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SUPERSYMMETRY VERSUS EXPERIMENT

Jorge C. ROMÃO

Centro de Física da Matéria Condensada/INIC Av. Gama Pinto, 2 1699 Lisboa Codex - Portugal

In these lectures we review the implications of supersymmetry at a low energy scale, the electroweak scale. The minimal supersymmetric standard model is described and its phenomenological consequences are discussed. The possibilities of discovery of supersymmetric particles at the accelerators presently being built or proposed are reviewed in detail.

1. INTRODUCTION

At a School where so much emphasis has been put on superstrings, it seems most appropriate to ask the question "Is there any experimental evidence in favour of the idea of supersymmetry (SUSY)?"

Although the answer up to present day energies is "No", it is worthwhile to review the status of SUSY versus experiment. The reason for this, is that there are some theoretical arguments indicating that SUSY might be of relevance for physics below the O(1TeV) energy scale and the accelerators being built or proposed for the near future can explore this energy region.

The theoretical arguments for SUSY are of many sorts. The most commonly invoked are:

i) Interrelates matter fields (leptons and quarks) with force fields, (gauge and/or higgs bosons)

ii) As local supersymmetry implies gravity it could provide a way to unify gravity with the other interactions.

iii) As supersymmetry and supergravity have fewer divergences than conventional field theories the hope is that it could provide a consistent (finite) quantum gravity theory.

iv) Supersymmetry can help understand the mass problem, in particular solve the naturalness problem (and in some models even the hierarchy problem) if supersymmetric particles have masses $\lesssim$O(1TeV).

As it is the last argument that makes SUSY particularly attractive for the experiments being proposed for the next decade or so, let us explain the idea in more detail.

As the standard model is not asymptotically free, at some energy scale Λ, the interactions must become strong indicating the appearance of new physics. Candidates for this scale are, for instance, $M_X \simeq O(10^{15}$ GeV) in GUTs or more fundamentally the Planck scale $M_P \simeq O(10^{19}$ GeV). This alone does not indicate that the new physics should be related to supersymmetry. But the so-called mass problem does. The only consistent way to give masses to the gauge bosons is through the Higgs mechanism involving at least one spin-0 particle, the higgs particle. Although the higgs mass is not fixed by the theory, a value much bigger than $< H^0 > \sim G_F^{-1/2} \sim 250$ GeV would imply that the higgs sector would be strongly coupled making it difficult to understand why we are seeing an apparently sucessful perturbation theory at low energies Now,the one-loop radiative corrections to the

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higgs mass would give

$$\delta m_H^2 = O(\frac{\alpha}{4\pi}) \Lambda^2 \quad , \qquad (1.1)$$

which would be too large if Λ is identified with Λ_{GUT} or Λ_{PLANCK}. SUSY cures this problem in the following way. If SUSY were exact, radiative corrections to the scalar masses squared would be absent because the contribution of fermion loops exactly cancels against the boson loops.

Therefore if SUSY is broken, as it must, we have

$$\delta m_H^2 = O(\frac{\alpha}{4\pi}) \mid m_B^2 - m_F^2 \mid \quad , \qquad (1.2)$$

showing that the SUSY-breaking mass squared difference between fermions and bosons acts as an effective cut off Λ^2. We conclude that SUSY provides a solution to the naturalness mass problem if

$$\mid m_B^2 - m_F^2 \mid \lesssim O(1TeV^2) \quad , \qquad (1.3)$$

because in this case we have $\delta m_H^2 = O(M_W^2)$ and not much bigger.

From the above argument we see that SUSY only solves the naturalness problem if the masses of the superpartners are less than O(1TeV). This is the main reason behind all the phenomenological interest in SUSY.

In the following sections we will give a brief review of the main aspects of the supersymmetric extension of the Standard Model, and describe how one can look for SUSY in the accelerators being built or proposed for the near future. Almost all the material is covered in the many excellent reviews that exist in the literature[1] . In most cases we will refer the review article instead of the original work. A very complete list of references can be found there.

2. THE SUPERSYMMETRIC STANDARD MODEL

In this section we will review the main properties of a supersymmetric extension of the Standard Model (SSM). As there are several possibilities we will stress the model independent properties instead of insisting upon the differences among them.

2.1 What is Supersymmetry (SUSY)?

Supersymmetry[2] is a symmetry that exchanges bosons with fermions, i.e.

$$Q_\alpha \mid \text{fermion}\rangle = \mid \text{boson}\rangle$$

$$Q_\alpha \mid \text{boson}\rangle = \mid \text{fermion}\rangle \quad , \qquad (2.1)$$

where Q_α are SUSY generators. These generators act as step operators for 1/2 unit of spin. In usual quantum field theories there are two Poincaré group invariants, $P^2=m^2$ and $W^2=W_\mu W^\mu$, where W^μ is the Pauli - Lubanski vector. These means that the irreducible representations of the Poincaré group are labelled by two numbers, the mass and the spin of the particle.

In supersymmetric quantum field theories the relevant group is the Super-Poincaré group which is obtained by extending the Poincaré group using the SUSY generators Q_α. One can then show[1] that

$$[Q_\alpha \ , P^2] = 0$$

$$[Q_\alpha \ , W^2] \neq 0 \quad , \qquad (2.2)$$

which implies that the irreducible representations, the so-called supermultiplets, are labelled just by one number, the mass. This means that particles in the same supermultiplet have the same mass and different spins (in intervals of 1/2 unit).

There could be a number N of generators Q_α^i , (i=1,...,N) of supersymmetric transformations carrying additional quantum numbers. However if $N \geq 2$ there would exist in the same multiplet the left-handed electron and a right-handed one which contradicts the $SU_L(2)$ assignments. Therefore at low energy SUSY should be broken to N=1. There is therefore no loss of generality in just considering the N=1 case.

Finally we would like to mention the connection of SUSY with gravity. The generators Q_α obbey the anticommutation relation

$$\{Q_\alpha , \bar{Q}_\beta\} = -2(\gamma^\mu)_{\alpha\beta} P_\mu , \qquad (2.3)$$

which implies that local SUSY is related to local translations, that is gravity.

2.2. The SUSY SU(3)xSU(2)xU(1) model .

The Lagrangian for the SSM is given by a sum of different terms :

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{leptons}} + \mathcal{L}_{\text{quarks}} + \mathcal{L}_{\text{higgs}} + \mathcal{L}_{\substack{\text{Yukawa}\\ \text{+higgs pot.}}} + \mathcal{L}_{\substack{\text{SUSY}\\ \text{breaking}}} . \qquad (2.4)$$

The easiest way to insure that the lagrangian is supersymmetric is to use the superfield construction[1]. Essentially there are two types of superfields, the real (or vector) superfield and the chiral superfield. Supersymmetric lagrangians are obtained by taking the D-component of vector superfields or the F-component of chiral superfields, which is a well known procedure[1]. Vector superfields describe the gauge supermultiplet consisting of a spin-1 gauge field plus a spin-1/2 majorana fermion, the gaugino. Chiral superfields describe matter supermultiplets consisting of a chiral fermion plus a spin-0 complex scalar, the scalar fermion. Both chiral and vector superfields also have the so-called auxiliary fields. These are fields that have no kinetic term in the lagrangian. Therefore they do not correspond to physical propagating degrees of freedom. Their equations of motion are algebraic and can be used to eliminate them from the lagrangian in favour of the physical fields.

In the following we will give the expressions for the various pieces in eq.(2.4) in terms of the superfields and of their component fields.

2.2.1. Gauge fields.

We need three vector superfields, $\mathbb{V}_i$ (i=1,2,3) for the three gauge groups $U_Y(1)$, $SU_L(2)$ and $SU_c(3)$, respectively. We then construct chiral superfields by taking superspace derivatives ,

$$\mathbb{W}_{i\alpha} = \bar{D}_R D_R \left(e^{+2g_i \mathbb{V}_i} D_{L\alpha} e^{-2g_i \mathbb{V}_i} \right) , \qquad (2.5)$$

where for $SU_L(2)$ and $SU_c(3)$ we use a matrix notation

$$\mathbb{V}_2 = \mathbb{V}_2^a \frac{\sigma_a}{2} \ ; \ \mathbb{V}_3 = \mathbb{V}_3^a \frac{\lambda_a}{2} , \qquad (2.6)$$

and σ_a , λ_a are, respectively, the Pauli and Gell-Mann matrices for those groups. The field content of these superfields is, in an obvious notation, $\mathbb{V}_1=(B_\mu, \tilde{B}, D_1)$, $\mathbb{V}_2=(W_\mu, \tilde{\omega}, D_2)$ and $\mathbb{V}_3=(A_\mu, \tilde{g}, D_3)$, where B_μ, W_μ and A_μ are the gauge fields, $\tilde{B}$, $\tilde{\omega}$ and $\tilde{g}$ the gauginos and D_1, D_2 and D_3 the auxiliary fields for $U_Y(1)$, $SU_L(2)$ and $SU_c(3)$ respectively.

