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Scalar particles in superstring models *

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Abstract

The role played by scalar fields in superstring models is reviewed, with an emphasis on recent developments. The case of the dilaton and moduli fields is discussed in connection with the issues of spacetime duality and supersymmetry breaking. Constraints on the Higgs sector are reviewed in the different classes of models.

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1 INTRODUCTION.

Although the superstring tidal wave seems to have somewhat receded, important progress has been made recently in the detailed knowledge of these theories. After the days of the "superstring revolution" where the possible relevance of strings to the fundamental structure was recognized, came the time when the basic problems of this approach were identified: in particular the painfully large number of possible candidates and the issue of supersymmetry breaking. However, since one has started to address these problems, some definite progress has been made.

There are many different ways to present these advances but a particularly convenient one is to consider the status of scalar particles in these models. This is due to a striking property of strings: *in string theory, all relevant dimensionless parameters are vacuum expectation values of some scalar fields.* There is indeed only one fundamental string scale M_S: M_S is given in terms of the string coupling α' (which sets the normalisation of the 2-dimensional metric) by

$$M_S = \alpha'^{-1/2}. \tag{1.1}$$

All other relevant scales in the theory (Planck scale M_{Pl}, compactification scale M_{comp}, grand unification scale M_{GUT}, scale of supersymmetry breaking...) are expressed in terms of M_S and the vacuum expectation values (*vevs*) of associated scalar fields.

In the following Section, I will review some aspects of superstring models with an emphasis on the progress achieved in the last years. This is of particular relevance for phenomenologically-oriented studies since most of the so-called superstring-inspired phenomenology has dealt with the first models proposed and ignored later developments.† Section 3 reviews the structure of the scalar sector of string models and describes how the dilaton and the so-called moduli fields are invoved in the basic issues of string theories. Finally Section 4 deals with the subject of this Meeting: the Higgs field. This together with the spectrum of supersymmetric particles, is certainly the place where most of the superstring-related phenomenology will take place in the future. There is no definite prediction yet from superstring but, as we will see, rather stringent bounds can be derived in most models, which put the Higgs sector within the reach of the next generation of accelerators. If nothing was found, this would point towards some fairly pathological models among the many present candidates. By the time these experimental searches are completed, one may hope that theoretical studies would have themselves pointed these models out. Otherwise string models would be in serious trouble.

2 SUPERSTRINGS MODELS.

2.1 The original E_6 model.

It is important to recall what is the content of the original model and how it was

†One should stress however that some of these developments have in fact strengthened basic hypotheses made in the early days, which were based at that time on little more than prejudice.

derived because it still is the unavowed basis of many superstring-related phenomenology discussions. It was obtained by Candelas, Horowitz, Strominger and Witten[1] by compactification of the 10-dimensional heterotic string theory[2] on a 6-dimensional compact Calabi-Yau manifold (the Calabi-Yau property ensuring that the 4-dimensional theory incorporates one supersymmetry charge).

The spectrum of this class of models is rather simple. It consists of:

(i) a supergravity multiplet which includes the graviton, its supersymmetric partner the gravitino, an antisymmetric tensor field, the dilaton and its fermionic partner.

(ii) chiral supermultiplets describing the "shape" of the compact manifold, of which more later.

(iii) a gauge non-singlet sector associated with gauge group $E_6 \otimes E_8'$ or one of its subgroups $\mathcal{H} \otimes \mathcal{H}'$ ($\mathcal{H} \subset E_6$ and $\mathcal{H}' \subset E_8'$). This sector divides itself into two:

- an observable sector where one finds the usual quarks, leptons, gauge fields and Higgs; it consists indeed of gauge supermultiplets (spin 1, spin 1/2) associated with $\mathcal{H}$ and matter chiral supermultiplets (spin 0, spin 1/2) singlet under $\mathcal{H}'$ and charged under $\mathcal{H}$: in the case of E_6, which we will take for illustration purpose, they are in full representations 27 and $\overline{27}$; otherwise, they are in representations of $\mathcal{H}$ contained in 27 and $\overline{27}$.

- a sector of gauge supermultiplets associated with $\mathcal{H}'$; these fields are neutral under $\mathcal{H}$ and thus interact with the observable sector only through gravitational interactions. They are said to form a *hidden sector*.

It proved to be a more difficult task to derive the couplings of these fields. The first attempt was an educated guess by Witten[3], who found his inspiration in a toy model of torus compactification supposed to reproduce the basic features of Calabi-Yau compactification. Witten started with the Lagrangian of 10-dimensional supergravity which, to a first approximation, is the field theory limit of the heterotic string and truncated it so as to obtain a 4-dimensional model with a $E_6 \otimes E_8'$ gauge supermultiplet and *one* matter supermultiplet in a 27 of E_6, coupled to supergravity.

Quite generally, in a supergravity model, all the couplings are determined by only three field-dependent functions[4]:

(a) the normalisation of the Yang-Mills kinetic term f_{ab},

$$\mathcal{L}_{YM} = -\frac{1}{4} f_{ab} F^{a\mu\nu} F^{b}_{\mu\nu} + \cdots \tag{2.1}$$

where the indices a and b run over all the gauge group generators,

(b) the Kähler potential K which describes the normalisation of the kinetic term for the scalar fields:

$$\mathcal{L} = \frac{\partial^2 K}{\partial \Phi^A \partial \bar{\Phi}^{\bar{B}}} \partial^\mu \Phi^A \partial_\mu \bar{\Phi}^{\bar{B}} + \cdots \tag{2.2}$$

(c) the superpotential W, an analytic function, which yields the scalar potential (F-term):

$$V = e^K \left[K_{A\bar{B}} \left(\frac{\partial W}{\partial \Phi^A} + W \frac{\partial K}{\partial \Phi^A}\right)\left(\frac{\partial \bar{W}}{\partial \bar{\Phi}^{\bar{B}}} + \bar{W} \frac{\partial K}{\partial \bar{\Phi}^{\bar{B}}}\right) - 3|W|^2 \right] \tag{2.3}$$

Witten's result is expressed in terms of the multiplets introduced earlier: the matter fields $\Phi^i, i = 1,\ldots,27$, the dilaton φ and the squared radius T of the compact manifold: [‡]

$$
\begin{aligned}
f_{ab} &= S\,\delta_{ab}, &(2.4)\\
K &= -\ln(S+\bar{S}) - 3\ln(T+\bar{T}-\sum_i \Phi^i\Phi^{i*}), &(2.5)\\
W &= d_{ijk}\Phi^i\Phi^j\Phi^k, &(2.6)
\end{aligned}
$$

where

$$S+\bar{S} = (T+\bar{T})^3\varphi^{-3}. \qquad (2.7)$$

Eqs.(2.4-2.7) fully determine the couplings in the model.

