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SUPERSTRING INSPIRED PHENOMENOLOGY

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ABSTRACT

Recent progress in superstring model building is reviewed with an emphasis on the general features of the models obtained. The problems associated with supersymmetry breaking and intermediate gauge symmetry breaking ($M_W < M_I < M_{GUT}$) are described. Finally, the phenomenology of these models is summarized, with a discussion of the role that new experimental results could play to help clearing up the above difficulties.

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Recent progress in superstring model building is reviewed with an emphasis on the general features of the models obtained. The problems associated with supersymmetry breaking and intermediate gauge symmetry breaking ($M_W < M_I < M_{GUT}$) are described. Finally, the phenomenology of these models is summarized, with a discussion of the role that new experimental results could play to help clearing up the above difficulties.

1. INTRODUCTION

Some progress has been made recently in the study of the effective particle physics models obtained from superstring theories : computation of Yukawa couplings, determination of the symmetries of the theory ... But this allows to specify the theory only at the Planck scale : there is basically only one scale in superstring models[1] which is fixed to be of order M_{Pl} because strings are interpreted as a theory of gravity. Even when string models are formulated in more than 4 dimensions, the compactification process that leads to 4-dimensional fields occurs at a scale M_{comp} which remains close to M_{Pl}. Indeed at this scale M_{comp} all 4-dimensional gauge interactions have, as a rule, the same coupling[2], which likens superstring unification to a standard grand unification at a scale $M_{GUT} \approx M_{comp}$. Hence

$$M_{string} \approx M_{Pl} \approx M_{comp} \approx M_{GUT} \tag{1}$$

where $\approx$ means equality up to say 2 orders of magnitude.

For many superstring models, it is now possible (notwithstanding the technical difficulties) to determine the basic couplings between the massless zero modes of the string (after compactification to four dimensions) at the common scale given by (1). As is well-known, the difficulty at this stage arises from the large number of models that one obtains. In Section 2, I will try to review where we stand and what are the general features of such models.

There is of course a long way from the region described by (1) down to the region below the TeV scale which can be tested experimentally. When bringing the models down to these energies, two interesting phenomena can occur :

(i) some further breaking of the gauge symmetry by the Higgs mechanism.

(ii) if we started with a supersymmetric model, the breaking of supersymmetry.

There are several good reasons to start with a supersymmetric theory :

- at the string level, supersymmetric theories are almost singled out when one requires the absence of tachyons and a zero cosmological constant[3].
- supersymmetry is the only known way to give a technical solution to the hierarchy problem : why should the low energy effective theory be stable under radiative corrections when very heavy fields (with mass of order M_{comp} or M_{Pl}) are present ?
- it is much easier to obtain the right prediction for $\sin^2\theta_W$ in a supersymmetric grand unified theory.

Because in phenomenology we are only interested in the massless modes of the string (massive modes have masses of order M_{str} or M_{comp}), supersymmetry breaking is the only way we have to generate a non-zero mass for these fields. Therefore, the determination of a mechanism for supersymmetry breaking is a key issue for phenomenological studies. An outstanding constraint on such a mechanism is that it should yield the typical scale of $SU(2) \times U(1)$ breaking, in a theory where the only scale available is given by (1). This is where most candidates for breaking supersymmetry fail.

Indeed, the fact that only supersymmetry breaking can generate non-zero masses for the fields of the effective theory, connects it closely with gauge symmetry breaking: once non-zero squared masses have been generated for Higgs fields at the supersymmetry breaking scale, radiative corrections turn them into scale-dependent–running–masses (their evolution is summarized in the renormalization group equations). For some of the fields, they become negative at a smaller scale, which induces a non-zero vacuum expectation value (vev) for the corresponding field and a breaking of the gauge symmetry if it is a gauge nonsinglet. Hence points (i) and (ii) above are strongly connected.

These two complementary aspects of superstring models in the range $M_W < \mu < M_{Pl}$ will be reviewed in Section 3. Finally, in Section 4, I will describe the phenomenology of these models at existing and future accelerators by putting the emphasis on how new experimental results could give precious indications on the issues discussed in Sections 2 and 3.

2. SUPERSTRING MODELS AT SCALES OF ORDER M_{Pl}

All the realistic models which have been proposed are within the framework of the heterotic string theory[4)*)]. I will therefore start with a brief and oversimplified account[6)] of the basic concepts of this theory.

It is well-known that the generic bosonic closed string is consistent only in 26 dimensions whereas the addition of fermionic degrees of freedom brings the number of critical dimensions down to 10. On the other hand, for a closed string, left moving and right moving oscillations are completely independent (at least in some gauges). The idea of the heterotic string is to describe left movers with 26 bosonic coordinates and right movers with 10 bosonic <u>and</u> fermionic coordinates. Of course, a subset of each of these two sets corresponds to the D dimensions of space-time in which the string propagates. The rest, that is 26-D and 10-D, corresponds to internal degrees of freedom. To see how this works, take for example one of the 26-D coordinates X and assume periodic boundary conditions : $X \equiv X + L$. As usual, the corresponding momentum is quantized: $P = 2\pi k/L$ and one generates this way a quantum number k, which for some values of L, can be interpreted as a gauge quantum number. With the 26-D such coordinates, one can therefore generate at most 26-D gauge quantum numbers that is a gauge group of rank[**)] 26-D. The heterotic string theory was originally proposed with D=10 and we will first review this special case, where most of the important concepts can be discussed.

2.1. Compactification of the 10-Dimensional Heterotic String

Taking D=10, we expect a gauge group of rank 16. The lattice described by the 16 coordinates obeying periodic boundary conditions must have some special properties (connected with modular invariance) to yield a self-consistent string theory. This, together with the requirement of supersymmetry, singles out the gauge groups SO(32) or

*) No model based on type II superstrings has been found which accomodates the particle spectrum of the standard model[5)].

**) The rank of a group is, loosely speaking, the maximum number of quantum numbers that one can define.

$E_8 \times E_8$. The former is not phenomenologically interesting because it does not yield chiral fermions. Hence, our starting point is a 10-dimensional theory with :

- N=1 supersymmetry,
- a gauge symmetry $E_8 \times E_8$;
- the field content is very simple : the graviton and its companion fields necessary for supersymmetry (gravity supermultiplet) ; the gauge fields in the adjoint representation of $E_8 \times E_8$, ($\underline{248}$, $\underline{1}$) + ($\underline{1}$, $\underline{248}$), and their supersymmetric partners, the gauginos (gauge supermultiplet).