The gauge field lagrangian is then

$$\mathcal{L}_{\text{gauge}} = \frac{1}{128\, g_3^2} \mathrm{Tr} \left[\mathbb{W}_3^T C\, \mathbb{W}_3 + \text{h.c.} \right]_F + \frac{1}{128\, g_2^2} \mathrm{Tr} \left[\mathbb{W}_2^T C\, \mathbb{W}_2 + \text{h.c.} \right]_F + \frac{1}{256\, g_1^2} \left[\mathbb{W}_1^T C\, \mathbb{W}_1 + \text{h.c.} \right]_F$$

$$= -\frac{1}{4} G^a{}_{\mu\nu} G^{a\mu\nu} - \frac{1}{4} F^a{}_{\mu\nu} F^{a\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} + \frac{i}{2} \bar{\tilde{g}}^a \gamma^\mu D_\mu^{ab} \tilde{g}^b + \frac{i}{2} \bar{\tilde{\omega}}^a \gamma^\mu D_\mu^{ab} \tilde{\omega}^b + \frac{i}{2} \bar{\tilde{B}} \gamma^\mu \partial_\mu \tilde{B} + \frac{1}{2} D_3^a D_3^a + \frac{1}{2} D_2^a D_2^a + \frac{1}{2} D_1 D_1 , \qquad (2.7)$$

where

$$G_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu - i g_3 [A_\mu , A_\nu]$$

$$F_{\mu\nu} = \partial_\mu W_\nu - \partial_\nu W_\mu - i g_2 [W_\mu , W_\nu]$$

$$B_{\mu\nu}=\partial_\mu B_\nu-\partial_\nu B_\mu \quad . \qquad (2.8)$$

Besides the usual couplings of the gauge fields in the standard model we have aditional couplings between the gauginos and the gauge fields.

2.2.2. Leptons.

We need the following chiral superfields

$$\mathbb{L}_i\ (1,2,-1/2)\ ;\ \boldsymbol{\ell}^c_i(1,1,1)\quad (i=1,2,3)\ , \qquad (2.9)$$

where we have indicated the quantum numbers under $SU_c(3)\times SU_L(2)\times U_Y(1)$, and the subscript $\underline{i}$ is a generation index. For convenience the right handed $SU_L(2)$ singlet is taken as the charge conjugated of a left handed singlet. In this way all the basic chiral superfields are left handed. To write the lagrangian in terms of the component fields we make the following definitions:

$$\mathbb{L}_i=(\tilde{L}_i,L_i)\quad ;\quad \boldsymbol{\ell}^c_i=(\tilde{\ell}^c_i,\ell^c_{iL})\quad , \qquad (2.10a)$$

where $\tilde{L}_i$ and L_i are the $SU_L(2)$ doublets of scalar leptons and leptons, respectively, which we write as follows:

$$\tilde{L}_i=\begin{pmatrix}\tilde{\nu}_i\\ \tilde{\ell}_i\end{pmatrix}\quad ;\quad L_i=\begin{pmatrix}\nu_{iL}\\ \ell_{iL}\end{pmatrix}\quad . \qquad (2.10b)$$

We also denote by z_A (y_A), A=1,...,n , the scalars (auxiliary fields) of the $\underline{n}$ chiral lepton superfields (n=9 for 3 generations).

With these definitions the minimally coupled kinetic lagrangian is :

$$\mathcal{L}_{\text{leptons}} =$$

$$=\left[\mathbb{L}_i^+ e^{(-2g_2\mathbb{V}_2-g_1\mathbb{V}_1)}\mathbb{L}_i+\boldsymbol{\ell}^{c*}_i e^{2g_1\mathbb{V}_1}\boldsymbol{\ell}^c_i\right]_D$$

$$=i\bar{L}_i\gamma^\mu D_\mu\gamma_L L_i+i\bar{\ell}_i\gamma^\mu D_\mu\gamma_R\ell_i$$

$$+\left(D_\mu\tilde{L}_i\right)^+D^\mu\tilde{L}_i+\left(D_\mu\tilde{\ell}^c_i\right)^*D^\mu\tilde{\ell}^c_i$$

$$+ig_2\sqrt{2}\left(\frac{\sigma_a}{2}\right)_{\alpha\beta}\left(\tilde{L}^*_{i\alpha}\bar{\tilde{\omega}}^a\gamma_L L_{i\beta}-\tilde{L}_{i\beta}\bar{L}_{i\alpha}\gamma_R\tilde{\omega}^a\right)$$

$$-ig_1\sqrt{2}\left(-\frac{1}{2}\tilde{L}^*_{i\alpha}\bar{\tilde{B}}\gamma_L L_{i\alpha}+\frac{1}{2}\tilde{L}_{i\alpha}\bar{L}_{i\alpha}\gamma_R\tilde{B}\right.$$

$$\left.+\tilde{\ell}^{c*}_i\bar{\ell}_i\gamma_L\tilde{B}-\tilde{\ell}^c_i\bar{\tilde{B}}\gamma_R\ell_i\right)$$

$$-\sum_{i=1}^{2}g_iD^a_i\sum_{(A,B)=1}^{n}z^{A*}\left(T^a_i\right)_A{}^Bz_B+\sum_{A=1}^{n}y^*_Ay_A, \qquad (2.11)$$

where we have defined the chiral projectors by

$$\gamma_L=\frac{1-\gamma_5}{2}\quad ;\quad \gamma_R=\frac{1+\gamma_5}{2}. \qquad (2.12)$$

From this lagrangian it is clear that the couplings of the leptons are the usual ones, plus a new coupling that involves the gauginos and the scalar leptons.

2.2.3. Quarks

For the quark sector one needs the chiral superfields

$$\mathbb{Q}_i\,(3,2,\tfrac{1}{6})\quad ;\quad \mathbb{d}^c_i(3,1,\tfrac{1}{3})\ ;\ \mathbb{u}^c_i(3,1,-\tfrac{2}{3}). \qquad (2.13)$$

In terms of the component fields we have

$$\mathbb{Q}_i=(\tilde{Q}_i,Q_i)\ ;\ \mathbb{d}^c_i=(\tilde{d}^c_i,d^c_{iL})\ ;\ \mathbb{u}^c_i=(\tilde{u}^c_i,u^c_{iL}) \qquad (2.14a)$$

where $\tilde{Q}_i$ and Q_i are the $SU_L(2)$ doublets

$$\tilde{Q}_i=\begin{pmatrix}\tilde{u}_i\\ \tilde{d}_i\end{pmatrix}\quad ;\quad Q_i=\begin{pmatrix}u_{iL}\\ d_{iL}\end{pmatrix}\quad . \qquad (2.14b)$$

The lagrangian is then:

$$\mathcal{L}_{\text{quarks}}=\left[\mathbb{Q}_i^+e^{(-2g_3\mathbb{V}_3-2g_2\mathbb{V}_2+\frac{1}{3}g_1\mathbb{V}_1)}\mathbb{Q}_i\right.$$

$$+\mathbb{u}^{c*}_i e^{(-2g_3\mathbb{V}_3-\frac{4}{3}g_1\mathbb{V}_1)}\mathbb{u}^c_i$$

$$\left.+\mathbb{d}^{c*}_i e^{(-2g_3\mathbb{V}_3+\frac{2}{3}g_1\mathbb{V}_1)}\mathbb{d}^c_i\right]_D$$

$$= i\bar{Q}_i\gamma^{\mu}D_{\mu}\gamma_L Q_i + i\bar{d}_i\gamma^{\mu}D_{\mu}\gamma_R d_i$$

$$+ i\bar{u}_i\gamma^{\mu}D_{\mu}\gamma_R u_i + \left(D_{\mu}\tilde{Q}_i\right)^{+}D^{\mu}\tilde{Q}_i$$

$$+ \left(D_{\mu}\tilde{d}_i^c\right)^{*}D^{\mu}\tilde{d}_i^c + \left(D_{\mu}\tilde{u}_i^c\right)^{*}D^{\mu}\tilde{u}_i^c$$

$$+ i g_3\sqrt{2}\left(\frac{\lambda_a}{2}\right)_{\alpha\beta}\left(\tilde{Q}^*_{i\,\alpha}\bar{\tilde{g}}^a\gamma_L Q_{i\beta} - \tilde{Q}_{i\beta}\bar{Q}_{i\alpha}\gamma_R\tilde{g}^a\right)$$

