DEUTSCHES ELEKTRONEN-SYNCHROTRON DESY

DESY 79/67
October 1979

SUPERSYMMETRIC PARTICLES AT LEP

ECFA/LEP Specialized Study Group 9 "Exotic Particles"

G. Barbiellini - *INFN, Frascati and CERN*

G. Bonneaud - *Strasbourg and CERN*

G. Coignet - *LAPP, Annecy-le-vieux*

J. Ellis - *CERN*

M. K. Gaillard - *LAPP, Annecy-le-vieux*

C. Matteuzzi - *CERN*

B. H. Wiik - *DESY*

NOTKESTRASSE 85 · 2 HAMBURG 52

To be sure that your preprints are promptly included in the
HIGH ENERGY PHYSICS INDEX ,
send them to the following address (if possible by air mail) :

DESY
Bibliothek
Notkestrasse 85
2 Hamburg 52
Germany

DESY 79/67
October 1979

SUPERSYMMETRIC PARTICLES AT LEP

ECFA/LEP Specialized Study Group 9 "Exotic Particles"

G.Barbiellini - INFN, Frascati and CERN
G.Bonneaud - Strasbourg and CERN
G.Coignet - LAPP, Annecy-le-vieux
J.Ellis - CERN
M.K.Gaillard - LAPP, Annecy-le-vieux
C.Matteuzzi - CERN
B.H.Wiik - DESY

1. Introduction

Supersymmetry is (perhaps) the most beautiful irrelevant theory[1] we have today.We are familiar, from other branches of physics as well as high energy physics, with symmetries that inter-relate particles with the same spin (nuclear SU(2), flavour SU(3), colour SU(3), electroweak SU(2) x U(1)), or with spins differing by integers (nuclear SU(4), non-relativistic SU(6)). Be they global or local gauge symmetries they all have the common feature that they group into multiplets bosons with bosons, and fermions with fermions. Supersymmetry goes beyond this level by introducing symmetry transformations between fermions and bosons, and therefore possessing multiplets containing both fermions and bosons. Furthermore this connection between fermions and bosons is established by gauging the spin, which seems like a natural thing to do after the great success obtained by gauging isospin. And what else is there left to do ?

Supersymmetric theories thus have the aesthetic attraction of introducing a higher degree of symmetry into our description of nature. They may also have mundane calculational advantages, in that many higher order divergences cancel between fermions and bosons, rendering theories involving large spins more convergent[2] and possibly even finite. As an example, extended supergravity theories which combine general relativity with the principle of supersymmetry seem to be less infinite than conventional gravity coupled to ordinary matter[2].Since a computable theory of gravity is the holy grail of theoretical physics, this suggests that supersymmetry should eventually play a role in the dynamics of elementary particles.

Unfortunately, there is no phenomenological sign of supersymmetry in the spectrum of particles and interactions observed to date, which is the basis for our overly harsh opening sentence. This by no means implies that supersymmetry will always be irrelevant. Our present theories of particle interactions contain many widely different mass scales - a strong interaction scale $\sim$ 1 GeV, fermion mass scales $O(10^{-3}$ to 10^{+2}?)GeV, a weak interaction scale $\sim$ 100 GeV, a possible grand unification scale $\sim 10^{15}$ GeV, and of course the Planck scale of gravitation $\sim 10^{19}$ GeV. The argument based on the greater convergence of supergravity suggests that supersymmetry should manifest itself at or before the Planck scale, but we have no theoretical insight into which of the above interaction scales - if any - may be supersymmetrized. All we know phenomenologically[3] is that "naive" supersymmetric partners of some of our known particles (leptons, quarks, gluons) do not exist with masses $\lesssim O(1)$ GeV.

The object of this report is to analyze in more detail the phenomenological signatures of the possible supersymmetric particles mentioned in table 1, with a view to their production and detection at LEP. Our objective will be to see whether the supersymmetrization of nature at a mass scale up to O(100) GeV can be confirmed or excluded by experiments with LEP. We

find that the scales at which supersymmetry is probed can be refined by one or two orders of magnitude.

In section 2 we review the qualitative features of the new spectroscopy in table 1 suggested by supersymmetric theories. There might be scalar coloured partners of quarks. There might be scalar uncoloured partners of spin 1/2 leptons which would look like Higgs particles[4], with the exception that their couplings to light fermions would be larger than those of conventional Higgs particles. There might also be supersymmetric heavy leptons which would be partners of the $W^{\pm}$ and Z^0 and perhaps have weak isospin $\neq$ 1/2, unlike conventional sequential heavy leptons. There might be spin 1/2 coloured partners of the gluons, called gluinos[3]. There might be new massless neutral spin 1/2 leptons analogous to the neutrinos - called collectively nuinos - perhaps associated with the photon for example (photinos)[3]. In addition to briefly describing the possible properties of each of these conjectured particles, we give approximate phenomenological limits on the possible masses of such objects, relying heavily on the pioneering phenomenological work of Farrar and Fayet[3,5,6].

The rest of this report consists of sections discussing in turn possible production rates and means of detection at LEP of each of the classes of supersymmetric particles outlined above. Where appropriate, we also make remarks about how other projects for future high energy physics machines may also be useful in the quest for supersymmetry. For example, a high energy ep machine[7] might be a good place to look for the indirect effects of gluinos, or for the direct production of scalar leptons.

We will see that LEP may find, or exclude, supersymmetric particles in wide ranges of masses which differ according to species. Section 8 contains a more detailed exposition of our conclusions about the prospects for supersymmetric phenomenology.

2. Particles suggested by supersymmetry

2.1 Introduction

The fundamental feature of supersymmetry[1] is its introduction of symmetry transformation α with 1/2 integral spin which may generate fermions from bosons, or vice versa. These transformations α may (extended supersymmetry[8]) or may not (simple supersymmetry[9]) have some internal (e.g. SU(3) colour or SU(2) electroweak) symmetry index as well: $\alpha \to \alpha_i$. In the extended supersymmetry case α_i the supersymmetric partners of conventional particles will not in general have the same internal symmetry (SU(3) or SU(2)) transformation properties, whereas the supersymmetric partners would have identical internal quantum numbers in simple supersymmetry theories. Applying to our familiar fundamental particles the operations of raising and lowering spins by 1/2, we find the particles listed in table 1 [10]. Not all possibilities have been listed. For example we have discarded

"vector quarks" and "vector leptons" because at present the only way of including vector particles *in* renormalizable field theory is to make them gauge bosons, rôles that are already fulfilled by the $W^{\pm}$, Z^0, γ and gluons. We mention only briefly particles with spin $\geq 3/2$ [11]. No renormalizable theory of such particles is known, and they are avoided in conventional supersymmetry schemes, though not in the more sophisticated supergravity schemes[12]. The non-renormalizability of the couplings of such particles is reflected in the fact that the production of charged 3/2 particles in e^+e^- by a direct channel γ or Z^0 would not have the simple scaling form $\sigma \sim 1/Q^2$, but would instead grow with Q^2. If any such particles existed with masses $<$ 100 GeV there should be no problem in detecting them[13] ! We will now review in turn some of the possible properties of the new supersymmetric particles mentioned in Table 1.

The contents of Table 1 are very similar to the particles discussed by Farrar and Fayet[3,5,6,11] in their important work on supersymmetry phenomenology. In their work they assumed a simple supersymmetry in which the internal quantum numbers of particles in the same horizontal line of table 1 were identical. They also had a new discrete quantum number R = 0 for conventional particles and R = $\pm$ 1 for their supersymmetric partners[10]. Much more general patterns of internal quantum numbers are possible, and one should not be too fixed on these simplest possibilities. However, the simple supersymmetry quantum number assignments often realize the most basic internal quantum numbers which are not regarded as conventional for elementary particles. For

Table 1 - Possible supersymmetric particles

Type of conventional particle	Spin 1	Spin 1/2	Spin 0
Matter		quarks q	scalar quarks $\tilde{q}$
		leptons ℓ	supersymmetric $\tilde{\ell}$ (scalar leptons)
Massive gauge bosons	$W^{\pm}$, Z^{o}	supersymmetric heavy leptons $\tilde{W}^{\pm}$, $\tilde{Z}^{o}$, $\tilde{H}$	Higgs scalar H
Massless gauge bosons	photon γ gluon g	photino $\tilde{\gamma}$ + other nuinos $\tilde{\nu}$ gluinos $\tilde{g}$	

The placing of conventional Higgs scalars on the same line as the $W^{\pm}$, Z^{o} does not imply the suggestion that they lie in a common supermultiplet, but just that their supersymmetric heavy lepton partners would have similar properties. A Higgs scalar may also have a supersymmetric partner. We generally use the notation X, $\tilde{X}$ for a particle and its supersymmetric partner with the same internal quantum numbers, except that we use the historical notation $\tilde{\nu}$ for nuinos.

example, simple supersymmetric heavy leptons would have weak isospin I = 1, corresponding to the isospin of the $W^{\pm}$, whereas a "conventional" heavy lepton has I = 1/2. As another example, simple supersymmetric gluinos would belong to an $\underline{8}$ of colour: if they had a $\underline{3}$ of colour we would surely call them conventional quarks. For this reason our phenomenological discussion will often closely parallel that of Farrar and Fayet[3], although we remain open to extended supersymmetry schemes with different spectroscopies.

We now discuss in turn the possible properties and present experimental limits on supersymmetric particles in the classes listed in Table 1.

2.2 Scalar quarks

These have spin 0 and non-trivial colour for which the simplest possibilities are 3 and 8. Let us first consider the 3 possibilitiy which is realized in simple supersymmetry models. Possible couplings of such scalar quarks are illustrated in Fig. 1a,b,c,d.