The next step is compactification : we interpret the 10 space-time dimensions X^M, M=1...10 as four dimensions describing our ordinary space-time (X^μ, μ=1...4) and 6 dimensions living on a compact manifold K (x^k, k=5...10). The point is to determine the manifolds K which yield N=1 supersymmetry in four dimensions and a realistic four-dimensional particle spectrum (in particular non-singlet scalar fields which can play the role of Higgs fields).

Calabi-Yau Models[7)]

These were the first models to be proposed that satisfy the above constraints. For this reason, their phenomenological aspects have been more extensively studied than other candidates introduced later. Their key property is that the gauge symmetry is $E_6 \times E_8'$ or some of its subgroups ("E_6-based models"). It is important to have an idea how this E_6 gauge symmetry which plays an important role in superstring-inspired phenomenology, arises.

To get a hint, let us consider a gauge field of E_8 in 10 dimensions : $A^a{}_M$ where a is a gauge index (a=1...10)*). Accordingly, we can decompose it into $A^a{}_\mu$, μ=1...4, and $A^a{}_k$ k=5...10. Because of its μ index, $A^a{}_\mu$ is a 4-dimensional vector field (under Lorentz transformations) whereas $A^a{}_k$ is a 4-dimensional scalar. Conversely, $A^a{}_k$ is a vector on K whereas $A^a{}_\mu$ is a scalar on K. Now take a closed curve in K and parallel transport a vector (such as $A^a{}_k$) around this loop. If K is curved, the vector will have undergone a rotation when we return to the starting point. A classical example is the case of the sphere which is described in Fig. 1. In this example, one can generate all the rotations in the tangent plane. Hence the group of such transformations, called the holonomy group H, is precisely SO(2). For a 6-dimensional space such as K, H is generally SO(6) = SU(4).

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*) We have singled out here one of the two E_8. From now on, the other one will be referred to as E_8' and will occasionally be dropped.

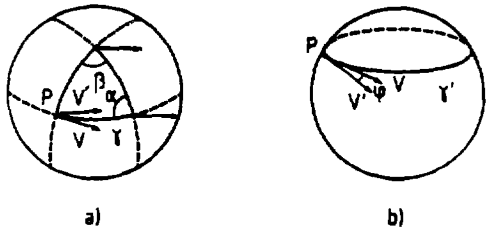

Fig. 1 : Parallel transport on the sphere S_2 along a closed curve
a) a triangle made of 3 great circles b) a small circle.

Let us now try to devise some simple compactification prescriptions:
i) one could try to require invariances under rotations in H (the physics in the effective 4- dimensional theory should not depend on whether we parallel transport the fields in the compact manifold). Then $A^a{}_\mu$ is present in the effective theory but not $A^a{}_k$ and we do not find any gauge non-singlet scalars in the effective theory.

Also, we started with N=1 supersymmetry in 10 dimensions, that is one supersymmetry charge Q_A, A=1...8 (Q_A, a 10-dimensional Majorana-Weyl spinor, has 8 degrees of freedom), which satisfies the usual supersymmetry algebra :

$$\{Q_A, Q_B\} = 2(\Gamma^M)_{AB} P_M \tag{2}$$

where P_M is the 10-dimensional momentum and Γ^M the 10-dimensional gamma-matrices. In 4 dimensions, the supersymmetry charge Q_α is a Majorana (i.e. real) spinor and has two degrees of freedom. Hence one supersymmetry in 10 dimensions yields in general through compactification N = 8/2 = 4 supersymmetries in 4 dimensions.

This is not tolerable because it is not possible to implement chiral fermions in N>1 supersymmetric theories. However, under the SO(6) = SU(4) group that we just discussed, Q_A transforms as a $\underline{4}$. In the special case where H = SU(3), then $Q_A = \underline{1} + \underline{3}$ and our prescription selects N=1 supersymmetry in 4 dimensions (the part of Q_A singlet under SU(3)). Therefore requiring N=1 supersymmetry in 4 dimensions singles out *(Kähler) manifolds with H = SU(3) holonomy, the Calabi-Yau manifolds.*

ii) one of the problems of the prescription *i)* is that it ignores the gauge degrees of freedom (a=1...248). To remedy this, take $H = SU(3)$ and identify H with part of the gauge symmetry. Indeed pick up $E_6 \times SU(3)_{YM} \subset E_8$; under this group

$$\underline{248} = (\underline{78}, \underline{1}) + (\underline{27}, \underline{3}) + (\overline{\underline{27}}, \overline{\underline{3}}) + (\underline{1}, \underline{8}) \tag{3}$$

First consider the gauge fields A_μ ; they are invariant under H, hence under $SU(3)_{YM}$ and therefore in the $(\underline{78}, \underline{1})$ piece : we find gauge fields in the adjoint representation of E_8.

Similarly the A_k, k=5...10, being in representations $\underline{3} + \overline{\underline{3}}$ of $H = SU(3)$ ($A_1 = A_5 + iA_6$, $A_{\bar{1}} = A_5 - iA_6$, $A_2 = A_7 + iA_8$ ...), we find them, through the identification $H \equiv SU(3)_{YM}$, in $\underline{27}$ and $\overline{\underline{27}}$ of E_6 : this prescription yields gauge nonsinglet scalar fields.

A parallel analysis can be performed on the 10 dimensional gauginos, giving the supersymmetric partners of these fields. We thus find :

- a gauge symmetry $E_6 \times E_8'$: the gauge fields and gauginos in $\underline{78}$ of E_6 form a 4-dimensional gauge supermultiplet (we have not touched E_8' and the corresponding gauge supermultiplet).
- non gauge neutral scalar fields in $\underline{27}$ and $\overline{\underline{27}}$ of E_6 and their supersymmetric partners, which form matter supermultiplets (and will hopefully describe quarks, leptons, Higgs...).

Finally if there is a hole in the compact manifold K, there is the possibility to have non contractible loops γ around the hole along which the circulation of A_k is nonzero in the vacuum :

$$< \oint_\gamma A_k \, dx^k > \neq 0 \tag{4}$$

From our interpretation of A_k as a Higgs field, it is clear that (4) will induce some breaking of the E_6 symmetry into a subgroup G of E_6 (*breaking by Wilson loops*)[2,8,9]. Matter supermultiplets will fall into representations of G ($\subset \underline{27}$, $\overline{\underline{27}}$ of E_6).