$$+ i g_2\sqrt{2}\left(\frac{\sigma_a}{2}\right)_{\alpha\beta}\left(\tilde{Q}^*_{i\,\alpha}\bar{\tilde{\omega}}^a\gamma_L Q_{i\beta} - \tilde{Q}_{i\beta}\bar{Q}_{i\alpha}\gamma_R\tilde{\omega}^a\right)$$

$$- i g_1\sqrt{2}\left(\frac{1}{6}\tilde{Q}^*_{i\,\alpha}\bar{\tilde{B}}\gamma_L Q_{i\alpha} - \frac{1}{6}\tilde{Q}_{i\alpha}\bar{Q}_{i\alpha}\gamma_R\tilde{B}\right.$$

$$- \frac{2}{3}\tilde{u}_i^{c*}\bar{u}_i\gamma_L\tilde{B} + \frac{2}{3}\tilde{u}_i^c\bar{\tilde{B}}\gamma_R u_i$$

$$\left.+\frac{1}{3}\tilde{d}_i^{c*}\bar{d}_i\gamma_L\tilde{B} - \frac{1}{3}\tilde{d}_i^c\bar{\tilde{B}}\gamma_R d_i\right)$$

$$-\sum_{i=1}^{3} g_i D_i^a \sum_{(A,B)=1}^{m} z^{A*}\left(T_i^a\right)_A{}^B z_B + \sum_{A=1}^{m} y^*_A y_A, \quad (2.15)$$

where $z_A(y_A)$, $A=1,\ldots m$ are the scalars (auxiliary fields) for the m quark chiral superfields.

This lagrangian has a very similar structure to $\mathcal{L}_{\text{leptons}}$ only differing in the hypercharge assignements for quarks and leptons.

2.2.4. Higgs

For a technical reason connected with the Yukawa mass terms of the quarks (see section 2.2.5) one needs at least two higgs doublets. In some models a singlet is also introduced. The chiral superfields are :

$$\mathbb{H}_1\ (1,2,-\tfrac{1}{2})\ ;\ \mathbb{H}_1\ (1,2,\tfrac{1}{2})\ ;\ \mathbb{s}\ (1,1,0). \quad (2.16)$$

In terms of the component fields we have

$$\mathbb{H}_1=(H_1,\tilde{H}_{1L})\ ;\ \mathbb{H}_2=(H_2,\tilde{H}_{2L})\ ;\ \mathbb{s}=(s,\tilde{s}_L) \quad (2.17a)$$

where $\tilde{H}_i$ and H_i are the $SU_L(2)$ doublets,

$$H_1=\begin{pmatrix} H_1^o \\ H_1^- \end{pmatrix}\ ;\quad \tilde{H}_{1L}=\begin{pmatrix} \tilde{H}_{1L}^o \\ \tilde{H}_{1L}^- \end{pmatrix},$$

$$H_2=\begin{pmatrix} H_2^+ \\ H_2^o \end{pmatrix}\ ;\quad \tilde{H}_{2L}=\begin{pmatrix} \tilde{H}_{2L}^+ \\ \tilde{H}_{2L}^o \end{pmatrix}. \quad (2.17b)$$

The Lagrangian is then :

$$\mathcal{L}_{\text{higgs}} = \Big[\mathbb{H}_1^+ e^{(-2g_2\mathbb{V}_2 - g_1\mathbb{V}_1)}\mathbb{H}_1$$

$$+\mathbb{H}_2^+ e^{(-2g_2\mathbb{V}_2 + g_1\mathbb{V}_1)}\mathbb{H}_2 + \mathbb{s}^*\mathbb{s}\Big]_D$$

$$= i\bar{\tilde{H}}_1\gamma^{\mu}D_{\mu}\gamma_L\tilde{H}_1 + i\bar{\tilde{H}}_2\gamma^{\mu}D_{\mu}\gamma_L\tilde{H}_2$$

$$+ i\bar{\tilde{s}}\gamma^{\mu}D_{\mu}\gamma_L\tilde{s} + \left(D_{\mu}H_1\right)^{+}D^{\mu}H_1$$

$$+ \left(D_{\mu}H_2\right)^{+}D^{\mu}H_2 + \partial_{\mu}s^*\,\partial^{\mu}s$$

$$+ i g_2\sqrt{2}\left(\frac{\sigma_a}{2}\right)_{\alpha\beta}\left(H^*_{1\,\alpha}\bar{\tilde{\omega}}^a\gamma_L\tilde{H}_{1\beta} - H_{1\beta}\bar{\tilde{H}}_{1\alpha}\gamma_R\tilde{\omega}^a\right.$$

$$\left.+ (H_1\leftrightarrow H_2)\right)$$

$$- i g_1\sqrt{2}\left(-\frac{1}{2}H^*_{1\,\alpha}\bar{\tilde{B}}\gamma_L\tilde{H}_{1\alpha} + \frac{1}{2}H_{1\alpha}\bar{\tilde{H}}_{1\alpha}\gamma_R\tilde{B}\right.$$

$$\left.- (H_1\leftrightarrow H_2)\right)$$

$$-\sum_{i=1}^{2} g_i D_i^a \sum_{(A,B)=1}^{p} z^{A*}\left(T_i^a\right)_A{}^B z_B + \sum_{A=1}^{p} y^*_A y_A, \quad (2.18)$$

where $z_A(y_A)$, $A=1,\ldots p$ are the scalars (auxiliary fields) for the p higgs chiral superfields.

2.2.5. Yukawa terms and Higgs potential.

In supersymmetric theories the Yukawa terms and higgs potential are obtained from a superpotential $\mathbb{W}$ through the relation

$$\mathcal{L}_{Y+HP} = \left[\ \mathbb{W} + h.c.\ \right]_F , \quad (2.19)$$

where $\mathbb{W}$ is a chiral superfield . $\mathbb{W}$ is a polynomial of dimension ≤3 in the chiral fields of the theory with the same chirality. It is this last requirement, that forces us to have at least two higgs doublets to be able to give masses to the up and down quarks. In the standard model the Yukawa terms are

$$\mathcal{L}_Y = g_i \bar{L}_i \gamma_R \ell_i H + g_{d_i} \bar{Q}_i \gamma_R d_i H + g_{u_i} \bar{Q}_i \gamma_R u_i \hat{H} + \text{h.c.} \quad , \tag{2.20}$$

where $\hat{H} = i\,\sigma_2\, H^+$. In this way we have

$$\langle H \rangle = \begin{pmatrix} 0 \\ v \end{pmatrix} \; ; \; \langle \hat{H} \rangle = \begin{pmatrix} v \\ 0 \end{pmatrix} \quad , \tag{2.21}$$

and with the same higgs doublet one can give masses to both the up and down quarks. This construction can not be used in the supersymmetric case because the superfields $\mathbb{H}_1$ and $(\mathbb{H}_1)^+$ have opposite chirality.

If we write the superpotential as

$$\mathbb{W} = \mathbb{W}_Y + \mathbb{W}_{HP} \quad , \tag{2.22}$$

then in the minimal model, the Yukawa part of the superpotential reads :

$$\mathbb{W}_Y = \sum_{i=1}^{3} \left(g_i\, \mathbb{H}_1 \mathbb{L}_i \ell_i^c + g_{u_i}\, \mathbb{Q}_i \mathbb{H}_2 u_i^c + g_{d_i}\, \mathbb{H}_1 \mathbb{Q}_i d_i^c \right) \quad , \tag{2.23}$$

where the $SU_L(2)$ doublets are contracted to form a singlet according to the rule

$$\mathbb{H}_1 \mathbb{L}_i = \mathbb{H}_{1\alpha}\, \varepsilon_{\alpha\beta}\, \mathbb{L}_{i\beta} \; . \tag{2.24}$$

For the higgs potential the minimal choice is

$$\mathbb{W}_{HP} = \mu\, \mathbb{H}_1 \mathbb{H}_2 \quad , \tag{2.25}$$

which breaks SU(2)x U(1) only at one loop level through radiative corrections. An alternative choice where the breaking occurs at tree level is

$$\mathbb{W}_{HP} = \beta\, \mathbb{S}\, \mathbb{H}_1 \mathbb{H}_2 - \varepsilon\, \mathbb{S} \quad . \tag{2.26}$$

The advantage of the first choice over the second is the smaller number of new fields that one has to introduce. It has also been claimed[3] that the vacuum obtained from eq. (2.19) not only breaks SU(2)x U(1) but also breaks the color and electric charges. Although this is a possibility that one has to take in account, we have shown[4] that there is a large region in parameter space where that does not occur.

