Superstring Phenomenology: An Overview

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The different aspects of superstring phenomenology before and after 1995 are briefly reviewed.

1. INTRODUCTION

String phenomenology has been under intense study for more than twelve years already. The main subtopics we can distinguish on this field are: model independent results, model building, low-energy effective actions and supersymmetry breaking. Here we will briefly mention the current status of each of these subtopics, comparing the situation before 1995 and after 1995.

2. SUPERSTRING PHENOMENOLOGY BEFORE 1995

Before 1995 there were five consistent string theories, all existing in ten dimensions, namely, type I open strings with gauge symmetry $SO(32)$, type IIA and IIB closed strings and the two closed heterotic strings with gauge symmetries $E_8 \otimes E_8$ and $SO(32)$. Out of these five theories, the one that attracted most of the attention was the heterotic $E_8 \otimes E_8$ theory because it was the most promising for phenomenology: upon compactification to 4D it gives rise to chiral $N = 1$ supersymmetric models with an observable sector, coming from the first $E_8$ which contains the standard model symmetry $SU(3) \otimes SU(2) \otimes U(1)$ and several families of matter fields. The second $E_8$ gives rise to a hidden sector, which fits perfectly with the attempts of supersymmetric model building prior to string theory, there a hidden sector was proposed that breaks supersymmetry at an intermediate scale $\sim 10^{12}$ GeV and gravity plays the role of messenger of supersymmetry breaking to the observable sector which feels the breaking of supersymmetry at the electroweak scale $\sim 10^2$ GeV.

The other four string theories were much less interesting, there was even a no-go theorem proving the impossibility to obtain the standard model out of type II theories.

At that time, we could only extract perturbative information in discussing phenomenological aspects of string theory because the theory itself was formulated only at the perturbative level. Therefore nonperturbative effects, such as those induced by the stringy versions of instantons, monopoles etc, were not under control. The issues of supersymmetry breaking and dilaton potentials needed to be addressed at the nonperturbative level and the only possibility was to consider field theoretical nonperturbative effects such as the condensation of gauginos in the hidden sector.

The tools used to study string phenomenology were threefold:

(i) Conformal field theory (CFT) with all its technology was heavily used to build ‘exact’ solutions of string theory giving rise to 4D string models with the phenomenologically desired properties, also CFT was applied to explicitly compute the important couplings such as gauge and Yukawa couplings directly from string theory and moreover, to extract general properties of string vacua which are independent of the model,
although are perturbative in nature.

(ii) Powerful topological techniques were used in the geometrical interpretation of string compactifications in terms of compact 6D Calabi-Yau manifolds, in order to construct compactifications, compute (cubic) Yukawa couplings and scalar kinetic terms and study the 'moduli space' of these solutions.

(iii) Finally, some 'macroscopic' techniques were used in order to extract information about the 4D effective action. Mostly based on general symmetry arguments. This allowed to prove remarkable results such as the non-renormalization of the superpotential and of the gauge kinetic function beyond one-loop.

The combination of these three different approaches allows us to have a general understanding of model building, effective actions and model independent results in string perturbation theory, as well as supersymmetry breaking from non-perturbative field theoretical effects. We will briefly discuss these subjects here, for more complete treatments including the relevant references we refer the reader to [1], [2], [3]. In particular the next subsections are based on the long review [4] and I will refer to this paper not only for a longer exposition but also for a detailed guide to the references. See also [5].

2.1. Model Building

The geometrical approach of 4D string model building can be described à la Kaluza-Klein in the sense of viewing the 10D spacetime as a product of 4D Minkowski spacetime and a small 6D compact manifold. This manifold is constrained by the requirement of having \( N = 1 \) supersymmetry to be a Calabi-Yau manifold, which is a 3D complex space with holonomy \( SU(3) \). There are many such string 'vacua', first because of the many possible Calabi-Yau manifolds, which are not completely classified yet, but also because, for each compactification, we still have the freedom of choosing the embedding in the gauge degrees of freedom (of the heterotic string for instance). The particular embedding identifying the spin connection and the gauge connection gives rise to a \( (2, 2) \) supersymmetric CFT, where the entries refer to the number of supersymmetries for left and right moving string modes, respectively. The generic embedding however has only \( (0, 2) \) supersymmetry. The \( (0, 2) \) Calabi-Yau compactifications are not yet well understood and most of the work on this area has been done on the \( (2, 2) \) models.

