Recent Developments in Superstring Phenomenology

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Abstract

Recent developments in superstring phenomenology are summarized on a non-technical level.

Talk presented at the XXVIIth Rencontre de Moriond on Electroweak Interactions and Unified Theories.
Superstring theory is still today the only candidate theory incorporating a consistent quantum gravity. That is, it provides a prescription of how to sensibly compute loop corrections to graviton–graviton scattering. Unfortunately, these corrections are far too small to be measured in the near future and hence this aspect of the theory cannot be verified directly at present. However, superstring theory also accommodates a spin 1 gauge theory with chiral fermions in the fundamental representation, which makes it a candidate for a theory of all known interactions. Clearly, it should be possible to check the validity of this assertion. To do so, one needs to extract the predictions of string theory at the experimentally accessible energy scale, the weak scale $M_W$. Being a consistent quantum gravity the characteristic scale of string theory is the Planck mass ($M_{Pl} \sim 10^{19}$ GeV) and thus the physics at $M_W$ is described by some low-energy effective field theory. At first sight, this effective field theory resembles rather closely (some extension of) the Standard Model (SM). It has a gauge group which comfortably includes $SU(3) \times SU(2) \times U(1)$ and ‘roughly’ the right spectrum: family replication of light chiral fermions in the fundamental representation as well as Higgs-like states inducing gauge symmetry breaking. Only when one tries to make this low-energy limit more precise does one encounter a number of technical and phenomenological difficulties. This is (partly) due to the fact that string theory is only known in a rather incomplete fashion: as a perturbative expansion. Consequently we have at present no handle on its non-perturbative properties and the physical effects hidden in this regime. Thus it might well be that once non-perturbative effects are taken into account all phenomenological shortcomings disappear. Being so close to having a consistent theory of all interactions with a light spectrum not unlike the SM we hesitate to abandon string theory too quickly, even if some of the more detailed predictions are not yet exactly matched in the particle world as we observe it. After all, the fermion masses agree to leading order (they are zero), and the subleading corrections are a tiny effect with respect to $M_{Pl}$ and require an enormous computational precision. There are two other aspects that I would like to mention in favour of string theory. Firstly, string theory does have the potential to answer some of the questions we have (so far) no way of addressing in the context of the SM. In principle it explains the origin of the gauge group and the fermion mass

* In this talk I exclusively focus on the so-called ‘heterotic string’, which is phenomenologically the most promising string theory.
spectrum. Even if we are not yet able to obtain the answers from string theory, it seems a major advance to be able to formulate these questions in a physical theory. Secondly, string theory does inspire our imagination about possible physics ‘beyond the SM’. Some of the physical effects I mention below could also occur in any other theory operating at some high-energy scale.

In this talk I describe some of the technical and phenomenological problems of superstring theory, and indicate the recent propositions of how non-perturbative physics might evade them. Let me first outline how one obtains the effective field theory which is the low-energy limit of string theory. The particle spectrum consists of a finite number of light fields (compared with $M_{\text{Pl}}$) and an infinite number of heavy states whose masses are proportional to $M_{\text{Pl}}$. The low-energy effective Lagrangian depends only on the light fields (these will be identified with the quarks and leptons) and the heavy modes contribute to their effective interactions. This process of ‘integrating out’ the heavy fields depends on which ground state (or vacuum) of the string theory has been chosen. For every ground state there can be a different effective Lagrangian, a different low-energy spectrum with different effective couplings.

Unfortunately, there exists an enormous vacuum degeneracy of every known string theory. That is, there are a large number of consistent classical ground states and there is no physical mechanism known at present which distinguishes between them. As I already indicated it is commonly believed that a mechanism for lifting this degeneracy is hidden in the non-perturbative structure of string theory. This state of affairs forces two distinctly different strategies onto us. On the one hand, we need to understand the non-perturbative properties of string theory. Here we have seen a lot of progress in simpler ‘toy’ string theories where the non-perturbative properties appear to be under control. Hopefully further progress in this area will occur. In this talk I will focus on the other, more phenomenological strategy and ask if there exists any classical ground state which contains the SM at the weak scale. Instead of randomly choosing a string ground state, one selects instead those which have a chance of containing the SM. Even in this phenomenological approach it would be

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† Being a consistent quantum gravity it also teaches us about space–time and its singularity structure.

‡ In fact, one is not even able to enumerate or classify the space of all consistent string vacua.

∗ In the terminology of ‘model building’ each string ground state corresponds to a different ‘model’.
a major success if we were able to compute the top-quark mass along with all the other fermion masses. A vacuum that reproduces all measured parameters of the SM clearly has some predictive power, since it would tell us what kind of ‘new physics’ we can expect, at what scale. Of course this second approach makes the assumption that string theory is weakly coupled at $M_{Pl}$ so that we can trust this perturbative analysis. Indeed, we know that a theory unifying the known interactions is weakly coupled at its characteristic scale (which is $M_{Pl}$ in the case of string theory). In practice one invokes a self-consistent analysis in which weak coupling is assumed and then shown to arise.

