The second string (phenomenology) revolution

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Abstract

In the past three years our views on how the standard model of particle physics could be embedded into string theory have dramatically changed. The heterotic string is no longer the only possibility for such an embedding and other perturbative (or non-perturbative) corners of M-theory, like Type I or Type II strings seem now possible. It has also been realized that the string scale $M_s$ is not necessary close to the Planck scale and could be much smaller, of order the intermediate scale $\sqrt{M_W M_p}$ or even close to the weak scale. In addition, semi-realistic three generation models have recently been constructed starting with Type IIB compact orientifolds. I briefly discuss some of these developments which represent a revolution in our understanding of string phenomenology.

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1 Introduction

Breakthroughs in string theory are often named string revolutions. The first took place around 1984 when the cancellation of anomalies in Type I string theory was shown by Green and Schwarz and the two hererotic strings were constructed \[1\]. The second revolution around 1994 when the importance of weak/strong coupling dualities \[2\] were generally appreciated, leading to a unification of all five supersymmetric string theories with 11-D supergravity \[3\]. If string theory is correct, it should contain as a certain (low energy) limit the observed standard model (SM) of particle physics. String phenomenology \[4\] is the study of the possible avenues by which the $SU(3) \times SU(2) \times U(1)$ SM could be embedded in string physics. The string revolutions also imply string phenomenology revolutions, which typically take place a few years later. This is because it takes some time to construct explicit semi-realistic string vacua with the new techniques. We are now living the second string phenomenology revolution. Progress is sometimes realizing what we do not know but thought we knew. This is now the case: we have realized that we do not even know what is the fundamental energy scale of string theory! To see what has changed in the last three years or so, let me remind you the pre-1995 orthodoxy in string phenomenology \[4\]:

i) Only the heterotic strings (mostly $E_8 \times E_8$) where considered as contenders for the unified theory of all interactions. Four dimensional string vacua with one unbroken SUSY are considered.

ii) The string scale $M_s$ is tied up to the Planck scale $M_p$ by the heterotic string expression $M_s = g M_p$. Thus string physics has its realm at energies of order $10^{17}$ GeV.

iii) The weak scale $M_W$ is generated by some non-perturbative effect like gaugino condensation in a hidden sector of the theory. If $\langle \lambda \rangle \propto M_W M_p^2$ the required SUSY-breaking soft terms $\propto M_W$ would be produced for the SUSY SM particles.

iv) Gauge coupling constants of the MSSM unify at a scale $M_X = 2 \times 10^{16}$ GeV, only slightly below the string scale.

A few examples of semi-realistic four-dimensional string models \[4\] with three generations of quarks and leptons have been constructed using different techniques ( $Z_N$ orbifolds, free fermionic constructions, Calabi-Yau compactifications etc.). This is by itself an important achievement since these models represent the first unified models of all interactions including gravity. There is however a loss of predictivity in the models because in order to make contact with the SM (or the MSSM) one has to abandon the string theory techniques and analize at the level of the effective field theory the flat directions in each model. Thus possible predictions like quark/lepton masses or proton stability become dependent on the flat field direction chosen. And of course, the questions of vacuum degeneracy and dilaton/moduli stabilization (not to mention the cosmological constant) remain to be solved.

Let us enumerate now some of the important points (from the phenomenological point of view) which have changed in the last few years:

i) End of the heterotic monopoly: other string theories like Type I or Type II or techniques like F-theory provide for new classes of $D = 4$, $N = 1$ string vacua which can lead to
new avenues to embed the observed physics.

ii) It has been realized that the number of extra dimensions felt by gauge fields and gravity fields may be in general different. Indeed, gauge fields in e.g., Type I theory live on the world-volume of D-branes \[5\] which span a number of dimensions often smaller than the full ten dimensions felt by the gravity fields \[6, \ 7, \ 8, \ 9\].

iii) As a consequence of ii), the string scale \(M_s\) has no theoretical bound and we only have the phenomenological bound \(M_s \geq 1\) TeV \[7, \ 8, \ 9\].

Concerning the first point, we are only starting to scratch the space of vacua offered by the new techniques. Indeed, with the new techniques available, Type I and Type IIB strings are as good as the heterotic models from the point of view of model building (see in particular ref. \[10\]). On the other hand, any of the string theories are just perturbative limits of the underlying unique M-theory. So the question is whether some corner of the M-theory moduli space sits sufficiently close to the observed SM physics. I am going to concentrate here on the case of the \(D = 4\) vacua obtained from Type I theory (or equivalently, Type IIB \(D = 4\) orientifolds \[11, \ 12, \ 13\]). These are interesting because they are examples of explicit, perturbative string vacua in which many of the new features in model building (like e.g., the possibility of a reduced string scale, large gauge groups, multiple pseudoanomalous \(U(1)\)'s etc.) are already present. Although Type I strings were discovered well before the heterotic strings, little effort has been devoted in the past to the construction of Type I four-dimensional string vacua. Of course, smooth Calabi-Yau compactifications of Type I are possible consistent solutions but they have no phenomenological interest and, anyway, present no obvious advantage over the heterotic \(SO(32)\) theory. On the other hand the concept of D-brane \[5\] has provided us with a new understanding of Type I string theory.