The simplest way to construct \( (0, 2) \) models is by constructing explicit CFTs such as orbifold compactifications. An orbifold is a twisted torus that happens to be a particular singular limit of a Calabi-Yau manifold. The advantage of working with orbifolds is that, similar to flat space or torus compactifications they correspond to free CFTs, their only complication relies in the choice of boundary conditions which is understood. Similarly, by either bosonizing or fermionizing all the coordinates of the extra space (the orbifold) and choosing different boundary conditions for each of the fields we arrive at the 'fermionic' and 'lattice' constructions of 4D models. Furthermore, there are more complicated constructions based on coset CFTs and known as Gepner-Kazama-Suzuki models, they attracted substantial attention especially after the realization that the \( (2, 2) \) versions of these models were also particular points of a Calabi-Yau manifold despite their original nongeometrical formulation.

Out of all these different formulations of models we learned that there is a twofold degeneracy of string models. One is discrete, given by say topologically different Calabi-Yau compactifications with different numbers of generations and antigenerations of quarks and leptons (for (2, 2) Calabi-Yau's these numbers are topological, in particular the net number of generations minus antigenerations is given by the Euler number of the manifold.). The second degeneracy is continuous, i.e. for each compactification, there are a number of free parameters or 'moduli' which can be varied freely, these moduli can describe the size and shape of the internal space and in the 4D effective theory they correspond to massless scalar fields \( 'T' \), having vanishing potential. Therefore the continuous degeneracy is given by the moduli space of the \( T \) fields defined by the existence of flat directions. Due to the existence of non-renormalization theorems, the flatness of
the potential stays to all orders in perturbation theory.

This enormous degeneracy limits the predictive power of the theory. We might expect that the continuous degeneracy may be lifted by the existence of non-perturbative string effects which could create a nonvanishing scalar field potential, but the discrete degeneracy seems more serious. Before 1995 it seemed impossible to be able to compare two different compactifications differing for instance on the total number of generations since, having topological origin, it was impossible to deform one space into another piece of information we can extract about this degeneracy of string models is by taking a phenomenological approach and search among all of these models those which resemble the standard model of particle physics with 3 families of quarks and leptons and standard model gauge group \( SU(3) \otimes SU(2) \otimes U(1) \).

Some quasi-realistic models have been studied with much detail using this approach. This includes models with three families and standard model gauge group \( SU(3) \otimes SU(2) \otimes U(1), SU(4) \otimes SO(4) \) as well as a version of \( SU(5) \otimes U(1) \) known as flipped SU(5). All of these models also reproduce many of the nice features of the standard model (proton stability, doublet-triplet splitting, pattern of fermion masses etc.). Nevertheless, as in the case of orbifolds, there is not a totally realistic model yet. In particular there is no model yet with just the spectrum of the supersymmetric standard model.

An unsolved issue in most of these models is the fact that after breaking to the standard model group, the number of extra light doublets is not well determined since their mass is not generally protected by symmetries and the actual calculation of the remaining superpotential, after solving the \( D \) and \( F \) flatness conditions, is done only up to operators of a given dimension. The presence of extra light doublets in orbifold models was understood in terms of the selection rules established for those models, but they are precisely the main source of problems for those models also, because of their contribution to the running couplings and therefore to the string unification scale and Weinberg’s angle. The extra light doublets is then an obstacle for obtaining realistic models in all the different approaches studied so far. Furthermore, without understanding the breaking of supersymmetry, we cannot confront directly these models with physics.

The more recent progress in this area was the successful construction of 4D string GUTs, i.e. string models with simple Grand Unified groups (such as \( E_6 \) and \( SO(10) \)) and an adjoint representation that can break the group to the standard model. The presence of adjoints requires special string constructions named ‘higher level Kac-Moody models’ which are a particularly difficult- to-construct subclass of the \((0,2)\) string models discussed above. These constructions are expected to share the phenomenological virtues of supersymmetric GUTs, especially in what refers to the unification of the gauge couplings. The models obtained so far have not been studied in great detail but they still suffer from some difficulties, such as fine tuning and the possible survival of extra states at low energies that could modify the running of the gauge couplings. Therefore it is still an open question to construct a fully realistic string model.

### 2.2. Effective Actions

The general Lagrangian coupling \( N = 1 \) supergravity to gauge and chiral multiplets depends on three arbitrary functions of the chiral multiplets:

1) The \( \text{Kähler potential} \ K(z, \bar{z}) \), which is a real function. It determines the kinetic terms of the chiral fields

\[
\mathcal{L}_{\text{kin}} = K\bar{z}\partial_{\mu}\bar{z}\partial^{\mu}z
\]

with \( K_{\bar{z}z} \equiv \partial^2 K/\partial z\partial\bar{z} \). \( K \) is called Kähler potential because the manifold of the scalar fields \( z \) is Kähler, with metric \( K_{\bar{z}z} \).