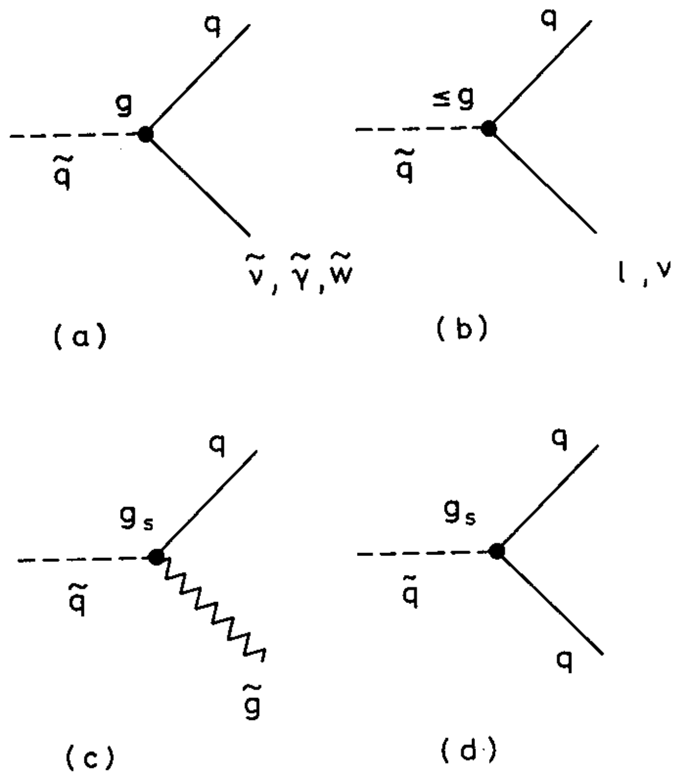

Fig. 1
Possible vertices involving scalar quarks $\tilde{q}$ and gluinos $\tilde{g}$

We might expect the vertex in Fig. 1a involving uncoloured particles to be of order the semi-weak gauge coupling g with $g^2/4\pi = O(\alpha)$, whereas the vertex in Fig. 1b might be the same magnitude or smaller, if it is analogous to cenventional Higgs scalar couplings, and the vertices in Figs. 1c,d involving only coloured particles should be strong $O(\alpha_s)$. If there were a coupling to conventional quarks and leptons of the type shown in Fig. 1b, then the scalar quark would need a very large mass $\gg M_W$. Otherwise it could mediate via the diagrams of Fig. 2 four fermion weak interactions of strength

$$O(g^2/m_{\tilde{q}}^2) \gtrsim O(G_F) \qquad (2.1)$$

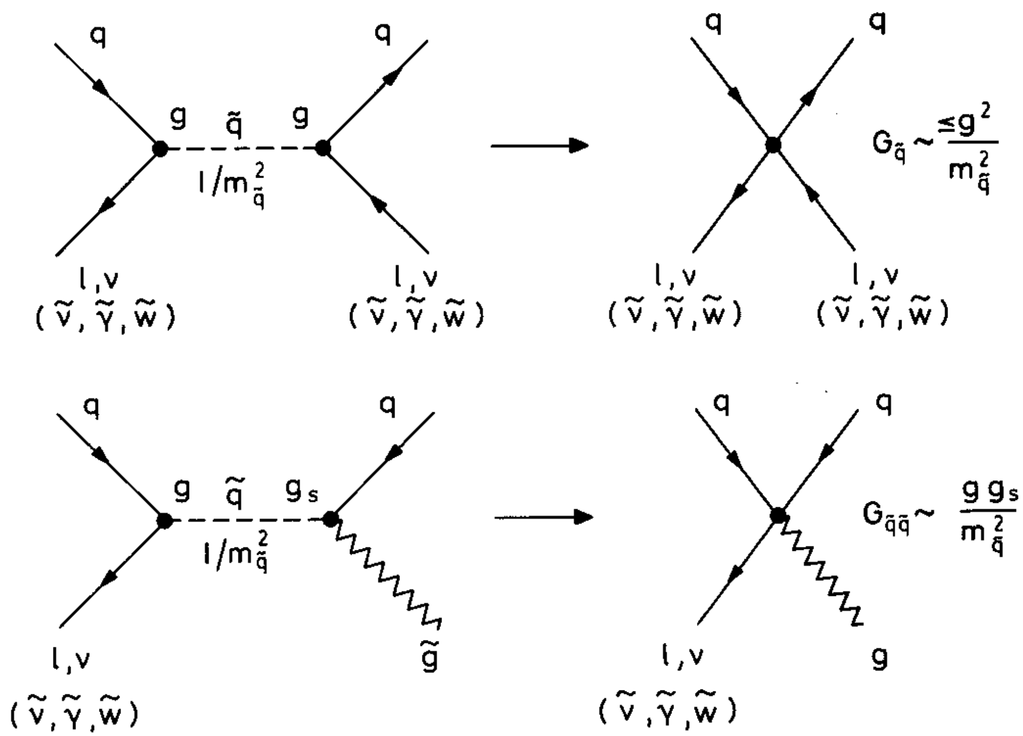

Fig. 2 - Possible effective low energy four particle interactions generated in supersymmetric theories.

which would be in conflict with the standard (V-A) form. This would not yet matter phenomenologically if the only quark q to couple to $\tilde{q}$ as in Fig. 1b were very heavy. Since even LEP cannot observe particles with masses $>> m_W$, we will assume that couplings of the type in Fig. 1b are absent, as in the simple models considered by Farrar and Fayet[3], or only exist for very heavy quarks q. There is then no stringent limit on the masses of possible scalar quarks. Limits on the production of hadrons containing scalar quarks are very weak. The discussion of scalar quark onia in section 3 suggest that the non-observation in e^+e^- annihilation of $J^{PC} = 1^{--}$ $(\tilde{q}\bar{\tilde{q}})$ onia only excludes $m_{\tilde{q}} \lesssim 1/2$ GeV. The best limits may come from the beam dump experiments[14]. If we assume that any exotic supersymmetric signal can be no larger than the observed prompt signal which is identified as being due to charm, then we have

$$\frac{\sigma(hh \to \tilde{q}\bar{\tilde{q}}X)}{\sigma(hh \to c\bar{c}\ X)} \cdot \frac{B(\tilde{q} \to \tilde{\nu} \text{ or } \tilde{\gamma} + X)}{B(c \to \nu + X)} \cdot \frac{\sigma(\tilde{\nu} \text{ or } \tilde{\gamma} + h \to X)}{\sigma(\nu + h \to X)} \lesssim 1 \qquad (2.2)$$

where the $B(\tilde{q} \to \tilde{\nu}$ or $\tilde{\gamma} + X)$ and $B(c \to \nu + X)$ are branching ratios whose ratio is probably $\gtrsim 1$. From the vertices of Fig. 1a we would deduce that a four-fermion diagram mediated by $\tilde{q}$ exchange (see Fig. 2) would yield

$$\frac{\sigma(\tilde{\nu} + h \to \tilde{\nu} + X)}{\sigma(\nu + h \to \nu + X)} \geq 0\left(\frac{m_W^4}{m_{\tilde{q}}^4}\right) \qquad (2.3)$$

and we might guess that phenomenologically for $m_{\tilde{q}}/m_c$ of 0(1)

$$\frac{\sigma(hh \to \tilde{q}\bar{\tilde{q}} + X)}{\sigma(hh \to c\bar{c} + X)} = 0\left(\frac{m_c^2}{m_{\tilde{q}}^2}\right) \qquad (2.4)$$

Naively putting (2.2, 2.3 and 2.4) together we would get

$$m_{\tilde{q}} \geq (m_W^4\, m_c^2)^{1/6} >> 2 \text{ GeV} \qquad (2.5)$$

but we are limited by phase space in the production of $\tilde{q}\bar{\tilde{q}}$ so that formula (2.4) is not appropriate for such heavy quarks at present energies. The present limit on $m_{\tilde{q}}$ from this consideration is probably a few GeV.

However, this limit would not apply to a scalar quark which only coupled to very heavy quarks via the couplings in Fig. 1a[3,6]. Nor would the limit apply to scalar quarks which only had couplings of the type in Fig. 1c. In this case the only way to avoid the gluino being stable would be to invoke a vertex of the type in Fig. 1d, which does not exist in simple supersymmetry models of the type discussed by Farrar and Fayet[3]. If such a vertex were to exist, one would deduce from the beam dump experiments that

$$\sigma(hh \to \tilde{q}\bar{\tilde{q}} + X) \leq \sigma(hh \to c\bar{c} + X) \qquad (2.6)$$

and so probably

$$m_{\tilde{q}} > 2 \text{ GeV} \qquad (2.7)$$

If a vertex of the type in Fig. 1d did not exist, then the lighter of the scalar quark and gluino would be stable. But there are limits on the masses of the stable hadrons that would be made of such particles[15], implying that

$m_{\tilde{q}}$ > 5 GeV (from high energy physics searches for stable particles (2.8)

> 16 GeV (from cosmology, for hadrons making anomalous Hydrogen)

> 38 GeV (from cosmology, for hadrons making anomalous Oxygen)

We conclude[16] that

$$m_{\tilde{q}} \geq 1/2 \text{ GeV (model independently)} \qquad\qquad (2.9)$$
$$\geq (2 \text{ to } 38) \text{ GeV (model dependently)}$$

and hope to do better with LEP.

2.3 Scalar leptons

These would be colourless scalar particles $\tilde{\ell}$ which might have couplings to $f\bar{f}$ as in Fig. 3a.

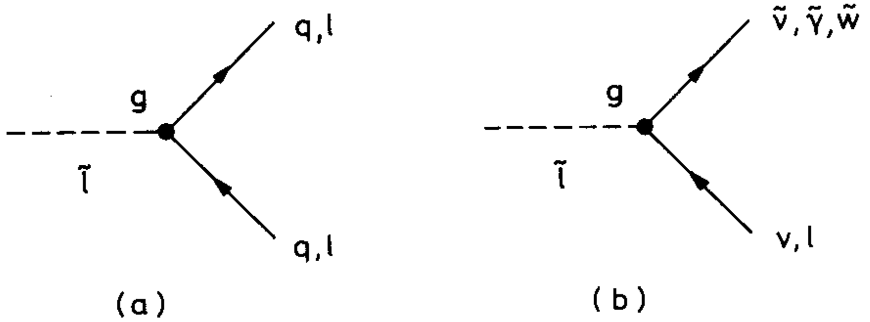

Fig. 3
Possible vertices involving scalar leptons $\tilde{\ell}$

$$g_{\tilde{\ell} f \bar{f}} = 0\,(g) \qquad\qquad (2.10)$$

rather than the usual order of magnitude of Higgs couplings

$$g_{\tilde{\ell} f f} = 0\left(\frac{g m_f}{m_W}\right) \qquad\qquad (2.11)$$

If such an $\tilde{\ell}$ had couplings of the type (2.10) to light fermions, then it would need to have

$$m_{\tilde{\ell}} > 100 \text{ GeV} \qquad\qquad (2.12)$$

Otherwise, for example, its contribution to (g-2) of the muon would be too large:

$$\Delta(g-2) \simeq 10^{-8}\left(\frac{m_W^2}{m_{\tilde{\ell}}^2}\right)\left(\ln(m_{\tilde{\ell}}^2/m^2) - \frac{7}{6}\right) \qquad\qquad (2.13)$$

which must be smaller than the present experimental error of $0(10^{-7})$. Furthermore a light scalar lepton $\tilde{\ell}$ would dominate the observed (V-A) structure of weak interactions among light fermions. On the other hand, the bound (2.12) would not apply to any supersymmetric scalar leptons which were only coupled to heavy quarks. For such a supersymmetric scalar lepton the best present limits may come from the heavy $q\bar{q}$ onium decays. Substitution of the coupling (2.10) instead of the usual coupling (2.11) in the standard formula[17] for $V \to H + \gamma$ yields

$$\frac{\sigma(V \to \tilde{\ell} + \gamma)}{\sigma(V \to \mu^+\mu^-)} = \frac{G_F\, m_q^2}{\sqrt{2}\,\pi\alpha}\left(1 - \frac{m_{\tilde{\ell}}^2}{m_V^2}\right) \to \frac{G_F\, m_W^2}{\sqrt{2}\,\pi\alpha}\left(1 - \frac{m_{\tilde{\ell}}^2}{m_V^2}\right) = 0(1) \qquad (2.14)$$

We therefore conclude that

$$m_{\tilde{\ell}} \gtrsim 3 \text{ GeV (for } \tilde{\ell} \text{ coupled to } c\bar{c}) \qquad\qquad (2.15)$$
$$\gtrsim 9 \text{ GeV (?) (for } \tilde{\ell} \text{ coupled to } b\bar{b})^*$$

The bound (2.12) is also invalidated if the $\tilde{\ell}$ has no diagonal couplings to conventional particles, as is in fact the case in the simple supersymmetric models of Farrar and Fayet[3] where all supersymmetric particles have a conserved quantum number R and only vertices of the type in Fig. 3b may occur. In this case no model-independent limit on $m_{\tilde{\ell}}$ can be obtained if they are neu-

---

* From DORIS experiments[18] looking at the inclusive γ spectrum in Υ decays.

tral, but the latest PETRA data[19] probably impose a limit close to 15 GeV on charged $\tilde{\ell}^{\pm}$, but not on neutral $\tilde{\ell}^{o}$.
We conclude that

$$m_{\tilde{\ell}} > 0 \text{ (model independently)} \quad > (3 \text{ to } 100) \text{ GeV (model dependently)} \qquad (2.16)$$

and hope to do better with LEP.