Among the models that incorporate $SU(3) \times SU(2) \times U(1) \subset G$, one finds two minimal models :

Model A : minimum number of generations : $n_g = 3$.
The corresponding manifold was built by Yau[10]. So far, this is the only Calabi-Yau model with 3 generations (at least as far as the particle content is concerned). The gauge symmetry is of rank 6. Greene et al.[11] have studied the case where $G = SU(3)_C \otimes SU(3)_L \otimes SU(3)_R$ for which the particle content is : $7(\underline{3}, \underline{3}, \underline{1}) + 4(\bar{\underline{3}}, \bar{\underline{3}}, \underline{1})$; $7(\bar{\underline{3}}, \underline{1}, \bar{\underline{3}}) + 4(\underline{3}, \underline{1}, \underline{3})$; $9(\underline{1}, \bar{\underline{3}}, \underline{3}) + 6(\underline{1}, \underline{3}, \bar{\underline{3}})$. In each case, one has 3 chiral fields $(n_g = 3)$ + $n=4$ or 6 mirror fermions. This gives a net number of 9 quarks, antiquarks and 27 leptons respectively.

Model B : minimal gauge symmetry $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\eta$.
The corresponding manifold was built by Strominger and Witten[12]. The rank is 5 and therefore there is, in the approach that we have described, at least one extra gauge quantum number[2]. In this minimal model, it is the charge Y_η associated with $U(1)_\eta$, called the "extra U(1)". As we will see, this "extra U(1)" is not a characteristic of superstring models and we have no guarantee that it will survive down to low energies (see Section 3.1). The number of generations is $n_g = 4$ for this model.
iii) it is possible in some cases[13] to identify a $SU(4)_{YM}$ subgroup picked up in E_8 ($E_8 \supset SO(10) \times SU(4)_{YM}$) with a SU(4) group which contains $H = SU(3)$. In this case, the gauge symmetry is $SO(10) \times E_8'$. This gauge symmetry can be further broken by Wilson loop breaking and in principle a model with the minimal gauge symmetry $H = SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ could be constructed.

Such models are often referred to as (0,2) models*). Their phenomenology has not been studied in much detail yet because of a somewhat troubled history : soon after they were proposed[13], they were shown to be unstable under non-perturbative effects on the world-sheet[14] ; but a class of them does not suffer from this unstability[15].

Before leaving Calabi-Yau models, let us stress that the complexity of Calabi-Yau manifolds prevents from computing explicitly all the couplings, although many Yukawa couplings can be obtained by algebraic geometry arguments. Recently however, Gepner has constructed explicit string solutions which seem to describe compactification

*) (0,2) refers to the supersymmetry charges at the 2-dimensional level (i.e. on the world-sheet). We saw earlier that left and right movers are completely independent in a closed string. Correspondingly, one can define a (p,q) supersymmetry with p left-moving and q right-moving supersymmetry charges in 2 dimensions. It turns out that the Calabi-Yau models described in ii) correspond to a (2,2) supersymmetry and, using the same terminology, are also called (2,2) models.

on some Calabi-Yau manifolds[16]. This would allow in principle an exhaustive computation of all physical quantities.

<u>Orbifolds</u>[17]

They represent the second large class of spaces which can be chosen to compactify a 10-dimensional string. Some of the corresponding models are simple enough to allow detailed computations of the couplings between the different fields.

To have a feeling of how orbifolds work, let us take the simplest example. Consider a torus which we represent by a square $A_1A_2A_3A_4$ whose we identify the opposite sides (Fig. 2). One constructs an orbifold $\mathcal{O}$ by further identifying two points which are symmetric with respect to the center O of the square. We are left with the lower triangle $A_1A_3A_4$ (the upper triangle is redundant) for which we should identify some of the boundaries : by symmetry with respect to O, $A_1B = A_3B'$ and because we are on a torus $A_3B' = A_4B$; hence $A_1B = A_4B$; similarly $A_1O = A_3O$ and $A_4C = A_3C$. Indeed, if we fold along the dotted lines to realize these identifications, we end up with a mere tetrahedron, which is therefore the simplest example of orbifold.

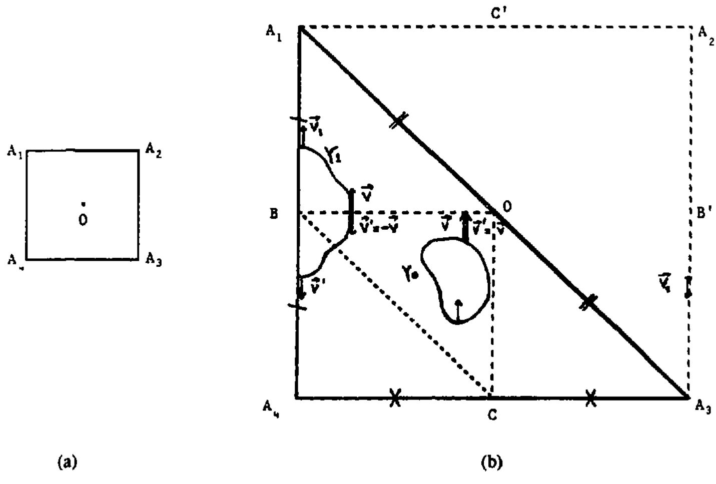

Fig. 2 : $\mathbb{Z}_2$-orbifold (b) obtained from a torus (a).

It is straightforward to compute the holonomy group. First consider the loop γ_0 in Fig. 2-b ; because the region enclosed is flat, a vector parallel-transported along γ_0 is not rotated. On the other hand, along the curve γ_1 which is closed in $\mathcal{O}$ (but not in the original torus Fig. 2-a), a vector undergoes a sign flip, due to the identifications : $\mathbf{BA_1} = \mathbf{B'A_3} = \mathbf{BA_4}$. [Note that in the tetrahedron formulation, γ_1 is a curve around the apex B where there is a singularity of curvature]. We conclude that the holonomy group is : $H = \{-1, 1\} = \mathbb{Z}_2$, a discrete group.

We could try to identify H with part of the gauge symmetry as in the Calabi-Yau case. To do so, one needs to isolate a discrete subgroup inside a gauge group. Consider for example SU(2) : the two matrices $\mathbf{1}_2$ and $-\mathbf{1}_2$ ($\mathbf{1}_n$ being the n-dimensional identity matrix) are elements of SU(2) which commute with all the others ; they form the center of SU(2) $C_2 = \{\mathbf{1}_2, -\mathbf{1}_2\} = \mathbb{Z}_2$. We therefore pick up a SU(2) subgroup inside E_8 ($E_8 \supset E_7 \times SU(2)$) and identify H with C_2. The gauge symmetry is broken to $E_7 \times SU(2)$*) but we have N=2 supersymmetry : the four supersymmetry charges that Q_A represents (see above) transform as $\underline{2} + \underline{1} + \underline{1}$ under H and therefore two 4-dimensional charges are invariant.