2) The superpotential \( W(z) \) which is a holomorphic function of the chiral multiplets (it does not depend on \( \bar{z} \)) \( W \) determines the Yukawa couplings as well as the \( F \)-term part of the scalar potential \( V_F \):

\[
\epsilon^{K/M_p^2} \left\{ D_{\bar{z}}W K_{\bar{z}z}D_{z}\bar{W} - 3 \frac{|W|^2}{M_p^2} \right\},
\]

where

- \( D_{\bar{z}} \) denotes the covariant derivative,
- \( K_{\bar{z}z} \) is the Kähler metric,
- \( W \) is the superpotential,
- \( M_p \) is the Planck mass,
- \( |W|^2 \) is the squared superpotential,
- \( \epsilon \) is a running coupling constant.
with $D_z W \equiv W_z + W K_z / M_p^2$. Here and in what follows, the internal indices labelling different chiral multiplets $z_i$ are not explicitly written.

The gauge kinetic function $f_{ab}(z)$ which is also holomorphic. It determines the gauge kinetic terms

$$\mathcal{L}_g = \text{Re} f_{ab} F^a_\mu F^{\mu b} + \text{Im} f_{ab} F^a_\mu \tilde{F}^{\mu b} \quad (3)$$

it also contributes to gaugino masses and the gauge part of the scalar potential.

$$V_D = (\text{Re} f^{-1})_{ab} (K_z, T^a z) \left( K_S, T^b \tau \right)$$
$$V = V_F + V_D \quad (4)$$

These three functions are arbitrary for a generic $N = 1$ supersymmetric model, but in string theory we should be able to compute them for each model. General non-renormalization theorems can be applied to the holomorphic functions $W$ and $f$. The main reason is that the dilaton field $\phi$ is always present in string models and its vev is the loop counting parameter. In $N = 1$ 4D models it joins an axion-like field $a$ having a Peccei-Quinn symmetry $a \to a + \text{constant}$, to form a complex (chiral) multiplet $S = \phi + i a$, because of the existence of the Peccei-Quinn symmetry $W$ cannot depend on $a$ and, because $a$ only appears through $S$, then $W$ cannot depend on $S$ and so it does not depend on the loop-counting parameter, therefore it is not renormalized. A similar argument applies to $f$ beyond one loop. The Peccei-Quinn symmetry is usually broken by nonperturbative effects therefore they will contribute to corrections to $W$ and $f$. For $K$ there are no simplifications and very little is known beyond some tree level calculations.

The general structure of these functions of matter multiplets $Q$, moduli $T$ and dilaton $S$ is the following:

$$W(S, T, Q) = W_{\text{tree}}(T, Q) + W_{np}(S, T, Q)$$
$$f(S, T, Q) = S + f_{1-\text{loop}}(T, Q) + f_{np}(S, T, Q)$$

$$K(z, \tau) = (K_{\text{tree}} + K_{1-\text{loop}} + \cdots) + K_{np} \quad (5)$$

Therefore we can see that the lack of control on the perturbative corrections to $K$ is the main source of ignorance of the full perturbative 4D effective actions.

2.3. Model Independent Results

Let us recapitulate in this section what we can say about string theories which is independent of the model. This is the closest we can get to string predictions and help us in approaching general questions, differentiating the generic issues from those of a particular model. Since the full non-perturbative formulation of string theory is not yet available, we have to content ourselves mostly with predictions of string perturbation theory, assuming that the corresponding string model is given by a CFT.

(i) String models predict the existence of gravity and gauge interactions. This is a point that cannot be overemphasized since it is the first theory that makes those fundamental predictions for interactions we experience in the everyday life.

(ii) The dimension of spacetime is dynamical and $D \leq 10$ raising the hope that eventually we could explain if a 4D spacetime is in some way special, although at present. Also the rank of the gauge group is bounded $r \leq 22$

(iii) There are other fields which survive at low energies: charged matter fields $Q$, candidates to be basic building blocks of matter but also the dilaton field $S$ and the moduli $T$. We have to mention that, although as yet there is no 4D model without moduli fields, there is not a general theorem implying their existence. In that sense the dilaton is the most generic modulus field, with a flat potential in perturbation theory.

