New Perspectives in String Phenomenology from Dualities

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Abstract

After a review of some topics concerning the phenomenological applications of perturbative string theory, I discuss to what extent all of it is affected by the recent developments in string dualities.

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1 Introduction

The excitement of the theoretical physics community after the so-called first string revolution in 1984 had in principle a phenomenological basis. The cancellation of anomalies for the $D = 10$, $N = 1$ superstring theories with gauge symmetries made in principle possible the idea of a unification of a chiral gauge theory like the standard model (SM) with gravity into a finite theory. String phenomenology is just the study in detail of that general idea: how is embedded the SM (or perhaps the minimal supersymmetric standard model MSSM) into string theory? A good deal of effort has been devoted to this field in the last ten years and some of its results have inspired other string theory areas as well as physics beyond the standard model in general. After the second (1995) string revolution (for reviews see ) a natural question appears. How is the general scheme of string phenomenology affected by the discovery of the non-perturbative string dualities? It is possibly too soon to make any definite statement. In spite of this, some qualitative features of the new physics appearing in the new non-perturbative and perturbative string vacua can already be extracted.

2 Some general aspects of perturbative string phenomenology

Circa 1985 the scheme for embedding the known standard model (SM) interactions within the string scheme were thought by many practitioners to be relatively simple and unique. If one starts from the $E_8 \times E_8$ heterotic string and compactifies on a Calabi-Yau manifold one lands on an $N = 1$ theory with gauge group $E_6 \times E_8$. Such a construction gives rise to a number of massless chiral $E_6$ generations given by $n_g = \chi/2$, where $\chi$ is the Euler characteristic of the Calabi-Yau manifold. It was soon realized that 1) it is not clear whether such an $E_6$ structure is phenomenologically viable (problems with unwanted extra massless matter, proton stability, neutrino masses, gauge coupling unification etc.); 2) there are many, many more ways to construct consistent $N = 1$, $D = 4$ string vacua leading to a variety of gauge groups (including directly the SM group) and massless particle content; 3) the $SO(32)$ heterotic is equally good from the point of view of model-building. In the last ten years four-dimensional string models based on toroidal orbifolds, free-fermion 4-D strings or Gepner type of models have been constructed. Some of them
have a massless spectrum tantalizingly close to the particle content of the minimal
supersymmetric standard model (MSSM). I think that this, by itself, is already an
achievement since they represent the first unified theories of all interactions including
gravity. Two broad classes of $D = 4$, $N = 1$ string models have been constructed,
those which involve heterotic constructions where the algebra of the gauge group is
realized at Kac-Moody level $k = 1$ and those with $k > 1$. As we will recall below,
the first class of models (by far the most studied up to now) gives rise to theories
without adjoint Higgs fields and hence one gets models with gauge groups like e.g.,
$SU(3) \times SU(2) \times U(1)^n$, $SU(4) \times SU(2) \times SU(2)$ or $SU(5) \times U(1)$ in which one can
make the symmetry breaking down to the SM without any adjoint Higgs. A number
of three generation models of these characteristics have been constructed (see e.g.,
[6, 7, 8, 9]). Models with $k > 1$ have only been considered in the last five years
or so [10]. Three and four generation models with the massless spectra of $SU(5)$,
$SO(10)$ or $E_6$ GUTs have been constructed.

In spite of the above successes, all of the realistic perturbative string vacua
constructed up to now have the general property of yielding extra unwanted massless
chiral fields beyond those present in the MSSM. One has to abandon the string
techniques and analyze the effective field theory. Then one has to assume that, after
SUSY is broken, some particular direction in the scalar field space is taken so that
(via Yukawa couplings) all of the unwanted massless particles disappear from the
low energy spectrum. Important phenomenological properties like the quark-lepton
spectrum and proton stability depend on the choice made for the pattern of gauge
symmetry breaking. That, of course, leads to a partial loss of predictivity.

There are a number of properties of the above perturbative heterotic vacua which
are general and appear in all different classes of constructions. They are particularly
interesting because they may be considered as generic predictions of perturbative
string unification. Let us briefly recall some of these properties:

a) In perturbative heterotic vacua the gauge coupling constants are unified with
the gravitational couplings:

$$G_{\text{Newton}} = \frac{1}{4} k_i \alpha_i^2 \alpha'$$

where $\alpha_i = g_i^2 / 4\pi$ ($k_i$) is the coupling constant (KM level) of the gauge group factor
$G_i$ and $\alpha'$ is the inverse string tension squared. Thus in string theory one has
unification of gauge coupling constants even in the absence of a GUT group. The
$k_i$ are integers $k_i \geq 1$ for non-Abelian factors ($k_i = 1$ in most models constructed)
and fractional normalization factors for the $U(1)$’s.
b) There is an upper bound on the rank of the gauge group in perturbative models. It comes from imposing the cancellation of conformal anomalies (vanishing of the total central charge, \( c = 0 \)) on the string world-sheet. It is easy to see that it must be \( \text{rank}(G) \leq 22 \). This 22 may be understood as coming from the rank of the gauge group in \( D = 10 \) before compactification \( \text{rank}(E_8 \times E_8) = \text{rank}(SO(32)) = 16 \) plus the maximum rank (six) of the gauge group one may obtain from Kaluza-Klein chiral \( N = 1 \) compactification on the extra six dimensions.

c) There are general restrictions on the possible gauge quantum numbers of massless fields in this class of theories. For example, for an \( SU(N) \) group realized at KM level \( k \), chiral fields which transform under the group as a representation with Young-tableaux wider than \( k \), cannot be present in the spectrum of the theory\(^\text{[10]}\). This implies, for example, that models with adjoint Higgs fields (like GUTs) can only be obtained for \( k > 1 \). Furthermore, the (left-moving gauge) conformal weight \( h_w \) associated to a massless field must obey \( h_w \leq 1 \). This implies that very large representations of the gauge group cannot possibly be in the massless spectrum of the theory, since one can see that \( h_w \) grows as the dimension of representation grows\(^\text{[10]}\). This is a very nice property of string theory since, from the point of view of gauge field theories, there is no reason at all to prefer lower dimensional representations like those appearing in the SM.

d) There are no exact continuous global symmetries in perturbative string theory\(^\text{[12]}\). Whenever they seem to be present they really correspond to local symmetries. This implies that continuous global symmetries appearing at low energies (like e.g., baryon or lepton numbers) can only be approximate (accidental) symmetries if seen from the string point of view.

e) Typically (although not always), when there are models with \( U(1) \) gauge symmetries, there is one linear combination which is apparently anomalous. In fact the anomaly is canceled by a four-dimensional version of the Green-Schwarz mechanism. At the same time a dilaton-dependent Fayet-Iliopoulos term proportional to the trace of the \( U(1) \) charge over the massless spectrum is created.

f) In all \( D = 4, N = 1 \) heterotic vacua there is a massless complex chiral multiplet whose complex scalar field (denoted \( S \)) contains the dilaton \( \text{Re} S = 4\pi/g^2 \). At the tree level one has\(^\text{[13]}\):

\[
 f_a = k_a S ; \quad K(S, S^*) = -\log(S + S^*)
\]

(2.2)

where \( f_a \) is the \( N = 1 \) gauge kinetic function and \( K(S, S^*) \) is the dilaton dependent piece of the Kahler potential.
g) $D = 4, N = 1$ string models with different numbers of chiral generations are (perturbatively) disconnected. The same is true for models with different numbers of supersymmetries.