It is useful to see more closely how the field T appears because it provides a first example of the fields described in (ii) above, the so-called moduli fields. The radius R of the compact manifold defined for instance from the metric g_{IJ} (I, J are 6-dimensional real indices describing the compact manifold) by[3]

$$g_{IJ} = R^2\,\delta_{IJ} \qquad (2.8)$$

is precisely given in string units in terms of the real part of the T field:

$$R = (ReT)^{1/2}M_S{}^{-1} \equiv M_{comp}{}^{-1}. \qquad (2.9)$$

The imaginary part of T, on the other hand, is deduced from the antisymmetric tensor contribution ($\epsilon_{12} = -\epsilon_{21} = \epsilon_{34} = -\epsilon_{43} = \epsilon_{56} = -\epsilon_{65} = 1$).

$$ImT = B, \quad B_{IJ} = B\,\epsilon_{IJ}. \qquad (2.10)$$

These two equations provide the full complex scalar component of the unique moduli superfield present in this toy model (in general, there are as many such scalar fields as there are 27 representations in the matter sector[5]).

Eq.(2.9) is the first example that we encounter of the property stated in the Introduction: the compactification scale, which turns out in these models to be also the grand unification scale, is given in string units by the vev of the T field. Similarly[6], the Planck scale is expressed in terms of the vev of the S field introduced in (2.7): [§]

$$M_{Pl} = (ReS)^{1/2}M_S. \qquad (2.11)$$

Another important parameter is similarly derived by considering Eqs.(2.1,2.4); the gauge coupling is obtained from the vev of the S field:

$$1/g^2 = ReS. \qquad (2.12)$$

[‡] All scalar fields will be taken to be dimensionless.

[§] The vev symbol $< \cdots >$ which should be present whenever we discuss mass scales is suppressed throughout.

2.2 Problems with the original model.

The model described in the last subsection covers more or less what was known in 1985 and served as a starting point for many superstring-inspired phenomenology papers. It has however some severe drawbacks which had to be addressed.

First of all the number of generations (only one 27) is not realistic and nothing is known of the couplings of matter fields in $\overline{27}$.

Secondly, the truncation procedure used in Ref.3, which consists in setting the heavy fields to zero, is a very dangerous and misleading procedure. The standard procedure would consist in integrating the heavy modes out but this is very difficult to implement in the case of Calabi-Yau compactification. Indeed, in order to do so, one should know the metric of the compact manifold; the choice (2.8) is obviously an oversimplification. Actually, there is no Calabi-Yau manifold for which the metric is known explicitly (this is in fact a long-standing mathematical problem). This seemed to lead to a dead end and the possibility of reliably computing the matter couplings in a realistic case seemed infinitely remote.

Finally, the compactification procedure used above, which starts from the field theory limit of the string, makes sense only if the heavy string modes lie at a scale much larger than the scale of compactification, i.e. $M_S \gg M_{comp}$. From (2.9), this means that ReT which measures the radius in string units has to be large in order for the whole approach to be consistent. But the grand unification scale given by (2.9) and (2.11)

$$M_{GUT} = M_{comp} = \frac{M_{Pl}}{(ReS)^{1/2}(ReT)^{1/2}} \tag{2.13}$$

together with the value (2.12) of the gauge coupling at grand unification are constrained by present data. This in turn sets the *vev* of ReT to be of order 1, in contradiction with our consistency condition[6,7].

2.3 The modern formulation.

In the modern formulation, the point of view changes completely. Instead of considering the problem at the level of a many-dimensional spacetime and having to deal with the intricacies of compactification, one looks at it from the 2-dimensional point of view of *Conformal Field Theory.*

A closed string which propagates in a D-dimensional space describes a 2-dimensional surface known as the *world-sheet.* Because the string is itself a one-dimensional object, the oscillations propagating along the string in one direction, say to the left, do so independently of the oscillations propagating in the other direction i.e. to the right (conformal transformations are precisely the reparametrisations of the world-sheet which leave this distinction untouched). There is therefore no difficulty in having a different number of left- and right-moving coordinates. On the other hand, well-known arguments of internal consistency impose this number to be 26 or 10 (if combined with fermionic degrees of freedom).