(iv) There is only one arbitrary parameter $\alpha'$ fixed to be close to the Planck scale $M_p$. All other parameters of the effective action are determined by expectation values of fields such as the dilaton and the moduli. In particular the gauge coupling is given at tree level by the vev of $S$. 
(v) The existence of spacetime supersymmetry is needed for consistency, although $N = 1$ is selected for phenomenological reasons. There is a general requirement for a CFT to lead to $N = 1$ spacetime supersymmetry: It has to have $(0, 2)$ supersymmetry in the worldsheet (2D) (plus a quantization condition on the charges of the $U(1)$ group mixing the two supersymmetries).

(vi) There are no global internal symmetries in 4D string models, besides the already mentioned Peccei-Quinn symmetry of the $S$ field and some accidental global symmetries (like baryon and lepton numbers in the standard model). This puts very strong constraints to string models compared with standard field theory models.

(vii) There are generically some discrete symmetries in string models. Some infinite dimensional such as $T$-duality which in the simplest version takes the form of an $SL(2, \mathbb{Z})$ transformation

$$T \rightarrow \frac{aT + b}{cT + d}$$

with $a, b, c, d$ integers satisfying $ad - bc = 1$.

There are also finite dimensional discrete symmetries, such as those inherited from the twist defining orbifold constructions, which are seen as discrete gauge symmetries in the 4D effective theory. These can in principle be useful for model building, hierarchy of masses etc. There are however some couplings that vanish in string theory and cannot be explained in terms of symmetries of the effective 4D theory, these are called ‘string miracles’ since from the 4D point of view they seem to break the criterium for naturalness. $T$-duality symmetries restrict very much the form of the effective action and quantities such as Yukawa couplings have to be modular forms of a given duality group. These symmetries are valid to all orders in string perturbation theory and are thought to be also preserved by nonperturbative effects. Matter fields $Q^I$ are assigned special quantum numbers, the modular weights $n$, according to their transformation properties under the duality group. For a $SL(2, \mathbb{Z})^m$ group for instance we have:

$$Q^I \rightarrow (ic T_1 + d_I)^{n_I} Q^I, \quad I = 1, \ldots, m.$$  (7)

Since fermions transform nontrivially under these symmetries, there may be ‘duality anomalies’ which have to be cancelled for consistency. This imposes strong constraints on the possible spectrum of the corresponding string model.

(viii) There is unification without the need of a GUT. If the gauge group is a direct product of several groups we have for the heterotic string:

$$k_1 g_1^2 = k_2 g_2^2 = \cdots = \frac{8\pi}{\alpha'} G_N \equiv g_{\text{string}}^2.$$  (8)

Where $k_i$ are special stringy constants known as the Kac-Moody levels of the corresponding gauge groups (for the standard model groups it is usually assumed that $k_2 = k_3 = 1, k_1 = 5/3$), $g_i$ are the gauge couplings and $G_N$ is Newton’s constant. We can see there is a difference with standard GUTs in field theory for which we compute the unification scale by finding the point where the different string couplings meet. In heterotic string theory, the unification scale is given in terms of the string coupling $g_{\text{string}}$ and the Planck scale $\pi$. More precisely: $M_{\text{string}} \approx 5.27 \times 10^{17} g_{\text{string}} \text{ GeV}$. For $g_{\text{string}} \sim O(1)$ this shows a discrepancy with the ‘observed’ value of the unification scale given by the experiments $M_{\text{GUT}} \sim 2 \times 10^{16} \text{ GeV}$. Also the Weinberg angle gives $\sin^2 \theta_W = 0.218$ differing from the experimental value of $\sin^2 \theta_W = 0.233 \pm 0.0008$.

Therefore the string ‘predictions’ are very close to the experimental value, which is encouraging, but differ by several standard

\footnote{For type I strings the gravitational and gauge couplings are independent, so we have the freedom to adjust the unification scale with experiment as in usual GUTs.}
deviations from it. This is the string unification problem. The situation looks much better for simple GUT’s which have good agreement with experiment. Several ideas have been proposed to cure this problem, including large values of threshold corrections, intermediate scales, extra particles, changing the values of Kac-Moody levels etc, with no compelling solution yet (for recent discussions of this issue see for example the review of K. Dienes in ref. [8]).

(ix) There are usually fractionally charged particles in 4D string models. In fact it can be shown that we cannot have simultaneously \( k_2 = k_3 = 3k_1/5 = 1 \) in the standard model and only integer charged particles, because if that is the case the standard model gauge group would be enhanced to a full level-one \( SU(5) \). This ‘problem’ can be evaded in models where the fractionally charged particles are heavy string states, it has also been proposed that those particles could confine at intermediate energies and be unobservable.