2.4 Supersymmetric Heavy Leptons

Since conventional heavy leptons have I = 1/2 for us the distinguishing feature of a supersymmetric heavy lepton $\tilde{W}^{\pm}$ will be its having isospin different from 1/2. This would be the case for a simple supersymmetric partner of the $W^{\pm}$ and Z^{0}. We know that the $e^{\pm}$ and $\mu^{\pm}$ leptons have $|I| = 1/2$, from their neutral and charged current interactions. We do not know the isospin of the $\tau^{\pm}$, which could well be 1. If it were, there should be an $I = 0$ neutral lepton τ^{o} which could not have been detected with present experiments because of the low $e^{+}e^{-} \to \tau^{o}\bar{\tau}^{o}$ cross-section at present energies. From higher energy $e^{+}e^{-}$ measurements we know there is no heavy lepton with mass $\leq$ 4 GeV, and this limit is probably increased to O(15) GeV by data from PETRA. There is no restriction on indirect effects of a supersymmetric heavy lepton $\tilde{W}$.

We conclude that either

$$m_{\tilde{W}} = 1.78 \text{ GeV (if it is the } \tau) \quad \text{or} \quad \gtrsim O(15) \text{ GeV (if it is not the } \tau) \qquad (2.17)$$

and hope to do better at LEP.

2.5 Gluinos

For the reasons discussed earlier, we will restrict ourselves to possible 8 gluinos, as found in simple supersymmetric theories[3,5,10]. Naively, one might expect such objects to be massless like the gluons, but this is not necessarily the case[20], and the gluino mass should probably be regarded as a free parameter. The gluinos can bind with ordinary quarks, gluons and gluinos to form very exotic looking hadrons, e.g.

$$\text{spin } 1/2 : g\tilde{g} \;,\quad q\bar{q}\tilde{g} \;;\quad \text{spin } 0 : \tilde{g}\tilde{g} \qquad (2.18)$$

As mentioned earlier[15], there are limits on the production of stable charged hadrons in hadron-hadron collisions which would suggest

$$m_{\tilde{g}} > 5 \text{ GeV} \qquad (2.19)$$

and cosmological constraints on the existence of stable particles which would give

$$m_{\tilde{g}} > 16 \text{ or } 38 \text{ GeV} \qquad (2.20)$$

if the gluino were stable. However, in general one expects the gluino to decay into final states with a nuino via the vertices shown in Fig. 1, so that the limits[15] do not apply.

One limit on unstable gluinos comes again from the beam dump experiments[5,17]. Let us assume that

$$\frac{\sigma(hh \to \tilde{g} + X)}{\sigma(hh \to c\bar{c} + X)} = O\left(\frac{m_c^2}{m_{\tilde{g}}^2}\right) \qquad (2.21)$$

and that the gluino then yields a nuino in the final state with probability O(1). The principle uncertainty now comes from

$$\frac{\sigma(\tilde{\nu} + h \to X)}{\sigma(\nu + h \to X)} \qquad (2.22)$$

which depends on the strength of the four-fermion interactions of Fig. 2 mediated by spin-0 quark $\tilde{q}$ (or leptons $\tilde{\ell}$) exchange. If $m_{\tilde{q}}$ (or $m_{\tilde{\ell}}$) $\sim m_W$, then the ratio (2.22) would be O(1) and one would deduce from (2.21) that

$$m_{\tilde{g}} \gtrsim 2 \text{ GeV} \qquad (2.23)$$

But this is clearly a highly model-dependent estimate.

There are some ways in which gluinos might show up indirectly, just as indirect evidence for gluons is now being found[21]. One possibility (see section 6) might be in onium decay, e.g. $1^{++} \to g\tilde{g}\tilde{g}$, but our present level of experimental understanding of the charmonium system is probably not precise enough to establish a firm limit[22]. A more promising signal seems to be in deep inelastic scattering. We know that at present energies about 50% of the nucleon momentum are carried by neutral particles, and these are usually identified as gluons. Different field theories make different predictions about the asymptotic values of the ratio of the gluon to quark momentum fractions, based on diagrams of the type shown in Fig. 4.

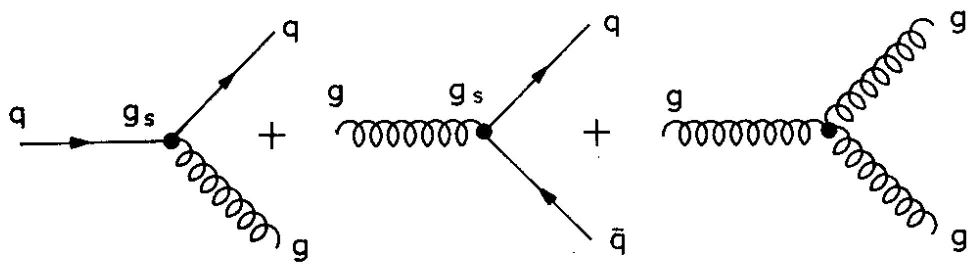

Fig. 4 - Important vertices for scaling violations in conventional QCD

The present data are in qualitative agreement[23] with the asymptotic prediction of QCD which is

$$\int_0^1 dx\, F_2^N(X, Q^2) \to \frac{3f}{16+3f} = 3/7 \text{ for } f = 4 \qquad (2.24)$$
$$= 9/17 \text{ for } f = 6$$

The prediction (2.24) gets dramatically modified if there are gluinos, because of the diagrams shown in Fig. 5.

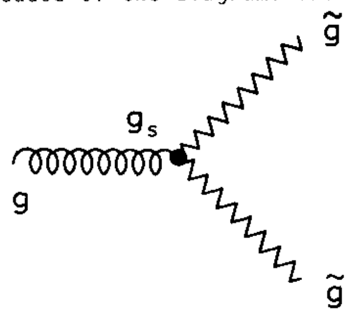

Fig. 5 - New vertex in QCD with gluinos

We now have

$$\frac{3f}{16+3f} \to \frac{3f}{16+3f+18} = \frac{3f}{34+3f} = \frac{6}{23} \text{ for } f = 4 \qquad (2.25)$$
$$= \frac{9}{26} \text{ for } f = 6$$

The transition from an asymptotic approach to (2.24) or (2.25) will presumably occur for $Q^2 \gtrsim 4m_{\tilde{g}}^2$ as illustrated in Fig. 6.

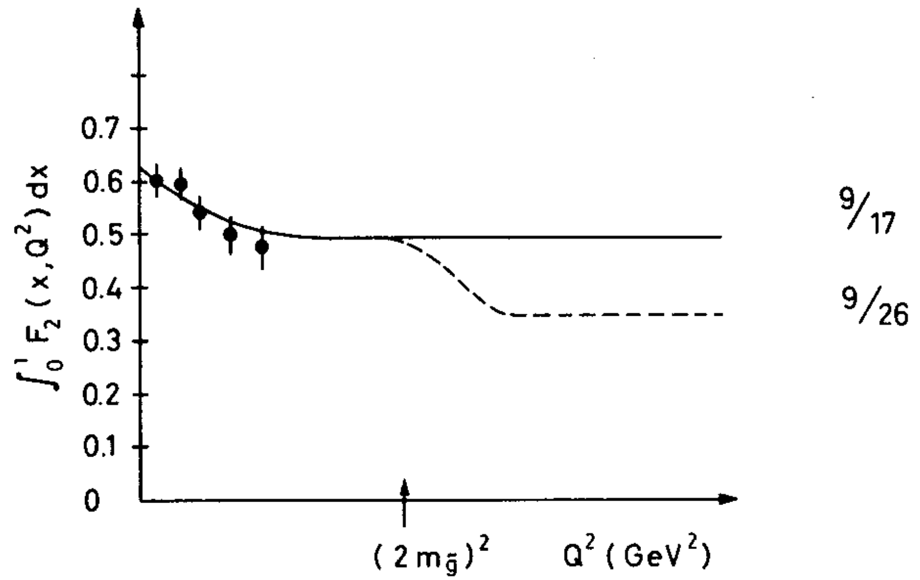

Fig. 6 - Effect of gluinos on $\int_0^1 dx\, F_2(X,Q^2)$

If we accept that the present data are inconsistent with gluino effects below $Q^2 = 40\ GeV^2$, then we deduce that

$$m_{\tilde g} > 3\ GeV \qquad (2.26)$$

This limit could clearly be dramatically improved with one of the projected ep colliding ring machines[7]. If one accepts that $\int_0^1 dx\, F_2(X,\, Q^2)$ can be reliably measured up to $Q^2 = 10000\ GeV^2$, two orders of magnitude better than at present, then the indirect effects of $\tilde g$ with masses up to 50 GeV could be detected. But this is for the future. For the moment we conclude that

$$\left.\begin{aligned} m_{\tilde g} &\gtrsim 3\ GeV \quad \text{(model independently)} \\ &\gtrsim (5 \text{ to } 38)\ GeV \text{ (if they are stable)} \end{aligned}\right\} \qquad (2.27)$$

and hope to do better with LEP.

2.6 Nuinos and Gravitinos

This is the collective name applied to a variety of spin 1/2 massless (or essentially massless) colourless neutrino-like objects which occur in supersymmetric theories. They are partners of particles of various other spins and include photinos $\tilde\gamma$, goldstinos, etc. They characteristically have weak interactions (see the vertices Figs. 1 and 3), but the strength of their effective four-fermion interactions is rather uncertain because of the unknown masses of the mediating particles (e.g. $\tilde g$ and $\tilde \chi$ in Figs. 2 and 3).