All the above are generic features of $D = 4, N = 1$ perturbative string vacua. In addition a good amount of work on different aspects of string phenomenology has been done in the last ten years. Some of the directions particularly pursued in the literature are the following [1]:

i) **Effective low-energy action**

A general $N = 1, D = 4$ supergravity Lagrangian is determined (up to two derivatives) by its particle content and three functions of the scalars in the theory: the kinetic function $f_a$ ($a$ runs over the different gauge groups), the Kahler potential $K$ and the superpotential $W$. In principle those functions can be perturbatively computed for any given $D = 4, N = 1$ heterotic vacuum. In practice this has only been done for some classes of Abelian toroidal orbifolds [14] and some fermionic models. Some information is also known for some specific Calabi-Yau compactifications. In particular, the dependence of the one-loop (threshold) corrections on the moduli ($T_i$) and complex structure ($U_j$) scalar fields has been computed. These scalars characterize the size and shape of the compactifying manifold. One finds for the Wilsonian action a result of the form [15]:

$$f_a = k_a S + \Delta_a(T_i, U_i)$$  \hspace{1cm} (2.3)

The first piece is the tree level result which we already mentioned above. The $\Delta_a$ piece is the moduli-dependent one-loop correction and is a holomorphic function. As I said it has been computed for a large class of toroidal orbifold compactifications. One of the most interesting properties found is that it has definite properties under the $T$ -- dualities of the underlying torus. In particular $\exp(\Delta_a(T_i))$ behaves as a modular form with respect to the $SL(2, \mathbb{Z})$ symmetries corresponding to those dualities. The tree level Kahler potential and superpotential $W$ are also known for some classes of orbifold and fermionic models (see [1] for references).

ii) **Gauge coupling unification**

As we said, in string unification all gauge couplings meet at scale that should be close to the string scale which has been computed [16] to be of order $M_{string} = 3 \times 10^{17}$ GeV. It is nice that this unification occurs but an $N = 1$ extrapolation of low energy experimental results seems to indicate a lower unification scale of order $2 \times 10^{16}$ GeV. Several explanations have been proposed for this discrepancy. We direct the reader to ref.[17] for a review and references on the subject.
iii) **SUSY-breaking and soft terms**

One of the phenomenological problems of $D = 4$ string models (large rank of the gauge group) may be a virtue. These models often contain (hidden) gauge interactions which do not directly couple to the observable particles. If the gauginos $\lambda_a$ of those hidden groups condense ($\langle \lambda_a \lambda_a \rangle \neq 0$) supersymmetry will in general be broken \([18]\). Although this is a nice possibility which may generate naturally a hierarchy of scales, it is not free of problems. In particular, working at the effective Lagrangian level one finds that the scalar potential has a qualitative dependence on the dilaton $S$ and overall modulus $T$ of the form:

$$V(S, T) \propto \frac{1}{\text{Re} S \text{Re} T^3} \exp(3\text{Re} S)/2|\beta|^2$$  \hspace{1cm} (2.4)

The exponential comes from the condensate (recall $\text{Re} S = 4\pi/g^2$) and $\beta$ is the one-loop beta function of the confining gauge interaction. Here the scalar $T$ correspond to the overall modulus which measures the size of the compact manifold ($\text{Re} T = R^2$, $R$ being the compactification radius). From this equation one concludes that the perturbative vacuum lies at $\text{Re} S \rightarrow \infty$ and/or $\text{Re} T \rightarrow \infty$ which correspond to a non-interacting and/or a decompactification limit. These are the so called run-away problems. It has been argued \([20]\) that the large $T$ decompactification problem may disappear in some particular models if the one-loop threshold corrections $\Delta_a(T_i)$ are taken into account. Indeed, if this is done the above formula gets multiplied by a factor of the form $1/|\Delta(T)|^6$ in such a way that two interesting things happen \([20]\): 1) the scalar potential becomes invariant under the $SL(2, \mathbb{Z})$ modular invariance associated to $T$-duality. 2) An additional T-dependence appears in the potential in such a way that for large $T$ the potential grows and the $T$ vev is stabilized around the string scale (one has to be careful though to check that one remains in the perturbative regime). A similar stabilization seems difficult for the $S$ field, at least if one remains within perturbation theory. In fact this possibility was one of the motivations in ref.\([21]\) to introduce the concept of $S$-duality in string theory. An interesting new proposal for the stabilization of the dilaton field has been put forward in \([19]\).

Another approach to the problem of SUSY breaking at the level of the effective Lagrangian was discussed in refs.\([22, 23]\) (for a recent review see \([24]\)). The idea rests on the assumption that SUSY is predominantly broken by the vevs of the auxiliary fields $F_S$, $F_{T_i}$ of the dilaton/moduli fields present in large classes of $D = 4$, $N = 1$ heterotic vacua (particularly in toroidal orbifolds). This simple assumption (plus that of a vanishing cosmological constant) leads to specific relationships among
the different soft terms possible in the effective Lagrangian. Here is an example. Consider an $N = 1$, $Z_N$ or $Z_N \times Z_M$ heterotic orbifold. Any such a model (it does not matter whether it is a $(2,2)$ or a $(0,2)$ compactification) has three sets of matter chiral fields associated to the three complex compact dimensions (this is the untwisted matter sector of these models). Then one can show that if SUSY is broken by an arbitrary combination of $F_S$ and $F_T^i$ ($i = 1, 2, 3$) the following relationship between SUSY-breaking soft terms exist \[23\]:

$$m_1^2 + m_2^2 + m_3^2 = M^2 ; \quad M = A_{123}$$

(2.5)

Here $m_i$ are the soft masses of the scalars in the three untwisted sectors, $M$ is the gaugino mass and $A_{123}$ is the trilinear soft term associated to the Yukawa coupling which relates the three types of untwisted sectors. A particular case is that in which only the dilaton auxiliary field $F_S \neq 0$ (dilaton dominated limit) contributes. In that case one gets the simple expression $\sqrt{3}m_i = M = -A$. Perhaps the most interesting aspect of this kind of relationships is that they give rise to specific constraints on the supersymmetric spectra if applied to the MSSM. Thus one can hope to test this kind of ideas if SUSY is found at LHC. Another interesting point regarding these relationships is their behaviour with respect to field theory finiteness. Indeed, it turns out that boundary conditions for soft terms of the type shown in eq.(2.5), when applied to $N = 1$ two-loop finite theories, preserve finiteness \[25\].

iv) Anomalous $U(1)$’s

Perturbative heterotic vacua in which there are gauged $U(1)$’s often have one of those $U(1)$’s anomalous. In fact the anomaly is canceled by a $D = 4$ generalization of the Green-Schwarz mechanism \[20\]. This works as follows. The gauge kinetic term in the Lagrangian is $k_a SW_a W_a$, where $W_a$ is the field strength superfield. Under an anomalous $U(1)_X$ gauge transformation one has $\text{Im } S \to \text{Im } S - \delta_{GS} \Lambda(x)$, where $\Lambda(x)$ is the gauge $U(1)_X$ parameter and $\delta_{GS}$ is a constant model-dependent coefficient. Thus one can compensate an anomalous transformation from the standard triangle graphs by an appropriate shift of the axion-like field $\text{Im } S$. For this to work in a theory with gauge group $U(1)_x \times \prod_a G_a$, $a = 2, \cdots, n$ the mixed anomalies $A_a$, $a = 1, \cdots, n$ of the $U(1)_X$ with all group factors must be in the ratio of the KM level coefficients \[27\]:

$$A_1 : A_2 : \cdots : A_n = k_1 : k_2 : \cdots : k_n$$

(2.6)

This has an interesting phenomenological consequence. Since the $k_a$ give the normalization of the gauge coupling constants at the string scale, one can relate those
normalizations to the mixed anomalous of $U(1)_X$. Consider for example a situation in which we have as part of the gauge group the SM one, $SU(3) \times SU(2) \times U(1)_Y$. If there is an additional anomalous $U(1)_X$ one obtains [27]:

$$\sin^2 \theta_W = \frac{1}{1 + k_1/k_2} = \frac{1}{1 + A_1/A_2}$$

(2.7)

where $A_1$, $A_2$ are the mixed anomalies of $U(1)_X$ with $U(1)_Y$ and $SU(2)_L$ respectively. Thus the value of the weak angle can be computed in terms of anomaly coefficients, independently of any grand unification symmetry. This possibility has been also used in the last few years to construct models which predict definite patterns(textures) for fermion mass matrices [28, 29].

The presence of an anomalous $U(1)_X$ has other dynamical consequences. A (one-loop) dilaton dependent Fayet-Iliopoulos (FI) term proportional to $\text{Tr} Q_X$ is also generated. In particular, the $U(1)_X$ D-term contribution to the scalar potential has the form [30]:

$$V_X = \frac{g^2}{2} \left( \sum_i q_i |\phi_i|^2 + \text{Tr} Q_X \frac{g}{192\pi^2} \right)^2$$

(2.8)

where $\phi_i$ denotes the scalar fields which have charge $q_i$ with respect to $U(1)_X$. The presence of the second (FI) term in this formula forces some of the $\phi_i$ scalars to get a non-vanishing vev. Thus the classical vacuum is unstable but in all cases studied up to now there is a nearby supersymmetric vacuum.

There are several other phenomenological aspects of $D = 4$, $N = 1$ perturbative vacua with have also been studied in the last few years including the possible role of the $\text{Im} T_i$ fields as invisible axions, cosmological constraints on the dilaton/moduli sector etc. We refer the reader to the reviews [1].