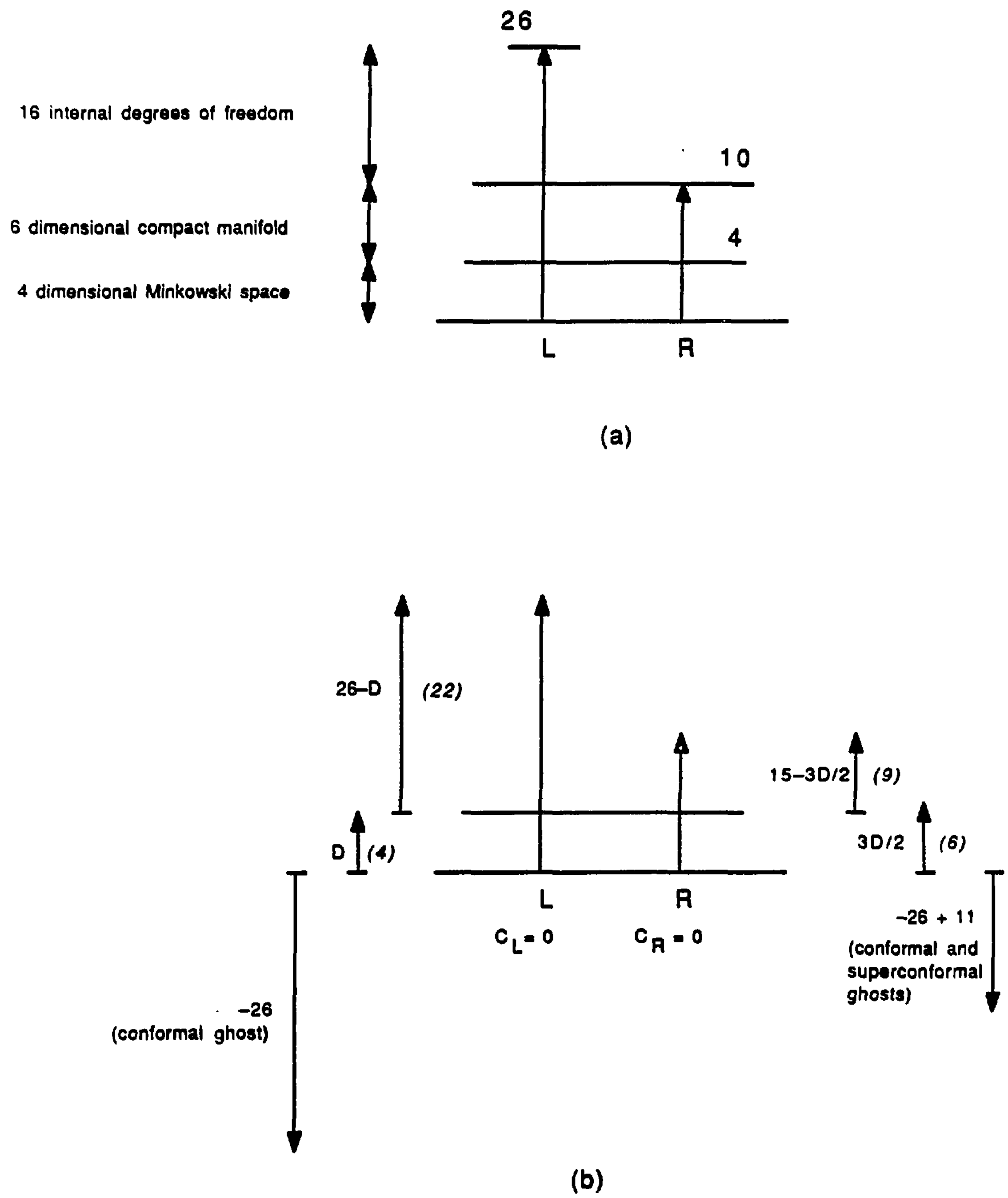

Figure 1: (a) Left- and right-moving degrees of freedom of the heterotic string theory. (b) Corresponding chart for the central charges.

In the heterotic string case, which was our starting point, space is 10-dimensional; out of the 26 left-moving degrees of freedom, 16 are internal (i.e. quantum numbers¶) and the remaining 10 describe, together with the 10 right-moving modes, the ten spacetime dimensions, six of which are compact (see Fig.1(a)).

The situation is somewhat simpler from the 2-dimensional point of view because conformal theories are distinguished only by their value of the central charge c and their field (representation) content. The central charge measures the conformal anomaly (i.e. the violation of conformal invariance which appears at one loop) but, for our purpose here, it will be merely used as a counter of degrees of freedom. Indeed, a bosonic degree of freedom, whether internal (quantum number), compact or continuous contributes one unity to c, whereas a fermionic degree of freedom contributes 1/2.

We can now represent the same system in Fig.1(b) where, this time, no difference is made between internal and compact indices. One thus bridges the gap with the 4-dimensional string approach where only 4 degrees of freedom are interpreted as spacetime. The key idea behind Fig.1(b) is the requirement of conformal invariance, even at one loop, which imposes an overall central charge of zero. Let us restrict our attention to the left-movers for a second. Because conformal symmetry is local, we must fix the gauge in a way very similar to the Faddeev-Popov procedure familiar in gauge theories and thus introduce a ghost (the conformal ghost) which contributes (-26) to the central charge[8]. Hence the famous 26 to cancel this charge. Similarly, each right-moving degree of freedom contributes 1+1/2=3/2 and ten such degrees are necessary to cancel the charges of the conformal ghost (26) and its partner the superconformal ghost (-11).

Supersymmetry can also be described at the 2-dimensional level. Again a little counting will take us a long way along the road. Quite generally in order to have chiral fermions, as we have in the standard model, we are looking for $N = 1$ supersymmetry in 4 dimensions i.e. one supersymmetry charge Q_α satisfying

$$\{Q_\alpha, Q_\beta\} = 2(\Gamma^\mu)_{\alpha\beta} P_\mu. \tag{2.14}$$

Since Q_α is a Majorana spinor, it counts for 2 degrees of freedom. Now, in 2 dimensions, one can similarly define a supersymmetry charge q_a such that

$$\{q_a, q_b\} = 2(\gamma^m)_{ab} p_m \tag{2.15}$$

where a, b are 2-dimensional spinorial indices, γ^m is a 2-dimensional gamma matrix and p_m is the momentum operator on the world-sheet. This time, q_a is a Majorana-Weyl spinor and corresponds to a single degree of freedom. Hence just by counting the number of degrees of freedom, one concludes that we need at least two supersymmetry charges on the world-sheet. These charges can be defined independently for left and right movers. Since in the heterotic string, fermionic coordinates are introduced only in the right-moving sector, we conclude that the minimal model which incorporates space-time supersymmetry has (0,2)

¶In the Calabi-Yau case, two of these numbers are identified with the local properties of the curved compact manifold -remember that general relativity is a gauge theory- and the remaining 14 are the 6 + 8 quantum numbers of $E_6 \supset E_8'$.

world-sheet supersymmetry i.e. 0 left-moving supersymmetry charges and 2 right-moving ones[9,10].

In the case of Calabi-Yau compactification, it turns out that the identifications made enhance the system of 2-dimensional supersymmetry charges to (2,2) world-sheet supersymmetry. ‖ Hence in modern language, Calabi-Yau compactification is described with the help of (2,2) superconformal systems (with 9 units of central charge to describe the internal or compact degrees of freedom; see Fig.1(b))[11].