2 Type IIB \(D = 4\), \(N = 1\) compact orientifolds

Let us describe a bit how these orientifolds \[14, \ 13, \ 16\] are constructed. In a \(D = 4\) Type IIB orientifold, the toroidally compactified theory is divided out by the joint action \[14, \ 13, \ 17\] of a discrete symmetry group \(G_1\), like \(Z_N\) together with a world sheet parity operation \(\Omega\), exchanging left and right movers. The \(\Omega\) action can be accompanied by extra operations thus leading to a generic orientifold group \(G_1 + \Omega G_2\) with \(\Omega h \Omega h' \in G_1\) for \(h, h' \in G_2\). We will consider here the cases \(G_1 = G_2\) and \(G_1 = Z_N\) and such that \(D = 4\) \(N = 1\) theories are obtained, when the twist \(\Omega\) is performed on Type IIB compactified on \(T^6/G_1\). The allowed orbifold groups, acting crystallographically on \(T^6\) leading to \(N = 1\) unbroken supersymmetry were classified in \[18\]. The finite list of twists is \(Z_3, Z_4, Z_6, Z_6', Z_8, Z_8', Z_{12}\) and \(Z_{12}'\). Here the primed twists correspond to different implementations of discrete rotations of the given example.

Orientifolding the closed Type IIB string introduces a Klein-bottle unoriented world-

\[1\]See talk by B. Ovrut for an alternative which embeds the SM in strongly coupled heterotic theory.
sheet. Amplitudes on such a surface contain tadpole divergences. In order to eliminate such unphysical divergences Dp-branes must be generically introduced. In this way, divergences occurring in the open string sector cancel up the closed sector ones and produce a consistent theory. For $Z_N$, with $N$ odd, only D9-branes are required. They fill the full space-time and six dimensional compact space. For $N$ even, $D5_s$-branes, with world-volume filling space-time and the $k^{th}$ complex plane, may be required. This is so whenever the orientifold group contains the element $\Omega O_i O_j$, for $k \neq i,j$. Here $O_i \ (O_j)$ is an order two twist of the $i^{th} \ (j^{th})$ complex plane.

One denotes (see ref.[12] for conventions ) open string states by $|\Psi, ab\rangle$, where $\Psi$ refers to world-sheet degrees of freedom while the $a, b$ Chan-Paton indices are associated to the open string endpoints lying on D$p$-branes and D$q$-branes respectively. These Chan-Paton labels must be contracted with a hermitian matrix $\lambda^{pq}_{ab}$ which parametrize the gauge indices. The action of an element of the orientifold group on Chan-Paton factors is achieved by a unitary matrix $\gamma_{g,p}$ such that $g: \lambda^{pq}_{ab} \rightarrow \gamma_{g,p} \lambda^{pq}_{ab} \gamma_{g,p}^{-1}$. We denote by $\gamma_{k,p}$ the matrix associated to the $Z_N$ orbifold twist $\theta^k$ acting on a Dp-brane. A generic matrix $\gamma_{1,p}$ can be written as $\gamma_{1,p} = (\tilde{\gamma}_{1,p}, \tilde{\gamma}_{1,p}' )$ where $*$ denotes complex conjugation. $\tilde{\gamma}$ is a $N_p \times N_p$ diagonal matrix given by

$$\tilde{\gamma}_{1,p} = \text{diag}(\cdot \cdot \cdot, \alpha^{N_p}I_{n^p_1}, \cdot \cdot \cdot, \alpha^{N_p}I_{n^p_p})$$

(1)

with $\alpha = e^{2i\pi/N}$. $V_j = \frac{i}{N}$ with $j = 0, \ldots, P$ corresponds to an action “with vector structure” $(\gamma^N = 1)$ while $V_j = \frac{2j-1}{2N}$ with $j = 1, \ldots, P$ describes an action “without vector structure” $(\gamma^N_{1,p} = -1)$. The gauge fields living on the world-volume of a D$p$-brane have associated Chan-Paton factors $\lambda^p$ corresponding to the gauge group $G_p$ with $G_9 = SO(2N_9)$ and $G_5 = Sp(2N_5)$. In Cartan-Weyl basis such generators are organized into charged generators $\lambda_a = E_a, a = 1, \ldots, \dim G_p - \text{rank} G_p$, and Cartan algebra generators $\lambda_I = H_I, I = 1, \ldots, \text{rank} G_p$, such that $[H_I, E_a] = \rho_I^a E_a$ where the $(\text{rank} G_p)$-dimensional vector with components $\rho_I^a$ is the root associated to the generator $E_a$.