3 String theory dualities

A lot has been learned about the non-perturbative structure of string theories in the last three years or so. There is a good number of reviews on the subject [3] and I will not try to describe here these developments. Let me just describe the different connections which have emerged among the different types of supersymmetric string theories. First of all, it was known that, upon toroidal compactification of one dimension, Type IIA string is T-dual to IIB if we simply exchange the radius of compactification $R \rightarrow \alpha'/R$ [31]. This means that these two theories are perturbatively equivalent. The same happens with the two type of heterotic strings: $E_8 \times E_8$ is T-dual to the $SO(32)$ heterotic upon the same exchange (in this case a
particular Wilson line breaking both groups to the common subgroup $SO(16)^2$ is also needed). These relationships are perturbative in nature and were known since the mid-eighties [31]. New non-perturbative dualities have been found in the last three years [3]: Type I, $SO(32)$ string theory is S-dual to the $SO(32)$ heterotic string. This means that the weak coupling of the former is equivalent to the strong coupling limit of the latter. Furthermore it has been found that Type IIA theory and the $E_8 \times E_8$ heterotic may be both obtained from an underlying 11-dimensional theory termed M-theory. Type IIA string is obtained from M-theory upon compactification on a circle, the string coupling constant being given by the compactification radius. Thus when one sends the 10-dimensional string coupling constant to infinity one recovers the eleventh dimension of M-theory. The $E_8 \times E_8$ string is obtained upon compactification of M-theory on a segment. The two $E_8$ factors are associated to the two boundaries of the segment. Combining all these connections one obtains the remarkable result that all known $D = 10$ string theories are different perturbative limits of a unique underlying 11-dimensional structure, M-theory. Our knowledge of M-theory is still very limited but we know that it contains membranes and fivebranes as fundamental ingredients and that its low-energy limit gives rise to 11-dimensional supergravity.

These duality connections between the different $D = 10$ supersymmetric strings gives rise upon compactification to lower dimensions to a complicated network of non-trivial connections between the different theories. The apparent consistency of all the different derived connections in lower dimensions is one of the strongest arguments in favour of all this unique construction. For example, one expects a duality between Type-IIA compactified on a Calabi-Yau manifold and heterotic $E_8 \times E_8$ compactified on $K3 \times T^2$ [32]. Both compactifications yield $N = 2$, $D = 4$ theories and abundant evidence has been found for this equivalence. New kinds of perturbative and non-perturbative Type IIB vacua have also been constructed. In particular a new non-perturbative way to obtain Type IIB vacua has been termed F-theory [33]. This theory lives in twelve dimensions, although the extra two dimensions are dynamically frozen. Compactifications of this F-theory on complex (elliptic) Calabi-Yau n-folds yield $D = (12 - 2n)$-dimensional theories with $N = 1$ supersymmetry. For example, F-theory compactified on a complex Calabi-Yau four-fold is expected to be dual to heterotic compactifications on Calabi-Yau 3-folds. Thus one expects to extract non-perturbative information about heterotic vacua from the F-theory dual.
New perturbative Type I vacua have also been recently constructed. This has renewed interest since its duality to the $SO(32)$ heterotic implies that one should be able to obtain non-perturbative information on the latter by working with Type I perturbative vacua. In particular, one can obtain Type I theory as a sort of world-sheet parity orbifold (orientifold) of Type IIB theory [34, 35]. In this construction of Type I, the open strings appear as the twisted sector of the Type IIB theory modded with respect to world-sheet parity reversal (see below). New consistent perturbative $N = 1, D = 4$ Type I vacua can be obtained [36, 37] by combining the above orientifold action with discrete $Z_N$ twists analogous to those appearing in standard heterotic orbifold models.

As we will discuss below, all this rich structure leads to a number of important implications for the possible phenomenological applications of string theory. These implications can be classified as follows:

1) New non-perturbative phenomena.
2) New classes of perturbative and non-perturbative vacua.
3) Extraction of non-perturbative information about previously known vacua.

The study of all these aspects is still in its infancy so I will limit myself to discuss some expected general new features in what follows.

4 New general features

a) Chirality-changing transitions

Some of the new phenomena appearing in $D = 4, N = 1$ string vacua may be more easily understood in terms of the simpler $D = 6, N = 1$ case. Let me then first review a few aspects of the latter type of field theories [32]. In $N = 1, D = 6$ theories the relevant (non-gravitational) SUSY multiplets are of three types: 1) Vector multiplets. They contain a vector field and its gaugino partner; 2) Hypermultiplets. Contain a couple of complex scalars and their fermionic partners. In general they transform under the gauge group of the vector multiplets of the given theory; 3) Tensor multiplets. They contain a two index antisymmetric field $B^{\mu\nu}$, a real scalar $\phi$ and a fermionic partner. They do not carry quantum numbers with respect to the gauge group. Perhaps the simplest type of $D = 6, N = 1$ string vacua may be obtained by compactifying the heterotic string on a Calabi-Yau complex 2-fold (i.e., the K3 manifold). Consider for example the $E_8 \times E_8$ heterotic. If we do the standard embedding of the spin connection into the gauge connection we obtain
a $E_7 \times E_8$ gauge symmetry and hypermultiplets transforming like $10(\overline{56}) + 65(\mathbf{1})$. To obtain an anomaly free result one can check that the background contains a total of 24 instantons. This kind of perturbative vacua have only one tensor multiplet whose real scalar $\phi$ is the string dilaton. It is directly obtained after dimensional reduction of the $D = 10$ supergravity multiplet. Now, it was found [18] that there can be non-perturbative transitions in the theory under which a particular (gauge anomaly free) combination of 29 hypermultiplets transforms into a (singlet) tensor multiplet:

$$1 \text{ Tensor} \leftrightarrow \frac{1}{2}(56) + (\mathbf{1})$$  \hspace{1cm} (4.1)

This type of transitions occur when the size of one the 24 underlying background instantons is set to zero. An analogue of this kind of transitions for the $D = 4$, $N = 1$ case is also expected on general grounds [38, 40]. If we compactify the $E_8 \times E_8$ heterotic on a Calabi-Yau (complex three-fold) with standard gauge embedding we get instead as gauge group $E_6 \times E_8$ and we get 27-dimensional chiral multiplets transforming like $E_6$ fundamentals. Since we know that the $E_7$ fundamental branches as $56 \rightarrow 27 + 2\overline{7} + 2\mathbf{1}$, non-perturbative (in general, chiral) transitions of the type

$$\text{Singlets} \leftrightarrow 2\overline{7}$$  \hspace{1cm} (4.2)

are expected. More generally, anomaly-free combinations of charged chiral fields may be transmuted into chiral singlets and vice versa. This applies, of course, to other groups and not only $E_6$. This seems to imply that string vacua with different numbers of e.g., chiral $E_6$ (or $SU(5)$, or SM) generations can be non-perturbatively connected. This would have far reaching consequences since it opens the way to a dynamical determination of the number of quark-lepton generations. Indeed, in perturbative string theory vacua with different numbers of chiral generations are disconnected and it seems difficult a dynamical choice of the number of generations, since the different vacua with different number of generations could not be dynamically preferred. Of course, it is still not clear what could eventually determine dynamically the chiral particle spectrum. Furthermore this kind of non-perturbative transitions in $D = 4$, $N = 1$ have still to be better understood.

Apart from the above variation of the number of chiral generations other possibilities appear to be (in principle) opened. It is well known that both in GUTs or in string models whenever there are extra unwanted particles in the massless spectrum there is essentially always the same idea to get rid of them: look whether the model has appropriate Yukawa couplings so that, by giving vevs to appropriate scalars, the
unwanted fields become superheavy. The new type of non-perturbative transitions seem to provide another mechanism by which one can get rid of unwanted charged fields. An important difference is that whereas standard Yukawas can only give masses to vector-like, non-chiral combinations of fields, the above non-perturbative transitions can also make to disappear chiral (anomaly free) collections of charged fields. Let me emphasize however that these are expectations which still have to be realized in specific models.

b) Rank of the gauge group

I already remarked how \( D = 4, N = 1 \) heterotic compactifications have the rank of the gauge group bounded, rank \( G \leq 22 \). It turns out that non-perturbative effects can give rise to new gauge interactions so that, in practice, the rank of the gauge group in string vacua is essentially unbounded above! Take again for a start the case of heterotic vacua in \( D = 6 \) with \( N = 1 \) SUSY. Consider now the case of the \( SO(32) \) heterotic compactified on \( K3 \) with standard embedding in the gauge group. The \( SO(32) \) gauge symmetry is generically broken to \( SO(28) \times SU(2) \) and there are hypermultiplets transforming as \( 10(28, 2) + 65(1, 1) \). As in the \( E_8 \times E_8 \) case, this corresponds to the presence of a background of 24 instantons (which turns out to be required for anomaly cancellation). Each of these instantons have parameters which govern e.g., their size. It was found by Witten \[41\] that if \( n_I \) of those instantons is put to zero size some interesting non-perturbative phenomena occur (at finite coupling constant): 1) Additional gauge interactions with the simplectic \( Sp(n_I) \) gauge group appear and 2) Hypermultiplets transforming as the antisymmetric representation of \( Sp(n_I) \) and \( n_I \) vectorials of \( SO(28) \) appear. If the instantons are located at some singularity of the \( K3 \) manifold, gauge groups \( U(N) \) and \( SO(2N) \) with a variety of hypermultiplet representations may appear. In these more general cases extra tensor multiplets may also appear in the spectrum \[32\]. Very singular configurations of the instanton moduli space may yield very large gauge groups. For example, if we consider a \( E_8 \times E_8 \) compactification on \( K3 \) and we locate all the 24 instantons at the same point and coinciding at a \( K3 \) singularity of the so called \( E_8 \) type, one gets the gauge group \[42\]:

\[
E_8^{17} \times F_4^{16} \times G_2^{32} \times SU(2)^{32}
\] (4.3)

and 193 tensor multiplets! The same kind of phenomena appear in \( D = 4, N = 1 \) non-perturbative vacua. A particularly impressive gauge group is obtained with certain heterotic compactification of \( E_8 \times E_8 \) (whose dual may be obtained as an
F-theory compactification on certain Calabi-Yau four-fold). It yields a group (4.3):

\[ E_8^{2561} \times F_4^{7576} \times G_2^{20168} \times SU(2)^{30200} \]  