If we want N=1 supersymmetry in 4 dimensions, we must construct orbifolds with $H = \mathbb{Z}_3$, the cubic roots of unity. One then picks up a SU(3) subgroup of E_8 and identify H with its center $C_3 = \{\mathbf{1}_3, e^{2i\pi/3}\,\mathbf{1}_3, e^{4i\pi/3}\,\mathbf{1}_3\}$. The gauge symmetry is then $E_6 \times SU(3)$. This example is instructive for two reasons :

– it shows that the gauge symmetry can be larger than E_6,

– but for a large class of models (often called with "standard embedding"), E_6 plays an important role in connection with N=1 supersymmetry ; we saw in the Calabi-Yau case the importance of SU(3) holonomy and how it selects E_6 gauge symmetry by identification but our orbifold example shows that this class is larger than Calabi-Yau manifolds**) . Indeed all phenomenologically viable models which have been presented so far incorporate E_6 gauge symmetry in one way or another***).

*) Note here the difference with the Calabi-Yau case : because only the center of SU(2) has been identified with H, the SU(2) gauge fields still participate in the local gauge symmetry.

**) One can show however that there are no orbifolds with standard embedding leading to 3 generations[17].

***) So far we have disregarded the "other" E_8 group. Because H is discrete, it is possible to have an identification partly in E_8, partly in E_8'. One obtains models such as $[E_7 \times U(1)] \times [SO(14) \times U(1)]'$ which are truly non-E_6-based[18].

Of course, the E_6 gauge symmetry can be broken by Wilson loops. One obtains this way models[19] with $n_g = 3$ and gauge symmetry $SU(3) \times SU(2) \times U(1)^n$ (n large).

An interesting feature of orbifolds is the presence of twisted sectors i.e. massless modes of strings that wind around the singularities of curvature (similar to γ_1 in Fig. 2-b). The fields in these sectors do not have the 10-dimensional "field theory" interpretation that could be given for untwisted sectors (say $A^a{}_k$ in the Calabi-Yau case). But a mechanism similar to Wilson loop breaking can be implemented in twisted sectors[20] and has recently been used to construct truly minimal models[21,22], i.e. models with $n_g = 3$ and $SU(3)_C \times SU(2)_L \times U(1)_Y$.

2.2. Four-Dimensional Strings[23]

An alternative way to compactification is to start directly to work in D=4 space-time dimensions. This corresponds to the so-called 4-dimensional string theories which have received much attention in recent times.

There are several possible constructions of these theories:

- one is the method that we sketched in the introduction to Section 2 : it involves a (22,6) dimensional lattice (bosonic formulation).
- the other one uses the well-known equivalence between one boson and two fermions in 2 dimensions to replace the $22 + 6 = 28$ bosonic internal degrees of freedom by 56 free fermion fields - to which one should add the 6 fermionic degrees of freedom present in the right-moving sector (fermionic formulation).
- it is also possible to combine these constructions with the orbifold construction presented in Section 2.1.

This leaves us with numerous possible models. The gauge group is a priori of rank $26 - 4 = 22$, for example SO(44) but it is possible to reduce the rank. However, at least in the fermionic formulation, no model with gauge symmetry $SU(3) \otimes SU(2) \otimes U(1)^n$, n=1,2, has been found.

Due to the huge number of possible models, not much attention has been paid to the phenomenological aspects which could be specific to the 4-dimensional string models. In this respect an interesting and somewhat different approach has been adopted by I. Antoniadis et al.[24] which led them to revive the so-called flipped $SU(5) \times U(1)$ model[25].

Flipped SU(5) x U(1) model

The idea is to try to obtain a SU(5) model similar to the familiar grand unified theory. The immediate problem that one faces is that in grand unified SU(5), there are Higgs fields in the adjoint representation $\underline{24}$, which is the representation of the gauge fields. It turns out to be very difficult in string models to have a given representation present simultaneously in the gauge field and scalar spectrum. It is easy to realize why in the Calabi-Yau case (the discussion is similar for 4-dimensional strings). We saw that both gauge fields and gauginos originate from the 10-dimensional gauge field $A^a{}_M$ of E_8. Restrict a to the 24 indices $\underline{a}$ that correspond to the $\underline{24}$ of SU(5) $\subset E_8$. Depending on the transformation properties of $A^{\underline{a}}{}_M$ under the holonomy group H (through the identification of H with part of the gauge symmetry), the effective 4-dimensional fields will be either of the $A^{\underline{a}}{}_\mu$ type or of the $A^{\underline{a}}{}_k$ type but not both.

To remedy this problem, Antoniadis et al.[24)] have looked for a generalisation of SU(5) that involves only Higgs fields in representations 10, $\overline{10}$, 5 and $\overline{5}$: this is the flipped SU(5)' x U(1) model[25)]. The matter field content is (the second entry is the U(1) charge)

$$(\underline{10}, 1) = (u_i, \nu^c, d_i, d^c{}_i) \qquad (\underline{\bar{5}}, -3) = (u^c{}_i, \nu, e^-) \qquad (\underline{1}, 5)$$

to be compared with the standard SU(5) assignment:

$$\underline{10} = (d_i, e^c, u_i, u^c{}_i) \qquad \underline{\bar{5}} = (d^c{}_i, e^-, \nu).$$

It is clear that the SU(5)' group does not coincide with the standard SU(5). For example, the electric charge is a combination of SU(5)' generators and of the extra U(1) generator. One loses this way the property of quantization of the electric charge that the standard SU(5) model incorporated (indeed $\sum_i Q_i \neq 0$ for a given representation)*).
The Higgs content is, as advertised above,

$$(\underline{10}, 1) \quad (\overline{\underline{10}}, -1) \quad (\underline{5}, -2) \quad (\underline{\bar{5}}, 2) \quad (n_g+1) \times (\underline{1}, 0).$$

This model has been constructed by compactification of 10-dimensional strings ((1,2) model)[26)] as well as from a 4-dimensional string model (fermionic formulation)[27)].

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*) But this might not be a problem for superstring models since SU(5)' x U(1) is incorporated in a "ghost" E_6 or E_8, of which the electric charge is a generator.

To conclude this section, one should stress that real progress has been made recently in the detailed knowledge of specific models. Still a lot needs to be done however to devise general criteria that would allow to discriminate between the numerous different models. In the range of scales that we have discussed so far (of order M_{GUT} to M_{Pl}), one can single out the following remarks:

- the only minimal candidates (n_g=3 and SU(3) x SU(2) x U(1) gauge symmetry) which have been explicitly written are of the orbifold type.
- as for a gauge symmetry larger than SU(3) x SU(2) x U(1), it is a generic property of Calabi-Yau models (type (2,2)) and it is a property of all 4-dimensional models which have been explicitly constructed so far (fermionic formulation). It remains to be seen how this extra symmetry survives when we go down the energy scale.

3. THE ROAD FROM M_{Pl} TO M_W

We now wish to take the models that were described rather sketchily in the last section down to the TeV scale where they can be tested experimentally.

