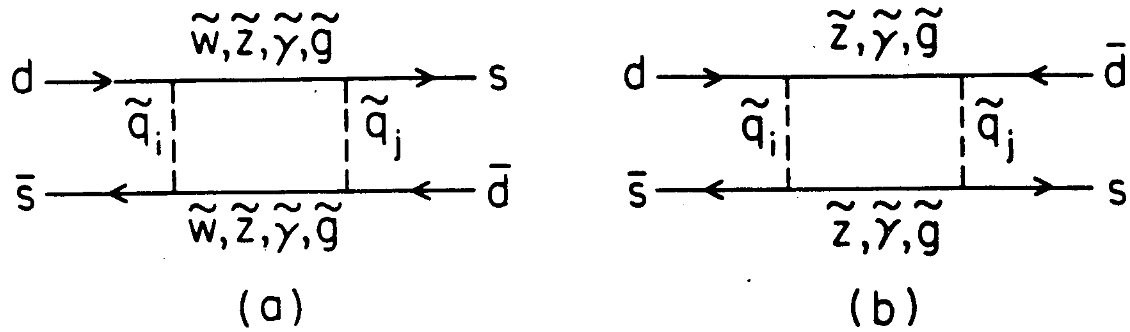


FIG 50

While the calculational techniques are somewhat different here, the experimental techniques are, of course, the same as for ordinary particle exchanges.

14. Conclusions

Experiments are beginning to set good limits on the masses of supersymmetric particles. However, none have yet reached m_W. If supersymmetry is to explain the origin of the electroweak scale, then it is natural to assume that this is the relevant scale for supersymmetric masses. Therefore, supersymmetric signals may be just around the corner. If we are lucky, the signals could be spectacular. If we are unlucky, much hard work will be necessary to confirm supersymmetry (at high-energy hadron colliders we will need to know how to scale up lower-energy calculations). As a result, I conclude by noting: Our luck is *model dependent.*

ACKNOWLEDGMENTS

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