When we compute the F-term corresponding to eq.(2.19) we get terms linear in the auxiliary fields D_3^a , D_2^a , D_1 and y_A . As these auxiliary fields appear quadratically in $\mathcal{L}_{gauge}$, $\mathcal{L}_{leptons}$, $\mathcal{L}_{quarks}$ and $\mathcal{L}_{higgs}$,their equations of motion are purely algebraic. We can solve for them and substitute back in the lagrangian. The result of this procedure can be stated as follows. Let us denote the N chiral fields of the theory by $\phi_A=(z_A,\psi_{AL},y_A)$, A=1,...N, where z_A are the scalars, ψ_{AL} the spin 1/2 fermions and y_A the auxiliary fields. In our previous notation N=n+m+p. Then the sum of the D-terms plus the F-terms is[1]:

$$\begin{aligned} \mathcal{L}_{D+F} &= \frac{1}{2} D_3^a D_3^a + \frac{1}{2} D_2^a D_2^a + \frac{1}{2} D_1 D_1 + \sum_{A=1}^{N} y^*_A y_A \\ &\quad - \sum_{i=1}^{3} g_i D_i^a \sum_{(A,B)=1}^{n} z^{A*} \left(T_i^a\right)_A{}^B z_B + [\, \mathbb{W} + \text{h.c.}\,]_F \\ &= -\frac{1}{2} D_3^a D_3^a - \frac{1}{2} D_2^a D_2^a - \frac{1}{2} D_1 D_1 \\ &\quad - \sum_{A=1}^{N} \left| \frac{\partial f}{\partial z_A} \right|^2 \\ &\quad - \frac{1}{2} \sum_{(A,B)=1}^{N} \left[\frac{\partial^2 f}{\partial z_A \partial z_B} \bar{\psi}_A \gamma_L \psi_B + \text{h.c.} \right] \end{aligned} \tag{2.27}$$

where

$$D_3^a = g_3 z^{A*} \left(T_3^a\right)_A{}^B z_B$$
$$D_2^a = g_2 z^{A*} \left(T_2^a\right)_A{}^B z_B$$
$$D_1 = g_1 z^{A*} \left(T_1\right)_A{}^B z_B \qquad (2.28)$$

and $\left(T_i^a\right)_A{}^B$ are the generators of the i^{th} gauge group in the representation to which the matter fields belong. The f-function in eq.(2.27) is obtained from the superpotential $\mathbb{W}$ by the substitution of the superfields by their scalar components. For instance for the minimal case, eq.(2.25), we have

$$f = g_i\, \varepsilon_{\alpha\beta} H_{1\alpha} \tilde{L}_{i\beta} \tilde{\ell}_i^c + g_{d_i}\, \varepsilon_{\alpha\beta} H_{1\alpha} \tilde{Q}_{i\beta} \tilde{d}_i^c + g_{u_i}\, \varepsilon_{\alpha\beta} \tilde{Q}_{i\alpha} H_{2\beta} \tilde{u}_i^c + \mu H_{1\alpha} \varepsilon_{\alpha\beta} H_{2\beta} \,. \qquad (2.29)$$

2.2.6. SUSY breaking.

As it has been said before, the fact that one has not yet observed any supersymmetric partner of the known particles forces SUSY to be broken in realistic models. This is the only way to account for particles in the same supermultiplet with different masses. This is the less well defined part of the models. In all models, the SUSY breaking Lagrangian contains tree level masses for the gauginos, plus some polynomial in the scalars of the theory. This function of the scalars, not only gives tree level masses for the scalar superpartners, but also modifies the higgs potential. Then, in general we have :

$$\mathscr{L}_{\substack{SUSY\\ breaking.}} = -\frac{1}{2} M\, \tilde{\omega}^{aT} C \tilde{\omega}^{a} - \frac{1}{2} M' \tilde{B}^T C \tilde{B} - m_0^2 \sum_{A=1}^{N} z^*_A z_A + \left(G(z) + h.c.\right) \,, \qquad (2.30)$$

where M and M′ are the tree level masses of the $SU_L(2) \times U_Y(1)$ gauginos, m_0 is a common mass scale for the scalars and G(z) is a function of the scalars of the theory. In the models[1] in which the SUSY breaking is induced via the coupling to N=1 supergravity, the function G(z) is specified, up to some arbitrary parameters. Before radiative corrections we have

$$G(z) = m_0 \left[(A-3)\, f(z) + \sum_{A=1}^{N} \frac{\partial f}{\partial z_A} z_A \right], \qquad (2.31)$$

where A is a numerical constant and f(z) is the function defined before,eq.(2.29).

In the models where $SU_L(2) \times U_Y(1)$ is broken via radiative corrections, the various couplings and masses are modified according to the renormalization group equations when we go down from some high energy scale to the electroweak scale. In this case, the very simple situation described above only occurs at the high scale. At the electroweak scale the mass parameters and couplings are modified[1].

2.2.7. Particle content.

In Table 2.1 we summarize the particle content of the minimal SUSY extension of the standard model. There is no uniform convention either for symbols or names. Here, we follow the conventions used by the LEP I and LEP II study groups [5,6]. The following comments are in order:

i) In the column of "known" particles we have included particles not yet discovered. Of these, only the top and one of the neutral higgs are predicted in the Standard Model. The other neutral higgs and the charged higgs, are a consequence of having two higgs doublets in the minimal SUSY extension of the Standard Model.

ii) For each helicity state of the fermions there is a scalar partner. For instance, there are two partners for the electron while there is only one for the neutrino.

iii) If SUSY were exact all particles belonging to the same supermultiplet would be degenerate in mass. As the supersymmetric partners of the known particles have not yet been discovered SUSY must be broken. This breaking makes, as we will see, supersymmetric partners with the same quantum numbers to mix. Therefore, for these particles the mass eigenstates are not the interaction eigenstates.

"Known" Particles	Supersymmetric Partners: Weak Inter. eigensates		Supersymmetric Partners: Mass eigenstates	
$\ell=(e,\mu,\tau)$	$\tilde{\ell}_L, \tilde{\ell}_R$	s.-lept.	$\tilde{\ell}_1, \tilde{\ell}_2$	s.-lept.
$\nu=(\nu_e \nu_\mu \nu_\tau)$	$\tilde{\nu}$	scalar-neutrino	$\tilde{\nu}$	scalar-neutrino
$q=(uds\ cbt)$	$\tilde{q}_L, \tilde{q}_R$	scalar-quarks	$\tilde{q}_1, \tilde{q}_2$	scalar-quarks
g	$\tilde{g}$	gluino	$\tilde{g}$	gluino
$W^\pm$	$\tilde{\omega}^\pm$	wino	$\tilde{W}$	wino
H_1^+	$\tilde{H}_1^+$	higgsino	$\tilde{W}_h$	heavy wino
H_2^+	$\tilde{H}_2^+$	higgsino		
γ	$\tilde{\gamma}$	photino	$\tilde{\gamma}$	photino
Z^0	$\tilde{Z}$	zino	$\tilde{Z}$	zino
H_1^0	$\tilde{H}_1^0$	higgsino	$\tilde{Z}_h$	heavy zino
H_2^0	$\tilde{H}_2^0$	higgsino	$\tilde{H}$	higgsino

TABLE 2.1
Particle content of the supersymmetric minimal extension of the standard model.

2.3. R-Parity.

In most supersymmetric theories one can introduce a multiplicatively conserved quantum number R, (R-parity) defined by

$$R=(-1)^{2J+3B+L}. \tag{2.32}$$

One can immediatly check that

$$R=\begin{cases}+1 \text{ for conventional particles}\\ -1 \text{ for supersymmetric partners}\end{cases}$$

The multiplicative conservation of the R-parity has very important phenomenological consequences:

i) Supersymmetric particles are always produced in pairs

$$\begin{aligned} e^+e^- &\to \tilde{\mu}^+\,\tilde{\mu}^- \\ &\to \tilde{e}^\pm\, e^\mp\, \tilde{\gamma} \\ &\dots \end{aligned} \tag{2.33}$$

ii) Every supersymmetric particle decays into another supersymmetric particle

$$\begin{aligned} \tilde{e} &\to e\,\tilde{\gamma} \\ \tilde{W} &\to q\,\bar{q}'\,\tilde{\gamma} \end{aligned} \tag{2.34}$$

iii) The lightest supersymmetric particle (LSP) is absolutely stable since it has no allowed decay mode.

The last property has important cosmologial consequences, which seem to exclude a strongly or electromagnetically interacting LSP. Therefore the candidates for LSP are the scalar neutrino $\tilde{\nu}$, the photino $\tilde{\gamma}$ (or the lightest of the neutralinos) and the gravitino $\tilde{G}$. The most favoured candidate is the photino and in the following we will assume that.