One might wonder what is the use of introducing this whole lot of new concepts and whether this is not an attempt of hiding old problems under a new phraseology. It turns out that this new point of view allows to attack many tasks that seemed formidable in the old approach. Consider for example the problem of computing the couplings of matter in a realistic Calabi-Yau model (hence not the simple toy model presented above). If we look at it from the point of view of compactification, it looks like an impossible task: indirect arguments allowed to compute some of the couplings but certainly not all of them since even the metric of the compact manifold is not known. On the other hand, in the 2-dimensional approach, each of the physical states is represented by an operator (*vertex operator*) whose explicit form is known; computing the couplings between given states amounts to computing the correlation functions of the corresponding operators on the 2-dimensional surface –a task familiar in statistical physics–. Thus within a given model it is in principle possible to compute all the couplings of interest (one is of course limited by the huge number of possible models).

The next Section describes the main results that were obtained using these methods.

3 THE STRUCTURE SCALARS.

By "structure scalars" we mean (i) the dilaton and (ii) the moduli fields which were introduced in the compactification models of subsection 2.1 as the fields describing the shape of the compact manifold. These fields interact only gravitationally with the matter scalars to be described in the next Section and are therefore ***hidden*** from them. They nevertheless play an important role in issues of key importance for low energy predictions, such as supersymmetry breaking and the discremination between the large number of models.

3.1 The dilaton field.

We have seen in subsection 2.1 that the dilaton appears in the 4-dimensional supermultiplet spectrum as the scalar field S of Eq.(2.7). It is this field that we call here dilaton and it is as such that it provides the ratio M_{Pl}/M_S as given in Eq.(2.11).

‖One might wonder how one can generate a left-moving supersymmetry when we have introduced only bosonic coordinates in the left-moving sector. This is due to an equivalence between bosons and fermions which appears in two dimensions. The 16 left-moving internal bosonic degrees of freedom (see Fig.1(a)) are equivalent to 32 left-moving fermionic degrees of freedom $\lambda^I, I = 1\dots 32$. The solution found by Candelas *et al.*[1] yields an extra supersymmetry between the transverse degrees of freedom $X_L^I, I = 1\dots 8$ and $\lambda^I, I = 1\dots 8$.

Indeed one finds that the properties of S obtained in the toy model are quite general. In particular, it remains true at the tree level of string theory that the Yang-Mills normalisation function f_{ab} is still simply given by S as in (2.4): the value of the gauge coupling at grand unification is obtained from the *vev* of S

$$1/g^2 = Re\, S. \tag{3.1}$$

Also as before S does not appear in the superpotential, which means that the corresponding term in the scalar potential is zero (cf. (2.3)):

$$V(S) = 0. \tag{3.2}$$

Thus the S field corresponds to a flat direction of the potential and its *vev* remains undetermined at this stage. This is certainly not a desirable feature of the model since it prevents us from determining M_{Pl} as well as the value of the gauge coupling.

There is however an important property which relates the presence of flat directions of the potential with unbroken supersymmetry. One can prove the following theorem[12]: *as long as supersymmetry is not broken, flat directions are not lifted to all orders of perturbation theory.* Thus, the lifting of the degeneracy associated with a flat direction is connected with supersymmetry breaking and it must occur through non-perturbative effects. We will come back to this in a moment but let us first introduce other flat directions of the scalar potential: the moduli fields.

3.2 The moduli fields.

We have already encountered an example of such a field in our example of subsection 2.1: the T field whose *vev* was connected with the radius of the compact manifold through (2.9). This is the representative of one class of moduli which we may call of the radius type and generically represent by T. They somehow determine the topological class of the manifold[9]. There is however another class of moduli which we will denote by the letter C: these determine the complex structure of the manifold.

Let us illustrate the difference on the simplest of the compact manifolds: the torus. As is well-known, the torus can be represented as a rectangle whose opposite sides are identified. We choose to represent it in the complex plane by a square of unit length (Fig.2(a)). Changing the radius of this torus amounts to multiply all its dimensions by a factor λ (Fig.2(b)). On the other hand, the case of Fig.2(c) where only the "imaginary" direction is dilated cannot be described in the complex plane by a holomorphic transformation $z \to \lambda z$: it corresponds to a change in the complex structure of the torus.

Using (2,2) superconformal invariance, one can prove the following theorem[5]: in a given model, there are as many T fields as there are 27 representations in the matter sector and there are as many C fields as there are $\overline{27}$. One can check this result on our toy model (respectively one and zero).

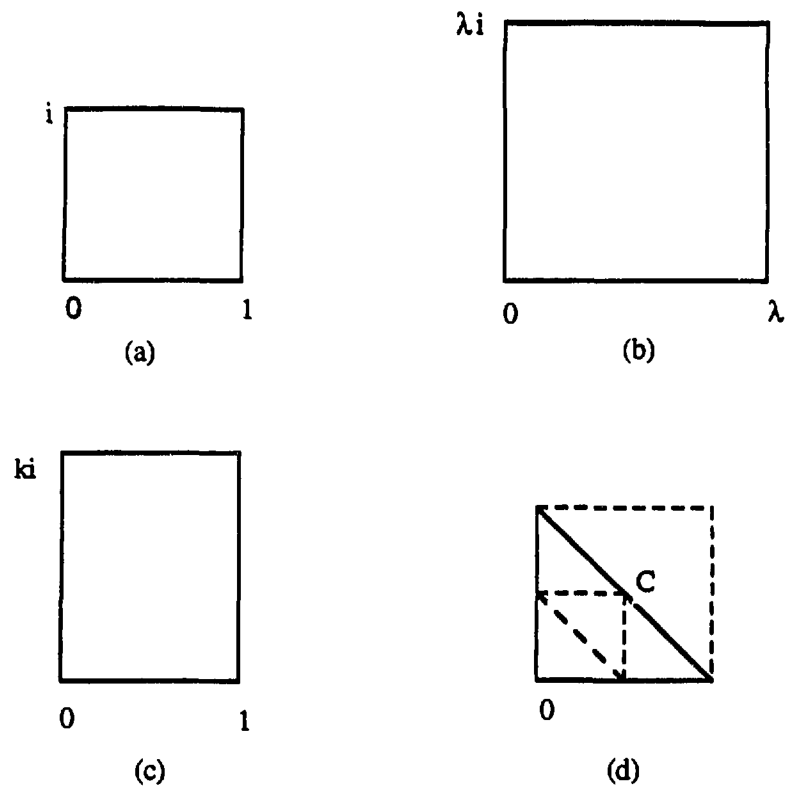

Figure 2: (a) Torus represented by a square of unit length in the complex plane: opposite sides are identified. (b) Another torus with the same complex structure. (c) Another torus with a different complex structure. (d) Orbifold obtained by symmetrizing torus (a) with respect to its center C.