(x) There are ‘anomalous’ \( U(1) \) groups in most of the models, but there is also a counterterm in the action cancelling the anomaly and generating a Fayet-Iliopoulos kind of term:

\[
\frac{1}{S + \bar{S}} \left| \frac{Trq^a}{48\pi^2} \frac{1}{(S + \bar{S})^2} + \sum q^a_i |Q_I|^2 \right|^2, \tag{9}
\]

where \( q^a_i \) are the anomalous charges of the scalar fields \( Q_I \). This term is responsible to break the would be anomalous group by fixing the value of a combination of the matter fields \( Q_I \), breaking the would be anomalous \( U(1) \) and usually other gauge groups, but not supersymmetry (although this has not been shown in general). A combination of the fields \( Q_I \) and the dilaton \( S \) still remains massless and plays the role of the new dilaton field.

There are further model independent results which refer to nonperturbative string effects and will be discussed next.

### 2.4. Supersymmetry Breaking

As we discussed previously there are two main problems of string perturbation theory namely, the enormous vacuum degeneracy and supersymmetry breaking. There were two lines of research towards attacking these problems.

(i) To use a particular field-theoretical nonperturbative effect which is the condensation of gauginos induced after the asymptotically free hidden sector becomes strongly interacting at lower energies (∼ \( 10^{12} \) GeV).

(ii) To consider the supersymmetry breaking sector as a black box and study its possible implications assuming that a combination of the moduli fields \( T \) and the dilaton \( S \) are responsible for breaking supersymmetry, but without specifying how.

The second approach has been recently reviewed in [3] and I refer the reader to that review for the general results. As for gaugino condensation, it was also reviewed in [4],[10], we will only mention here that the combination of \( T \) duality with the inclusion of several condensing groups with matter in the hidden sector (very generic in string theory) can give rise to interesting results, namely, supersymmetry is broken at the phenomenologically desired scale (∼ \( 10^2 \) Gev), with the moduli \( T \) and dilaton \( S \) fixed at the desired values \( T \sim 1, S \sim 2 \) in Planck scale units. This value of \( S \) gives the expected value of the gauge coupling at the unification scale. However there were several problems:

(i) The cosmological constant is very large and negative.

(ii) The dilaton potential is runaway and there has to be some fine tuning in order to obtain a nontrivial minimum besides the minimum at zero coupling and \( S \rightarrow \infty \).

(iii) There are at least two serious cosmological problems for the gaugino condensation scenario. First, it was found under very general grounds, that it was not possible to get inflation with the type of dilaton.
potentials obtained from gaugino condensation \[1\]. Second is the so-called ‘cosmological moduli problem’ which applies to any (non-renormalizable) hidden sector scenario including gaugino condensation \[12\]. In this case, it can be shown that if the same effect that fixes the vev’s of the moduli, also breaks supersymmetry, then: the moduli and dilaton fields acquire masses of the electroweak scale (\(\sim 10^2\) GeV) after supersymmetry breaking \[12\]. Therefore if stable, they overclose the universe, if unstable, they destroy nucleosynthesis by their late decay, since they only have gravitational strength interactions. At present there is no satisfactory explanation of this problem and it stands as one of the unsolved generic problems of string phenomenology.

The runaway behaviour of the dilaton has been argued to be a generic problem for string models \[13\]. The reason for this is that being \(S\) the string coupling, we know that for \(S \rightarrow \infty\) the theory becomes free and then the scalar potential has to vanish. This was used in \[13\], to argue that strings have to be strongly coupled in order to develop another minimum, unless some parameters conspire to fix \(S\) at weak coupling. This argument has been revised recently in \[14\]. There it was argued that even if the scalar potential vanishes at \(S \rightarrow \infty\) in the full theory, in an effective theory, after integrating out some of the massive fields the remaining potential for \(S \rightarrow \infty\) could blow up for \(S \rightarrow \infty\). A simple example of this is the \(\lambda \phi^4\) theory since upon minimizing the potential

\[ V = -\frac{1}{2}m^2 \phi^2 + \frac{1}{4}g^2 \phi^4 \]

we find \(\phi = \pm m/g\) and substituting back into \(V\) gives \(V(\phi) = -m^4/4g^2\) which blows up as the coupling \(g \rightarrow 0\). A particular mechanism to achieve this could be to consider nonasymptotically free models which appear very often in string theories. If the potential blows up at \(S \rightarrow \infty\), then the cosmological problems discussed in \[11\] are not necessarily there. Therefore this is an interesting way of fixing the dilaton that deserves further investigation.