As emphasized in the introduction, we need only consider massless string (and compactification) zero modes. But to be able to discuss in the next section the phenomenology associated with these fields, we must give them a non-zero mass and thus break supersymmetry.

3.1. Supersymmetry-Breaking

The standard way[28] to break supersymmetry in any supergravity theory is by giving large vevs to scalar fields in a hidden sector, i.e. a sector that interacts only gravitationally with observable matter (leptons, quarks...). The sign of supersymmetry-breaking in this sector is a non-zero mass for the gravitino : $m_{3/2} \neq 0$. The information is then sent to the observable sector through gravitational interactions : supersymmetry is broken there at a scale $\tilde{m}$, which is a function of $m_{3/2}$, but it is only broken softly (no quadratic divergences are generated)[29]. More explicitly, supersymmetry-breaking terms which appear in the Lagrangian describing the interactions of the observable sector are scalar masses m_s, A-terms (terms in the scalar potential which are proportional to the superpotential W) and gaugino masses m_λ :

$$\mathcal{L} = -\frac{1}{2} m_s^2 |\varphi|^2 - A\,(W(\varphi)+W^*(\varphi)) - \frac{1}{2} m_\lambda \bar{\lambda}\lambda. \qquad (5)$$

Each of these supersymmetry-breaking parameters is of order $\tilde{m} = f(m_{3/2})$. Since this is true in particular for the mass of the Higgs of SU(2) x U(1) which we know to be of order M_W (up to some couplings), we obtain the following constraint:

$$\tilde{m} = f(m_{3/2}) = 0(M_W). \tag{6}$$

Several mechanisms have been proposed to break supersymmetry in superstring models but none of them has succeeded so far in satisfying Eq.(6). Indeed, more than a constraint, (6) represents the generation of a scale, M_W, which is some 16 orders of magnitude below the common scale of the theory, Eq.(1). Among the mechanisms that have been proposed, one counts :

- condensation of the gauginos of E_8'[30].
  By construction, the fields non-singlet under E_8' (i.e. gauge fields and gauginos) form a hidden sector. When the E_8' interaction becomes strong, one can expect the gauginos to form condensates[31]. This breaks supersymmetry but generates a large cosmological constant. There exist however some companion mechanisms which cancel the cosmological constant at the tree level[30,32] and sometimes even at one loop[33]. *There is a hope that in this case a hierarchy is generated between* $m_{3/2}$ *and* $\tilde{m}$[33,34], but no proof that it is large enough to accomodate (6).
- generation of supersymmetry-breaking terms by compactification of the Green-Schwarz anomaly-cancelling terms[35].
- supersymmetry-breaking terms by a Fayet-Iliopoulos D-term (connected with the presence of an anomalous U(1))[36].
- an interesting proposal[37] involves breaking supersymmetry at the string level using an analogue of the Scherk-Schwarz mechanism in supergravity. In its existing formulation however, this requires to write the string theory in an intermediate (D>4, say D=5) number of dimensions.

Without waiting for some progress in the theoretical understanding of these different mechanisms, one can take a more pragmatic and phenomenologically-oriented approach to discreminate between them. The idea is that usually only one type of supersymmetry breaking term is generated at first : the others are then induced by radiative corrections. It is therefore possible to analyse them according to which type of term appears first :

- A-terms as the dominant source of supersymmetry breaking usually lead to catastrophic consequences[33] : non-zero vevs are generated for colored scalars (squarks), which breaks $SU(3)_C$.
- in the case of m_s, it is often difficult to reproduce $SU(2) \times U(1)$ breaking and to break the extra gauge symmetry, if any, at large enough scales (see the constraints on extra neutral gauge bosons in the next section)[38].
- hence the favourite source of supersymmetry breaking are gaugino masses. In fact, if there are no large Yukawa couplings in the theory, this hypothesis can even be tested experimentally because masses evolve according to their gauge interactions and appear in fixed ratios. One obtains for the "minimal" Calabi-Yau model (**Model B**)[39]:

$$m_{\tilde{q}} : m_{\tilde{g}} : m_{\tilde{e}_L} : m_{\tilde{e}_R} : m_{\tilde{\omega}} : m_{\tilde{\gamma}} = 1.9 : 1 : 0.7 : 0.4 : 0.28 : 0.14.$$

3.2. Grand Desert Scenario vs. Intermediate Scale

We stressed earlier the connection between supersymmetry-breaking and gauge symmetry breaking. This is indeed how we could derive the constraint (6) on the supersymmetry-breaking scale, by looking at the specific breaking of $SU(2) \times U(1)$. A similar analysis can be performed for the extra gauge symmetries that are usually present in the models discussed in Section 2. We can therefore try to answer the following question particularly relevant for phenomenology : is the extra gauge symmetry broken, partly or completely, at scales intermediate between 1 TeV and M_{GUT}, or is it still present at low energies (grand desert scenario)?

If we want to break extra gauge symmetry at a scale $M_I >> 1$ TeV, we must use $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge singlet scalars to do it. It turns out that in E_6-based models, there are only two possible candidates ($\in \underline{27}$; see Table 2 at the beginning of Section 4) which we will call ν^c and N, together with the fields with opposite quantum numbers $\overline{\nu^c}$, $\bar{N}$ ($\in \overline{\underline{27}}$).

What could generate a large vev for these fields? For simplicity, we will restrict our analysis to one N and one $\bar{N}$, and determine the corresponding potential. Quite generally in a supersymmetric model, the scalar potential is given by

$$V = \sum_i \left|\frac{\partial W}{\partial \phi_i}\right|^2 + \frac{1}{2}\sum_a (D^a)^2 \qquad D^a = \phi^{i*} T^a{}_i{}^j \phi_j \tag{7}$$

where the $T^a{}_i{}^j$ are the gauge generators. The superpotential term does not yield any self-coupling for N, $\bar{N}$ (see eq.(12) below), and the D term is proportional to $(|N|^2 - |\bar{N}|^2)^2$ since N and $\bar{N}$ have opposite quantum numbers. Hence $V(N = \bar{N}) = 0$: *the scalar potential is flat in the direction* $N = \bar{N}$.

Flat directions of the potential are a characteristic of supersymmetric theories ; indeed one can show that they are not lifted by radiative corrections (as they would in any non-supersymmetric theory)[40]. There is a plethora of them in superstring models (for instance, one counts in **Model A** nine N, ν_c and six $\bar{N}$, $\overline{\nu_c}$!) and they often play an important role : no-scale property[41] of some models[42]; flat directions associated with the limit *Calabi-Yau manifold → orbifold*[43] : the orbifold limit corresponds to a zero vev of the corresponding scalar fields, called blowing-up modes ; other flat directions are associated with gauge symmetry breaking by Wilson loops.