The matrices $\gamma_{1,p}$ and its powers represent the action of the $Z_N$ group on Chan-Paton factors, and they correspond to elements of a discrete subgroup of the Abelian group spanned by the Cartan generators. Hence, we can write $\gamma_{1,p} = e^{-2i\pi V^p_H}$. Thus, this equation defines a $(\text{rank} G_p)$-dimensional vector $V^p$ with coordinates corresponding to the $V_j$'s defined in (1) above. In such a description the massless states are easily found $\gamma_{1,p}$. Let us consider the case in which all 5-branes sit at the origin. In the $(pp)$ sector the gauge group is obtained by selecting the root vectors satisfying $\rho^a \cdot V^p = 0 \text{ mod } \mathbb{Z}$ while matter states correspond to charged generators with $\rho^a \cdot V^p = v_i \text{ mod } \mathbb{Z}$. Here the $v_i$ are the eigenvalues of the $Z_N$ rotation of the three complex compact dimensions (see ref.[12]). In the (95) sector the subset of roots of $G_9 \times G_5$ of the form $P_{(95)} = (W_{(9)}; W_{(5)}) = (\pm 1, 0, \ldots, 0; \pm 1, 0, \ldots, 0)$ must be considered. Matter states are obtained from the projection $P_{(95)} \cdot V^{(95)}_{(95)} = (s_j v_j + s_k v_k) \mod \mathbb{Z}$ with $s_j = s_k = \pm \frac{1}{2}$, plus (minus) sign corresponding to particles (antiparticles) and $V^{(95)} = (V^9; V^5)$.

Twisted tadpole cancelation turns out to be quite restrictive in these models. In particular, it was shown in ref.[12] that the only tadpole-free $Z_N$ orientifolds are $Z_3, Z_6, Z_6', Z_7$. 
and \( Z_{12} \). Furthermore, for all those models (except \( Z_6 \) and \( Z_{12} \)) tadpole conditions fix completely the gauge group and massless spectrum (modulo Wilson lines and/or moving of 5-branes). All these models lead to \( N = 1, D = 4 \) consistent vacua with a chiral anomaly-free spectrum.

Instead of \( \Omega \) one can use other \( Z_2 \) modings which are still consistent with \( N = 1 \) SUSY in \( D = 4 \). Thus, for example one can use as orientifold projector \((-1)^{F_L} \Omega O_i \) or \((-1)^{F_L} \Omega O_i O_j O_k \), \( i \neq j \neq k \neq i \). Here \( F_L \) is the world-sheet left-handed fermion number. In this case tadpole cancellation conditions will require in general the presence in the vacuum of 7-branes and 3-branes respectively. There may be three different types of 7-branes, \( 7_i \), \( i = 1, 2, 3 \) depending what complex dimension \( X_i \) is transverse to the 7-brane world-volume. Thus we see that, depending on the orientifold generators, one can deal with 3-branes, 5-\( i \)-branes, 7-\( i \)-branes and 9-branes. Not all types may be present simultaneously if we want to preserve \( N = 1 \) in \( D = 4 \). For a given \( D = 4, N = 1 \) vacuum with D-p-branes and D-p’-branes one must have \( (p - p') = 0, \pm 4 \). The number of each type of p-brane in each case is dictated by tadpole cancellation constraints. These in turn guarantee the cancellation of gauge anomalies in the effective \( D = 4, N = 1 \) theory. T-dualities relate the different types of p-branes present in each given vacuum \([5]\). Consider for simplicity the 6-torus as the product of three two-tori, \( T_6 = T^2 \times T^2 \times T^2 \) each with compact radii \( R_i \), \( i = 1, 2, 3 \). Now, it is well known that a duality transformation \( R_i \rightarrow \alpha' / R_i \) transforms Neumann boundary conditions on the \( X_i \) coordinate into Dirichlet boundary conditions and vice versa \([4]\). This means that e.g., a 9-brane will turn into a 7-\( i \)-brane and vice versa under this transformation. Thus given any configuration with certain distribution of p-branes in the vacuum, there are a number of equivalent configurations which are obtained from T-dualities.

Given a p-brane in a background with six compact dimensions, open strings ending on that p-brane will only have Kaluza-Klein (KK) states along the compact dimensions with Neumann boundary conditions. On the contrary, it will have winding states only in those compact directions with Dirichlet boundary conditions. On the other hand, closed strings can have both KK and winding modes in all compact dimensions. This turns out to be important in order to study the structure of mass scales in the theory.