One must emphasize that these are not typical examples of gauge groups, of course, well on the contrary they correspond to very, very particular configurations in the moduli space of thousands of scalar fields. But they certainly provide examples showing how the perturbative heterotic bound rank \( G \leq 22 \) is badly violated at the non-perturbative level. In fact one does not need to invoke non-perturbative effects to obtain large gauge groups (not so enormous as the above !) in some particular (non-heterotic) string constructions. They appear easily, as we will discuss below, in Type I perturbative string constructions. However this fact was unknown till the advent of the D-brane technology developed after Polchinski’s paper in 1995 [44].

From the phenomenological side the above fact suggests a number of comments. To start with, where should we embed the SM group? It is all of it non-perturbative, perturbative or some factor (e.g. QCD) is non-perturbative and the rest perturbative? Due to the S-dualities probably it does not make sense to say that the SM group is fully non-perturbative or fully perturbative since the whole idea of strong/weak coupling duality is the equivalence of those two regimes (at least for relatively small gauge groups; certainly there is no perturbative model yielding such extreme spectrum as the ones above). But it could well be that part of the gauge group of the SM (or a GUT) or/and chiral particle content could have non-perturbative origin. For example, the chiral multiplet spectrum of the MSSM is suspiciously asymmetric: quarks and leptons come in chiral representations whereas the Higgs sector is vector-like. Perhaps the Higgs sector has non-perturbative and the three generations perturbative origins (or viceversa).

What seems clear is that the string perturbation theory used up to now to explore string vacua misses most of the moduli space. We will have to study the new possibilities opened up for embeddings of the SM into string theory.

c) Multiple dilaton-like fields

In perturbative \( D = 4, N = 1 \) heterotic vacua there is a single complex scalar field \( S \) whose real part is the dilaton and which couples in a universal manner to all gauge groups and matter fields. In non-perturbative vacua there are in general more than one field with similar characteristics. Let us again start with six dimensions. Consider the compactification of the \( E_8 \times E_8 \) heterotic on K3 with standard gauge embedding. As we said, the perturbative background must include 24 instantons, the gauge group is broken to \( E_7 \times E_8 \) and there are hypermultiplets transforming
like $10(56) + 65(1)$. This model has only one $D = 6$ tensor multiplet which includes one real scalar, the dilaton. This unique tensor multiplet comes from dimensional reduction of the unique $D = 10$ gravitational multiplet. We saw in the previous section that in the $SO(32)$ case, when $n_I$ of the instantons in the background are put at the same point in the $K^3$ and with zero size, some non-perturbative gauge group plus hypermultiplets appear. When we do the same thing in the present $E_8 \times E_8$ case something quite different happens \[32\]. What happens is that $n_I$ tensor multiplets appear in the massless sector (and some hypermultiplets disappear). This has a clear M-theory interpretation which we refrain from explaining here \[38\]. This is an example of a generic phenomenon. While perturbative heterotic vacua in $D = 6$ have just one single tensor multiplet which contains the dilaton, non-perturbative vacua contain an indefinite number of such tensor multiplet. The appearance of such type of tensor multiplets proliferates (both in $E_8 \times E_8$ and $SO(32)$ vacua) if the zero size instantons are located at some $K^3$ singularity. We already mentioned that in the $D = 6$ extreme example in previous section there were 193 tensor multiplets! Each of them contains one scalar with dilaton-like couplings to all the different gauge groups in the theory. When compactifying further to $D = 4$, $N = 1$ theory one generically gets from all these tensor multiplets complex scalar fields $S_i$ with behaviour similar (though not identical) to that of the usual perturbative $S$ field. Furthermore some of the dualities exchange the $T_i$ moduli fields of the compactification with the dilaton. All in all, if the string vacuum has a gauge group $\prod_a G_a$ one finds a general structure for the different gauge kinetic functions

$$f_a = c_a^1 S_1 + c_a^2 S_2 + \cdots + c_a^n S_n .$$

(4.5)

where the $S_i$ represent here general massless scalars with dilaton behaviour but may also include moduli type of fields $T_i$ in specific cases. The $c_a^i$ are model dependent coefficients. This implies that in general the gauge coupling constants of the different gauge group factors are going to be different. Unlike perturbative heterotic vacua in which the tree level $f$-function is proportional to a universal $S$ field, here things are more complicated. Gauge coupling unification is no longer automatic. Actually this is qualitatively not so different from the perturbative case: we already saw in eq.(2.3) that one-loop corrections to the $f$-function involve non-universal gauge couplings for different gauge groups, but one may hope those corrections to be smaller than the tree-level result. Here the dependence on the different $S_i$ fields of the $f_a$ functions is not expected to be particularly suppressed. In particular, if one embeds the SM group directly into a string vacuum (without an intermediate
GUT structure) one runs into the risk of getting different boundary conditions for
the three gauge coupling constants. This could perhaps be an argument in favour
of string GUT's [45], although one cannot exclude the existence of models in which
the different $SU(3) \times SU(2) \times U(1)$ couplings are equal without a GUT symmetry
(perturbative models are an example).

Let us also point out that the proliferation of massless dilaton-like fields makes
also harder an analysis of soft terms along the lines discussed in chapter 2.

d) Anomalous $U(1)$'s

We already remarked that perturbative $D = 4, N = 1$ heterotic vacua often
contain one anomalous $U(1)_X$. The anomaly is actually canceled by the four-
dimensional version of the Green-Schwarz mechanism. Since there is only one
complex dilaton field $S$ to do the trick, there can only be one anomalous $U(1)_X$
and its mixed anomaly with all the gauge groups have to obey eq.(2.6).) In non-
perturbative heterotic vacua (or in perturbative Type I vacua) there are more than
one field which can help in canceling $U(1)$ anomalies, there is a generalized $D = 4,$
$N = 1$ Green-Schwarz mechanism 1 at work [37]. There are two consequences of
this : 1) there may be more than one anomalous $U(1)$ in the models and 2) even
in cases in which there is a single anomalous $U(1)$, its mixed anomaly with respect
to the different group factors may be non-universal. One also expects the presence
of several Fayet-Iliopoulos terms depending on the different $S_i$ fields for each of the
different anomalous $U(1)$'s.

The above facts may have implications on the phenomenological use of anomalous
$U(1)$’s in order to construct fermion mass matrices [28, 29]. For example, one can
consider new classes of models with two anomalous $U(1)$’s or else models with only
one anomalous $U(1)$’s but non-universal mixed anomalies with the different gauge
groups etc. Of course one loses in this case the simplicity of the perturbative models.

5 New $D = 4, N = 1$ String vacua

Most of the new perturbative and non-perturbative string vacua constructed using
the new duality and D-brane techniques are higher dimensional and with extended
supersymmetry. Much work remains to be made in the more phenomenologically
interesting $D = 4, N = 1$ case. Let me briefly review some of the latter construc-

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1 An analogous generalized mechanism was in fact known to exist in $D = 6$ [46].
5.1 F-theory compactifications on 4-folds

This is not the place to present a review of F-theory [3]. Let me just describe a few main points. F-theory [33] is a new non-perturbative method to get consistent Type IIB vacua in a variety of dimensions (including $D = 4, N = 1$ theories). The bosonic massless fields in $D = 10$ Type IIB theory include two scalars, the dilaton $\phi$ (from the NS-NS sector) and a scalar $a$ (from the Ramond-Ramond sector). The particular complex field combination $\tau = a + i e^{-\phi/2}$ turns out to be specially relevant. Indeed, Type IIB theory present a $SL(2,\mathbb{Z})$ S-duality which is generated by the modular transformations in the $\tau$ field. The usual perturbative vacua take $\tau$ to be a constant. F-theory vacua allow $\tau$ to vary in a particular form consistent with the $SL(2,\mathbb{Z})$ S-duality of the theory. One identifies $\tau$ with the complex structure of a torus $T^2$ living in some unphysical 11-th and 12-th dimensions. The idea is then to compactify this 12-dimensional F-theory on a manifold $M$ which locally looks like $M \propto T^2 \times CY$ (i.e., an ‘elliptical fibration’), $T^2$ being in the 11-th and 12-th dimensions. This gives rise to non-perturbative compactifications of IIB theory on the CY submanifold in the fibration.