What could lift the degeneracy in the direction $N = \bar{N}$? First, as in any effective theory where heavy modes (mass scale given by (1)) have been integrated out, non-renormalizable interactions should be present

$$W(N, \bar{N}) = \lambda \frac{(N\bar{N})^n}{M_{Pl}^{2n-3}} . \qquad (8)$$

This is something that we are familiar with in weak interactions i.e. the Fermi coupling with a constant G_F of order $1/M_W^2$.

Also, as for $SU(2) \times U(1)$ breaking, supersymmetry-breaking induces a squared mass of order $\tilde{m}^2$ for N and $\bar{N}$ (cf. eq.(5)). Radiative corrections turn them into running masses $m^2(\mu)$ but because renormalisation group equations are homogeneous, $m^2(\mu)$ remains of order $\tilde{m}^2$. Below some scale μ_0, $m^2(\mu) \leq 0$ and the potential acquires a non-trivial minimum. More explicitly, the potential reads in the direction $N = \bar{N}$ using (5) and (8)

$$V(N = \bar{N}) = - \tilde{m}^2 N^2 + \lambda^2 \frac{N^{4n-2}}{M_{Pl}^{4n-6}} \qquad (9)$$

where we have neglected quantities of order one. The ground state is, using (6):

$$\frac{<N>}{M_{Pl}} = \frac{<\bar{N}>}{M_{Pl}} \cong \left(\frac{\tilde{m}}{\lambda M_{Pl}}\right)^{1/2n-2} \cong \left(\frac{M_W}{\lambda M_{Pl}}\right)^{1/2n-2} \tag{10}$$

and some gauge symmetry is broken at[44])

$$M_I \cong g <N> = 0 \left[M_{Pl} \left(\frac{M_W}{\lambda M_{Pl}}\right)^{1/2n-2} \right] \tag{11}$$

where g is the corresponding gauge coupling. One obtains $M_I \cong 10^{10,11}$ GeV for n=2 and $M_I \cong 10^{14,15}$ GeV for n=3*).

For example, in **Model A**, Greene et al.[11]) have exhibited flat directions which allow a two step breaking that takes $SU(3)^3$ down to $SU(3)_C \times SU(2)_L \times U(1)_Y$. As for the flipped $SU(5) \times U(1)$ model[24]), there is a unique flat direction ($|\underline{10}> = |\overline{\underline{10}}> \neq 0$) which breaks $SU(5) \times U(1)$ to $SU(3)_C \times SU(2)_L \times U(1)_Y$; the $\underline{5}$ and $\overline{\underline{5}}$ naturally provide doublets through the missing partner mechanism, which ensures the final breaking to $SU(3)_C \otimes U(1)_{e.m.}$.

There is a running debate on the question whether intermediate scale scenarios are phenomenologically viable or not. Arguments for and against are summarized in Table 1. The conclusion is that an intermediate scale can exist if it is larger than 10^{16} GeV. If there is no intermediate scale, the gauge symmetry that was found in Section 2 remains untouched until one reaches the TeV region and should provide striking new phenomenology in future accelerators.

*) There exist cases where no non-renormalizable terms are generated in the flat direction considered[13]). One can then set λ to zero which sends M_I to infinity. This tells us that in fact $M_I \cong M_{comp}$. For these models, everything happens as if N (or ν^c) acquired a non-zero vev through compactification. Since $<N>\neq 0$ breaks E_6 to SO(10), one obtains SO(10)-based models at compactification. These are actually the (0,2) models that we have discussed earlier.

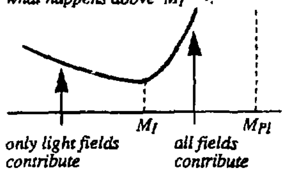
For		Against
Allows to prevent baryon-violating interactions by giving a mass of order M_I to the fields responsible for it.		Proton stability $\Rightarrow M_I > 10^{16}$ GeV
This large number allows to push M_I to larger values (by increasing n) [11].	←	Loss of predictive power due to the large number of flat directions.
Allows to give masses of order M_I to many fields (usually models have too many color triplets to have asymptotic freedom).	→	*True but one needs to be careful with what happens above M_I* [45].  *$\Rightarrow M_I > 10^{16}$ GeV.*
Model-dependent. Favourite supersymmetry breaking mechanism avoid this problem[47].	←	Cosmological problem (late release of entropy)[46].
		Difficult to prevent baryon number violating interactions (dimension 4 operators) to appear below M_I if there is no extra gauge symmetry to prevent them[48].

Table 1 : Arguments for and against the intermediate scale scenario.

4. LOW-ENERGY PHENOMENOLOGY

In the absence of a breakthrough in the theoretical understanding of string theory which would either invalidate the theory or single out a unique model, new experimental results will be needed to further constrain the models and give better indications on how to solve the problems described in the previous sections. Let us now review what this new physics could be, and what information it could give.

New Gauge Bosons

We saw earlier that gauge symmetry is broken either at an intermediate scale $M_I > 10^{16}$ GeV or at the TeV scale. If some extra gauge symmetry survives below 10^{16} GeV, there is a good hope to find new gauge bosons in the range of energy which will be explored by future accelerators. In the case of E_6-based models, there are at most 2 extra neutral gauge bosons (the rank is 6, that is 2 units more than the rank of the standard model). In terms of the $SU(3)_C \otimes SU(2)_L \otimes U(1)_{Y_L} \otimes U(1)_{Y_R} \otimes U(1)_{T_{3R}}$ decomposition of E_6 indicated in Table 2, the corresponding charges are linear combinations of Y_L, Y_R, T_{3R} and T_{3L} (Table 2 gives the value of these charges for all the members of a 27-plet ; $Y = Y_L + Y_R$).

The existing limits on the masses of these new gauge bosons Z' are discussed in P. Langacker's talk[49]. Let me summarize the situation[50,51] by saying that they range from 130 GeV (**Model B**) to a few hundred GeV : neutral current data usually give better limits than direct searches at $Sp\bar{p}S$ (except for **Model B**). Let us note that these limits are substantially higher[51] (in the 200 GeV region) if one makes some reasonable assumption on the way gauge symmetry and/or supersymmetry is broken. On the other hand, if one allows for some rescaling of the couplings at the compactification scale (relaxing therefore grand unification *stricto sensu*), one generates a much larger class of models and the limit goes down : for example, since $2Y_L - Y_R$ is zero for the leptons (see Table 2), the corresponding neutral gauge boson couples to leptons only through its mixing with Z^0 and the limit on its mass is practically inexistant[52].