### 3 The string scale and the Planck mass

Let us study the relationship between string, Planck and compactification scales in Type I \( D = 4 \) strings of the type described above. We consider the dimensional reduction down to four dimensions obtained by compactification on an orbifold with an underlying compact torus of the form \( T^2 \times T^2 \times T^2 \). The three tori are taken with volumes \((2\pi R_i)^2\), \( i = 1, 2, 3 \) respectively. The relevant piece of the bosonic action of the \( D = 4, N = 1 \) effective Lagrangian for a generic distribution of D-branes has the form:

\[
S_4 = -\int \frac{dx^4}{2\pi} \sqrt{-g} \left( \frac{R_2^2 R_2^2 R_3^2 M_s^8}{\lambda_f^4} \right) R + \frac{R_2^2 R_2^2 R_3^2 M_s^6}{\lambda_f} \frac{1}{4} F^{(2)}
\]
where $\lambda_I$ is the $D = 10$ Type I dilaton, $M_s = 1/\sqrt{\alpha'}$ is the Type I string scale and we have displayed the kinetic terms for gauge bosons of the different groups which may come from the different p-branes, $p = 9, 7, 5, 3$. As discussed above, not all the different p-brane sectors should be present in the vacuum if we want to respect $N = 1$ SUSY. From the above equation one obtains for the gravitational coupling

$$G_N = \frac{1}{M_{\text{Planck}}^2} = \frac{\lambda_I^2 M_s^4 M_3^2}{8 M_s^8}$$

and for the gauge couplings $\alpha_p$ for the different p-branes:

$$\alpha_9 = \frac{\lambda_I M_1^2 M_2^2 M_3^2}{2 M_s^6} \quad ; \quad \alpha_7 = \frac{\lambda_I M_j^2 M_k^2}{2 M_s^4}, \quad i \neq j \neq k \neq i$$

$$\alpha_5 = \frac{\lambda_I M_i^2}{2 M_s^2} \quad ; \quad \alpha_3 = \frac{\lambda_I}{2}$$

where $M_i = 1/R_i$. From the above formulae we observe that, unlike what happens in the heterotic case, $M_{\text{Planck}}$ and $M_s$ do not need to be of the same order of magnitude [$6$.]

Consider for example the simple isotropic case in which all compactification radii are taken equal, $R_i = R = 1/M_c$. Then one gets [$3, 7, 8, 9, 20, 29$]

$$M_{\text{Planck}}^2 = \frac{8 M_s^8}{\lambda_I^2 M_c^6} \quad \alpha_p = \frac{\lambda_I}{2} (\frac{M_c}{M_s})^{p-3}, \quad p = 9, 7, 5, 3$$

that combined give the following relationship

$$\frac{M_c^{(p-6)}}{M_s^{(p-7)}} = \frac{\alpha_p M_{\text{Planck}}}{\sqrt{2}}$$

Notice that in principle these equations give us a certain freedom to play with the values of the Type I string scale $M_s$ and the compactification scale $M_c$. This is to be compared to the analogous equation in the perturbative heterotic case where the relation $M_{\text{string}} = \sqrt{\alpha_X/8} M_{\text{Planck}}$ fixes the value of the string scale independently of the compactification scale.

Consider for example the case of a set of 3-branes in an isotropic compactification. One then has: $M_s^4 = \frac{\alpha_X}{\sqrt{2}} M_c^3 M_p$. Thus one can e.g. lower the string scale down to e.g. 1 TeV (which is the lower phenomenological bound) by choosing $M_c = 10$ MeV [$8, 9$]. Notice that the compactification scale $M_c$ may be this low because the charged fields living on the 3-branes have no KK excitations and hence there is no charged threshold at the $M_c$ scale. Thus the lesson that we learn from these considerations is that we do not really know what the string scale is! This is quite surprising because one of the most widely spread dogmas about string theory is that its natural scale is the Planck scale [$8$].

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$^2$Recently an interesting alternative has been suggested [$21$] in which gravitons are trapped on domain walls and the extra dimensions are non-compact.
So, what is the string scale? There are a number of natural options for the string scale $M_s$:

**i)** $M_s \approx M_{\text{Planck}}$. As we said, this is the option which is forced upon us in perturbative heterotic vacua [4]. In this case $M_c \approx M_I$ and gauge coupling unification should take place also about the same scale. There is a slight problem here since plane extrapolation of the MSSM gauge couplings indicate unification at $2 \times 10^{16}$ GeV, which is about a factor 20 too small compared to the string scale. This is naturally solved identifying $M_s$ with the unification scale, which is option ii):

**ii)** $M_s \approx M_X$ [6]. Here $M_X$ is the GUT scale or the scale at which the extrapolated gauge couplings of the minimal supersymmetric standard model (MSSM) join. Numerically this is of order $10^{16}$ GeV. This corresponds (for the 3-brane case) to choices for $M_c$ only slightly below $M_s$.