In order to reach a $D = 4, N = 1$ theory one has to compactify F-theory on a complex Calabi-Yau four-fold which is an elliptical fibration [14]. The latter kind of manifolds are quite complicated and are not so well known as the complex 3-fold Calabi-Yau manifolds abundantly used for heterotic compactifications in the last fourteen years. Furthermore extracting the spectra of these models requires substantial expertise in algebraic geometry. Still some particular $D = 4, N = 1$ models have been studied, although up to now the specific examples obtained seem to be non-chiral. In spite of their complication, F-theory vacua are particularly well suited to extract non-perturbative information of string vacua. For example, superpotentials for some particular $D = 4, N = 1$ F-theory vacua have been obtained which have interesting modular behaviour under dualities [18].

5.2 Brane cooking

Dirichlet-branes (D-branes) have played a dominant role in many of the recent duality developments. Again, this is not the place to review this subject (see ref.[49] for a nice pictorial review). D-p-branes are solitonic states of Type II string theory which have $p+1$-dimensional worldvolume and carry Ramond-Ramond charges [44]. An important property of D-branes (which may serve as a definition) is that open
strings are allowed to end and start on them. D-branes have gauge fields living on their worldvolume and when sets of them are located at the same point they give rise to non-Abelian symmetries like e.g. $SU(N)$ on the worldvolume. Now, one can cook certain combinations of Type IIA D-branes plus Neveu-Schwarz fivebranes distributed in such a way that on the worldvolume of some particular D-branes lives a $D = 4, N = 1$ theory of interest to us. This provides us with a new geometrical tool to study some perturbative and non-perturbative aspects of the gauge theory living on the worldvolume. The simplest such configuration \[50\] corresponds to $N_c$ Type IIA D-fourbranes stretching in between two Neveu-Schwarz fivebranes \[.\] On the worldvolume of the D-fourbranes lives a $D = 4, N = 2$ gauge theory with $SU(N_c)$ gauge group.

Chiral models with $D = 4, N = 1$ can also be obtained by acting with some orbifold actions (see J. Lykken’s contribution to these proceedings). Typically one gets gauge groups of the form $U(N)^m$ with chiral matter in bifundamental representations of the type:

$$(N, \overline{N}, 1, \cdots, 1) + (1, N, \overline{N}, 1, \cdots, 1) + \cdots + (\overline{N}, \cdots, 1, N) \quad (5.1)$$

Many other possibilities are however possible. One particularly nice point about this type of constructions is that they can often be embedded into M-theory and one can obtain important non-perturbative information (like the Seiberg-Witten curves in the $N = 2$ case).

It is important to remark that these kind of theories are not compactifications, gravity still lives in $D > 4$ dimensions, so they cannot be considered as they stand as $D = 4$ unified theories of all interactions. They are however an important tool to study the properties of pure gauge theories (possibly with matter) leaving on the worldvolume. It would be interesting to try to obtain in this way a brane configuration with an $N = 1$ model living on the worldvolume resembling the MSSM. It will not be a unified model of all interactions but perhaps could give us some interesting hints.

### 5.3 Type IIB, $D = 4$, $N = 1$ orientifolds

These are new classes of perturbative $N = 1$ four-dimensional strings whose structure was considerably developed with the advent of the D-brane technology in the

\[2\]The latter are not, strictly speaking, D-branes but solitonic fivebranes associated to the the massless $B^{\mu \nu}$ field coming the NS-NS sector of the theory.
last few years. Actually they may equivalently be considered as Type I vacua, as we will discuss below.

Type I vacua with $D = 4, N = 1$ were essentially ignored before 1995 (see however ref. [34, 35] for early work). Type I in $D = 10$ has gauge group $SO(32)$, and fully fledged four-dimensional strings based on it were ignored because of the technical difficulties in ensuring anomaly cancellation. The latter has to be checked case by case. On the other hand the heterotic $SO(32)$ has also low energy group $SO(32)$ but it is trivial to obtain general conditions for absence of anomalies in $D = 4, N = 1$ models. Anomaly cancellation is a direct consequence of world-sheet modular invariance of the closed heterotic string and it is very easy to obtain general conditions on string (e.g., orbifold) backgrounds guaranteeing absence of anomalies. This is why no serious attempt to do $D = 4, N = 1$ string model building based on Type I was made before: 1) it is technically cumbersome to ensure absence of anomalies and 2) another $SO(32)$ (the heterotic) string exists in which one can obtain equally interesting models with less effort!

There are a few reasons to reconsider Type I vacua in the light of recent developments. First of all, before 1995 the existence and relevance of Dirichlet branes in the context of Type I theory [44] was unknown. In fact their existence may be considered as a consequence of T-dualities in the context of open strings. Secondly, as we remarked above, there is evidence of an S-duality between Type I $SO(32)$ string theory and the $SO(32)$ heterotic. This means that strongly coupled heterotic string is equivalent (dual) to weakly coupled Type I. Thus one might expect to obtain information about non-perturbative heterotic vacua by studying the perturbative (weakly coupled) Type I theory. Actually, the strong and weakly coupled regimes get more entangled when compactifying both theories to lower dimensions. Doing a dimensional reduction in both theories one finds that the mapping between the dilatons in the different dimensions is [51, 52]:

$$\Phi_I = \frac{6 - D}{4} \Phi_H - \frac{D - 2}{16} \log \det G_H^{(10 - D)}$$

(5.2)

where $D$ is the dimensionality of space-time and $\Phi_I$ ($\Phi_H$) is the Type I (heterotic) dilaton (recall $\exp(\Phi)$ yields the strength of e.g., gauge interactions in both theories). $G_H$ is the metric of the $(10 - D)$ compact dimensions in the heterotic frame. Notice that for $D = 10$, $\Phi_I = - \Phi_H$, so that indeed a strongly coupled heterotic corresponds to a weakly coupled Type I string. The case $D = 6$ is special since the Type I dilaton is in fact mapped to the overall modulus (compact variety size). Now in the phenomenologically interesting $D = 4$ case the coefficient of $\Phi_H$ is positive so
that there are different regimes. For moderate compactification radius there is a weak↔weak coupling duality between Type I and heterotic in $D = 4$. For large heterotic compactification radius there is a weak↔strong duality. The two regimes can yield complementary information in particular vacua [32].

Let us discuss now in somewhat more detail how Type IIB $D = 4$, $N = 1$ orientifolds are constructed [34, 35, 36, 37]. The idea is analogous to that of toroidal orbifolds in heterotic strings [53]. One starts from Type IIB string and compactifies on $T^6/G$, where $G$ is some discrete Abelian group ($Z_N$ or $Z_N \times Z_M$) which acts rotating the lattice vectors defining the 6-torus $T^6$. This means that there will be an ‘untwisted sector’ in the spectrum obtained by just a projection under $G$ of the $IIB/T^6$ spectrum and ‘twisted sectors’ which correspond to IIB strings which are closed modulo an element of $G$. This is in general not enough to get $N = 1$ in $D = 4$ since (unlike the heterotic case which has only one) there are two $D = 10$ supersymmetries in Type II strings (one coming from right-moving closed string oscillations and the other from the left-moving ones). To further reduce the number of supersymmetries one can mode the theory under the $Z_2$ operation termed world-sheet parity $\Omega$ [34]. Let $(\sigma, \tau)$ be the two (space-like and time-like) worldsheet coordinates of the string. Defining the complex world-sheet coordinate $z = \exp(\tau + i\sigma)$, one then has $\Omega z = \bar{z}$. Thus $\Omega$ transforms left-moving and right-moving vibrations of the string into each other, so that the result of a projection of IIB string under $\Omega$ is a closed unoriented string with only one $D = 10$ SUSY.

In addition it turns out that the consistency of the theory requires the addition of twisted sectors with respect to $\Omega$. These are nothing but Type I open strings which have to be added to the closed unoriented strings discussed above. All in all this is a compactification of Type IIB on $T^6/\{G, \Omega\}$ yielding in fact a Type I $D = 4$ string theory with $N = 1$.