Some of the T fields play an important role. We have already encountered the overall radius of the manifold and will see its role in connection with duality in the next subsection. In the case of orbifold models, one also finds among the T fields the ***blowing modes***.

Crudely speaking, orbifolds are obtained by symmetrizing tori (or more generally manifolds)[13]. If we take for example the torus of Fig.2(a) and symmetrize it with respect to its center, we need to fold it along the dotted lines in Fig.2(d) in order to make the correct identifications (we can also disregard half of its surface, say the half above the diagonal). We thus obtain a tetrahedron whose vertices correspond to the fixed points of the symmetrization procedure. It is possible to obtain a Calabi-Yau manifold from an orbifold by blowing up the singularities of curvature (vertices): one replaces each vertex by a Calabi-Yau manifold. It turns out that, for each fixed point, there is a T field whose *vev* is the radius of the Calabi-Yau manifold attached to it. Hence $< T >= 0$ corresponds to the orbifold limit whereas $< T > \neq 0$ corresponds to the Calabi-Yau case. This provides a simple way for blowing up the singularities of orbifolds and studying the Calabi-Yau models in their vicinity[5,13,14].

Most of the progress made in the last year has been in the determination of the properties of moduli. Remember that it is sufficient to know the dependence of the Kähler potential and of the superpotential (cf. (2.2) and (2.3)). One finds[15] that in the case of (2,2) models, the same holomorphic function determines the superpotential of the matter fields and the Kähler potential of the moduli (until now, the only connection we had made between the two was that they appear in equal numbers) More explicitly, both the Kähler potential $K_1(T)$ for the T fields and the superpotential W for the matter fields in the 27 are fixed by the same function $\mathcal{F}_1(T)$:

$$K_1(T) = -\ln Y_1, \tag{3.3}$$

$$Y_1 = \sum_{a=1}^{N_1}(\partial_a\mathcal{F}_1 + \partial_{\bar{a}}\overline{\mathcal{F}}_1)(T^a + \bar{T}^{\bar{a}}) - 2(\mathcal{F}_1 + \overline{\mathcal{F}}_1), \tag{3.4}$$

$$W = \frac{1}{3}\,\partial_a\partial_b\partial_c\mathcal{F}_1\;\;\Phi^a\Phi^b\Phi^c, \tag{3.5}$$

where N_1 is the number of T fields (i.e. the number of matter fields Φ in a 27).

Note that for $N_1 = 1$ and $\mathcal{F}_1 = T^3$ one recovers the couplings originally found by Witten (Eqs.(2.5,2.6)) **. Also, setting the matter fields Φ to zero, one automatically gets a vanishing potential. Thus, as promised, the moduli fields correspond to flat directions of the scalar potential.

A similar relation exists between the Kähler potential $K_2(C)$ for the C fields and the superpotential for the matter fields in $\overline{27}$. Finally, the normalisation of the kinetic term (2.2) for the matter fields can be expressed in terms of $K_1(T)$ and $K_2(C)$.[15]

All this might seem rather academic from a phenomenological point of view since the moduli remain *hidden* from the observable sector. Nevertheless, we will see in the next subsection that these fields play an important role in issues of key importance for superstring phenomenology. Moreover, even though their couplings are of a gravitational type, they could be of direct relevance for present day experiments. Take for example $N_1 = 1$ and

$$\mathcal{F}_1 = \frac{1}{8}\lambda T^4/M_{Pl}, \tag{3.6}$$

which yields a superpotential for the matter fields:

$$W = \frac{\lambda}{M_{Pl}}T\Phi^3. \tag{3.7}$$

Such a term typically gives a coupling in the Lagrangian of the form

$$\frac{\lambda}{M_{Pl}} < H > \overline{\Psi}\Psi\, T \tag{3.8}$$

where H is a Higgs field which we set to its *vev* (say 300 Gev) and the Ψ are quarks or leptons.

**A note on notations: in Eq.(2.6), there is only one 27 and the index i runs over the components of this 27. In (3.3-3.5), there are N_1 T fields, hence N_1 27 and $a = 1\ldots N_1$; the index i is suppressed.

The coupling (3.8) induces a force between the constituents of ordinary matter which is mediated by the T field. Its range depends on the mass of this modulus. Quite generally, the C fields acquire a rather large mass through supersymmetry breaking because this breaking determines a particular choice for the complex structure (hence it induces a structure in the otherwise flat direction C). The case of the T fields is somewhat different: some acquire a mass (for example[16], in some cases, the T field corresponding to the overall radius acquires a mass $m_{3/2}M_{GUT}/M_{Pl}$) but other directions remain flat. If the T field considered in (3.6) is one of them, it will eventually obtain a non-zero mass at one loop through the couplings (3.8). Cvetič[17] evaluated it to be possibly as small as 10^{-18}Gev. Then the above force mediated by T would be long-ranged (< 1000m) and could be tested in fifth force experiments.

3.3 The importance of structure scalars.

If superstring theory has not come yet with definite predictions, it is due mainly to two obstacles: the huge number of possible models and the uncertainties associated with the supersymmetry-breaking mechanism. In both instances, moduli play a very special role.

In principle, a criterion that would discreminate between the different string models and determine the ground state among them would require to fully develop a quantum field theory of strings. We are right now quite far from having such a criterion at our disposal. On the other hand, one has gradually become aware of the important role played by the symmetry known as *duality* in comparing different string theories.