It can be shown that the cross section for $\tilde{\gamma}$ interacting with matter is not much bigger than the cross section for ν in matter[1]. This means that the normal signature for SUSY is missing $\not{p}$ carried away by the LSPs, that escape detection, e.g.

$$e^+e^- \to \tilde{\mu}^+ \tilde{\mu}^- \to \mu^+\mu^-\tilde{\gamma}\tilde{\gamma} \quad . \qquad (2.35)$$

2.4. Couplings.

Contrary to the masses, that are largely unknown (see section 2.5), the couplings of supersymmetric particles are precisely determined and model independent. Except for some numerical factors (spin Clebsch-Gordon coefficients) they are just the normal gauge couplings, g_s, g_2 and e, or Yukawa couplings. To have an idea of the vertices, one takes each vertex in the standard theory and converts it into several new ones by replacing particles by their supersymmetric partners in pairs (to conserve R-parity).

For instance

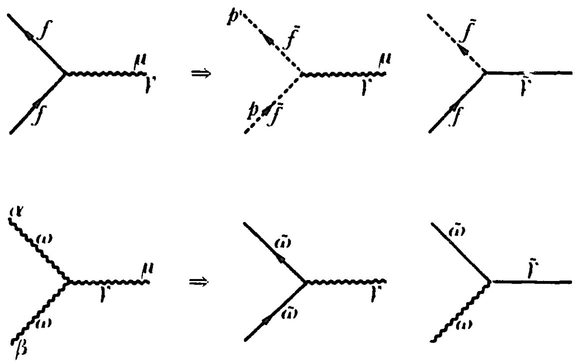

FIGURE 2.1
Examples of vertices with SUSY particles.

In these examples we have neglected the mixings between the winos and the higgsinos. As we will see below, the mass eigenstates are different from the weak interactions eigenstates which complicates matters .

2.5. SUSY breaking and mass matrices.

As it was said above if supersymmetry were an exact symmetry of Nature, particles and their superpartners would have the same mass. As no evidence of supersymmetric particles has been found so far, supersymmetry has to broken to lift this degeneracy.

There are many possibilities for the SUSY breaking, but in most current models the breaking is generated by supergravity via the so-called super-Higgs mechanism. In this analogue of the Higgs phenomena, the massless Goldstino (Goldstone fermion) of spin 1/2 is eaten by the massless spin 3/2 gravitino to give a massive gravitino. Just as quarks and leptons acquire mass when the gauge symmetry is spontaneously broken, the supersymmetric particles get extra contributions to their masses when SUSY is spontaneously broken. The precise contributions for the supersymmetric mass matrices are largely unknown and will be treated in the following as free parameters. The only constraint is that, if SUSY is going to play a role in explaining the mass problem, the splitting should obbey

$$| m_B^2 - m_F^2 | \lesssim O(1\text{TeV}^2) . \qquad (2.36)$$

2.5.1. Scalar-fermions mass matrix.

The generic form for the scalar fermions (leptons and quarks) mass matrix is

$$(\tilde{f}_L^* \tilde{f}_R^*) \begin{pmatrix} L^2 m_0^2 + m_f^2 & A m_0 m_f \\ A m_0 m_f & R^2 m_0^2 + m_f^2 \end{pmatrix} \begin{pmatrix} \tilde{f}_L \\ \tilde{f}_R \end{pmatrix} . \qquad (2.37)$$

The parameter m_0 is a typical scalar fermion mass (which must be $\gtrsim 20$ GeV), and m_f is the usual fermion mass. A, L and R are dimensionless model dependent parameters. For any given model it is straightforward to calculate the eigenvalues and eigenstates. There are two particular cases of interest. If L=R then the mixing is maximal, i.e.

$$\tilde{f}_{1,2} = \frac{\tilde{f}_R \pm \tilde{f}_L}{\sqrt{2}} \quad , \qquad (2.38)$$

and

$$m^2_{1,2} = m_o^2 + L^2 m_o^2 \pm A m_o m_f \quad , \qquad (2.39)$$

which indeed shows that the mass scale is set by m_o. On the other hand, if $L \neq R$ and $m_f << m_o$ (which is true for all cases except, possibly, for the top) then the mixing between $\tilde{f}_L$ and $\tilde{f}_R$ is negligible and we can consider $\tilde{f}_{L,R}$ as the mass eigenstates.

2.5.2. The charginos mass matrix.

We define the following Dirac spinors

$$\begin{cases} \tilde{W}^+ = -\tilde{\omega}^+ \\ \tilde{H}^+ = i\gamma_L \tilde{H}_2^+ - i\gamma_R \tilde{H}_1^{-c} \end{cases} \qquad (2.40)$$

In terms of these, the chargino mass matrix reads

$$(\bar{\tilde{W}}, \bar{\tilde{H}}) \begin{pmatrix} M & g v_2 \\ g v_1 & \mu \end{pmatrix} \begin{pmatrix} \tilde{W} \\ \tilde{H} \end{pmatrix}, \qquad (2.41)$$

where M is a SU(2) gaugino mass, μ comes from the higgs mixing term and $v_{1,2}$ are the vacuum expectation values of the two higgs scalars. Using $M_W^2 = 1/2\, g^2(v_1^2 + v_2^2)$ and taking $v_1/v_2=1$, which is a value favoured in most models ,we have

$$(\bar{\tilde{W}}, \bar{\tilde{H}}) \begin{pmatrix} M & M_W \\ M_W & \mu \end{pmatrix} \begin{pmatrix} \tilde{W} \\ \tilde{H} \end{pmatrix} , \qquad (2.42)$$

which has as eigenvalues

$$\tilde{M}_\pm^2 = \frac{1}{2}\Big[M^2+\mu^2+2M_W^2 \pm\sqrt{(M^2-\mu^2)^2 + 4M_W^2(M+\mu)^2}\,\Big] \qquad (2.43)$$

The eigenvectors are given by

$$\begin{cases} \tilde{W} = \cos\phi\, \tilde{\omega} - \sin\phi\, \tilde{H} \\ \tilde{W}_h = \sin\phi\, \tilde{\omega} + \cos\phi\, \tilde{H} \end{cases} \qquad (2.44)$$

where W and $\tilde{W}_h$ have masses $\tilde{M}_-$ and $\tilde{M}_+$, respectively, and ϕ is a mixing angle easily calculated. Many times is referred that the lightest chargino has a smaller mass than the W. This is of course a model dependent statement, but just to see how it could happen take $\mu = 0$. Then from (2.12) one easily gets

$$\tilde{M}_+ \tilde{M}_- = M_W^2 \qquad (2.45)$$

impliying that $\tilde{M}_- < M_W$.

2.5.3.The neutralinos mass matrix.

There are in the minimal case four neutralinos, which can be chosen as $\tilde{B}, \tilde{\omega}, \tilde{H}_1^o$ and $\tilde{H}_2^o$ or $\tilde{\gamma}, \tilde{Z}, \tilde{H}_1^o$ and $\tilde{H}_2^o$. In the Lagrangian the relevant mass terms are:

$$\mathcal{L} = -\frac{1}{2}\tilde{\omega}^{3T}C\tilde{\omega}^3 - \frac{1}{2} M'\tilde{B}^T C\tilde{B} - \mu\tilde{H}_1^{1T}C\tilde{H}_2^2 \qquad (2.46)$$

where M, M′ and μ are the same as before. Choosing a basis

$$\psi^o = \left(\tilde{B} \;\; -\tilde{\omega}^3 \;\; -i\gamma_5\tilde{H}_1^o \;\; -i\gamma_5\tilde{H}_2^o\right) \qquad (2.47)$$

we can write the lagrangian as $\mathcal{L} = \dots - \frac{1}{2}\bar{\psi}\mathcal{M}\psi$ where the mass matrix $\mathcal{M}$ is

$$\begin{pmatrix} M' & 0 & -\frac{M_W}{\sqrt{2}} t_W & \frac{M_W}{\sqrt{2}} t_W \\ 0 & M & \frac{M_W}{\sqrt{2}} & -\frac{M_W}{\sqrt{2}} \\ -\frac{M_W}{\sqrt{2}} t_W & \frac{M_W}{\sqrt{2}} & 0 & -\mu \\ \frac{M_W}{\sqrt{2}} t_W & -\frac{M_W}{2} & -\mu & 0 \end{pmatrix} \qquad (2.48)$$

In the last expression $t_W = \tan\Theta_W$, and for simplicity we have already considered the special case $v_1/v_2 = 1$. Without extra assumptions the diagonalization has to be handled numerically. One such assumption[5] is to put first zero mass for the gauginos. In this case we get a massless photino plus a zino, an heavy zino and another neutral particle which we call higgsino. The zino and heavy zinc masses are related by:

$$M_{\tilde{Z}} M_{\tilde{Z}_h} = M_Z^2 \quad . \tag{2.49}$$

Then, as a second step, we take $M, M' \neq 0$ but treat them as small parameters calculating the first order corrections to the above result.