All this has been said using mostly nonperturbative field theoretical effects but we know that there should also be stringy nonperturbative effects that could play a role in fixing the dilaton and breaking supersymmetry. In general we should always consider the two types of nonperturbative effects: stringy (at the Planck scale) and field theoretical (like gaugino condensation). Four different scenarios can be considered depending on which class of mechanism solves each of the two problems: lifting the vacuum degeneracy and breaking supersymmetry.

For breaking supersymmetry at low energies, we expect that a field theoretical effect should be dominant in order to generate the hierarchy of scales (it is hard to believe that a nonperturbative effect at the Planck scale could generate the Weinberg-Salam scale). We are then left with two preferred scenarios: either the dominant nonperturbative effects solve both problems simultaneously, or there is a ‘two steps’ scenario in which stringy effects dominate to lift vacuum degeneracy and field theory effects dominate to break supersymmetry. The first scenario has been the only one considered so far, it includes gaugino condensation. The main reason this is the only scenario considered so far is that we can control field theoretical nonperturbative effects but not the stringy. In this scenario, independent of the particular mechanism, we have to face the cosmological moduli problem. In the two steps scenario stringy effects only lift the vacuum degeneracy and supersymmetry may be broken dynamically by field theoretical effects, such as discussed in a whole session at this conference \[15\]. We have to conclude that there is no prefer scenario yet.

3. SUPERSTRING PHENOMENOLOGY AFTER 1995

The recent progress in understanding nonperturbative issues of string theory \[16\], \[17\] has necessarily strong impact on the phenomenological questions, we are only starting to explore these implications which can be summarized as follows.

(i) **Unification of theories:** We mentioned in the introduction that there are five consistent superstring theories and each has thousands or millions of different vacua. It is now believed that the five string theories are related by strong-weak coupling dual-
ities and furthermore, they appear to be different limits of a single underlying fundamental theory, probably in 11D, the $M$ theory (probably related with membranes or higher dimensional objects such as five-branes), which is yet to be constructed. If this is true it may solve the arbitrariness in the number of fundamental string theories by deriving them from a single theory.

(ii) Unification of vacua (?): Recent work based on comparison of string compactifications with the Seiberg-Witten theory, has lead to the conclusion that many and probably all Calabi-Yau compactifications are connected. Then it seems that not only the five different theories are unified, but also all the vacua of these theories could also be unified: since, if they are all connected, we can foresee a mechanism that lifts the degeneracy and select one point in the set of compactifications, something it could not have been done before because they were thought to be disconnected vacua. These transitions occur in singular points of the corresponding moduli space where a particular state (massless black hole) or even an infinite tower of states (tensionless strings) become massless. They were partially understood for $N > 1$ compactifications, but recently extensions to the phenomenologically interesting $N = 1$ case have been found,[13] implying for instance that models with different number of families would belong to the same moduli space reducing in some sense the discrete degeneracy problem to the level of the continuous degeneracy problem and so we may expect that probably one particular number of families could eventually be selected dynamically.

(iii) Nonperturbative vacua: The fact that the strong coupling regime of a given string theory would simply be the weak coupling regime of another string theory would be very disappointing since that means that the problems present at weak coupling would remain at strong coupling. Fortunately this is not the case. For instance, the strong coupling limit of the $E_8 \otimes E_8$ string is believed to be given by $M$ theory compactified in the orbifold $S^1/\mathbb{Z}_2$ which is just a one dimensional interval. $M$ theory contains elementary membranes and their magnetic dual, 5—branes. The membranes can end at each of the two 10D ends of the interval (fixed points) which are 9—branes and generate an $E_8$ symmetry at each end. The distance between the two 9—branes $\rho$ is proportional to the heterotic coupling and when this is very small the two $E_8$'s collapse to a single 10D point which is the heterotic string. For any finite coupling the membrane is a cylinder between the two 9—branes with heterotic strings at the intersection. This reproduces the standard perturbative spectrum of heterotic strings. The new ingredient comes mostly from the 5—branes which for the $E_8 \otimes E_8$ case, carry two-index antisymmetric tensors, therefore introducing more than one of these fields in the spectrum after compactifications (in the $SO(32)$ version they may lead to extra vector fields depending on the compactification). In perturbative heterotic string there was a single antisymmetric tensor $B_{\mu\nu}$ that we saw is dual to an axion field. The appearance of several of those fields in the spectrum shows clearly that the corresponding vacuum is nonperturbative and may eventually create more possibilities for using these axion fields for solving the strong CP problem in string theory. There is even a model with zero tensor fields. This may be relevant because $B_{\mu\nu}$ is a supersymmetric partner of the dilaton and having a model without antisymmetric tensors would mean that somehow the dilaton was fixed, lifting the corresponding degeneracy, and acquired a mass (avoiding the cosmological moduli problem)! Furthermore for compact spaces with nontrivial 4—cycles, the corresponding 5—branes could wrap around those cycles giving rise to another string (different of course from the one obtained from the membrane). These nonperturbative strings...
will generically have their own nonperturbative gauge group, therefore enhancing the maximum rank required in perturbation theory \[1\] (the world record seems to be right now a group of rank of order \(10^9\) \[21\]). The physical relevance of the nonperturbative gauge fields is yet to be explored.