There are no important restrictions on their numbers or masses from high energy physics phenomenology. As discussed in ref. 24 the best such limit on the number of neutrinos, which would also be the correct order of magnitude for $\tilde\nu$, if they have normal strength weak interactions, is

$$N_{\nu+\bar\nu} \lesssim O(6000) \qquad (2.28)$$

from the upper limit on $K \to \pi(\nu\bar\nu)$ However, cosmology does place important restrictions on the number and masses of neutrino-like objects[25]. If they have masses $\lesssim O(1)$ MeV, then there can be at most one more of them (assuming that $m_{\nu e}$, $m_{\nu\mu}$ and $m_{\nu\tau}$ are small), and the sum of the masses of neutrino-like objects must be < 50 eV. The proliferation of nuinos in supersymmetric theories can only be consistent with this cosmological limit if most of the $m_{\tilde\nu}$ are > O(1) MeV or if their couplings are much weaker than the standard G_F[26]. In simple supersymmetric theories of the type studied by Farrar and Fayet[3] where there is a conserved R quantum number, at least one of the nuinos must be absolutely stable. Cosmology tells us that there can be no stable heavy leptons with masses between O(100) eV and O(1) GeV. Hence this type of model can only be consistent with all the cosmological constraints if it has exactly one very light $\tilde{\nu}$, and all the rest have masses $\geq$ O(1) MeV.

We conclude that if the conventional cosmology arguments are correct:

$$m_{\tilde{\nu}} \geq O(1)\text{ MeV, except for at most one} \quad (2.29)$$

which must have $m_{\tilde{\nu}} < 50$ eV

(exactly one in theories with R symmetry)

and hope to do better with LEP.

We conclude this section with some remarks about gravitinos. These are very light spin 3/2 particles $\tilde{G}$ which are supersymmetric partners of the graviton. A phenomenological restriction on their mass has been discussed by Fayet[11]. In a class of supersymmetric models

$$m_{\tilde{G}} \cdot \approx \frac{1}{m_{Planck}} \frac{\Delta m^2}{g} \quad (2.30)$$

where $m_{Planck} \sim 10^{19}$ GeV, g is a typical gauge coupling, and Δm^2 is a typical supersymmetric mass splitting, estimated to be $\Delta m^2 \lesssim 1/4\, m_W^2 \sim 2500\ \text{GeV}^2$. (2.31)

One finds from (2.30) that

$$m_{\tilde{G}} \sim \begin{cases} 2 \times 10^{-7}\text{ eV} & \text{for } \Delta m^2 \sim (15\text{ GeV})^2 \\ 2 \times 10^{-6}\text{ eV} & \text{for } \Delta m^2 \sim (40\text{ GeV})^2 \end{cases} \quad (2.32)$$

In such a theory the $\tilde{G}\tilde{\gamma}$ photon coupling has[11] a form

$$\tilde{V}_\mu \sim \frac{\gamma_\mu q^2}{m_{\tilde{G}} m_{Planck}} \quad (2.33)$$

The best restriction on the mass of such a gravitino comes from the non-observation of $J/\psi \to \tilde{\bar{G}}\tilde{\gamma} + \tilde{\bar{G}}\tilde{\gamma}$, which Fayet estimates to have a rate

$$\frac{\Gamma(J/\psi \to \tilde{\bar{G}}\tilde{\gamma} + \tilde{\bar{G}}\tilde{\gamma})}{\Gamma(J/\psi \to e^+e^-)} \approx O\left(\frac{m^2_{J/\psi}}{2\Delta m^2}\right) \quad (2.34)$$

From an experimental limit

$$\frac{\Gamma(J/\psi \to \text{unobserved neutrals})}{\Gamma(J/\psi \to e^+e^-)} < 1/10 \quad (2.35)$$

he[11] then deduces a limit on Δm^2 and hence on $m_{\tilde{G}}$:

$$m_{\tilde{G}} > 1.5 \times 10^{-8}\text{ eV} \quad (2.36)$$

which is comfortably consistent with the expected range (2.32). If a sensitivity of the type (2.35) could be established also for $\Upsilon \to$ unobserved neutrals, then one would be sensitive to

$m_{\tilde{G}} \sim (1 \text{ or } 2) \times 10^{-7}$ eV, and approaching the range (2.32).

We conclude that

$$m_{\tilde{G}} > 1.5 \times 10^{-8}\text{ eV}$$

and expect to do better at LEP.

3. Scalar Quarks

We will discuss the phenomenology of these particles at LEP by analogy with that for a conventional spin 1/2 heavy quark. The most distinctive features of their production in e^+e^- collisions are[27]:

- the appearance of vector onium peaks
- a threshold for high sphericity events
- a high energy 2-jet structure with distinctive angular distribution.

We start with the $\tilde{q}\bar{\tilde{q}}$ onia, which have a distinctively different spectrum from that of ordinary $Q\bar{Q}$ onia. There are no spin effects, and therefore no 4-fold degeneracy at each orbital angular momentum level, as shown in Fig. 7.

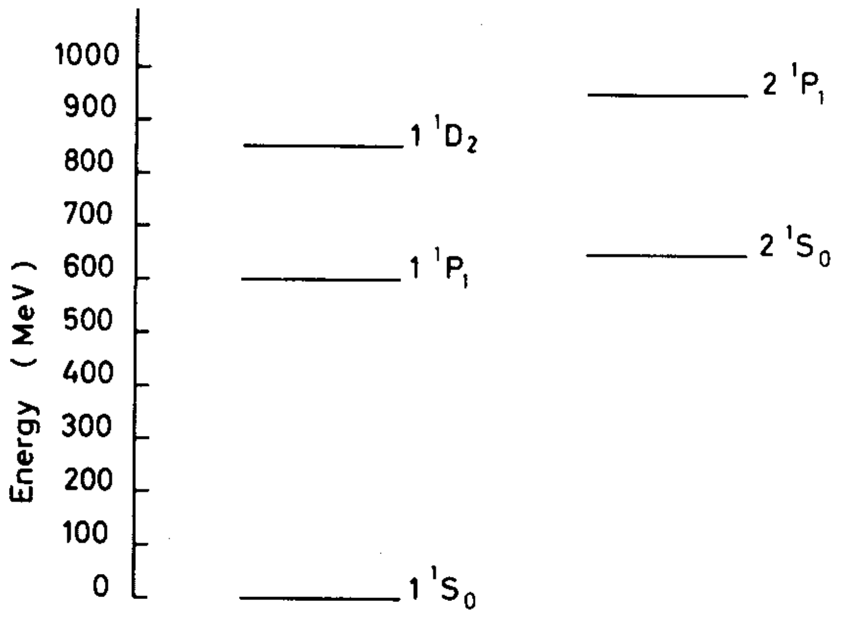

Fig. 7
Spectroscopy of scalar quark onia for $m_{\tilde{q}} \simeq 50$ GeV

Another important distinctive feature is that the lowest lying S-states are <u>not</u> produced directly in e^+e^- annihilation since they have $J^{PC} = 0^{++}$. The states $J^{PC} = 1^{--}$ most easily produced in e^+e^- annihilation are P-waves, which turns out to have dramatic implications for their observability[28]. The colour forces between scalar quarks are the same as those between spin 1/2 quarks, so that the level structure of the states should be identical with that of the centre-of-gravity of the corresponding spin 1/2 quark onium spin multiplets. The energy levels in Fig. 7 correspond to expectations for a nominal scalar quark of mass 50 GeV using the asymptotically free potential model of Krasemann and Ono[29]. We notice that the mass difference ($\sim$ 350 MeV) between the (in principle) observable $\tilde{q}\bar{\tilde{q}}$ $^1P_1$ onia is considerably smaller than that ($\sim$650 MeV) between the (in principle) observable $Q\bar{Q}$ $^3S_1$ onia in the same mass region.

We now turn to the production of these onia in e^+e^- annihilation, which is determined by

$$\int_{\text{Resonance}} d\mathcal{E}_{cm}\ \sigma(V \to \text{all}) = 6\pi^2 \frac{\Gamma(V \to e^+e^-)}{m_V^2} \qquad (3.1)$$

Using the standard width formula

$$\Gamma(V \to e^+e^-) = \frac{P_{c.m.}}{4m_V} \frac{N_c}{2J+1} \int \frac{d\Omega}{4\pi^2} \sum_{Pol} |M|^2 \qquad (3.2)$$

and the scalar quark annihilation matrix element

$$M = \frac{e^2 e_{\tilde{q}}}{(2E)^2} \bar{u}(p')\gamma_\mu V(q')\ \varepsilon^\mu\ \phi_p'\ (r = 0)\ 1/m_{\tilde{q}} \qquad (3.3)$$

where $\phi_p(r)$ is the P-wave wave function, Krasemann[28] has calculated

$$\Gamma(\tilde{q}\bar{\tilde{q}}(^1P_1) \to e^+e^-) = \alpha^4 \, e_{\tilde{q}}^2 \, \frac{|R_p'(0)|^2}{2m_{\tilde{q}}^4}$$

where $R_p(r)$ is the radial wave function. One can then, given a potential model, calculate the ratio

$$K \equiv \frac{\Gamma(\tilde{q}\bar{\tilde{q}}\,(^1P_1) \xrightarrow{\gamma} e^+e^-)}{\Gamma(Q\bar{Q}(^3S_1) \xrightarrow{\Upsilon} e^+e^-)} = \frac{|R_p'(0)|^2}{2m_{\tilde{q}}^2\,|R_s(0)|^2} \qquad (3.5)$$

and one finds the results listed in Table 2.

Table 2 - The Ratio K of Equation (3.5)

m_s ($=m_Q$) (GeV)	Harmonic Oscillator Potential	Krasemann-Ono[29] Asymptotically free potential
1	5%	2.2%
8.5	0.6%	0.2%
17	0.3%	0.08%

These results indicate that the ADONE and SPEAR searches for narrow onium resonances were not sufficiently sensitive to have detected any $\tilde{q}\bar{\tilde{q}}$ onia. Certainly higher energy e^+e^- machines have nothing to say on this matter.

The only limit on $\tilde{q}\bar{\tilde{q}}$ onia seems to come from the very low energy data ($E_{c.m.} \lesssim 1.5$ GeV) which are perfectly described by the ρ, ω and ϕ peaks[30]. We therefore arrive at the $m_{\tilde{q}} > 1/2$ GeV limit quoted in section 2.2. We also conclude that LEP will not be able to do any better looking for $\tilde{q}\bar{\tilde{q}}$ vector onia.

A slightly less depressing possibility is looking for the S-wave 0^{++} $(\tilde{q}\bar{\tilde{q}})$ onia in $\gamma\gamma$ collisions. An analogous calculation to the P-wave analysis above yields

$$\Gamma(\tilde{q}\bar{\tilde{q}}(^1S_o) \to \gamma\gamma) = \frac{3}{2}\,\alpha^4\, e_{\tilde{q}}^4 \,\frac{|R_s(0)|^2}{m_{\tilde{q}}^2} \qquad (3.6)$$

which is therefore only a factor 2 smaller than $\Gamma(Q\bar{Q}(^1S_o) \to \gamma\gamma)$. Recent $e^+e^- \to \gamma\gamma$ collision experiments[31] have seen the $\eta'(960)$, so the non-observation of a $(\tilde{q}\bar{\tilde{q}})$ peak in this range would again imply $m_{\tilde{q}} > \frac{1}{2}$ GeV. For $\gamma\gamma$ collisions one has the production rate

$$\sigma(e^+e^- \to e^+e^-X) = \frac{8\pi\alpha\;\Gamma(X \to \gamma\gamma)}{m_X^3}\;(\ln\,(t^{max}/t^{min}) + 1)$$

which when combined with (3.6) yields the rates for $\gamma\gamma$ production of $\tilde{q}\bar{\tilde{q}}$ (^1S_o) onia listed in Table 3.