Another simplifying assumption that it is made in some models is to relate M′ to M as in the minimal grand unified theory, that is,

$$\frac{M'}{M} = \frac{5}{3} \tan^2 \theta_W \quad . \tag{2.50}$$

In conclusion, one can say that due to absence of knowledge in the SUSY breaking parameters, the masses of the supersymmetric particles remain largely undetermined and should be kept as free parameters in any phenomelogical study. The major problem that arises is the mixing between the gauginos and higgsinos that modifies the couplings of the mass eigenstates.

3. WEAKLY INTERACTING SUPERSYMMETRIC PARTICLES

In this section we review for each weakly interacting SUSY particle the decay and production mechanisms. In almost all situations we will assume that the photino is the LSP.

3.1. Scalar leptons.

3.1.1. Decay.

If photinos are light the decay will be

$$\tilde{\ell}^{\pm} \to \ell^{\pm} \tilde{\gamma} \tag{3.1}$$

with an almost 100% branching ratio. The photinos will then go through the apparatus without being detected giving a missing $\not{p}$ signature.

3.1.2. Production.

There are several processes by which one can produce scalar leptons. Let us briefly review them:

3.1.2.1. Pair-production.

This is the typical situation at e^+e^- colliders. The processes is

$$e^+e^- \to \tilde{\ell}^+ \tilde{\ell}^- \quad . \tag{3.2}$$

For the case ($\tilde{\ell} \neq \tilde{e}$) only contribute the diagrams of Fig. 3.1 :

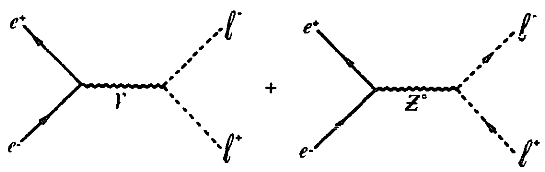

FIGURE 3.1
Diagrams contributing to $e^+e^- \to \tilde{\ell}^+ \tilde{\ell}^-$.

while for the scalar electron pair production we also have t-channel photino and zino exchanges as shown in Fig 3.2

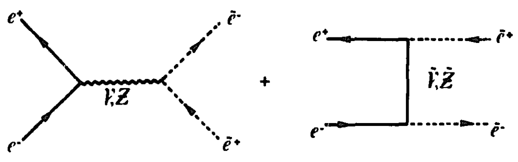

FIGURE 3.2
Diagrams contributing to $e^+e^- \to \tilde{e}^+ \tilde{e}^-$.

The presence of t-channel exchange makes the scalar electron more likely to go in the forward direction, giving a peak at small angles in the angular distribution of the electron coming from the scalar electron decay. These processes have been calculated[1,5]. They can set limits on masses $\lesssim E_{beam}$.

3.1.2.2. Associated production.

This a process for scalar electrons. The process is

$$e^+e^- \to e^+\tilde{e}^-\tilde{\gamma} \to e^+e^- \tilde{\gamma}\,\tilde{\gamma} \quad . \qquad (3.3)$$

The signal consists of an acoplanar e^+e^- pair and missing $\not{p}$. The diagrams and the expressions for the cross section can be found in the LEP I report[5].

3.1.2.3. Radiative processes.

This is an indirect way of getting limits on the masses of the scalar electrons. It has the advantage of being able to set higher limits on $m_{\tilde{e}}$. The process is

$$e^+e^- \to \gamma\,\tilde{\gamma}\,\tilde{\gamma} \quad , \qquad (3.4)$$

for which contribute the diagrams of Fig.3.3:

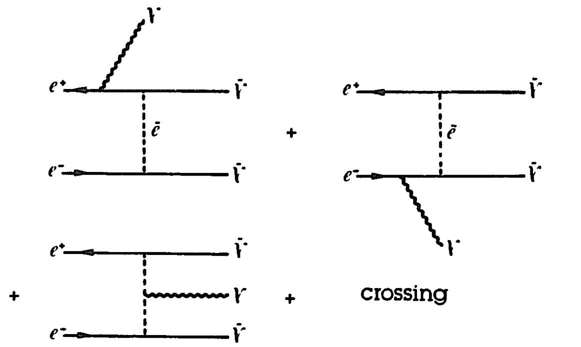

FIGURE 3.3

Diagrams contributing to $e^+e^- \to \gamma\,\tilde{\gamma}\,\tilde{\gamma}$.

This process[7,8] can set the highest limits on $m_{\tilde{e}}$ (assuming $m_{\tilde{\gamma}}=0$) or exclude a region on the $m_{\tilde{e}}$, $m_{\tilde{\gamma}}$ plane. In Fig. 3.4 we show the cross section for this process compared with the background, which is the neutrino counting experiment[9] $e^+e^- \to \gamma\,\nu\,\bar{\nu}$,

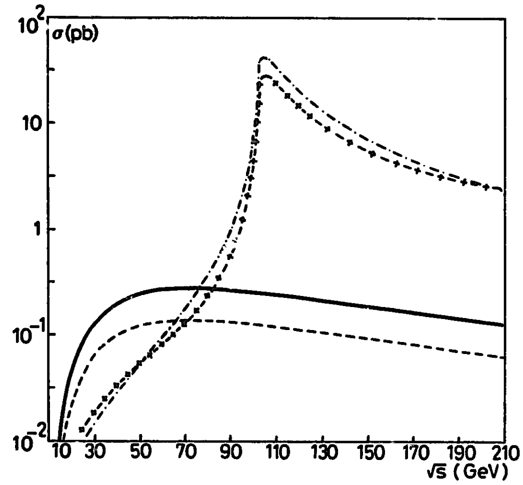

FIGURE 3.4

Photino cross section for unpolarized (dashed curve) and polarized (solid curve) beams. The dashed-crossed and dashed-dotted curves are the neutrino cross section for unpolarized and polarized beams, respectively. See Ref.9 for details.

In Fig. 3.5 we show the type of plot in the

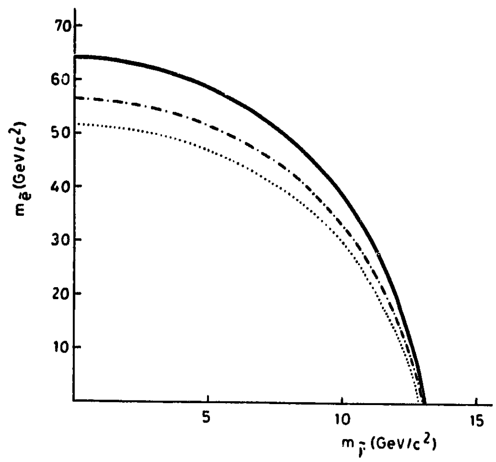

FIGURE 3.5

$m_{\tilde{e}}$, $m_{\tilde{\gamma}}$ plot. See Ref. 9 for details.

$m_{\tilde{e}}$, $m_{\tilde{\gamma}}$ plane that we can extract from a negative result.

3.1.2.3. Z^0 decay.

For SLC and LEP I this will be a very important process if $M_Z > 2m_{\tilde{e}}$. Then the Z^0 will decay in scalar-leptons pairs

$$Z^0 \to \tilde{e}^+ \tilde{e}^- \to e^+ e^- \tilde{\gamma} \tilde{\gamma} \quad , \tag{3.5}$$

with a rate[1]

$$\frac{\Gamma(Z^0 \to \tilde{e}^+ \tilde{e}^-)}{\Gamma(Z^0 \to e^+ e^-)} = \frac{1}{2}\left(1 - 4\, m_{\tilde{e}}^2 / M_Z^2\right)^{3/2} . \tag{3.6}$$

If $m_{\tilde{e}}$ is not too close to $M_Z/2$ this rate can be important.

3.1.2.4. Z^0 and $W^{\pm}$ decay at $p\overset{(-)}{p}$ colliders.

At $p\overset{(-)}{p}$ colliders colliders one could observe scalar-leptons as decaying products of Z^0 and $W^{\pm}$'s. The process would be

$$\begin{aligned} &W \to \tilde{e}\ \tilde{\nu} \\ &Z \to \tilde{e}^+ \tilde{e}^- \quad . \end{aligned} \tag{3.7}$$

Although these processes can be used to put limits on the scalar lepton masses normally they are not as good as the e^+e^- colliders.

3.1.3. Present limits.

The present limits on the masses of the scalar leptons are[6] :

$$\begin{cases} m_{\tilde{e}} \gtrsim 65 \text{ GeV} \\ m_{\tilde{\mu}} \gtrsim 20 \text{ GeV} \\ m_{\tilde{\tau}} \gtrsim 17 \text{ GeV} \end{cases} . \tag{3.8}$$

The limit on $m_{\tilde{e}}$ comes from $e^+e^- \to \gamma\, \tilde{\gamma} \tilde{\gamma}$. The other two come from pair production.