It is beyond the scope of this presentation to give a detailed description of duality and its many uses but one can give an idea of how it works Consider the compactification of one string coordinate on a circle of radius R. Remember that one can define independently left-movers and right-movers i.e.

$$\begin{aligned} X(\tau,\sigma) &= X_L(\tau+\sigma) + X_R(\tau-\sigma), \\ X_L(\tau+\sigma) &= x_L + 2\alpha' p_L(\tau+\sigma) + \text{oscillating modes}, \\ X_R(\tau-\sigma) &= x_R + 2\alpha' p_R(\tau-\sigma) + \text{oscillating modes}, \end{aligned} \tag{3.9}$$

where τ and σ are coordinates on the world-sheet, $x = x_L + x_R$ is the centre of mass of the string and we have written explicitly only the zero (non-oscillating) modes. Because the coordinate X describes a circle, the associated total momentum $p = p_L + p_R$ is quantized: $p = m/R$. On the other hand, since the string can wind around the circle n times (i.e. $X(\sigma+\pi) = X(\sigma) + 2n\pi R$), we have

$$\begin{aligned} p_L &= \frac{1}{2\alpha'} nR + \frac{m}{2R} + \text{oscillating modes}, \\ p_R &= -\frac{1}{2\alpha'} nR + \frac{m}{2R} + \text{oscillating modes}. \end{aligned} \tag{3.10}$$

As is well-known, the eigenstates of the string mass-squared operator

$$\frac{1}{2}M^2 = p_L^2 + p_R^2 + \text{oscillating mode contribution} \tag{3.11}$$

constitute the spectrum of particles of the theory. The key observation is that this spectrum is invariant under the transformation

$$\begin{aligned} n &\longleftrightarrow m \\ R &\longleftrightarrow \alpha'/R. \end{aligned} \tag{3.12}$$

This is the *duality transformation*[18]. Although its derivation is straightforward, it is rather surprising to uncover a symmetry which relates compact manifolds with very large radius to ones with very small radius. One can immediately note the role it could play in our case since, as was stressed at the end of subsection 2.2, the consistency of the Candelas *et al.*[1] approach required $R \gg 1$ whereas grand unification constraints seem to indicate that R is of order 1.

We have interpreted the radius as a modulus –the unique one in the case of compactification on a circle–. In more complicated models, duality generalizes to a discrete transformation on the whole space of moduli. To get an idea of the type of structure that is involved, we will describe briefly the case of 2-dimensional torus compactification which is very similar to the compactification described in subsection 2.1. Indeed Eqs.(2.8-2.10) can be translated *verbatim* (with now I,J 2-dimensional real indices describing the torus) and we see that torus compactification is described by a single complex modulus $T = R^2 + iB$ and the modulus space is half of the complex plane, which we write in coset space notation $\frac{SL(2,\mathbf{R})}{U(1)}$. In this case, the duality transformations are simply $SL(2,\mathbf{Z})$ transformations (also called modular transformations):

$$T \rightarrow \frac{aT - iB}{icT + d}, \qquad ad - bc = 1, \qquad a, b, c, d \in \mathbf{Z}. \tag{3.13}$$

This case is fairly representative of the general structure of the moduli space. Just for the sake of completeness, let us indicate that in the case of the heterotic string compactified on a D-dimensional torus, the moduli space has locally the structure of the coset space[19] $\frac{SO(16+D,D)}{SO(16+D)\otimes SO(D)}$. The general duality transformations form the discrete group[20] $SO(16+D, D, \mathbf{Z})$. The uses of duality transformations are many: they play an important role in comparing different string theories[21], in finding models with enhanced gauge symmetry (self-dual models)[22], in discussions of strings at high temperature[23], in statistical models[24]...

Another topic where the structure scalars play an important role is supersymmetry breaking. Both the dilaton S and the moduli fields correspond to flat directions of the potential and we have pointed out earlier that the lifting of the corresponding degeneracy is a sign of supersymmetry breaking through nonperturbative effects. It is clear that supersymmetry breaking is an important issue but all the more in superstring models since we need to determine the *vevs* of the structure scalars in order to fix the value of such important quantities as M_{Pl}, M_{GUT} or the value of the gauge coupling g (cf Eqs.(2.11)-(2.13)).

Another problem specific to string is the fact that all scales turn out to be of the same order. We mentionned in subsection 2.2 that constraints on M_{GUT} and g impose to consider *vevs* for ReS and ReT of order 1. Thus a realistic model would necessarily have[6,7]

$$M_S \simeq M_{Pl} \simeq M_{GUT} \simeq M_{comp}. \tag{3.14}$$

In these conditions, it remains to explain how a scale as small as M_W can ever be generated. We know that M_W is related to the mass term for the Higgs field of $SU(2) \otimes U(1)$, itself determined by supersymmetry breaking (all the fields present at low energy are found among the massless modes of the superstring; hence any mass in this sector has a supersymmetry-breaking origin). Thus supersymmetry breaking is the place where the string theory has to address the problem of hierarchy, in the hard way.

There are several ways of dealing with supersymmetry breaking in the models that we consider (for a recent review see for example Ref.25). To emphasize the issues at stake, I will concentrate on a specific mechanism based on gaugino condensation[26]. In models compactified on a Calabi-Yau manifold, we found a hidden sector of gauge supermultiplets corresponding to the gauge symmetry $\mathcal{H}' \subset E_8'$. One expects that, when the $\mathcal{H}'$ gauge coupling becomes strong, the corresponding gauginos form condensates which break local supersymmetry in this *hidden* sector: the gravitino acquires a non-zero mass $m_{3/2}$.

In order not to generate a huge cosmological constant, one usually accompany this with a second supersymmetry-breaking mechanism whose role is to cancel the vacuum energy. A favourite way is to generate a *vev* for the field strength of the antisymmetric tensor:

$$\begin{aligned} H_{IJK} &= \partial_I B_{JK} + \partial_K B_{IJ} + \partial_J B_{KI}, \\ < H_{IJK} > &= \mathbf{c}\, M_{Pl}^3\, \epsilon_{IJK} \end{aligned} \tag{3.15}$$

(this time I, J, K are 3-dimensional complex coordinates describing the 6-dimensional compact manifold; ϵ_{IJK} is the completely antisymmetric tensor). The end result is that a superpotential is generated

$$W = \mathbf{c} + he^{-\frac{3S}{2b_0}} \tag{3.16}$$

where the second term takes its origin from gaugino condensation. This induces a non-trivial potential in the S direction which fixes $< S >$; $< T >$ is then determined at one loop. One finds as expected that both of these *vevs* are of order 1; more precisely one generates at most two orders of magnitude between $m_{3/2} = e^{K/2}|W|$ and M_{Pl}.