**iii)** $M_s \approx \sqrt{M_W M_{\text{Planck}}}$. This is the geometrical intermediate scale $\approx 10^{11}$ GeV which coincides with the SUSY-breaking scale in models with a hidden sector and gravity mediated SUSY breaking in the observable sector. The interest of this choice has been recently pointed out in ref. [23] (see also ref.[22]) and recently explicit semirealistic D-brane models consistent with this structure of mass scales have been constructed [10]. An interesting point of this possibility occurs in the case of an isotropic compactification with $M_1 = M_2 = M_3 = M_c$ with the SM embedded into 3-branes. If there is a SUSY-breaking 3-brane sector which is away in transverse space from the 3-branes containing the SM, the non-SUSY brane sector acts as a standard hidden SUSY-breaking sector. In this case one expects $M_{\text{soft}} = M^2_c/M_p = \alpha^2_3/(M_c/M_s)^6$, where $M_{\text{soft}}$ is the scale of SUSY-breaking felt by the SM fields. Thus in order to have $M_{\text{soft}} \propto M_W$ it is enough to have $M_c/M_s \approx 0.01$, not very large relative scales are needed [23]. Thus the $M_W/M_p$ hierarchy may be naturally generated without any need for mechanisms like gaugino condensation.

**iv)** $M_s \approx 1$ TeV. This is the 1 TeV string scenario considered in refs. [7, 8, 24]. In this case it should be (in an isotropical 3-brane situation) $M_c/M_s \approx 10^{-5}$. This case is potentially very exciting since the string scale could perhaps be testable at accelerator energies. This case has been discussed by other speakers at this meeting [24].

We would like to argue that, within the context of D-brane models the third and fourth options are phenomenologically safer [23]. The reason is simple. A generic string theory background will contain in general D-brane systems with different amounts of supersymmetry. Thus the SM model may perhaps be embedded into a brane system which has $N = 1$ supersymmetry. This could guarantee that the standard gauge hierarchy problem is solved. However, generically there will be some other brane systems away in transverse directions which will have no supersymmetry ($N = 0$). If this is the case, the massless closed string states living in the bulk of extra dimensions couple to both the $N = 1$ brane sector (where the SM is contained) and the SUSY-breaking $N = 0$ brane sector. Thus these closed string states will transmit supersymmetry breaking to the SM sector supressed by the Planck mass and of order $M^2_c/M_p$. If we want the magnitude of these soft terms to be not bigger than
the weak scale $M_W$ (so that the hierarchy problem does not reappear) one needs to have:

$$M_s \leq \sqrt{M_W M_p} \propto 10^{11} \text{ GeV}$$  \hspace{1cm} (7)$$

Thus, from the phenomenological point of view is safer to have $M_s \leq 10^{11} \text{ GeV}$ in order to avoid too big SUSY-breaking effects from generic $N = 0$ brane sectors. Of course, this is not a theorem, but is quite suggestive. In fact, we will describe below a class of orientifolds which provide a realization of this constraint. Let me finally emphasize that in the cases iii) and iv) above one will have to eventually find a mechanism to explain why some of the compact dimensions became large compared to the string size. It is this large size of some dimensions which gives rise eventually to a $M_W/M_p$ hierarchy.

4 The gauge coupling unification problem

If options iii) or iv) above are correct, the string scale would be well below the standard grand unification scale $M_X = 2 \times 10^{16} \text{ GeV}$ where the unification of the $SU(3) \times SU(2) \times U(1)$ couplings takes place when the minimal SUSY standard model spectrum is assumed to hold. Thus we have to face the following two important problems:

1) Couplings should unify at a scale $M_s \leq 10^{11} \text{ GeV}$. How we make this compatible with the fact that the MSSM gauge couplings seem to unify at a much larger scale of order $10^{16} \text{ GeV}$?

2) Baryon and lepton number violating operators are suppressed only by inverse powers of $M_s$, which is now much lower than in the conventional heterotic scenario. Thus unless appropriate symmetries are present it is difficult to understand the level of proton stability indicated by the experimental limits.

Concerning the first point, one has to emphasize that a detailed knowledge of the gauge kinetic functions of the SM gauge groups is really required in order to check whether coupling unification is still possible. Concerning the second point, one has to study whether appropriate symmetries could be present in order to suppress sufficiently the operators violating proton stability (see ref. [25]). For these two questions it turns out to be relevant the study of the pseudo-anomalous $U(1)$ gauge symmetries which are generically present in the class of Type IIB orientifold models mentioned above.