Table 3 - Rates for $e^+e^- \to e^+e^-\;{}^1S_o\,(\tilde{q}\bar{\tilde{q}})$

m(GeV)	5	10	20	30	40
Events/day	570	71	9	3	1

These rates were evaluated assuming $\Gamma(X \to \gamma\gamma)$ = 10 keV and a luminosity of 5 x 10^{31} $cm^{-2}sec^{-1}$. The minimum momentum transfer squared t_{min} is given by $t_{min} = (m^4 \cdot m_e^2/s^2)$.

Taking into account the problems in picking out such an onium once produced, we conclude that LEP has a sensitivity to

$$m_{\tilde{q}} \lesssim 10 \text{ GeV} \qquad (3.8)$$

using this process.

The basic QED rate for $\tilde{q}\bar{\tilde{q}}$ pairproduction above threshold is

$$\sigma_o = \sigma(e^+e^- \to \gamma \to \tilde{q}\bar{\tilde{q}}) = \frac{\pi\alpha^2 e_{\tilde{q}}^2}{s} \quad \beta^3 \underset{s \to \infty}{\to} 66 \left(\frac{e_{\tilde{q}}^2}{s} \right) nb \qquad (3.9)$$

where $e_{\tilde{q}}$ is the charge (-1/3, 2/3), we have assumed as usual that the q is a colour triplet, and $s = Q^2 = E^2_{c.m.}$. The direct channel Z^o diagrams modify (3.9) by a factor

$$\frac{\sigma(e^+e^- \to \gamma, Z \to \tilde{q}\bar{\tilde{q}})}{\sigma_o} = 1 - \frac{2 s \chi v v_{\tilde{q}}}{(s/m_Z^2-1) + \Gamma_Z^2/(s-m_Z^2)} + \frac{s^2 \chi^2 (v^2 + a^2)(v_{\tilde{q}}^2)}{(s/m_Z^2-1) + \Gamma_Z^2/m_Z^2} \qquad (3.10)$$

where a, v are the electron weak couplings

$$a = -1, \quad v = 0 \quad \text{for } \sin^2\theta_W = 0.25 \qquad (3.11)$$

and $v_{\tilde{q}}$ is the scalar quark weak coupling

$$v_s = \begin{matrix} 1/3 \\ -2/3 \end{matrix} \quad \text{for } Q = \begin{matrix} 2/3 \\ -1/3 \end{matrix} \quad \text{for } \sin^2\theta_W = 0.25 \qquad (3.12)$$

The rate for $e^+e^- \to \tilde{q}\bar{\tilde{q}}$ production is shown in Table 4 as a function of centre-of-mass energy $\sqrt{s}$ and the scalar quark mass $m_{\tilde{q}}$. The rate was evaluated for a luminosity of $5 \times 10^{31} cm^{-2} sec^{-1}$.

We conclude that $e^+e^- \to \tilde{q}\bar{\tilde{q}}$ has an observable cross-section for

$$m_{\tilde{q}} < 50 \text{ GeV}. \qquad (3.13).$$

The production of scalar quarks near threshold will lead to high sphericity events.Scalar quark production can be separated from spin 1/2 quark production by observing:

Table 4 - Event rates for scalar quarks

$\sqrt{s}$ (GeV)	R(2/3 e)	Events(2/3e)/day	R(1/3e)	Events(1/3e)/day
50	0.34	51	0.09	13
60	0.34	35	0.12	12
70	0.37	28	0.22	17
80	0.67	39	1.41	82
86.4	12.30	614	47.80	2387
90	1.68	77	5.48	252
120	0.37	10	0.24	6
160	0.35	5	0.16	2
200	0.35	3	0.14	1

1) A cross section proportional to β^3/s.

2) A $\sin^2\theta$ angular distribution.

3) A smaller step than observed for spin 1/2 quark (factor 4).

4. Scalar Leptons

As discussed in section 2.3, the main difference between these and conventional Higgs particles is possibly their large coupling (2.10) to fundamental fermions. Thus their phenomenology at LEP has many points in common with that of the conventional Higgs particles discussed in our previous report[4]. We start with neutral $\tilde{\ell}$.

$V \to \tilde{\ell}^o + \gamma$[17]: This process is even more favourable than $V \to H^o + \gamma$ because of the possibly large coupling (2.10) which may yield the estimate

$$\frac{\Gamma(V \to \tilde{\ell}^o + \gamma)}{\Gamma(V \to \mu^+\mu^-)} = O(1) \qquad (4.1)$$

This process will yield any $\tilde{\ell}^o$ with mass less than any onium V accessible to LEP.

$Z^o \to \tilde{\ell}^o + (\ell^+\ell^-)$[32]: The $Z^oZ^o\tilde{\ell}^o$ coupling is probably the same order of magnitude as the $Z^oZ^oH^o$ coupling, unless it is forbidden by a selection rule such as the R quantum number, so that this reaction should have a comparable branching ratio

$$B(Z^o \to \tilde{\ell}^o + (\ell^+\ell^-)) \geq O\,(10^{-6}) \qquad (4.2)$$

$$\text{for } M_{\tilde{\ell}^o} < 50 \text{ GeV}.$$

As before[4], we conclude that an $\tilde{\ell}^o$ with mass less than 50 GeV could be found using this reaction.

$e^+e^- \to Z^o + \tilde{\ell}^o$[33]: Again the discussion parallels our previous discussion[4], the conclusion of which was that a Higgs with

m_H < (the maximum centre-of-mass energy)-100 GeV $\gtrsim$ 100 GeV could be found this way.

We continue with charged $\tilde{\ell}^\pm$.

$e^+e^- \to \tilde{\ell}^+\tilde{\ell}^-$: The cross section was evaluated for $\sin^2\theta_W = 0.2$ and the event rate is listed in Table 5 assuming a luminosity of $5 \times 10^{31}\text{cm}^{-2}\text{sec}^{-1}$, and calculating the cross section using just direct channel γ and Z^o diagrams. This is appropriate for scalar leptons which are not partners of the $e^\pm$. For such particles there are extra t-channel nuino exchange diagrams which give larger cross sections and a pronounced forward-backward asymmetry[6].

Table 5 - Events/day for $e^+e^- \to \tilde{\ell}^+\tilde{\ell}^-$

$\sqrt{s}$ GeV	70	90	93.9	97	100	120	140	160	180	200
$m_{\tilde{\ell}^+}$=10GeV	18	23.1	69	16	10.1	6	5	4	3	2
$m_{\tilde{\ell}^+}$=50GeV						1	2	2	2	2

This reaction can be used to search for scalar leptons with mass $\leq$ 50 GeV.

The next reaction is better suited to look for high mass scalar leptons, if one $W^\mp Z^o\tilde{\ell}^\pm$ coupling is not forbidden.

$e^+e^- \to W^\pm\tilde{\ell}^\mp$[34]: This cross-section rises rapidly above threshold particularly if the $\tilde{\ell}^\mp$ has a small mass, so that one may benefit from the nearby Z^o pole. Typical event rates evaluated for $L = 5 \cdot 10^{31}\text{cm}^{-2}\text{sec}^{-1}$ are shown in Table 6.

Table 6 - Events/day for $e^+e^- \to W^{\pm}\tilde{\chi}^{\mp}$

$\sqrt{s}$ (GeV)	100	120	140	160	180	200
$m_{\tilde{\chi}^+}$=10GeV	15392	1575	547	260	143	87
$m_{\tilde{\chi}^+}$=50GeV			257	192	122	77
$m_{\tilde{\chi}^+}$=83.6GeV					68	56

As before, we conclude that an $\tilde{\chi}^{\mp}$ could be found in this way if

$$m_{\tilde{\chi}^{\pm}} \leq 100 \text{ GeV} \quad (4.4)$$

Not all the above processes exist for theories with an R quantum number carried by supersymmetric particles[3]. However, two processes which do then exist are:

$e^+e^- \to \tilde{\chi}^o\bar{\tilde{\chi}}^o$: This proceeds by direct channel Z^0 exchange and has the cross section

$$\sigma = \left(\frac{\alpha^2\pi}{3s}\right) \beta^3 \frac{s^2 h^2 (v^2+a^2) v^2_{\tilde{\chi}^o}}{(s/m_Z^2-1)^2 + \Gamma_Z^2/m_Z^2} \quad (4.5)$$

where $v_{\tilde{\chi}^o} = \pm 1$ for an $I = 1/2$ $\tilde{\chi}^o$, and $h = 4.4 \times 10^{-5}$ GeV^{-2} [35]. If the $\tilde{\chi}^o$ has mass $\lesssim m_{Z^o/2}$ then at the Z^o peak

$$\frac{\Gamma(Z^o \to \tilde{\chi}^o\tilde{\chi}^o)}{\Gamma(Z^o \to e^+e^-)} \approx 1/4\ \beta^3 \quad \text{for } \sin^2\theta \sim 0.25 \quad (4.6)$$

We conclude that this reaction could be used to look for

$$m_{\tilde{\chi}^o} < 40 \text{ GeV} \quad (4.7)$$

$e^+e^- \to \tilde{\chi}^+\tilde{\chi}^-$: As discussed above, this should give access to

$$m_{\tilde{\chi}^{\pm}} \lesssim m_{Z^o}/2$$

We conclude that LEP gives us access to

$$m_{\tilde{\chi}^o,\tilde{\chi}^{\pm}} \lesssim \begin{cases} (40 \text{ to } 50) \text{ GeV (model independently)} \\ 100 \text{ GeV (model dependently)} \end{cases} \quad (4.8)$$

5. Supersymmetric Heavy Leptons

In this chapter we discuss the production rates and the signatures for the supersymmetric lepton $W^{\pm} \to \tilde{W}^{\pm}$. These leptons are pairproduced according to the usual Feynmann graph:

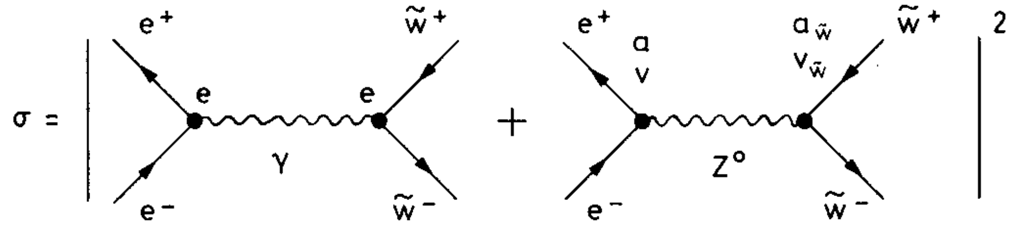

Fig. 8 - Diagrams for $e^+e^- \to \tilde{W}^+\tilde{W}^-$

Normalized to the point cross section this yields[35]:

$$R_{\tilde{W}^+\tilde{W}^-} = 1 - \frac{2sv_e v_{\tilde{W}} \cdot h}{(s/m_Z^2-1) + \Gamma_Z^2/s-m_Z^2} + \frac{s^2 \cdot h^2 (v_e^2 + a_e^2)(v_{\tilde{W}}^2 + a_{\tilde{W}}^2)}{(s/m_Z^2-1)^2 + \Gamma_Z^2/m_Z^2} \quad (5.1)$$

where $h = G_F/(8\sqrt{2}\pi\alpha) \approx 4.4 \times 10^{-5}$ GeV^{-2}.