A detailed analysis of radiative corrections[16] in the model of subsection 2.1 shows that no soft supersymmetry breaking terms are generated either at tree level or at one loop in the observable sector (this includes gaugino and scalar masses). The reason for this has recently been identified[27]: it is the presence of a $SL(2, \mathbf{R})$ symmetry for the full Lagrangian[28]. More precisely,

$$\begin{aligned} T &\rightarrow \frac{aT - iB}{icT + d}, \quad ad - bc = 1, \\ \Phi^i &\rightarrow \frac{\Phi^i}{icT + d}, \quad a, b, c, d \in \mathbf{R}. \end{aligned} \tag{3.17}$$

This symmetry forbids to send the information of supersymmetry breaking from the *hidden* sector (where $m_{3/2} \neq 0$) to the observable sector. Luckily, this symmetry has an anomaly which breaks nonperturbatively the invariance. An analysis[27] based on effective Lagrangian methods[29] shows that a scale $\tilde{m}$ of global supersymmetry breaking is thus generated in the observable sector which can easily account for the 16 orders of magnitude between M_W and M_{Pl} ($\tilde{m} \approx M_W$).

Of course, the similarity between (3.13) and (3.17) suggests that the whole scheme might be general. One should note however that the symmetry that we have just discussed is continuous whereas duality is discrete. It would be very interesting to see in the light of all this how the discrete character of duality appears at the level of field theory[30]. Of possible relevance are the quantization conditions that appear in discussing supersymmetry breaking (for example, topological arguments indicate that **c** is quantized, which in turn forces $m_{3/2}$ to take only discrete values).

4 THE MATTER SCALARS. THE HIGGS.

We now consider the scalars of the observable sector.

4.1 The low energy scalar potential.

In the type of model of subsection 2.1, there are cases where the gauge group $\mathcal{H}$ is smaller than E_6: they correspond to models where part of the gauge symmetry is broken by the compactification process (the so-called breaking by Wilson loops)[31]. In such instances, the scalars that we consider are in representations of the gauge group $\mathcal{H} \subset E_6$ which are contained in 27 and $\overline{27}$ of E_6. To be explicit, we give in Table 1 the field content of a complete representation 27. Fields in $\overline{27}$ have opposite quantum numbers. As we will stress later, one could expect in principle fields which do not belong to these representations.

Looking at Table 1, one realizes that there are two types of fields which can be used to break $\mathcal{H}$ down to $SU(3) \otimes U(1)$:

– fields with the quantum numbers of a Higgs field under $SU(3) \otimes SU(2) \otimes U(1)$; we will refer to them generically as H; as in any supersymmetric model, there are two types of such fields, H_1 and H_2, coupling respectively to the u-type and d-type (upper and lower component of a weak doublet) quarks;

–fields singlet under $SU(3) \otimes SU(2) \otimes U(1)$ referred to as N. In Table 1, there are actually two types of such fields: one has allowed trilinear couplings to the Higgs, i.e.

$$W = \lambda H_1 H_2 N, \tag{4.1}$$

whereas the other one, which we will call N', has couplings to the Higgs forbidden by E_6 gauge symmetry.

There are of course other trilinear couplings involving the other matter fields as well as possibly non-renormalisable terms (quartic terms in the superpotential...). Disregarding these for a moment, we obtain the following supersymmetric potential (cf Eq.(2.3) where we ignore all non-renormalisable terms):

$$V = \lambda^2[|N|^2(|H_1|^2 + |H_2|^2) + |H_1|^2|H_2|^2] + \text{D}-\text{terms}. \tag{4.2}$$

		$SU(3)_c \otimes SU(3)_L \otimes SU(3)_R$	$SU(3)_c \otimes SU(2)_L$	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
$(1,\bar{3},3)$	$H_1 = \begin{pmatrix} H_1^+ \\ H_1^0 \end{pmatrix}$	or $E^c = \begin{pmatrix} E^+ \\ N_E^c \end{pmatrix}$	$(1,2)$	$-1/3$	$4/3$	$1/2$
	$H_2 = \begin{pmatrix} H_2^0 \\ H_2^- \end{pmatrix}$	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		$(1,1)$	$2/3$	$4/3$	$1/2$
	N'		$(1,1)$	$2/3$	$-2/3$	$-1/2$
	N		$(1,1)$	$2/3$	$-2/3$	0

Table 1 : 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 $E_6 (Y = Y_L + Y_R)$.

Of course, supersymmetry breaking is needed in order to generate non-zero masses for the scalars. As explained in the last section, supersymmetry breaking mechanisms should yield a scale $\tilde{m}$ of global supersymmetry breaking in the observable sector which is typically of the order of M_W. When local supersymmetry is broken in a hidden sector, the most general supersymmetry breaking terms which appear in the Lagrangian of the observable sector are of only three different types: gaugino mass, scalar mass and A-term (term in the potential which is proportional to the superpotential)[32]

$$\mathcal{L}' = \frac{1}{2} m_\lambda \bar{\lambda} \lambda - m_s^2 |\Phi|^2 + A[W(\Phi) + W(\Phi)^*]. \qquad (4.3)$$

Hence supersymmetry breaking generates new terms in the scalar potential:

$$V' = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_N^2 |N|^2 + m_{N'}^2 |N'|^2 - A_\lambda [\lambda H_1 H_2 N + h.c.]. \qquad (4.4)$$

In principle, the complete scalar potential for the H and N fields alone,[††] is given by Eqs.(4.2),(4.4). In these expressions however, one should replace all couplings by running

[††] Including only renormalisable terms; see below.

ones $m(\mu), A(\mu), \lambda(\mu) \dots$ i.e. include large logarithmic ‡‡ radiative corrections of order $\ln(M_{GUT}/\mu)$. The equations of evolution of these running couplings are the well-known Renormalisation Group Equations (RGE) whose explicit form can be computed once the spectrum and the couplings of the model are determined. The supersymmetry breaking contributions that we discussed in Section 4 appear actually at compactification and provide boundary conditions for these RGE:

$$m_1(M_{GUT}) = m_1^{(0)}, m_2(M_{GUT}) = m_2^{(0)}, m_N(M_{GUT}) = m_N^{(0)}, \cdots \tag{4.5}$$

where $m_1^{(0)}, m_2^{(0)}, \dots$ are all of order $\tilde{m}$.