Going to Z' searches at future hadron colliders and restricting ourselves to signatures with known particles, the two most interesting decay channels are:

a) $Z' \to \ell^+\ell^-$

As in the energy range covered by $Sp\bar{p}S$, this channel has no serious background[53]. The width of the Z' is generally small[54] (for example $\Gamma_{Z'}/M_{Z'} \approx 7 \times 10^{-3}$ and $B(e^+e^-) \approx 4\%$ for **Model B**) and distinguishing among the models might require good

$SU(3)_c \otimes SU(3)_L \otimes SU(3)_R$		$SU(3)_c \otimes SU(2)_L$	Y_L	Y_R	T_{3R}
(3, 3, 1)	$Q = \begin{pmatrix} u \\ d \end{pmatrix}$	(3,2)	1/3	0	0
	D	(3,1)	-2/3	0	0
$(\bar{3}, 1, \bar{3})$	u^c	$(\bar{3},1)$	0	-4/3	-1/2
	d^c	$(\bar{3},1)$	0	2/3	1/2
	D^c	$(\bar{3},1)$	0	2/3	0
$(1, \bar{3}, 3)$	$H_u = \begin{pmatrix} H_u^+ \\ H_u^0 \end{pmatrix}$ or $E^c = \begin{pmatrix} E^+ \\ N_E^c \end{pmatrix}$	(1,2)	-1/3	4/3	1/2
	$H_d = \begin{pmatrix} H_d^0 \\ H_d^- \end{pmatrix}$ or $E = \begin{pmatrix} \nu_E \\ E^- \end{pmatrix}$	(1,2)	-1/3	-2/3	-1/2
	$L = \begin{pmatrix} \nu \\ e^- \end{pmatrix}$	(1,2)	-1/3	-2/3	0
	e^c	(1,1)	2/3	4/3	1/2
	ν^c	(1,1)	2/3	-2/3	-1/2
	N	(1,1)	2/3	-2/3	0

Table 2 : Value of the charges under the $SU(3)_C \otimes SU(3)_L \otimes SU(3)_R$ and $SU(3)_C \otimes SU(2)_L \otimes U(1)_{Y_L} \otimes U(1)_{Y_R} \otimes U(1)_{T_{3R}}$ subgroups of E_6 for the components of the 27 of E_6 ($Y = Y_L + Y_R$)

resolution detectors. In this respect, the forward-backward asymmetry of leptons pairs could make the difference between the different models[55,54] as well as with some composite models[56]. This would however need good lepton charge identification and is, in any case, less efficient to find a new Z' than the direct production of lepton pairs.

b) $Z' \rightarrow W^+W^-$

This channel is open if the Z' is heavy enough ($M_{Z'} > 2M_W$) and if there is a non-zero Z–Z' mixing (the branching ratio is significant even for small mixing angles[52,57]). The relative branching ratio $\Gamma(Z' \rightarrow W^+W^-)/\Gamma(Z' \rightarrow e^+e^-)$ is somewhat model-dependent ranging from a few hundred for the model with $U(1)_{2Y_L-Y_R}$ (see above) to 10^{-1} and below. The WW pair is best detected through jet-jet-$\ell\nu$. The process is much similar to $H \rightarrow W^+W^-$ (except that Z' is much narrower) and the fairly large background is reduced along the same methods[57]. But one loses this way a good part of the signal and in most cases the clean signal $Z' \rightarrow \ell^+\ell^-$ is more efficient.

Discovery limits[58] for extra Z' range from 0(250 GeV) at the Tevatron or 300 GeV to 500 GeV at LEP II to 0(3 TeV) for LHC and 0(4 TeV) for SSC (6 TeV if exotic fermion decay channels are inacessible). Forward-backward asymmetries should allow a good discrimination between the models for masses up to 1 TeV (1.5 TeV if exotic fermions suppressed).

Higgs particle

As any other supersymmetric model, a superstring model has at least two Higgs doublets, which yield one charged Higgs, two neutral scalars and one pseudoscalar (one should actually add at least one singlet of type N to generate a H_uH_d coupling in the potential, see Table 2 and eq.(12) below). A careful study of radiative $SU(2) \times U(1)$ breaking in the simplest models has allowed to put some upper limits on the masses of these fields[59] : in particular the mass of the lightest scalar should be smaller than 0(150 GeV). It remains to be seen how general this result is.

Exotic Leptons

As indicated in Table 2, some of the isodoublet supermultiplets can be interpreted as [Higgs, Higgsino] or [exotic leptons, sleptons] (if we allow intergenerational couplings, we need only one generation of Higgs but we have at least n_g such doublets). Present mass limits from PETRA and PEP are fairly low[60] : $m_{N_E,\nu_E} > 18$ GeV, $M_E >$ 23 GeV. Langacker and London[61] have recently been able to push these limits by looking at the constraints that arise from mixings with ordinary fermions but the new limits depend on the type of mixing chosen.

Since ν_E and N_E^c are isodoublets, they should give the standard neutrino contribution (up to phase space) to the Z^0 width. Because $E^\pm$ couples vectorially to Z^0, one should not expect any forward-backward asymmetry in $e^+e^- \to E^+E^-$[62,63].

Exotic Color Triplets

We find in Table 2 color triplet isosinglet fields D (D^c) which could lead to some striking phenomenology. Before going further let us write the most general superpotential in the case of E_6-based models :

$$\begin{aligned} W = {} & \lambda_u H_u Q u^c + \lambda_d H_d Q d^c + \lambda_e H_d L e^c \\ & + \lambda_\nu H_u L \nu^c + \lambda_D D D^c N + \lambda_H H_u H_d N \\ & + \mu_1 D^c L Q + \mu_2 D e^c u^c + \mu_3 D \nu^c d^c \\ & + \rho_1 D Q Q + \rho_2 D^c u^c d^c \end{aligned} \tag{12}$$

where generation indices are suppressed (in principle, intergenerational couplings are allowed but discrete symmetries that can be determined from the compact manifold chosen[2,11]) put constraints on them). Remember that the Yukawa couplings are obtained from the superpotential by:

$$\mathcal{L} = \sum_{ij} \frac{\partial^2 W}{\partial \phi_i \partial \phi_j} \bar{\psi}_i \psi_j + \text{h.c.} \tag{13}$$

We have already discussed some of these couplings : the λ_u (λ_d, λ_e) term is necessary in order that $\langle H_u^0 \rangle \neq 0$ ($\langle H_d^0 \rangle \neq 0$) generates a mass for u-type quarks (d-type quarks, charged leptons); the λ_H term allows a mixing $H_u H_d$. The μ_i couplings induce $D \to$ quark + lepton whereas the ρ_i couplings induce $D \to 2$ quarks. Because the couplings of D, D^c generally violate baryon number, one must discuss separately models with and without intermediate scale (in the former case, the λ_D coupling can generate a very large mass for $D : \langle N \rangle = 0(M_I)$).

a) Models without intermediate scale breaking.