Cancellation of anomalies is not guaranteed in this $D = 4$ models and has to be essentially imposed case by case [36, 37]. This cancellation may be reinterpreted as the vanishing of certain one-loop tadpole graphs. It is at this level that the introduction of the D-branes is forced upon us. For our purposes they may be defined as submanifolds of the full $D = 10$ space where the open strings can end. In the simplest type of models we are describing here only D-ninebranes and D-fivebranes turn out to be relevant. The world-volume of the ninebranes is the full $D = 10$ space. Thus the ends of open strings move freely in all ten dimensions. There is an index $i$ attached to each ninebrane as a label and open strings have then associated
matrices (the Chan-Paton matrices). For the case of the fivebranes the worldvolume is some six-dimensional submanifold and they also carry an index \(j\) as a label. It turns out that for fivebranes to be present in this class of models the orbifold discrete group \(G\) must contain an order-two twist. The fivebrane worldvolume spans the four uncompactified dimensions plus one of the three compact complex planes. Thus in this class of \(D = 4\) orientifolds there may be up to three different sets of fivebranes depending on the particular complex plane occupied by their worldvolume. There will also be open strings stretching between fivebranes which will also carry associated Chan-Paton matrices.

Let me describe how is the general structure of the massless spectrum of this kind of theories \([36, 37]\). As I said, they contain both closed and open strings. From the closed strings one gets the gravity multiplets as well as a number of moduli/dilaton singlet fields analogous to the \(S\) and \(T_i\) fields of the heterotic. From the open strings one gets both gauge fields and chiral multiplets transforming under them. Open strings stretching among ninebranes give rise to some gauge group \(G_9\) (typically with rank 16 or less). Each set of fivebranes yields some extra gauge group \(G_{5,i}\), \(i = 1, 2, 3\). It is thus obvious that the rank of the gauge group in Type IIB orientifold models exceeds in general the perturbative heterotic bound rank \(G \leq 22\). There are chiral multiplets \(C_9\) which are charged under \(G_9\) and chiral multiplets \(C_{5,i}\) which are charged with respect to the corresponding \(G_{5,i}\) group. In addition there are charged chiral fields \(C_{95,i}, C_{5,5j}\) which transform simultaneously under different sets of groups. They come from open strings stretching between different classes of D-branes.

Let us show as a first example \([54]\) a model which only contains fivebranes. In this example the Type IIB theory is compactified on \(T^6\) and modded by the orientifold group \(Z_6 = Z_3 \times Z_2\). Here \(Z_3\) is the standard \(Z_3\) action in \(D = 4\) which involves \(2\pi/3\) rotations on the three compact complex planes. The \(Z_2\) is generated by \(\Omega R\) (instead of simply \(\Omega\)), where \(R\) is a reflection of the first two complex coordinates. It turns out that due to the fact that \(\Omega\) is not a generator of the orientifold group, there are no ninebranes. The presence of the element \(\Omega R\) and tadpole cancellation requires the presence of 32 Dirichlet fivebranes whose worldvolume lives in the four non-compact dimension plus the 3-d compact plane. Now we will chose a particular configuration of the fivebranes on the fixed points under the \(Z_3\) action which obeys tadpole cancellation conditions. Eight fivebranes will be sitting at fixed point at the origin. The open strings stretching among these fivebranes give rise to a \(U(4)\) group
with three 6-plets. The remaining 24 fivebranes will be sitting at some other fixed point away from the origin. The corresponding open strings give rise to a gauge group $U(4)^3$ with chiral multiplets

$$3(4, \bar{4}, 1) + 3(\bar{4}, 1, 4) + 3(1, 4, \bar{4})$$  \hspace{1cm} (5.3)$$

The full chiral multiplet content (except for the dilaton field $S$) is shown in the table. This model is chiral and has the interesting property that the $SU(4)^3$ theory corresponding to the fivebranes away from the origin is a finite theory \[54\].

The table also shows the charges of the different particles with respect to the four $U(1)$’s. All of them are anomalous except for the linear combination $Q_1 + Q_2 + Q_3$ which does not couple to any chiral multiplet. In particular one finds $\text{Tr} Q_X = -36$ (gravitational anomalies) and also the mixed anomaly with respect to the last $SU(4)$ is $A_4 = -6$ (the mixed anomaly with the other three $SU(4)$’s vanishes). This provides as with an example of what we stated in the previous chapter: in Type I vacua the mixed anomalies of a $U(1)_X$ with respect to the different non-Abelian factors are not equal. In the present model there is a generalized Green-Schwarz mechanism in which some of the 27 twisted moduli participate. This model has a simple perturbative heterotic dual which has a similar (but not identical) massless particle spectrum.

Some of the simplest and more interesting types of $D = 4, N = 1$ orientifolds have both ninebranes and one set of fivebranes whose worldvolume spans the four non-compact dimensions plus (say) the 3-d compact complex plane \[36, 37\]. They are particularly interesting since, having in general a gauge group with rank larger than 22, they can only be dual to non-perturbative heterotic vacua. Thus the hope is to learn non-perturbative aspects of $D = 4, N = 1$ heterotic vacua by studying

\begin{table}
\centering
\begin{tabular}{|c|c|c|c|c|c|}
\hline
Sector & $SU(4)^3$ & $Q_1$ & $Q_2$ & $Q_3$ \\
\hline
Open Strings & $3(1, \bar{4}, 4, 1)$ & 0 & 0 & -1 & 1 \\
& $3(4, 1, \bar{4}, 1, 1)$ & 0 & 1 & 0 & -1 \\
& $3(\bar{4}, 4, 1, 1)$ & 0 & -1 & 1 & 0 \\
& $3(1, 1, 1, 1)$ & -2 & 0 & 0 & 0 \\
\hline
Closed Twisted Strings & $27(1, 1, 1, 1)$ & 0 & 0 & 0 & 0 \\
\hline
Closed Untwisted Strings & $9(1, 1, 1, 1)$ & 0 & 0 & 0 & 0 \\
\hline
\end{tabular}
\caption{Chiral multiplets in the $Z_2 \times Z_3$ Type IIB orientifold.}
\end{table}
the structure of their Type IIB orientifold duals. This class of models have a gauge group with a structure

\[ G_9 \times G_5 \]  

(5.4)

In these models both groups have (maximal) rank 16. If the fivebranes are all located at the fixed point lying at the origin in the first two complex dimensions, it turns out that one has \( G_9 = G_5 \) and there is an explicit symmetry under the exchange of the spectra coming from ninebrane and fivebrane sectors. It turns out that this symmetry is nothing but standard (Type I) T-duality with respect to the first two complex planes: \( R_i \leftrightarrow 1/R_i, \ i = 1, 2 \). One can see that T-duality in Type I theory exchange the roles of fivebranes and ninebranes and the above fivebrane configuration is self-dual with respect to T-duality. The untwisted closed string sector includes dilaton/moduli fields with compactification radius dependence

\[ S = e^{-\phi} R_1 R_2 R_3 + i \theta ; \ T_i = e^{-\phi} \frac{R_i}{R_j R_k} + i \eta_i \ (i \neq j \neq k) \]  

(5.5)

where \( R_i \) is the size of the \( i \)-th complex plane. These fields are the Type I duals of their well known heterotic counterparts. From the open string sectors one gets three sets of chiral fields \( C_9^i, \ i = 1, 2, 3 \) charged with respect to \( G_9 \), three sets of fields \( C_5^i, \ i = 1, 2, 3 \) charged with respect to \( G_5 \) and fields \( C_{95} \) charged with respect to both. One finds the remarkable result that the gauge kinetic functions of \( G_9 \) and \( G_5 \) are respectively:

\[ f_9 = S ; \ f_5 = T_3 \]  

(5.6)

This is to be compared to the heterotic perturbative result eq.(2.2). Looked from the heterotic dual point of view, \( G_9 \) is a perturbative gauge interaction whereas \( G_5 \) is non-perturbative. Equation (5.6) shows that some non-perturbative heterotic groups have gauge couplings governed by the compactification moduli \( T_i \) instead of the dilaton field \( S \). Notice that under T-duality in the first two complex planes one has

\[ R_1 \leftrightarrow \frac{\alpha'}{R_1} ; \ R_2 \leftrightarrow \frac{\alpha'}{R_2} \]  

\[ S \leftrightarrow T_3 \]  

\[ T_{1,2} \leftrightarrow T_{2,1} \]  

\[ C_9^i \leftrightarrow C_5^i \]  

\[ C_{95} \leftrightarrow C_{95} \]  

(5.7)
We observe that under T-duality the role of $S$ and $T^3$ are exchanged. Looked from the heterotic dual side this exchange between the dilaton $S$ field and the modulus $T^3$ field would look as a non-perturbative symmetry.