The weak axial and vector couplings are given by:

$$a = 2\,(I_3^L - I_3^R) \text{ and } v = 2\,(I_3^L + I_3^R) - 4e_{\tilde{W}} \sin^2\theta_W. \quad (5.2)$$

With $\sin^2\theta_W = 0.25$ we find that a supersymmetric lepton with $I_3^L = \pm 1$ has $a_{\tilde{W}^\pm} = \pm 2$ and $v_{\tilde{W}^\pm} = \pm 1$. The ratio $R_{\tilde{W}^+\tilde{W}^-}$ and the number of events per day evaluated assuming $L = 5 \times 10^{31} cm^{-2} sec^{-1}$ and $\beta = p_{\tilde{W}}/E_b \simeq 1$ are listed in Table 7 below.

Table 7 - Pair production of $e^+e^- \to \tilde{W}^+\tilde{W}^-$ and comparison to $e^+e^- \to \mu^+\mu^-$

$\sqrt{s}$	$R_{\tilde{W}^+\tilde{W}^-}$	Events/day	$R_{\tilde{W}^+\tilde{W}^-}/R_{\mu^+\mu^-}$	$A_{\tilde{W}^+\tilde{W}^-}/A_{\mu^+\mu^-}$
50	1.14	170	1.11	1.72
60	1.47	152	1.34	1.30
70	3.03	231	2.16	0.7
80	20.3	1178	4.18	0.02
86.4	703	35065	5.0	2.29
90	80.5	3703	4.76	1.80
120	3.33	86	2.27	1.10
160	2.08	30	1.71	1.32
200	1.86	18	1.59	1.40

The rates are sufficient to explore the energy range opened up by LEP provided a clean signature exist. For $m_{\tilde{W}^\pm} > m_{W^\pm}$ the dominant decay mode will presumably be $\tilde{W}^+ \to W^+ + \tilde{\nu}$ - i.e. $e^+e^- \to \tilde{W}^+\tilde{W}^- \to W^+W^- + \tilde{\nu}\bar{\tilde{\nu}}$.

The main background will thus come from $e^+e^- \to W^+W^-$ which is expected to have a larger cross-section[36]. For energies well above threshold the two reactions might be separated by a cut on the total energy in the final state. In general, the supersymmetric

lepton will show up as a sharp step, reflecting the spin 1/2 nature, on the yield curve for any typical W^+W^- signature like $e\mu$ or $e(\mu)$ hadrons. For $m_{W^\pm} > m_{\tilde{W}^\pm}$ the prominent decay modes will be:

$\tilde{W} \to \ell + \bar{\nu}_\ell + \tilde{\nu}$, $\tilde{W} \to q + \bar{q} + \tilde{\nu}$ and $\tilde{W} \to \tilde{q} + \bar{\tilde{q}} + \tilde{\nu}$ where $\tilde{q}$ denotes the supersymmetric scalar quark. The branching ratios into leptons will be given by:

$$B(\tilde{W} \to \ell + \bar{\nu}_\ell + \bar{\nu}_{\tilde{W}}) = \frac{1}{4 \cdot N_F + \frac{1}{2} N_{\tilde{F}} + N_{\tilde{W}}} \qquad (5.3)$$

Here N_F denotes the number of Fermion families, $N_{\tilde{F}}$ the number of scalar families ($q \to \tilde{q}$, $\ell \to \tilde{\ell}$) and $N_{\tilde{W}}$ the number of supersymmetric leptons. Hence the branching ratios into leptons and hadrons will presumably be comparable to those predicted for conventional leptons leading to final states $e^+e^- \to \tilde{W}^+\tilde{W}^- \to e(\mu)$ + hadrons + neutrino as shown in Fig. 9 with a probability of the order of 20%.

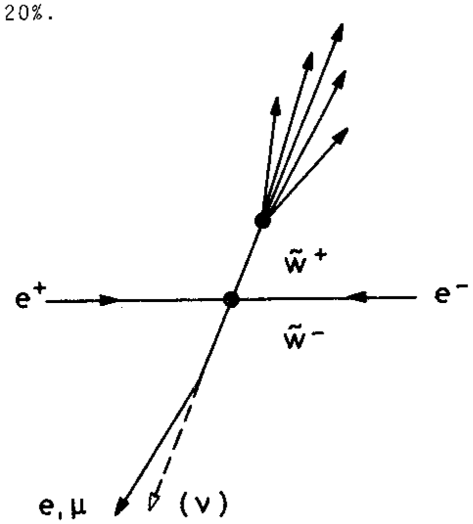

Fig. 9 - Signature for $e^+e^- \to \tilde{W}^+\tilde{W}^-$

Such event topologies can be used to separate lepton production from normal hadronic events. Conventional and supersymmetric lepton pairproduction can be separated by the larger production cross-section for the latter. The ratio $\sigma(e^+e^- \to \tilde{W}\,\tilde{\bar{W}})/\sigma(e^+e^- \to \mu^+\mu^-)$ is evaluated and listed in Table 7. It varies between 1.1 and 5 in the energy range 2E = 50 GeV to 200 GeV. Also the forward backward asymmetry[35] defined by

$$A = \frac{\int_0^1 d\cos\theta d\sigma - \int_{-1}^0 d\cos\theta d\sigma}{\int_{-1}^1 d\cos\theta d\sigma} \qquad (5.4)$$

can be used to separate conventional and supersymmetric leptons.

$$A = \frac{3/2\ \chi(-e\ a_e a_f + 2\ v_e a_e v_f a_f\ \chi)}{(e^2 - 2e\ v_e v_f + \chi^2 (v^2 + a^2)\ (v_f^2 + a_f^2))} \qquad (5.5)$$

with $\chi = h \dfrac{s\ m_Z^2}{s - m_Z^2}$.

The ratio $A_{\tilde{W}\tilde{\bar{W}}} / A_{\mu^+\mu^-}$ has been evaluated and is also listed in Table 7. It is clear that the angular distribution is quite different for the production of conventional and supersymmetric leptons.

6. Gluinos

Since gluinos do not have electric or weak charges, they cannot be produced directly in e^+e^- annihilation. This same problem is encountered in gluon hunting, and several ways of circumventing it have been proposed. Many of these have analogues in gluino hunting.

Gluino bremsstrahlung

It has been proposed[37] to look for 3 jet processes in e^+e^- annihilation generated by $e^+e^- \to q\bar{q}g$. In supersymmetric theories an analogous process may exist: $e^+e^- \to \tilde{q}\bar{q}\tilde{g}$ shown in Fig. 10.

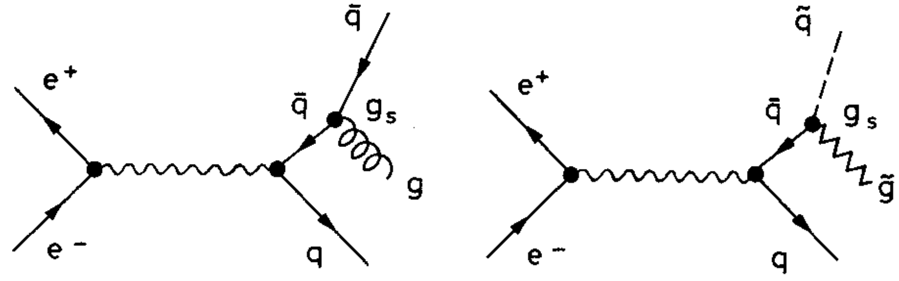

Fig. 10 - Gluino bremsstrahlung diagrams

The basic $\tilde{q}\bar{q}\tilde{g}$ vertex is expected to have normal QCD strength $O(\alpha_s)$, so that the two processes $q\bar{q}g$ and $\tilde{q}\bar{q}\tilde{g}$ should be competitive if there is no phase-space suppression of the supersymmetric process. If we denote the quark and gluon energies by

$$x_q = \frac{2E_q}{E_{c.m.}} \quad , \quad x_{\tilde{g}} = \frac{2E_{\tilde{g}}}{E_{c.m.}} \qquad (6.1)$$

then the cross-section from primordial $q\bar{q}$ production is of order (neglecting γ^* - Z interference)

$$\frac{1}{\sigma_{p_t}} \cdot \frac{d^2\sigma}{dx_q dx_{\bar{q}}} = \frac{7}{24\pi}\ \alpha_s\ \frac{x_{\tilde{g}}^2}{(1-x_q)(1-x_{\tilde{q}})} \qquad (6.2)$$

There is also a small contribution from primordial $\tilde{q}\tilde{q}$ production which we will neglect. This process is distinguishable from conventional[37] gluon bremsstrahlung (neglecting γ^* - Z interference)

$$\frac{1}{\sigma_{tot}} \frac{d^2\sigma}{dx_q dx_{\bar{q}}} = \frac{2\alpha_s}{3\pi} \frac{x_q^2 + x_{\bar{q}}^2}{(1-x_q)(1-x_{\bar{q}})} \qquad (6.3)$$

by virtue of the different energy distributions of the jets, or equivalently their relative angular distribution. For example, if in either case one works in the centre-of-mass of the two less energetic jets and defines $\tilde{\theta}$ to be the angle between the common axis of these jets and the axis of the most energetic jet, then one finds an angular distribution[38]

$$\frac{dN}{d(\cos\tilde{\theta})} \sim (1 + 0.2 \cos^2\tilde{\theta}) \qquad (6.4)$$

in the $q\bar{q}\tilde{g}$ case (6.2), and in the $q\bar{q}g$ (6.3)

$$\frac{dN}{d(\cos\tilde{\theta})} \sim (1 + 2 \cos^2\tilde{\theta}) \qquad (6.5)$$

(These distributions for the energy fraction T of the most energetic jet $\lesssim 0.9$). The two angular distributions (6.4) and (6.5) are almost as distinctive as the angular distributions $(1 \pm \cos^2\theta)$ for two-jet events $e^+e^- \to q\bar{q}$ or $\tilde{q}\bar{\tilde{q}}$. Future analysis of 3-jet candidate events at PETRA and PEP should bear in mind the possibility of a gluino bremsstrahlung contribution.

Finding a threshold in e^+e^- annihilation at high energy for the appearance of such events-signatured by exotic properties of the $\tilde{q}$ and $\tilde{g}$ jets such as high $\langle p_T \rangle$ and/or missing neutral energy from $\tilde{\gamma}$'s - would be nice evidence for gluinos. Unfortunately, a reliable estimate of the cross-section is impossible because of uncertainties in the masses of the $\tilde{g}$ and $\tilde{q}$.