In these models D is necessary light (of order $\lambda_D<N>$)*). Since one must in this case forbid to a first approximation any baryon number violating interaction, the μ and ρ couplings cannot exist simultaneously. The most interesting baryon and lepton number assignments are :

a1) D is a leptoquark $B = 1/3$ $L = 1$.

The ρ couplings must be set to zero. To be short, I will only discuss the phenomenology of the scalar component $\tilde{D}$: the phenomenology of the fermionic component is very similar[65].

The $p \overset{(-)}{p}$ production of $\tilde{D}\bar{\tilde{D}}$ pairs is similar to the production of squark pairs. $\tilde{D}$ can also be singly produced in ep collisions : the signature is very clean but depends strongly on the magnitude of the Yukawa coupling μ_2.

In e^+e^- annihilation, the ratio $\sigma(e^+e^- \rightarrow Z \rightarrow \tilde{D}\bar{\tilde{D}})/\sigma(e^+e^- \rightarrow Z \rightarrow \mu^+\mu^-)$ is $0(7\%)$ in the limit of light $\tilde{D}$; decays of a Z' could also be a copious source.

Among the possible decays, $\tilde{D} \rightarrow$ hard lepton or missing p_T + jet would provide a spectacular signature. A search for this decay has actually allowed UA1 to put a limit on the mass[66) : $m_{\tilde{D}} > 33$ GeV. Also of interest is $\tilde{D} \rightarrow D + (\tilde{\gamma}, \tilde{Z}, \tilde{Z}'...)$ if $m_{\tilde{D}} > m_D + ...$

Finally leptoquarks could be searched through their virtual effects on the Drell-Yan process[67).

As emphasized by Barger, Deshpande and Hagiwara[65), if $\tilde{D}$ is light enough, it might swamp conventional top and gluino signals :

– if $m_{\tilde{D}} < m_t$, $t \rightarrow \tilde{D}\ell^+ \rightarrow \ell^+\ell^- u$ or $\ell^+ \nu d$.

– if $m_D + m_{\tilde{D}} < m_{\tilde{g}} < m_{\tilde{q}}$, $\tilde{g} \rightarrow \bar{\tilde{D}} D$ or $\tilde{D}\bar{D}$.

Discovery limits for leptoquarks[64,65,68) are $0(150$ GeV$)$ at the Tevatron, $0(250$ GeV$)$ at Hera through direct production (with a coupling comparable with α_{em}), and $0(2$ TeV$)$ at LHC or SSC.

a2) D is a diquark $B = -2/3$ $L = 0$.

The μ couplings must be set to zero. In this case $\tilde{D}$ phenomenology is much less exciting due to the large QCD background. $\tilde{D}$ can be produced in $p\overset{(-)}{p}$ colliders

*) Existing limits are : $m_D > 23$ GeV from direct searches at PETRA, PEP[60) and from 20 GeV to 50 GeV from constraints on the mixing with ordinary fermions (depending on the type of mixing)[61). There exist also many constraints involving couplings which make it difficult to envisage a D much lighter than 100 GeV[64).

either in pair (similarly to squark production) or singly through the Yukawa coupling ρ_1. The decay of $\tilde{D}$ into 2 jets is difficult to isolate from the background.

In both cases **a1** and **a2**, there is a residual R-parity defined as usual by

$$R = (-)^{3B+L+2S} \tag{14}$$

which gives $R(D) = -1$, $R(\tilde{D}) = +1$. This R-parity might be spontaneously broken along with baryon or lepton number, but precisely for this reason, this should be a very small effect[69].

b) Models with intermediate scale breaking.

In this case, D can be made heavy through its coupling λ_D to N : $m_D = 0(M_I)$; and its couplings may violate baryon or lepton number. A leptoquark D for example is very similar to the Higgs triplet that mediates proton decay in standard SU(5) grand unification. We saw earlier (see Table 1) that, for this reason, one must have a rather large intermediate scale: $M_I > 10^{16}$ GeV. Because there now exist strong constraints on proton decay from IMB and KAMIOKA, it is in fact possible to obtain information on the low energy mass spectrum if M_I is near this critical value of 10^{16} GeV.

Indeed, Arnowitt and Nath[70] have studied the minimal **Model B** with a non-zero gaugino mass as a dominant supersymmetry breaking source. If $m_{\tilde{u}} = m_{\tilde{d}}$, the main contribution to the dominant mode $p \to K^+\bar{\nu}$ comes from Wino dressing of the effective baryon number violating operator $O(m_D^{-1})$. One obtains a decay rate which is a function of m_D, $m_{\tilde{q}}$ and $m_{\tilde{\gamma}}$ (due to the assumption about supersymmetry-breaking, all gaugino masses have a fixed ratio with respect to $m_{\tilde{\gamma}}$; see end of Section 3.1.). Its strong dependence on $m_{\tilde{q}}$ and $m_{\tilde{\gamma}}$ allows one to obtain stringent constraints. For example, Arnowitt and Nath find, for the typical value $M_I = 10^{16}$ GeV and for $m_{\tilde{\gamma}} = 1$ GeV (10 GeV) that $m_{\tilde{q}} > 500$ GeV (1 TeV). It would be interesting to see how model dependent this type of result is.

A very interesting possibility that appears with intermediate scale breaking models is the breaking of R-parity[71]. If R-parity is broken, the Lightest Supersymmetric Particle (LSP), say the photino, becomes unstable. For example, one could obtain $\Delta B \neq 0$ decays such as $\tilde{\gamma}d \to ud$ (using $\tilde{d}/\tilde{D}$ mixing and the ρ_1 coupling) or $\Delta L \neq 0$ decays such as $\tilde{\gamma} \to \ell^+\ell'^-\nu$ *). Dimopoulos and Hall[72] have recently discussed the

*) In a lepton number conserving and baryon number violating model, the proton would be stable since it decays into an odd number of leptons. If a superstring model of this type was found, the bound on M_I would not necessarily hold.

most striking phenomenological aspects of such a possibility. They stress in particular the following point : typical supersymmetric searches involve missing energy signatures (the LSP escapes the detector) whereas in this type of models, isolated leptons should be the prominent signature (cf. the $\Delta L \neq 0$ decay indicated above).

As we see, new physics observed at the future accelerators might give us ways to probe superstring models at much larger scales. This is one advantage of working with unified models which are a priori self-consistent. Another one is to make us think of exotic particles or exotic processes that have been disregarded until now : this can certainly be counted as the greatest (only?) success of superstring phenomenology until now.

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