(iv) Scales in \(M\) theory: It is interesting to analyze the different scales present in a 4\(D\) model built from \(M\)-theory. There are three relevant scales: the 11\(D\) Planck scale \(\kappa\), the length of the interval \(\rho\) and the overall volume of the compactified 6\(D\) space \(V\). In the 11\(D\) theory, the gauge and gravitational couplings can be written as:

\[
L = -\frac{1}{2\kappa^2} \int_{M^{11}} d^{11}x \sqrt{g} R - \sum_i \frac{1}{8\pi(4\pi\kappa)^2}{2/3} \int_{M^{10}_i} d^{10}x \sqrt{g} \text{tr} F_i^2. \tag{10}
\]

Where \(M^{11}\) is the 11\(D\) space (bulk) and \(M^{10}_i, i = 1, 2\) are the two 10\(D\) 9-branes at each end of the interval. We can see that after compactification, the 4\(D\) Newton constant and gauge couplings are given by

\[
G_N = \frac{\kappa^2}{8\pi^2 \kappa^2 \rho^2} \quad \text{and} \quad \alpha_{\text{GUT}} = \frac{(4\pi\kappa^2)^{2/3}}{2\pi^2}. \]

Notice that now \(M^2_{\text{GUT}} = V^{-1/3} = \frac{\alpha_{\text{GUT}}}{8\pi^2 G_N^2 \rho^2}\), since we have an extra parameter, \(\rho\), we can get \(M_{\text{GUT}} \sim 10^{16}\text{GeV}\) by setting \(\rho^{-1} \sim 10^{12-14}\text{GeV}\) something we could not have done in perturbative heterotic strings. This has been used by Witten to claim that it may be possible to solve the string unification problem by tuning the extra parameter as in standard GUTs \[21\]. We then get the following picture: at large distances the universe looks 4\(D\) at energy scales between \(10^{12-14}\text{GeV}\) and \(10^{16}\text{GeV}\) it looks 5\(D\) and at higher scales (smaller distances) it looks 11\(D\). This new intermediate scale \(\rho\) may play an interesting role for other phenomenological and cosmological questions. There is a complication that for \(\rho^{-1} \leq 10^{16}\text{GeV}\) the gauge coupling of one of the gauge groups blows up, this has been argued by Witten that could put a bound on Newton’s constant on a generic model. There are some specific models which avoid this problem which makes them more attractive. Also, the process of gaugino condensation can be reanalyzed in this picture \[22\]. \[24\]. A single condensate in the hidden \(E_8\) 9-brane, does not break supersymmetry in its vicinity nor in the 5\(D\) bulk but due to a topological obstruction it can break supersymmetry in the observable sector \[22\]. Note that in this picture the standard model lives at one of the ‘end of the world’ branes while gravity and the moduli fields live in the 5\(D\) bulk. The possible physical consequences of this new picture are only starting to be explored \[21\].