As we go down in scale, we reach a region where some of the mass-squared in (4.4) become negative; some of the scalar fields acquire a non-zero *vev*, part of the gauge symmetry is broken: one refers to this as a *radiative breaking of gauge symmetry*[34]. Solving for the ground state of the μ-dependent scalar potential $V_0 = V + V'$ yields vacuum expectation values

$$< H_1 >= v_1(\mu), < H_2 >= v_2(\mu), < N >= n(\mu), < N' >= n'(\mu). \tag{4.6}$$

4.2 The different types of models.

Until now, we have neglected the nonrenormalisable terms that may arise as in any effective theory obtained by decoupling heavy modes. Let us consider the N field alone for a while.

Suppose that there exists a term of the type $(27 \times \overline{27})^n$ in the superpotential

$$W = \frac{\tilde{\lambda}}{M_{Pl}^{2n-3}} N^n \overline{N}^n, \quad n \geq 2, \tag{4.7}$$

where $\overline{N}$ is the field in $\overline{27}$ with opposite quantum numbers as N. Then the supersymmetric part of the potential reads

$$V = \left(n \frac{\tilde{\lambda}}{M_{Pl}^{2n-3}}\right)^2 (|N|^{2(n-1)}|\overline{N}|^{2n} + |N|^{2n}|\overline{N}|^{2(n-1)}) + D-terms. \tag{4.8}$$

The D-terms vanish for $N = \overline{N}$ since these fields have opposite quantum numbers. In this direction, after we have included the supersymmetry breaking part V' in (4.4),

$$V = m_N^2 N^2 + 2\left(n \frac{\tilde{\lambda}}{M_{Pl}^{2n-3}}\right)^2 |N|^{2(2n-1)}. \tag{4.9}$$

When m_N^2 becomes negative, N acquires a non-zero *vev*.

$$< N >= (\sqrt{-m_N^2} M_{Pl}^{2n-3})^{\frac{1}{2(n-1)}}. \tag{4.10}$$

‡‡In principle, quadratic divergences appear in the hidden sector where one breaks supersymmetry but these divergences are softened by the gravitational coupling between hidden and observable sector (which contributes a factor $1/M_{Pl}^2$) and only logarithmic divergences are generated in this sector. This is precisely the reason of breaking supersymmetry in a hidden sector.[33]

Since $\sqrt{-m_N^2} \approx \tilde{m} \approx M_W$, one notes that $< N >$ is larger than 10^{10}GeV and the gauge symmetry under which N is charged is broken at a scale intermediate between M_W and M_{Pl}. This is referred to as "*intermediate scale*" breaking[35]. If all N-type fields present in the model obtain such an intermediate *vev*, then we are left only with H fields at low energy.

Let us note that, so far, we have suppressed a generation index. In fact, in realistic models, we might have 3 generations for each type of fields: $H_1^i, i \leq 3, H_2^j, j \leq 3, \cdots$. It is however desirable that only one H_1 and one H_2 get a non-zero *vev*.Indeed, if there are too many H_1^i, H_2^j with non-zero *vevs*, one runs into problems with the evolution of the gauge coupling constant: M_{GUT} becomes too small. This problem is actually generic to all models with a minimal gauge and field content[36]. One sometimes call "unHiggs" the H fields which do not acquire a *vev*.[37]

The prototype of the models with only H fields has therefore only 2 "real" Higgs H_1 and H_2 which survive at low energy. It is usually referred to as the Minimal Supersymmetric Model (MSM). Out of the eight scalar degrees of freedom, two are eaten in the Higgs mechanism and the physical scalar fields are:

- 2 neutral Higgs h^0, H^0,
- 1 pseudoscalar A^0,
- 1 charged Higgs $H^{\pm}$.

The phenomenological aspects of this model will be treated by H.Haber[38]. Let me just say here that the low energy mass spectrum in the scalar sector depends on only 2 parameters which we can choose to be m_{H^+} and β, $\mathrm{tg}\beta = v_2/v_1$. One thus obtains relations between the different masses which can be turned into mass limits:

$$\begin{aligned} m_{H^+} \geq M_W,\ m_{H_0} &\geq M_Z,\ m_{A^0} \geq m_{h^0}, \\ m_{h^0} &\leq M_Z|\cos 2\beta|. \end{aligned} \tag{4.11}$$

For a large class of orbifolds however, one can show[39] that there are no such non-renormalisable terms as in 4.7. In this case, if there are N or N' fields in the massless spectrum, they will only acquire a nonzero *vev* in the Tev region. The minimal model of this type has 3 scalars: H_1, H_2 and N (models with N' have phenomenological problems of their own[40]). The spectrum after gauge symmetry breaking includes 3 neutral Higgs, 1 pseudoscalar and 1 charged Higgs.

Another possibility arises with the (numerous) models which have a gauge symmetry smaller than E_6. It may happen in this case that N or N' are present in the massless spectrum but are gauge singlets. An example of this is provided by the so-called flipped $SU(5) \otimes U(1)$ model[41]. This model was obtained in an effort to reproduce the grand unified $SU(5)$ model. It turns out that in string models it is very difficult to have the same representation in the gauge sector and in the scalar sector (as for the 24 of grand unified $SU(5)$). The string model which comes closest symmetry is precisely the flipped $SU(5) \otimes U(1)$ model: one is obliged to enlarge the gauge symmetry and to "flip" the $SU(5)$ assignments of quarks and leptons. In the TeV region, one is left with H_1, H_2 and N which is precisely a gauge singlet in this model. A further difference with the previous cases is that now a renormalisable N^3 term is allowed in the superpotential. The phenomenology of all these models will be discussed in J. Gunion and H. Haber's talks.

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