Typically, the gauge group in this kind of models has the form

\[ G_9 = G_5 = U(M) \times U(N) \times U(P) \times \cdots \]  

(5.8)

whereas the chiral fields are usually bi-fundamental representations such as e.g.,

\[(M, \overline{N}, 1, \cdots, 1) + (1, N, \mathcal{P}, 1, \cdots, 1) + \cdots \text{ etc.},\]

although there can also be some antisymmetric representations among the $C_i^9$ and $C_i^5$ fields. This kind of representations are of similar type to the ones corresponding to the left-handed quarks in the SM, which are $(3, 2)$ representations under $SU(3) \times SU(2)$. Notice, however, that Type I vacua can never give rise \textit{at the perturbative level} to spinorial representations of $SO(2N)$ groups nor exceptional groups either (this is because in Type I theory the gauge group is actually $SO(32)$ and not $Spin(32)$). Thus Type I theory is not the simplest way to obtain e.g. $SO(10)$ GUT’s (which have generations in spinorial reps.) nor $E_6$ GUT’S (that have $SO(10)$ spinorial roots).

This kind of orientifold vacua have in general several anomalous $U(1)$’s [37]. The anomaly is canceled by a generalized Green-Schwarz mechanism in which not only $S$ but $T^3$ and other moduli fields get involved (their imaginary parts get shifted under anomalous $U(1)$ transformations). This is an example of the phenomenon we described in section 4-d.

As we said these are models with gauge group $G_9 \times G_5$ whose rank may be as high as 32. If they are S-dual to some heterotic vacua, certainly they must be non-perturbative heterotic vacua. It is easy to find heterotic vacua ($SO(32)$ heterotic on $Z_N$ or $Z_N \times Z_M$ orbifolds) whose gauge group and untwisted charged fields precisely match the $G_9$ gauge group and $C_i^9$ orientifold massless fields. However one finds that the candidate duals in fact violate the perturbative modular invariance constraints [37]. The heterotic duals of these class of orientifolds seem to correspond to some sort of non-perturbative $SO(32)$ heterotic orbifolds. Heterotic orbifolds of this type have been recently constructed in [40] (see Aldazabal’s contribution to these proceedings).

Realistic $D = 4, N = 1$ Type IIB orientifolds have not yet been constructed but it is certainly an interesting direction. The form of the Kahler potential for some orientifolds of the type discussed above have been obtained in ref. [37]. The models seem to have quite a different phenomenology to that of perturbative heterotic models in several respects, as the brief summary above shows. Furthermore they correspond by S-duality to non-perturbative heterotic vacua and may perhaps
teach us some non-perturbative secrets of perturbative heterotic models studied in the past.

5.4 M-theory compactifications on $\text{CY} \times S^1/\mathbb{Z}_2$

As we mentioned above, the strongly coupled limit of the heterotic $E_8 \times E_8$ string is M-theory compactified on a segment $(S^1/\mathbb{Z}_2)$ of length $\rho$. As $\rho \to 0$ one recovers the weakly-coupled heterotic string. The two boundaries of the segment correspond to the two $E_8$ factors of the heterotic which are purely ten-dimensional. Thus a way to obtain four-dimensional $N = 1$ vacua corresponding to strongly coupled heterotic models is to compactify M-theory on a Calabi-Yau $\times S^1/\mathbb{Z}_2$. In principle this is true but our lack of sufficient knowledge of the structure of M-theory at the moment only allow us to extract some qualitative (but nevertheless interesting) features. The low energy limit of M-theory is known to yield 11-dimensional supergravity. The massless sector of this theory is particularly simple, it includes a graviton $G_{MN}$, a gravitino $\Psi_M$ and an antisymmetric tensor $C_{MNP}$. When compactified on the segment $S^1/\mathbb{Z}_2$, $D = 10$, $E_8$ super-Yang-Mills fields are located at the two boundaries of the segment, $X_{11} = 0$ and $X_{11} = \rho$. The effective bosonic Lagrangian involving gravity and $E_8$ fields has then the general form

$$L = -\frac{1}{2\kappa^2} \int dX_{11} \sqrt{g} R - \sum_{i=1,2} \frac{1}{8\pi(4\pi\kappa^2)^{2/3}} \int dX_{10} \sqrt{g} \text{Tr} F_i^2 + \cdots \quad (5.9)$$

where $\kappa$ is the $D = 11$ gravitational constant and the sum corresponds to the two $E_8$ factors. Notice that the first integral extends over the full $D = 11$ space-time whereas the second is only ten-dimensional. If we compactify six dimensions on a Calabi-Yau space of volume $V$ and the 11-th dimension on a segment of length $\rho$, one gets a $D = 4$ action

$$L_4 = -\frac{1}{2\kappa^2 V\rho} \int dX_{11} \sqrt{g} R - \sum_{i=1,2} \frac{1}{8\pi(4\pi\kappa^2)^{2/3}} V \int dX_{10} \sqrt{g} \text{Tr} F_i^2 + \cdots \quad (5.10)$$

One can then identify the dependence of the 4-dimensional gravitational Newton constant $G_N$ and gauge fine-structure constant $\alpha_{GUT}$ on the different parameters and scales:

$$G_N = \frac{\kappa^2}{16\pi^2 V\rho} \quad ; \quad \alpha_{GUT} = \frac{(4\pi\kappa^2)^{2/3}}{2V} \quad (5.11)$$

We will see in the next section how the dependence of $G_N$ on the length $\rho$ may provide an interesting alternative to the problem of gauge coupling unification in
perturbative heterotic strings described in the second chapter. Doing a more careful dimensional reduction one can obtain the qualitative behaviour of this strongly coupled heterotic limit with respect to the equivalent to the dilaton $S$ and overall modulus $T$ in this approach. One finds a Kahler potential and gauge kinetic functions

$$K = -\log(S + S^* - \epsilon |C|^2) - 3\log(T + T^* - |C|^2) \quad (5.12)$$

$$f_{1,2} = S \pm \epsilon T$$

where the $C$ fields are charged matter scalars and the subindices 1, 2 correspond to the gauge groups $E_6$ and $E_8$ relevant for the standard gauge embedding. $\epsilon$ is a model-dependent constant which is not expected to be particularly small. Those familiar with effective Lagrangians in perturbative string theories probably have noted how similar to the perturbative result the above formulae are. There are a couple of differences though [57]: 1) the $C$-dependent piece in the $S$-term of the Kahler potential and 2) the $T$-dependent pieces in the gauge kinetic functions. Although similar corrections are in fact present [58] at the one-loop level in perturbative vacua (see eq. (2.3)), unlike in that case, here the coefficient $\epsilon$ may be of order one [59]. Furthermore, the $S$ and $T$ scalar fields above are defined as [57]

$$S = R^6 + i\sigma + \frac{\epsilon}{2}|C|^2; \quad T = \rho R^2 + i\eta + \frac{1}{2}|C|^2 \quad (5.13)$$

where $R = V^{1/6}$ is the overall CY compactification radius and $\sigma, \eta$ are axion-like fields. Notice that, whereas in perturbative heterotic vacua the size of the CY manifold was essentially given by $\text{Re} T$, in this non-perturbative limit it is given by the $S$ field. Notice also that this non-perturbative exchange of the roles of $S$ and $T$ appeared also when we discussed a class of $D = 4, N = 1$ orientifolds.

The full meaning of these results and their applications to different phenomenological questions like supersymmetry breaking are at present being the subject of intense research [60] by different groups and have been reported by several speakers at this conference. I forward the patient reader to their contributions for more details and complete lists of references.
6 Gauge coupling unification: a hint of non-perturbative dynamics?

I mentioned at the beginning of this talk the gauge unification problem of perturbative heterotic unification. Let me now be a bit more explicit. The gauge and gravitational bosonic terms in the $D=10$ effective Lagrangian from the perturbative heterotic string are

$$L_{10} = -\int dX_{10} \sqrt{g} e^{-2\phi} \left( \frac{4}{\alpha'^4} R + \frac{1}{\alpha'^3} \text{Tr} F^2 + \cdots \right)$$  \hspace{1cm} (6.1)$$

where $\phi$ is the dilaton field and $\alpha'$ is the inverse of the string tension. Upon dimensional reduction to four dimensions via CY compactification on a manifold of volume $V$ one arrives at an effective $D=4$ bosonic Lagrangian

$$L_4 = -\int dX_4 \sqrt{g} e^{-2\phi} V \left( \frac{4}{\alpha'^4} R + \frac{1}{\alpha'^3} \text{Tr} F^2 + \cdots \right).$$  \hspace{1cm} (6.2)$$

We then identify Newton’s constant and the gauge coupling \cite{56}:

$$G_N = \frac{e^{2\phi} \alpha'^4}{64\pi V}; \quad \alpha_{GUT} = \frac{e^{2\phi} \alpha'^3}{16\pi V}.$$  \hspace{1cm} (6.3)$$

Notice that one indeed has $G_N = \alpha_{GUT} \alpha'/4$, as we remarked in chapter 2. Consider now the mass of the massive gauge bosons in this theory. The lightest (massive) ones will have a Kaluza-Klein mass of order $M_{GUT} = 1/V^{1/6} = (\pi \alpha_{GUT}^4) / (4e^{2\phi} G_N^3)^{1/6}$. Now, if we want to remain in the perturbative regime we have to impose $e^{2\phi} \leq 1$, which then implies $M_{GUT}^2 \geq (\alpha_{GUT}^4)/G_N$. Plugging experimental numbers for $\alpha_{GUT}$ and $G_N$ one gets a lower bound \cite{56} $M_{GUT} \geq 10^{17}$ GeV, an order of magnitude larger than the value obtained by extrapolating low energy coupling data. More detailed analysis give results for the unification mass in perturbative heterotic unification around a factor 20 larger than one would expect on the basis of low energy data. This is the gauge coupling unification problem.