(v) Nonperturbative superpotential: It is quite remarkable that recently Witten and others have been able to extract information about superpotentials derived from stringy nonperturbative effects \[24\]. At the moment there have been found three classes of results, depending on the compactification:\(W = 0\), \(W \sim e^{-\Phi}\), \(W\) a modular form. Here \(\Phi\) is one of the moduli fields. The first case is interesting because it means there are compactifications for which the nonperturbative superpotential vanishes so the only source of superpotential could be strong coupling infrared effects such as gaugino condensation making the field theoretical discussion above more relevant. The second case gives the standard runaway behaviour of the scalar potential and the third possibility is a realization of the kind of duality invariant potentials proposed in the past \[25\]. \[16\] in this case there are nontrivial minima and it is yet to be studied in detail whether supersymmetry could be broken, in particular these models seem suitable for a realization of the two steps scenario alluded to before. We hope more progress will be made in this direction which is addressing the main problem of superstring phenomenology from a nonperturbative formulation.
(vi) Stringy $e^{-1/g}$ effects: Some time ago, Shenker proposed that in string theory, there would appear nonperturbative effects of the form $e^{-1/g}$ on top of the standard field theoretical effects of the form $e^{-1/g^2}$. These have been argued to correct the Kähler potential and contribute to the dilaton potential in such a way that the dilaton can be fixed even with a single exponential in $W$. Recently, these effects were explicitly computed for the heterotic string for a particular compactification.

We can see that many of the results from string perturbation theory are modified by the nonperturbative information obtained so far. Some of the other results are expected to be modified or need revision, for instance the nonexistence of global symmetries was proved using CFT techniques which are explicitly perturbative, it is expected that being string theory a theory of gravity, global symmetries will not be allowed (as usually found studying black holes and wormholes), but a general nonperturbative proof is not available yet. Also, the main problems such as supersymmetry breaking, are still open which is a good motivation to work on this field.

REFERENCES

1. M. Green, J. Schwarz and E. Witten, *Superstring Theory, volumes 1,2*, Cambridge University Press (1987); M. Kaku, *String Theory* Springer-Verlag (1988); D. Lüst and S. Theisen, *Lectures in String Theory*, Springer Lecture Notes in Physics, Vol. 346 (1989); P. Ginsparg, *Les Houches Lectures Elsevier* (1989); J. Polchinski, *Les Houches Lectures*, hep-th/9411028 and book to appear.

2. B. Schellekens, *Superstring Construction*, North-Holland (1989).

3. M. Dine, *String Theory in Four Dimensions*, North-Holland (1988).

4. F. Quevedo, hep-th/9603074 and references therein.

5. P. Nath and G. Kane contributions to this conference.

6. For a recent review see: A. Farggi, hep-ph/9707311.

7. Z. Kakushadze, contribution at this conference.

8. K. Dienes, hep-th/9602045 and references therein.

9. A. Brignole, L. Ibáñez and C. Muñoz, hep-ph/9707203.

10. F. Quevedo, hep-th/9511131.

11. R. Brustein, P. Steinhardt, *Phys. Lett.* B302 (1993) 196.

12. T. Banks, D. Kaplan, A. Nelson, *Phys. Rev.* D49 (1994) 779B; de Carlos, J.A. Casas, F. Quevedo and E. Roulet, *Phys. Lett.* B318 (1993) 447.

13. M. Dine and N. Seiberg, *Phys. Lett.* B162 (1985) 299.

14. C.P. Burgess, A. de la Macorra, I. Maksymyk and F. Quevedo, hep-th/9707062.

15. See the contributions of Nelson, Dine, Poppitz and Murayama at this conference.

16. A. Font, L.E. Ibáñez, D. Liist, F. Quevedo, *Phys. Lett.* 249B (1990) 35; S.-J. Rey, *Phys. Rev.* D43 (1991) 526; C. Hull and P.K. Townsend, *Nucl. Phys.* B438, (1995) 409; E. Witten, *Nucl. Phys.* B443 (1995) 85.

17. For recent reviews see: J. Polchinski, hep-th/9607050; J. Schwarz, hep-th/9607201; P.K. Townsend, hep-th/9612121; C. Vafa, hep-th/9702201.

18. S. Kachru and E. Silverstein, hep-th/9704185; G. Aldazabal, A. Font, L.Ibáñez, A. Uranga and G. Violero, hep-th/9706158.

19. E. Witten, hep-th/9507121.

20. P. Candelas, E. Perevalov and G. Rajesh, hep-th/9704097.

21. E. Witten, hep-th/9602070; T. Banks and M. Dine, hep-th/9605136.

22. P. Horava, hep-th/9608019.

23. I. Antoniadis and E. Dudas, contributions to this conference.

24. E. Witten, *Nucl. Phys.* B474 (1996) 343. See also the contribution of D. Liist to this conference.

25. S. Ferrara, D. Liist, A. Shapere and S. Theisen, *Phys. Lett.* B225 (1989) 363; E. Chun, J. Mas and H.P. Nilles, *Phys. Lett.* 233B (1989) 141.

26. Y.-Y. Wu, contribution to this conference.

27. E. Silverstein, hep-th/9611193.