3.2. Scalar neutrinos.

3.2.1. Decay.

First there is the possibility that $\tilde{\nu}$ is stable, in which it would be the LSP. This can not be ruled out, although the prefered choice is the photino. If $\tilde{\nu}$ is not the LSP then the decays will depend very much on the mass. For instance, if there exists a chargino or neutralino lighter than $\tilde{\nu}$ we could have

$$\begin{aligned} &\tilde{\nu}_e \to \tilde{\omega}^+ e^- \\ &\tilde{\nu} \to \tilde{Z}\ \nu \quad . \end{aligned} \tag{3.9}$$

If this is not the case, we could have

$$\tilde{\nu} \to \nu \tilde{\gamma} \quad . \tag{3.10}$$

Since there is no tree level coupling this decay mode would proceed via the 1-loop diagrams of Fig. 3.6 :

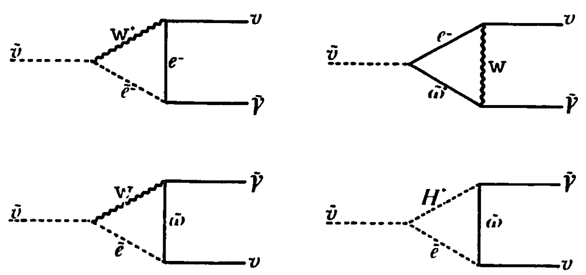

FIGURE 3.6

1-loop diagrams for $\tilde{\nu} \to \nu \tilde{\gamma}$.

3.2.2. Production.

Scalar neutrinos can be produced in the following processes:

3.2.2.1. Pair production.

The process is

$$e^+e^- \to \tilde{\nu}\ \tilde{\nu} \quad , \tag{3.11}$$

for which the diagrams are given in Fig. 3.7.

3.2.2.2. Associated Production.

Here the reaction is

$$e^+e^- \to e^+ \tilde{\omega}^- \tilde{\nu}_e \quad . \tag{3.12}$$

There are 12 contributing diagrams. They have been evaluated and can be found in the LEP I report[5].

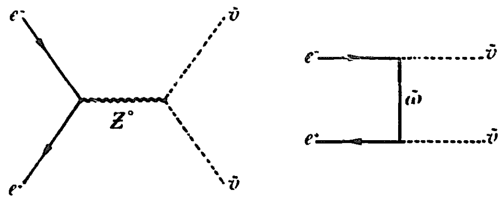

FIGURE 3.7
Diagrams for $e^+e^- \to \tilde\nu\ \tilde\nu$.

3.2.2.3. Radiative processes
The reaction here is

$$e^+e^- \to \gamma\tilde\nu\tilde\nu \quad , \tag{3.13}$$

with the diagrams shown in Fig 3.8:

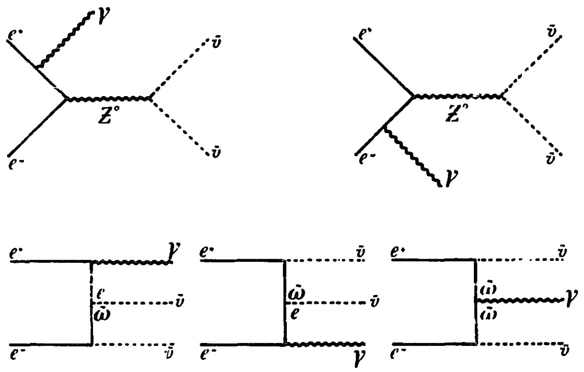

FIGURE 3.8
Diagrams for $e^+e^- \to \gamma\tilde\nu\tilde\nu$.

This is a process competing with $e^+e^- \to \gamma\tilde\gamma\tilde\gamma$ and $e^+e^- \to \gamma\ \nu\ \bar\nu$ (in general[8,9,10] $e^+e^- \to \gamma$ + missing neutrals).

3.2.2.4. Z^0 and $W^\pm$ decays.

The scalar neutrinos can be produced via the decays

$$Z^0 \to \tilde\nu\ \tilde\nu$$
$$W^\pm \to \tilde e^+\tilde\nu \quad . \tag{3.14}$$

If the decays are allowed ($M_Z > 2\,m_{\tilde\nu}$ and $M_Z > m_{\tilde e} + m_{\tilde\nu}$) we have[1]

$$\frac{\Gamma(Z^0\to\tilde\nu\tilde\nu)}{\Gamma(Z^0\to\nu\bar\nu)} = \frac{1}{2}\left(1-4\,m_{\tilde\nu}^2/M_Z^2\right)^{3/2} \tag{3.15}$$

and

$$\frac{\Gamma(W^+\to\tilde e^+\,\tilde\nu_e)}{\Gamma(W^+\to e^+\,\nu_e)} = \frac{1}{2}\left\{\left[\left(M_W^2-m_{\tilde\nu}^2-m_{\tilde e}^2\right)^2 - 4m_{\tilde\nu}^2 m_{\tilde e}^2\right] / M_W^4\right\} \tag{3.16}$$

which indicates that at LEP I we can have a large number of $\tilde\nu$'s.

3.2.3. Present limits.

Essentially no limits.

3.3. Charginos.

3.3.1. Decay.

There are many possibilities for the decays of the charginos depending on the mass

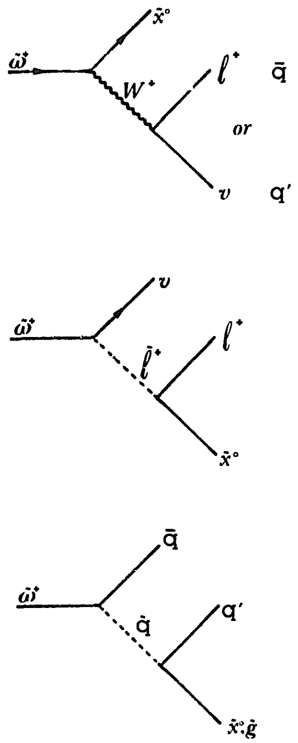

FIGURE 3.9
Examples of decay modes of the charginos.

and couplings. Some possibilities are given in Fig. 3.9 where $\tilde{\chi}^o$ is any of the neutralinos, depending on masses and mixings. The lifetimes are too short for the $\tilde{\omega}$'s to be observed directly.

3.3.2. Production.

3.3.2.1. Pair Production.

In e^+e^-colliders if kinematically allowed one can produce charginos (or the lightest of them) in the reaction[1,5]

$$e^+e^- \to \tilde{\omega}^+ \tilde{\omega}^- \quad , \tag{3.17}$$

which takes place via s-chamel exchange of γ and Z^0 and t-channel exchange of the scalar neutrino. At the $p\overset{(-)}{p}$ colliders the production can take place via the Drell-Yan mechanism. The relevant subprocesses are given in Fig. 3.10 :

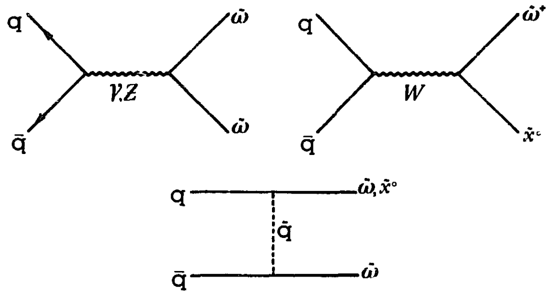

FIGURE 3.10
Sub-processes for chargino production.

3.3.2.2. Associated Production.

The process $e^+e^- \to e^+ \tilde{\omega}^- \tilde{\nu}_e$ discussed for the associated production of scalar neutrinos will also be relevant for the production of charginos.

3.3.2.3. Z^0 and $W^\pm$ decays.

At LEP I the chargino production via Z^0 decay could be important if kinematically allowed. At $p\overset{(-)}{p}$ colliders the production could happen through the W^+ decay

$$W^+ \to \tilde{\omega}^+ \tilde{\gamma} \quad , \tag{3.18}$$

if allowed.

3.3.3. Present limits.

Using the e^+e^- coliders the present limit is

$$m_{\tilde{\omega}} \gtrsim 23 \text{ GeV} \quad , \tag{3.19}$$

assuming $m_{\tilde{\nu}} = 0$.

3.4. Neutralinos.

In the minimal supersymmetric extension of the standard model there are four neutralinos (see section 2.5). The decay and production rates will depend very much on the actual values of the parameters that enter the mass matrix, Eq. (2.48).

3.4.1.Decay.