Onium decay into gluinos

The dominant decay of vector $Q\bar{Q}$ onium $^3S_1(1^{--})$ is supposed to be into 3g [39], whereas ^1S_o (0^{+-}), $^3P_o(0^{++})$ and $^3P_2(2^{++})$ are supposed to decay into 2g, and the $^3P_1(1^{++})$ is usually supposed to decay

into $(q\bar{q})g$, where the $(q\bar{q})$ are the lighter quarks[40]. Could gluinos disturb this picture ?

One could imagine decays such as

$^3S_1 \to \tilde{g}\tilde{g}g$	(cf Fig. 11a)	(6.6a)
^1S_o , ^3P_o , $^3P_2 \to \tilde{g}\tilde{g}$	(cf Fig. 11b)	(6.6b)
$^3P_1 \to g\tilde{g}\tilde{g}$	(cf Fig. 11c)	(6.6c)

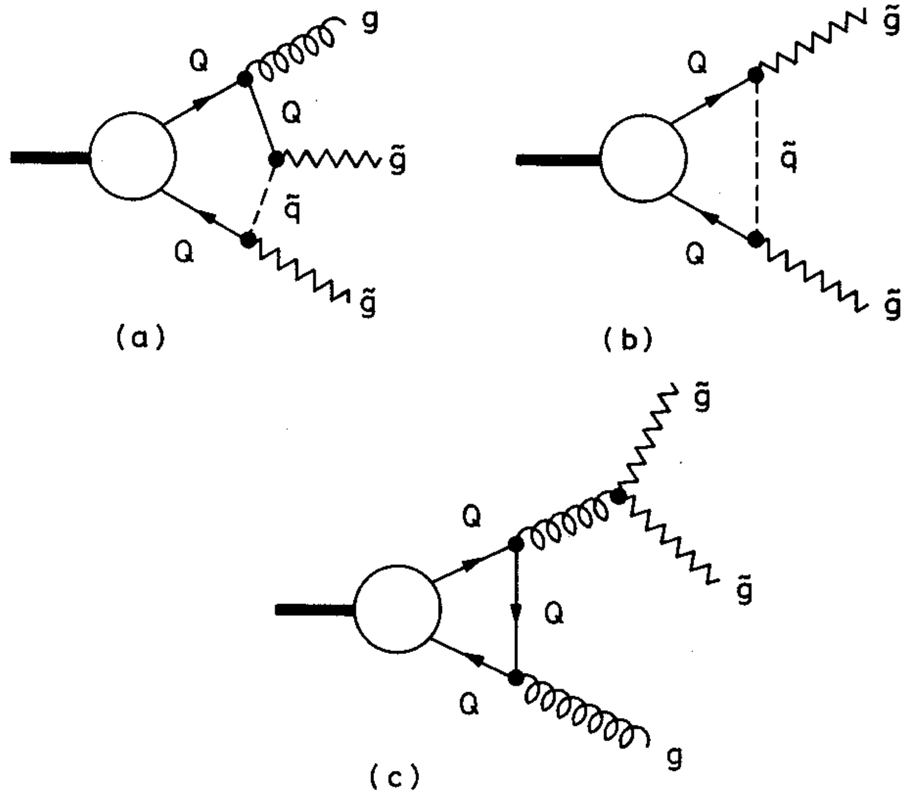

Fig. 11 - Onium decays involving gluinos

Of these processes, the rates for (6.6a) and (6.6b) are very uncertain because of the unknown mass of the scalar quarks $\tilde{q}$, but for the order of magnitude of masses to be expected it is difficult to see how they could be competitive with the conventional 3g or 2g decays. The situation is better for the process (6.6c) which should be competitive with the dominant[40] $g(q\bar{q})$ decay. The ratio between the two is just

$$\frac{\Gamma(^3P_1 \to g(\tilde{g}\tilde{g}))}{\Gamma(^3P_1 \to g(q\bar{q}))} = \frac{6}{f} \qquad (6.7)$$

where f is the number of light quark flavours: $m_q << m_Q$, and we have neglected any possible phase space suppression for the $g\tilde{g}\tilde{g}$ decay mode. If we consider a very heavy (LEP range) $Q\bar{Q}$ onium system, then we expect the lowest order "charmonium" predictions for their decays to work very well. Therefore we may expect that an extra $^3P_1$ decay mode as large as indicated in (6.7) should be readily distinguishable. This is particularly true since we expect gluino jets to exhibit novel features (large $\langle p_T \rangle$, missing neutral energy, stable hadrons ... ?) On the other hand we are not sufficiently confident in the "charmonium" picture for the $(c\bar{c})$ system to feel sure that present data[22] on $P_c(3.51)$ decays totally exclude the presence of a $\tilde{g}\tilde{g}g$ component In any case we had in section (2.7) the restriction from deep inelastic scattering that $m_{\tilde{g}} > 3$ GeV.

We conclude that LEP could improve this limit to

$$m_{\tilde{g}} \gtrsim m_Q \qquad (6.8)$$

where Q is any heavy quark whose onia are encountered.

7. Nuinos and Gravitinos

We expect nuinos to be manifested in many of the decays of supersymmetric particles. This is certainly the case when supersymmetric particles carry an extra R quantum number as considered by Farrar and Fayet[3], because nuinos are the lightest particles with $R \neq 0$. However, it is not clear that missing energy carried off by $\tilde{\nu}$'s is experimentally distinguishable from missing energy carried off by ν's. A good way to find evidence for $\tilde{\nu}$'s may be to look at the reactions that have been proposed for neutrino counting. In general these processes involve the coupling of the Z^0 to $\tilde{\nu}\tilde{\bar{\nu}}$, which we assume to be the same as that for Z^0 to $\nu\bar{\nu}$.

Onium $\to \tilde{\nu}\tilde{\bar{\nu}}$

The reaction

$$e^+e^- \to V' \to V + \pi\pi, \quad V \to \nu\bar{\nu} \qquad (7.1)$$

has been proposed[41] as a way to count unobserved neutrinos by looking at V decays into unvisible neutrals.

The ratio of branching ratios[24]

$$\frac{\Gamma(0 \to Z^0 \to \nu\bar{\nu})}{\Gamma(0 \to \gamma \to e^+e^-)} = \frac{G_F^2}{64\pi^2\alpha^2} \frac{m_V^4}{e_Q^2} (1-4|e_Q|\sin^2\theta_W)^2 N_\nu$$
$$\approx 0.2 \times 10^{-8} \times N_\nu \times M_V^4 \text{ for } e_Q = \frac{2}{3} \qquad (7.2)$$

indicates that this type of measurement using the J/ψ or Υ is not competitive with the $K \to \pi\nu\bar{\nu}$ limit of 6000, let alone the cosmological limit[25] of 4. In fact the quoted[11] limit

$$\frac{\Gamma(J/\psi \to \text{unobserved neutrals})}{\Gamma(J/\psi \to e^+e^-)} < \frac{1}{10} \to N_\nu \lesssim 5 \times 10^5 \qquad (7.3)$$

If we apply formula (7.2) to an onium in the LEP range then we find that

$$\frac{\Gamma(0^\gamma \to \nu\bar{\nu})}{\Gamma(0^\gamma \to e^+e^-)} = 0.2\% \times N_\nu \times \left(\frac{m_0}{30\ \text{GeV}}\right)^2 \gtrsim 0.2\% \times N_\nu \qquad (7.4)$$

so that if such an onium exists it should be possible to count all the ν's and $\tilde{\nu}$'s with masses $\lesssim m_0/2$.

<u>$Z^0 \to \tilde{\nu}\tilde{\bar{\nu}}$</u> - In the conventional Weinberg-Salam model with N generations of fermions,

$$\frac{\Gamma(Z^0 \to \nu\bar{\nu})}{\Gamma(Z^0 \to e^+e^-)} = 2N_\nu \text{ for } \sin^2\theta_W = 0.25 \qquad (7.5)$$

Studies[35] have shown that the total Z^0 width can be measured to an accuracy $O(\Gamma(Z^0 \to e^+e^-))$, so that a measurement of the total width of the Z^0 should be able to count ν's (and $\tilde{\nu}$'s) with masses $\lesssim \frac{1}{2}m_Z$ with a precision < 1.

One might therefore conclude that nuino counting at LEP should be no problem. The only complication that we can see arises if there is a gravitino[11]. Comparing equations (2.34) and (7.2) we see that the decay $0 \to \tilde{G}\tilde{\bar{\gamma}} + \tilde{\bar{G}}\tilde{\gamma}$ will dominate over the $\nu\bar{\nu}$ decays if

$$(\Delta m^2) < 10^4\ \text{GeV}^2 \qquad (7.6)$$

and so invalidate any attempt to count neutrinos with vector onia.

An even worse problem arises from $Z^0 \to \tilde{G}\tilde{\bar{\gamma}} + \tilde{\bar{G}}\tilde{\gamma}$ decays. One might expect in general (by analogy with equation (2.33)) a $Z^0 \tilde{G} \tilde{\gamma}$ coupling of the form

$$\frac{\gamma_\mu m_{Z^0}^2}{m_{\text{planck}}\, m_{\tilde{G}}} \qquad (7.7)$$

This then gives a decay rate

$$\frac{\Gamma(Z^0 \to \tilde{G}\tilde{\bar{\gamma}} + \tilde{\bar{G}}\tilde{\gamma})}{\Gamma(Z^0 \to e^+e^-)} \approx \left(\frac{m_Z^2}{\Delta m^2}\right)^2 \qquad (7.8)$$

where Δm^2 was introduced in equation (2.31) and is expected[11] (2.31) to be 0 (2500 GeV^2). If this expectation is true, the <u>dominant decay mode of the Z^0 will be to $(\tilde{G}\tilde{\bar{\gamma}}) + (\tilde{\bar{G}}\tilde{\gamma})$</u> ! In fact the Z^0 peak could even be completely washed out: the limit[11] (2.35) requires $\Delta m^2 > 15$ GeV^2 and hence only implies a very weak limit

$$\frac{\Gamma(Z^0 \to \tilde{G}\tilde{\bar{\gamma}} + \tilde{\bar{G}}\tilde{\gamma})}{\Gamma(Z^0 \to e^-e^-)} \lesssim 2500 \qquad (7.9)$$

One thing is clear: if a supersymmetry model with gravitinos of the type discussed by Fayet is correct, the Z^0 peak will differ greatly from the conventional expectations !