We already remarked that it is probably premature to say that this is indeed a problem and that several solutions have been put forward \cite{17}. But probably the most elegant one has been put forward by Witten in ref.\cite{56}. He has noted that lower unification scales may naturally appear in the context of strongly coupled heterotic compactifications. Let us first consider the case of the $SO(32)$ heterotic string. Semi-realistic models may be constructed with this heterotic string equally well as with the $E_8 \times E_8$ heterotic. Now, we know that the strongly coupled limit of the $SO(32)$ heterotic is the weakly coupled Type I string theory, whose gauge
group in $D = 10$ is also $SO(32)$. Let us now consider, as we did lines above, the
effective low energy bosonic Lagrangian but for the Type I case. It turns out that the
Lagrangian in $D = 10$ is identical to the heterotic case with a particular difference:
the gauge piece $F^2$ has an additional $\exp(\phi)$ factor in front. When going down to
four dimensions this extra factor gives rise to a $G_N$ analogous to the heterotic one
but $\alpha_{GUT}$ gets an extra factor $\exp(-\phi)$ compared to the heterotic result and hence
we get

$$G_N = \frac{e^2 \alpha_{GUT} \alpha'}{4}$$

(6.4)

instead of the perturbative heterotic result $G_N = \alpha_{GUT} \alpha'/4$. Now we have an extra
parameter to play and we can decouple the string scale $M_{string} = (\alpha')^{-1/2}$ from
the Planck scale $G_N^{-1/2}$. In particular one can simultaneously have e.g. $\alpha_{GUT} =
1/24$, $M_{Planck} = 10^{18}$ GeV and $M_{string} = 10^{16}$ GeV with a compactification scale
$M_c \propto 4 \times 10^{16}$ GeV. All this is possible remaining in the Type I perturbative regime
($\exp(\phi) \propto 10^{-3}$). In this case one identifies $M_{GUT}$ with $M_{string} \propto 10^{16}$ GeV, in
agreement with experiment. Thus within the context of perturbative Type I vacua
one can naturally evade the gauge coupling unification problem of perturbative
heterotic strings. Since weakly coupled Type I string is $S$-dual to strongly coupled
$SO(32)$ heterotic, one can equivalently say that one can solve this gauge unification
problem by going to the strongly coupled limit of that heterotic string.

The above discussion applied to models based on the $SO(32)$ heterotic (or their
Type I dual). What about $D = 4$, $N = 1$ models based on the $E_8 \times E_8$ heterotic? We
already recalled in the previous chapter that the strongly coupled limit of the $E_8 \times E_8$
heterotic was M-theory compactified on a $CY \times S^1/Z_2$. We found above (eq.(5.11))
the relationship in this scheme between $G_N$, $\alpha_{GUT}$ and the 11-dimensional gravita-
tional coupling $\kappa$, the CY volume $V$ and $\rho$. Identifying the unification mass with
the CY Kaluza-Klein scale one obtains $M_{GUT} = V^{-1/6} = (\alpha_{GUT}/8\pi^2 G^2 N^2 \rho^2/3)^{1/2}$.
Plugging the phenomenologically appropriated $G_N$, $\alpha_{GUT}$ values one can get the
correct GUT scale $M_{GUT} = 10^{16}$ GeV by stting $1/\rho \propto 10^{14}$ GeV. This also fixes the
fundamental M-theory scale $M_M = \kappa^{-2/9} \propto 2 \times M_{GUT}$. Thus the overall structure of
scales is as follows: Around $10^{10}$ GeV one has the M-theory scale and only slightly
below one has the $M_{GUT}$ (or CY) scale. In between that scale and $10^{14}$ GeV the
world is five-dimensional but with some peculiar characteristics: gauge interactions
and charged matter fields are purely four-dimensional and live at both boundaries
($X_{11} = 0, \rho$) of the 5-th dimension. On the other hand, gravitational fields, dilaton
and moduli live in the bulk of the 5-dimensional space. Below this intermediate scale
\[ M_1 = 1/\rho \propto 10^{14} \text{ GeV} \] 
there is a \( N = 1, D = 4 \) field theory. This structure of mass scales may have important consequences for different phenomenological aspects like supersymmetry breaking and structure of SUSY breaking as well as the possible role of moduli as invisible axions \[ [60]. \]

All in all, the possibility of an underlying strongly coupled heterotic string provides an attractive understanding of the unification of coupling constants within string theory. It is important to remark that in both strongly coupled \( E_8 \times E_8 \) and \( SO(32) \) schemes the GUT scale \( M_{\text{GUT}} \) is a scale at which also extended objects (M-theory in the first case, Type I strings in the second) appear. Thus in both schemes one expects \[ [60, 61] \] the appearance of operators of dimension > 4 suppressed by powers of \( (1/M_{\text{GUT}}) \) instead of \( (1/M_{\text{Planck}}) \) as was previously thought for perturbative heterotic unification. In particular, within the context of the MSSM one would expect the generation of baryon number violating dimension five operators like \[ [QQQL]_F. \] If these operators are only suppressed by \( (1/M_X) \), they would yield proton decay at rates excluded by present proton stability limits. Thus in the schemes discussed above there must be additional (e.g., discrete) symmetries which forbid this kind of dangerous operators.

### 7 Outlook

Our improved understanding of string theory from the different duality connections, give us both new insights into possible new dynamics as well as new possibilities for string vacua. There are new possibilities for obtaining potentially realistic \( D = 4, N = 1 \) string models embedding and unifying with gravity the standard model physics. Several general properties of perturbative heterotic unification which were thought to be general predictions of string theory turn out in fact to be predictions only of perturbative string theory. There may be new non-perturbative transitions which change the number of chiral generations. The gauge groups which can appear in string vacua may be quite large (for very particular configurations in the moduli space of the scalars in each model). \( D = 4, N = 1 \) vacua may have in general more than one anomalous \( U(1) \) and there may be more than one scalar field behaving like a dilaton.

The perturbative heterotic four-dimensional strings constructed in the last ten years constitute a particular class of string vacua in which conformal field theory techniques are applicable, but new non-perturbative vacua exist which remain to be
explored. In particular, Type II string theory which was essentially ignored in the past for model building purposes, may give rise to interesting $D = 4$, $N = 1$ vacua in terms of orientifolds or F-theory techniques. Although the underlying theory (M-theory) is unique, different techniques may be more appropriate to study different corners in the space of theories. For example, depending on the particular vacuum or coupling regime, the strongly coupled limit of a given perturbative $D = 4$, $N = 1$ heterotic vacuum may be reached in terms of a Type II orientifold or/and F-theory compactified on a CY four-fold or/and M-theory compactified on $CY \times S^1/\mathbb{Z}_2$. A systematic construction of this class of models is still to be done.

Apart from new vacua, the duality symmetries relate the strongly coupled limit of some theories to the weak coupling limit of other theories. This has been used in the last few years to extract non-perturbative information for some classes of $D = 4$, $N = 2$ vacua (like the exact prepotential). There is the hope to extend the use of dualities to gain knowledge of non-perturbative physics on some $N = 1$ (and perhaps $N = 0$) vacua. This could eventually give us a hint on how is broken the vacuum degeneracy (large number of apparently consistent perturbative string vacua). The fact that M-theory is a unique theory and all vacua seem non-perturbatively connected give us hope that the dynamics should determine things like the number of chiral generations (hopefully three!). In any event, a new epoch starts for string phenomenology and it will take us some time to figure out the full structure of the new possibilities which are open to us. The fact that gauge coupling unification takes place at $10^{16}$ GeV more naturally within the context of a strongly coupled heterotic string could be the first evidence in favour of a non-perturbative string unification of all interactions.

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