4.1 Anomalous $U(1)$’s and mirage unification

Indeed, if one computes the $U(1)$ triangle anomalies in Type IIB $D = 4$ orientifolds one finds that most of the $U(1)$’s are anomalous. This is not new in string theory: it is well known that in heterotic string vacua there are analogous $U(1)$ symmetries whose triangle anomalies are cancelled by a $D = 4$ version of the Green-Schwarz mechanism [26]. There is however a couple of important differences between the Type I and heterotic cases. In the heterotic case there is only one anomalous $U(1)$ and its mixed anomaly with all the non-Abelian gauge
groups is identical. This is because there is a single field (the complex dilaton $S$) giving rise to the GS mechanism. In addition a Fayet-Iliopoulos (FI) term of order $g^2 M_p^2 / 16\pi^2$ appears at one-loop. The latter are of order the string scale. In the Type IIB orientifold models the story is quite different. One finds that \cite{27}

i) There are multiple anomalous $U(1)$’s.

ii) The mixed anomalies of the $U(1)$’s with the different gauge factors is non-universal.

iii) It is the twisted moduli fields $M_k$ which participate in the GS mechanism, instead of the complex dilaton $S$.

iv) There appear FI-terms which are proportional to $<\text{Re}M_k>$, which are the ”blowing-up” fields of the orbifold singularities. Thus, unlike the heterotic case, the FI terms may be arbitrarily small \cite{27, 28}.

More specifically, cancellation of $U(1)$ anomalies results \cite{27} from the presence in the $D = 4, N = 1$ effective action of the term

$$\sum_k \delta^l_k B_k \wedge F_{U(1)_l}$$

where $k$ runs over the different twisted sectors of the underlying orbifold (see ref.\cite{27} for details) and $B^k$ are the two-forms which are dual to the imaginary part of the twisted fields $M_k$. Here $l$ labels the different anomalous $U(1)$’s and $\delta^l_k$ are model-dependent constant coefficients. In addition the gauge kinetic functions have also a (tree-level) $M_k$-dependent piece:

$$f_\alpha = S + \sum_k s^k_\alpha M_k$$

where the $s^k_\alpha$ are model dependent coefficients. Under a $U(1)_l$ transformation the $M_k$ fields transform non-linearly $\text{Im}M_k \rightarrow \text{Im}M_k + \delta^l_k \Lambda_l(x)$. This non-linear transformation combined with eq.\(8\) results in the cancellation of the $U(1)$ anomalies as long as the coefficients $C^l_\alpha$ of the mixed $U(1)_l$-$G^2_\alpha$ anomalies are given by

$$C^l_\alpha = - \sum_k s^k_\alpha \delta^l_k$$

Unlike the perturbative heterotic case, eq.\(10\) does not in general require universal mixed anomalies.

Now, equation \(8\) shows us an interesting point \cite{29, 30} : the gauge coupling constants at the string scale in this class of theories are only unified if one sits precisely at the orbifold points with $<\text{Re}M_k> = 0$. Also, if $<\text{Re}M_k> \neq 0$ the corrections are group dependent and not universal. Thus consider a simplified scenario in which we had only a single blowing up field $M$ so that $f_\alpha = S + s_\alpha M$. Consider now the renormalization group running of gauge couplings $g_\alpha$ from the weak scale to the string scale $M_s$:

$$\frac{4\pi}{g^2_\alpha(M_W)} = \text{Re} f_\alpha + \frac{b_\alpha}{2\pi} \log \frac{M_s}{M_W}$$

\(11\)
where $f_\alpha$ is the gauge kinetic function in eq. (9). We know that with the particle content of the MSSM coupling unification works nicely for a unification scale $M_X = 2 \times 10^{16}$ GeV. Thus if we had a model with:

$$s_\alpha = \gamma b_\alpha ; \quad \langle ReM \rangle = \frac{1}{\gamma 2\pi} \log(M_X/M_s)$$

(12)

we would nicely get (apparent) gauge coupling unification. This possibility may be named "mirage unification" because from a low-energy observer, everything looks like if there was just standard coupling unification at $M_X$ (for approaches similar in spirit see also refs.31). In fact what happens is that there are finite corrections to the gauge couplings at $M_s$ (which may be much smaller than $M_X$) which precisely mimic the effect. It turns out that there are indeed 29, 80 orientifolds in which in some simple cases one can have $s_\alpha \propto \beta_\alpha$ (e.g., the $Z_3$ and $Z_7$ orientifolds). However in those cases, the study of the scalar potential (including the FI terms) tells us that $< M > = 0$ at the minima with unbroken non-Abelian gauge group 30, 32. Thus in these orientifold examples the gauge couplings seem to unify at $M_X$, mirage unification does not occur. Nevertheless, odd orientifolds like these are very special and it could well be that in more general situations the vacua may sit at points with $< ReM > \neq 0$. Notice also that once SUSY is broken large non-vanishing D-terms will in general be allowed and $< ReM >$ may move from a vanishing value at the SUSY minimum to a non-vanishing one after SUSY-breaking effects are taken into account.