8. Conclusions

It is not possible to derive the mass scale of supersymmetric particles from first principles, although one might imagine that some should have a mass accessible with LEP. The mass region which can be explored with LEP for various

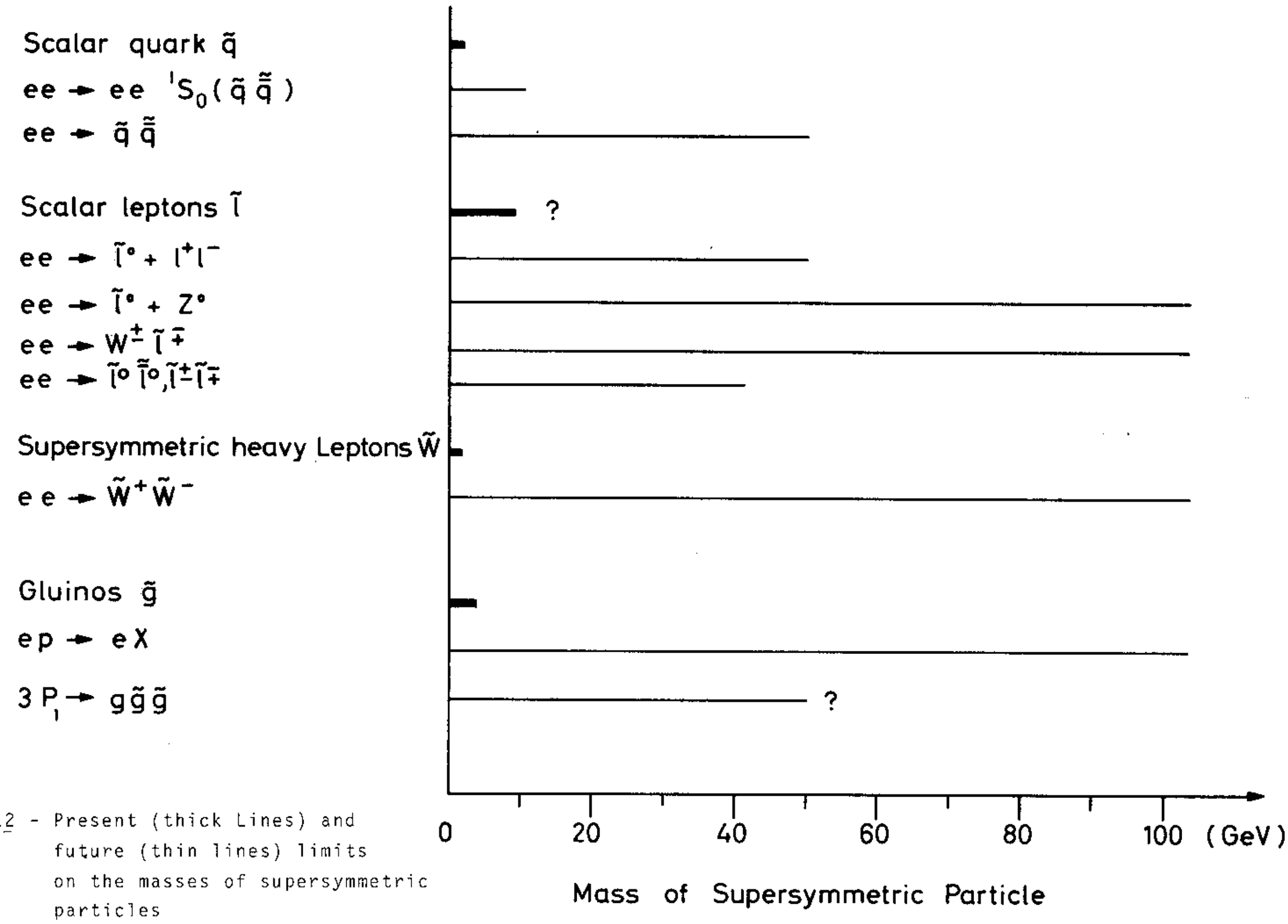

Fig. 12 - Present (thick Lines) and future (thin lines) limits on the masses of supersymmetric particles

kinds of supersymmetric particles is shown in Fig. 12 and compared with present limits. It is clear that LEP in general improves present limits by an order of magnitude of more to about 50 - 100 GeV. Because of the well defined production process in LEP the supersymmetric particles can easily be dis-entangled from more conventional particles. Needless to say, the search for such particles favours an e^+e^- ring with the largest energy possible.

References:

1. Yu.A.Gol'fand and E.P.Likhtman - JETP Letters 13, 323 (1971)
   D.V.Volkov and V.P.Akulov - Phys.Lett. 46B, 109 (1973)
   J.Wess and B.Zumino - Nucl.Phys. B70, 39 (1974)

2. For a view, see M.Grisaru and P.van Niewenhuizen-Deeper Pathways in High Energy Physics, Proc. Orbis Scientiae Coral Gables 1977 (Plenum Press, New York, 1977)

3. For review, see
   P.Fayet - Proc. Orbis Scientiae (Coral Gables, Florida, USA, Jan. 1978), New Frontiers in High Energy Physics, (Plenum, New York 1978), p. 413
   G.R.Farrar - Proc.Int.School of Subnuclear Physics (Erice, Italy, Aug. 1978)

4. G.Barbiellini et al., ECFA/LEP Specialized Study Group 9 DESY Preprint 79-17 (1979)

5. G.R.Farrar and P.Fayet - Phys.Lett. 76B 575 (1978) and Phys.Lett. 79B, 442 (1978)

6. G.R.Farrar and P.Fayet -"Searching for the spin-0 leptons of supersymmetry", Rutgers/Caltech preprint (1979)

7. CHEEP-an ep facility in the SPS - ed. J.Ellis, B.H.Wiik and K.Hübner - CERN Yellow Report 78-02 (1978)
   Colliding ring facilities for ep physics have also been actively considered at other laboratories, in particular at DESY at the present time

8. See J.Wess and B.Zumino, ref. 1

9. See R.Haag, J.T.Copuszanski and M.Sohnius Nucl.Phys. 88B, 257 (1975) and references therein

10. See Farrar and Fayet, ref. 3 which is largely based on the theoretical analysis of P.Fayet - Phys.Lett. 69B, 489 (1977)

references continued:

11. P.Fayet - Caltech preprint CALT-68-710 (1979)

12. S.Ferrara, D.Z.Freedman and P.van Nieuwenhuizen Phys.Rev. D13, 3214 (1976)
    S.Deser and B.Zumino - Phys.Lett. 62B, 335 (1976)
    P.Fayet - Phys.Lett. 70B, 461 (1977)

13. J.Kim, I.S.Ko and H.S.Song
    Preprint "A spin 3/2 consideration of heavy lepton" (1977)
    G.L.Kane and S.Raby - Stanford University preprint (1979)

14. P.Alibran et al., Phys.Lett. 74B, 134 (1978)
    T.Hansl et al., - Phys.Lett. 74B, 139 (1978)
    P.C.Bosetti et al., - Phys.Lett. 74B, 143 (1978)
    See also the talks by H.Wachsmuth and K.Winter at the 1979 Int.Symp. on Lepton and Photon Interactions at High Energies, FNAL, Batavia (Aug. 1979). There are still several puzzling phenomena in these experiments (the large cross sections, an apparent violation of ν_e - ν_μ universality, a copious source of low hadronic energy events, and events with missing p_T) which mean that one should not exclude non-charm contributions to the observed signals.

15. D.Cutts et al., Phys.Rev.Lett. 41, 363 (1978)
    R.Vidal et al., Fermilab preprint "A search for new massive particles" (1978)
    S.Wolfram - Phys.Lett. 82B, 55 (1979)
    C.B.Dover, T.K.Gaisser and G.Steigman Phys.Rev.Lett. 42, 117 (1979)
    R.Middleton, R.W.Zurmühle, J.Klein and R.V.Kollarits Phys.Rev.Lett. 43, 429 (1979)

16. Farrar and Fayet (ref. 6) get a limit of $3\frac{1}{2}$ GeV on scalar lepton masses from SPEAR data

references continued:

17. F.Wilczek - Phys.Rev.Lett. 39, 1304 (1977)

18. All three 1978 DORIS experiments looking at Υ decays had some sensitivity to $\Upsilon \to \gamma + X$ decays. In particular the PLUTO group looked at the inclusive γ spectrum, and can probably exclude the existence of $\Upsilon \to \tilde{\chi} + \gamma$ at a rate comparable to $\Upsilon \to e^+e^-$

19. The PLUTO group has no evidence for charged heavy leptons at PETRA, and quote a preliminary lower limit of 13 GeV (H.Meyer, private communication)

20. P.Fayet - Phys.Lett. 78B, 417 (1978)

21. H.L.Anderson, H.S.Matis and L.C.Myrianthopoulos Phys.Rev.Lett. 40, 1061 (1978)
P.C.Bosetti et al., Nucl.Phys. B142, 1 (1978)
J.G.H.De Groot et al., Zeit.Phys.C1, 143 (1979); Phys.Lett. 82B, 292, 456 (1979)
TASSO Collaboration, R.Brandelik et al., Phys.Lett. 86B, 243 (1979)
MARK J Collaboration, D.P.Barber et al., Phys.Rev.Lett. 43, 830 (1979)

22. For the latest results of the Crystal Ball Collaboration which suggest problems in understanding the $^3P(c\bar{c})$ hadronic decay widths in the context of lowest order QDC, see
E.Bloom - Talk at the 1979 Int. Symp. on Lepton and Photon Interactions at High Energies, FNAL, Batavia (Aug. 1979)

23. M.Glück and E.Reya - DESY preprint 79/13 (1979)

24. J.Ellis - Proc. 1978 SLAC Summer Institute on Particle Phys., ed. M.Zipf - SLAC-219 (1978)

references continued:

25. J.Yang, D. N.Schramm, G.Steigman and R.Rood-Astrophys.J. 227, 697 (1979)

26. G.Steigman, K.A.Olive and D.N.Schramm, Phys.Rev.Lett. 43, 239 (1979)

27. See for example
J.Ellis in Proc. ECFA/CERN Les Houches Summer Study on LEP - CERN Yellow Report 79-01 (1979)

28. H.Krasemann - ECFA/LEP Specialized Study Group 9 report SSG/9 (1979)

29. H.Krasemann and S.Ono - DESY preprint 79/09 (1979)

30. See particularly the talk on Novosibirsk results by V.Sidorov - Proc. 1979 Int. Symp. on Lepton and Photon Interactions at High Energies, FNAL, Batavia (Aug. 1979)

31. G.S.Abrams et al., Phys.Rev.Lett 43, 477 (1977)

32. J.D.Bjorken - Proc. 1976 SLAC Summer Institute on Particle Physics, ed. M.Zifp - SLAC-198 (1976)

33. J.Ellis, M.K.Gaillard and D.V.Nanopoulos, Nucl.Phys. B106, 292 (1976)

34. J.F.Donoghue and L.-F.Li
Carnegie-Melton University preprint (00-3066-113) (1978) and the discussion in G.Barbiellini et al., ref. 4

35. L.Camilleri et al., CERN Yellow Report 76-18 (1976), and the Les Houches proceedings, ref. 27.
The quantity we call h here was called g in these references

36. W.Alles, C.Boyer and A.J.Buras Nucl.Phys. B119, 125 (1977)

references continued:

37. J.Ellis, M.K.Gaillard and G.G.Ross
Nucl.Phys. B111, 253 (1976)

38. J.Ellis and I.Karliner
Nucl.Phys. B148, 141 (1979)

39. T.Appelquist and H.D.Politzer
Phys.Rev.Lett. 34, 43 (1975)

40. R.Barbieri, R.Gatto and E.Remiddi
Phys.Lett. 61B, 465 (1976)

41. J.Rich and D.R.Winn
Phys.Rev. D14, 1283 (1976) and J.Ellis, ref. 24.