Let us denote by $\tilde{\chi}_i^o$ (i=1,...4) the four neutralino states ($\tilde{\gamma}$, $\tilde{Z}$, $\tilde{Z}_h$ and $\tilde{H}$). Then if kinematically allowed possible decay modes are

$$\tilde{\chi}_i^o \to \tilde{\omega}^\pm \tilde{\ell}^\pm \bar{\nu}$$

$$\tilde{\chi}_i^o \to \tilde{\chi}_k^o f \bar{f}$$

$$\tilde{\chi}_i^o \to \tilde{\chi}_k^o H^o \quad . \tag{3.20}$$

Only knowing the mass matrix is possible to evaluate the various branching ratios.

3.4.2. Production.

3.4.2.1. Pair Production.

At e^+e^- colliders they can be pair produced via the s-channel Z^0 intermediate state and t-channel $\tilde{e}$ exchange. In particular the process

$$e^+e^- \to \tilde{Z}\tilde{\gamma} \quad , \tag{3.21}$$

can be very important at LEP I if $m_{\tilde{e}}$ and $m_{\tilde{\omega}}$ are greater than $M_Z/2$.

At $p\overset{(-)}{p}$ colliders the production occurs via the Drell-Yan mechanism.

3.4.2.2. Radiative production.

The neutralinos can be produced in radiative processes. An example of these processes is $e^+e^- \to \gamma \tilde{\gamma} \tilde{\gamma}$ already discussed in section 3.1.

3.4.2.3. Z^o and $W^\pm$ decays.

If masses allow, the neutralinos can be produced in the following reactions

$$Z^o \to \tilde{\chi}_i^o \tilde{\chi}_k^o$$
$$W^\pm \to \tilde{\omega}^\pm \tilde{\chi}_i^o \quad . \qquad (3.22)$$

3.4.3. Present limits.

The limits are very dependent on the masses. If we take $m_{\tilde{\gamma}} = 10$ GeV and $m_{\tilde{e}} > 40$ GeV we set[1] $m_{\tilde{Z}} \gtrsim 35$ GeV.

4. STRONGLY INTERACTING SUSY PARTICLES.

In this section we review the supersymmetric particles that have strong interactions.

4.1. Scalar quarks.

4.1.1.Decay.

The decay channels will depend on the masses. Normally it is assumed that winos and zinos are heavier than the scalar quarks. In this case the allowed channels are

$$\tilde{q} \to q\, \tilde{\gamma} \qquad (4.1a)$$
$$\tilde{q} \to q\, \tilde{g} \quad . \qquad (4.1b)$$

If $m_{\tilde{g}} < m_{\tilde{q}}$ one expects $\tilde{q} \to q\, \tilde{g}$ to dominate. On the contrary is $m_{\tilde{g}} > m_{\tilde{q}}$ then $\tilde{q} \to q\, \tilde{\gamma}$ will the be prefered decay mode. Even if not dominant the mode (4.1a) gives a much cleaner signature.

4.1.2. Production.

4.1.2.1.e^+e^-colliders.

At e^+e^- colliders the scalar quark can be pair produced in the process

$$e^+e^- \to \tilde{q}\, \bar{\tilde{q}} \quad , \qquad (4.2)$$

for which we have the s-channel contributions of the photon and of the Z^o. The signature is two acolinear and acoplanar jets plus $\not{p}$.

4.1.2.2. $p\overset{(-)}{p}$ colliders.

At hadron colliders scalar quarks are produced by the Drell-Yan mechanism using the following sub-processes:

$$gg \to \tilde{q}\, \tilde{q}$$
$$q\bar{q} \to \tilde{q}\, \bar{\tilde{q}}$$
$$qq \to \tilde{q}\, \tilde{q}$$
$$qg \to \tilde{q}\, \tilde{g}$$
$$qg \to \tilde{q}\, \tilde{\gamma} \quad . \qquad (4.3)$$

The diagrams for these processes are easily found and computed[1,3,5].

4.1.3. Present limits.

From e^+e^- collisons one has[6]

$$m_{\tilde{q}} \geqslant 21 \text{ GeV} \quad . \qquad (4.4)$$

Stronger limits can be obtained from the $p\bar{p}$ colliders (Sp$\bar{p}$S and Tevatron) and will be given at the same time as those for the gluinos (see section 4.2).

4.2. Gluinos.

4.2.1. Decay.

There are essentially two channels for gluino decay. They are

$$\tilde{g} \to q\, \bar{q}\, \tilde{\gamma} \qquad (4.5a)$$
$$\tilde{g} \to g\, \tilde{\gamma} \quad . \qquad (4.5b)$$

The first process although it is a 3-body decay is expected to be the dominant decay mode. This is because the decay (4.5b) is a C-violating decay, that only occurs at one-loop level if $\tilde{q}_L$ and $\tilde{q}_R$ are not degenerate. If one takes the photino massless and for one flavour of massless quarks one has[1]

$$\frac{\Gamma(\tilde{g} \to g\tilde{\gamma})}{\Gamma(\tilde{g} \to q\bar{q}\tilde{\gamma})} = \frac{3\alpha_s}{4\pi} \frac{\left(\tilde{m}_R^2 - \tilde{m}_L^2\right)^2}{\tilde{m}_R^4 + \tilde{m}_L^4} \quad , \qquad (4.6)$$

which confirms the above remarks.

4.2.2. Production.

Gluinos will be produced in the hadron-hadron colliders through the processes

$$p + \overset{(-)}{p} \to \tilde{g} + \tilde{q} + X \qquad (4.7a)$$

$$p + \overset{(-)}{p} \to \tilde{g} + \tilde{q} + X \qquad (4.7b)$$

$$p + \overset{(-)}{p} \to \tilde{g} + \tilde{\gamma} + X \quad . \qquad (4.7c)$$

These are Drell-Yan processes.

4.2.3. Present limits.

The best limits on scalar-quark and gluino masses come from the $p\bar{p}$ colliders, $Sp\bar{p}S$ at CERN and Tevatron at Fermilab. They are now[11] $m_{\tilde{g}}$, $m_{\tilde{q}} > 80$ GeV.

5. FUTURE PROSPECTS.

In the previous sections we have described the minimal supersymmetric extension of the Standard Model. The emphasis was put in describing the main production and decay mechanisms at the various accelerators. Limits on the masses of supersymmetric particles were also indicated. Here we want to indicate the limits that will be obtained at the accelerators presently being built and for those that are now proposed.

First, it will be useful to have an idea about those machines. In table 5.1 we summarize the relevant information. For each accelerator we indicate the center of mass energy and the proposed luminosity. We should mention, that there are also a number of lower energy machines (PEP, PETRA,...) that we have not include in our table, because they do not have enough center of mass energy to be relevant to the discovery of supersymmetric particles.

The limits that can be obtained at the accelerators shown in table 5.1 have been investigated by several study groups[5,6,12].

Acellerator	Type	$\sqrt{s}$	$\mathcal{L}$ ($cm^{-2}s^{-1}$)
SPPS	$p\bar{p}$	630 GeV	$5x10^{30}$
TRISTAN	e^+e^-	60 GeV	$8x10^{31}$
TEVATRON	$p\bar{p}$	1.8 TeV	$5x10^{30}$
LEP I	e^+e^-	100 GeV	10^{30}
HERA	ep	314 GeV	$2x10^{31}$
LEP II	e^+e^-	200 GeV	$5x10^{31}$
LHC	pp	17 TeV	10^{33}
SSC	pp	40 TeV	10^{33}-10^{34}
CLIC	e^+e^-	2 TeV	$4x10^{33}$

Table 5.1
Present and future accelerators.

The conclusions of these groups are compiled in table 5.2, where they are compared with the present day limits. We see that the prospects are very good to discover SUSY, if supersymmetric particles have masses up to O(1 TeV).

We can summarize our conclusions as follows:

i) Despite all the theoretical motivations, there is still NO experimental evidence for supersymmetry in Nature.

ii) The theoretical models, described in some detail in these lectures need input from the experiments to fix the parameters. Only then they can be predictive.

	TODAY	TEVA-TRON	LEP-I SLC	LEP-II	HERA	LHC	SSC	CLIC
$M_{\tilde{e}}$	67	-	70	90	-	300	300	850
$M_{\tilde{\mu}}$	20	-	50	85	-	300	300	850
$M_{\tilde{\tau}}$	17	-	45	75	-	-	-	850
$M_{\tilde{\omega}}$	23	-	60	80	-	450	450	850
$M_{\tilde{z}}$	36	-	60	90	-	-	-	-
$M_{\tilde{q}}, M_{\tilde{g}}$	80	200	45	85	-	1000	1500	850
$M_{\tilde{e}}+M_{\tilde{q}}$	-	-	-	-	150	-	-	-
$M_{\tilde{\nu}}+M_{\tilde{q}}$	-	-	-	-	150	-	-	-

TABLE 5.2
Discovery limits at the various accelerators

iii) In the near future, with the new accelerators, we will be able to test if the idea of supersymmetry is relevant to Physics at a scale $\lesssim$ O (1 TeV).

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