### 4.2 Precocious gauge coupling unification

If mirage unification as above does not occur and couplings join at the string scale, one has to find an explanation for "precocious" coupling unification at a scale $M_s << 10^{16}$ GeV. The simplest and most conservative possibility is to abandon the particle content of the MSSM and assume that there are extra massless charged particles beyond quarks, leptons and one set of Higgs fields. If e.g. $M_s \propto 10^{10} - 10^{12}$ GeV, it is easy to find extra sets of particles which can have this effect. In particular, a simple option is the addition of extra left-handed and right-handed leptons which were shown in ref.23 to be consistent with intermediate scale unification. This possibility could sound less natural than the MSSM paradigm with unification at $2 \times 10^{16}$ GeV but one should be more open minded and not look at that paradigm as the unique possibility. Let us remind that the MSSM structure is in some respects a bit artificial: whereas quarks and leptons come in three chiral copies, Higgsses come in only one (vector-like) copy. Furthermore, symmetries have to be imposed in order to insure sufficient proton stability. Interestingly enough it has been recently been found 10, 35 that in semi-realistic Type -I models there is a tendency to get extra massless leptons which could lead to gauge coupling unification at an intermediate scale (see below).
5 Standard-like models from Type I string vacua

It turns out to be quite difficult to construct semi-realistic $D = 4$ compact orientifolds with unbroken $N = 1$ supersymmetry. As we mentioned above, the tadpole cancellation constraints are so strong that there is little flexibility left for obtaining a realistic gauge group and three quark-lepton generations. One obvious direction in order to obtain more flexibility is to consider also non-supersymmetric vacua: after all the world is not (exactly) supersymmetric. This possibility was excluded in the past because of the hierarchy problem: if SUSY is broken at the string scale, with the latter close to the Planck scale, no scalar would survive radiative corrections and we would be left with no Higgs fields in order to break the $SU(2)_L \times U(1)$ symmetry.

Since we can now lower the string scale well below the Planck mass, the above exclusion of non-supersymmetric vacua must be reexamined. There are now two new possibilities opened: 1) Having a non-SUSY model with the string scale not much above the weak scale or 2) Having a non-SUSY model with the string scale of order the intermediate scale. In the latter case if the non-SUSY sector of the theory is only connected to the SM world by the exchange of bulk (closed string) fields, the hierarchy can in general be preserved.

Recently \cite{10,33} Type IIB, $D = 4$ compact orientifolds providing for explicit realizations of the above possibilities have been obtained for the first time. They are based on the observation \cite{33,34} that one can obtain tadpole- and tachyon-free configurations by adding brane-antibrane pairs to $N = 1$, $D = 4$ orientifolds. The presence of the anti-branes makes this kind of configuration non-supersymmetric. Let us present an specific orientifold example \cite{10} based on the construction in ref.\cite{34} yielding a semirealistic spectrum. It is based in the standard $Z_3$ orientifold constructed with 7-branes instead of 9-branes. One compactifies the Type IIB string on the standard $Z_3$ orbifold. The orientifold projector is given by $\Omega(-1)^F R_3$, where $R_3$ is the operation reflecting the third compact complex plane. Tadpole cancellation conditions require the presence of 32 7-branes with their worldvolume including the first two complex planes plus Minkowski space. Now, we embedd the $Z_3$ action into the 7-brane Chan-Paton factors by chosing \cite{10}:

$$V_7 = 1/3(1,1,1,-1,-1,0,0,0,1,1,1,1,1,1,1)$$ (13)

In addition a quantized Wilson line is added in the first complex plane given by:

$$W_7 = 1/3(1,1,1,1,1,1,0,0,0,0,0,0,0,0,0)$$ (14)

The gauge group from the $(77)$ sector will be $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times [U(1)^2 \times SO(4) \times U(7)]$. This contains the left-right symmetric extension of the SM which is a phenomenologically interesting model. One can also check that from the $(77)$ sector there are chiral fields transforming like:

$$3(3,2,1,1/3) + 3(\bar{3},1,2,-1/3) + 3(1,2,2,0)$$ (15)
under $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (plus extra fields transforming under the hidden group). These are three quark generations plus three sets of Higgs fields. Now, the nine fixed points under $Z_3$ in the first two complex planes split into three sets of three fixed points each which have associated twists $V, V + W$ and $V - W$ respectively. The corresponding value for $\text{Tr} \gamma_{\theta,7}$ are -4, -4 and -1. Now, for this orientifold tadpole cancellation conditions require $\text{Tr} \gamma_{\theta,7} = -4$. This means that we will have to add something on the three fixed points with shift $V - W$ in order to cancel tadpoles. It is easy to see that if we locate at each point two 3-branes with $\text{Tr} \gamma_{\theta,3} = -1$ all twisted tadpoles cancel. Now there are extra massless chiral fields from open strings extending between the 3-branes in the three fixed points and the 7-branes. They will have gauge quantum numbers under the $(77)$ gauge group:

\begin{equation}
(1, 2, 1, +1) + (1, 1, 2, -1) + (3, 1, 1, -2/3) + (\bar{3}, 1, 1, 2/3)
\end{equation}

plus extra fields transforming under the hidden group. These come in three copies (one per fixed point). Notice that in this sector three standard lepton generations appear. In addition there are three extra sets of vector-like colour triplets which turn out to get generically large masses (see [11, 35] for details). Thus we have easily constructed a three generation left-right symmetric model starting with the simplest $Z_3$ orientifold and adding appropriate numbers of 7-branes and 3-branes. This is, by the way, the simplest semi-realistic string model I have ever seen. Notice that this model has three generations of Higgs fields. This is a general trend in these constructions, there are typically extra massless weakly interacting fields. This is interesting because they lead to precocious gauge coupling unification, $SU(2)_L$ and $U(1)_Y$ interactions grow faster than in the MSSM and tend to join at an intermediate scale $10^8 - 10^{12}$ GeV with the $SU(3)$ coupling [10, 35]. Thus this class of models provide a natural alternative to the MSSM scenario in which couplings join close to the Planck mass.

This model is non-supersymmetric because there is an additional tadpole cancelation condition: the net-number of 3-branes minus anti-3-branes must be zero in this model. Thus there must be 6 anti-3-branes somewhere. Depending on the location of these extra anti-3-branes, the SUSY-breaking phenomenology is different. Let us locate for definiteness the 7-branes at the origin in the third complex dimension. Now, if the anti-3-branes are away from the origin in the third compact dimension, they have no overlap with the 7-branes and hence there are no massless chiral fields in the $(\bar{3} 7)$ sector. In this case the SUSY-breaking spectrum residing in this anti-3-branes can only communicate with the $(77)$ and $(3 7)$ sectors (which contain the observed physics) by the exchange of closed string fields which live in the bulk. The couplings of the latter are supressed by powers of the Planck mass. In this case the model behaves like the standard hidden sector SUSY-breaking models in which the role of hidden sector is played by the anti-3-branes. For this to work the string scale must be the intermediate scale. This can be made consistent with the observed Planck mass by e.g., choosing compactification scales $M_i$ along the three complex compact dimensions as follows: $M_1 \propto M_2 = \propto M_s \propto 10^{11}$ GeV and $M_3 \propto 1$ TeV. Alternatively, if the anti-3-branes are located at the origin in the third complex plane, their worldvolume will overlap with that of 7-branes and there will be a non-SUSY massless spectrum coupling to the SM gauge group.
In this case, if we do not want to have a hierarchy problem, one should lower the string scale down to $M_s \propto 1 - 10 \text{ TeV}$. This is again possible by choosing $M_1 \propto M_2 \propto M_s \propto 1 - 10 \text{ TeV}$ and $M_3 \propto 10^{-3} \text{ eV}$.

Since this class of models, although free of Ramond-Ramond tadpoles and tachyons, are non-supersymmetric, their stability should be farther studied. Notice however that this is a problem that we will have to face anyhow in any semirealistic model. In the traditional heterotic models one had to resort to field theory effects like gaugino condensation to break supersymmetry in a hidden sector, and this leads to the same questions that we face now in the non-supersymmetric Type IIB orientifolds. One of the advantages now is that those effects are produced by explicit anti-D-branes whose effects can in principle be better studied.

6 Outlook

We are witnessing at the moment something we could perhaps name (to follow the tradition) the second string (phenomenology) revolution. The heterotic string has lost its monopoly as the candidate for the unification of gravity and the standard model of particle physics. Although M-theory is supposed to be the unique underlying theory, one of the different perturbative limits like Type I, Type II and heterotics could perhaps be closer than the others to the observed physics. In the last fifteen years essentially only the heterotic string has been explored, with important (but limited) success. The first Type I semirealistic models are now starting to be built and they show some very interesting features compared to their heterotic precursors. One of them is the possibility that the string scale is much below the Planck mass, at the intermediate scale $\sqrt{M_W M_p}$ or even close to the weak scale. These are first steps in the search for realistic string vacua using D-brane techniques. Much work remains to be done both from the theoretical and more phenomenological sides.

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