In the process of selecting a ‘promising’ vacuum one first chooses a string ground state which corresponds to an effective field theory in 4 space–time dimensions. In addition one usually requires it to be supersymmetric at the Planck scale. The reason for this is twofold, non-supersymmetric vacua are often unstable at the loop level. Secondly, it is difficult to understand how a theory which operates at such different energy scales can exist without supersymmetry. (This is usually called the hierarchy problem.) The requirement of four-dimensional space–time supersymmetry constrains the couplings of the effective Lagrangian very strongly. In fact, in the two-derivative approximation, they can be expressed in terms of three arbitrary scalar functions: the Kähler potential $K$, the superpotential $W$, and the gauge kinetic function $f$. The $K$ governs the kinetic terms of the fields ($g_{ij} = \partial_i \partial_j K$, $g_{ij}$ being the $\sigma$-model metric) and is thus important for obtaining the proper field-normalization. The effective scalar potential $V_{\text{eff}}$, and thus the Yukawa couplings which eventually determine the fermion masses, are expressed in terms of $W$. Finally, the gauge couplings and the $\Theta$-angle are encoded in $f$. In some sense, string phenomenology is about computing $K, W, f$ for all vacua and to all orders in string perturbation theory, and if possible also non-perturbatively. Here one is helped by various non-renormalization theorems initially proved in supersymmetry and (partially) extended to string theory. The superpotential $W$ is not renormalized to all orders in string perturbation theory.

This is just a subset of all ground states and many exist in dimensions other than four. The hope is that string theory chooses dynamically a four-dimensional vacuum over the others, which would explain why we live in four rather than any other number of dimensions.

A caveat to this theorem at two loops with massless fields present has been pointed out in refs. [5,6].
whereas the gauge kinetic function $f$ is believed to receive corrections only at one loop $[8]$.

Let me now briefly describe how $K, W$ and $f$ are computed. The motion of the string itself is described by a two-dimensional (conformal) field theory (CFT) and the effective couplings are encoded in the correlation functions of this CFT. The two requirements of a four-dimensional string vacuum, which is also $N = 1$ space–time supersymmetric, translates into a condition on the CFT. It turns out that one needs a specific class of CFTs which are $N = 2$ supersymmetric (in two dimensions) and have a so-called central charge $c = 9$. The correlation functions of such a CFT have been studied extensively and were computed for a number of examples $[9–11]$. In calculating $K, W, f$, there have been two basic strategies. On the one hand, particularly promising string vacua have been analysed in detail $[10,11]$. On the other hand more generic properties, common to all or a large class of vacua, have been uncovered $[9]$. Unfortunately, for most string vacua we are currently not able to calculate the relevant correlation functions. Therefore, it could well be that for the vacuum which contains the SM we cannot compute any couplings.

Recently, there has been some progress in the technique of computing the effective Lagrangian. For a particular string vacuum (a compactification on a specific Calabi–Yau threefold) some of its couplings have been computed $[12]$ (at the string tree level) without relying on the underlying CFT, but instead employing methods of algebraic geometry and the recently observed mirror symmetry $[13]$. Even though the general lesson behind this specific example is not yet completely understood, it seems that the couplings can be computed without knowing the CFT correlation functions. Instead, the necessary information is encoded in the so-called Landau–Ginzburg potential from which one derives a set of coupled partial differential equations (Picard–Fuchs equations) $[14]$. The effective couplings are then obtained from the solutions of these differential equations. Similarly, some of the couplings also arise as correlation functions of some appropriate topological field theory which is a ‘twisted’ (and much simpler) version of the original CFT $[15]$. These new insights allow us to extend the computation of the effective Lagrangian to a new class of vacua and in a rather simple way.

Another property of string theory has proved useful in constraining the low-energy couplings. Most string vacua are believed to display a discrete symmetry
termed ‘duality’ \cite{16}. It appears as a symmetry of the string partition function and is due to the presence of so-called winding states, which arise as a consequence of the one-dimensional nature of the string. This symmetry has no field theory analogue and is a truly ‘stringy’ property of the theory. The explicit form of the symmetry group is so far only known for a few vacua but there (in conjunction with supersymmetry) it proved very powerful in constraining the couplings \cite{16–18}. (The methods employed in \cite{12} also allowed the determination of the duality group.) Apart from this technical merit this symmetry also implies an interesting physical phenomenon. Under the symmetry operation, small and large distances are being identified, which indicates the existence of a minimal length in nature.

So far I mentioned some of the technical problems (and the recent improvements) in computing the low-energy effective Lagrangian or equivalently $K, W, f$. Let me now turn to some of the phenomenological problems which are common to all four-dimensional supersymmetric string ground states. Generically, the gauge group $G$ in string theory is a product of factor groups $G_a$

$$G = \prod_a G_a.$$  

(1)

An example is the well-known $E_8 \times E_6$ but there are many vacua with other gauge groups.\footnote{However, the rank of $G$ cannot be arbitrarily large, string theory puts an upper bound on it.} Still, string theory chooses the same (unified) gauge coupling for every factor. Thus, there is unification in the sense of a universal gauge coupling. However, this does not imply a single gauge group as in conventional GUT theories.\footnote{From the string point of view, there seems to be almost no reason left for considering a conventional GUT theory.} Furthermore, the unification occurs at a different scale, the Planck scale.\footnote{In ref. \cite{19} it has been shown under what conditions this can be consistent with the LEP precision measurements and, vice versa, what conditions the LEP measurements put on the light spectrum of string theory.} In addition, this gauge coupling is dynamical, that is, it arises as the vacuum expectation value (VEV) of a scalar field $\phi$ called the dilaton

$$1/g_a^2 = \langle \phi \rangle,$$  

(2)
It can be shown that $\phi$ is a flat direction of the effective potential $V_{\text{eff}}$ to all orders in string perturbation theory. Thus $\langle \phi \rangle$ is undetermined and a free parameter of the theory, in contradiction with the SM. ‡ This is the first ‘generic’ problem.

If one chooses supersymmetric string vacua this supersymmetry has to be broken somewhere between the Planck scale and the weak scale. The reason for choosing a supersymmetric vacuum in the first place was the hierarchy problem. However, this requires not only supersymmetry at $M_{\text{Pl}}$ but also that it be broken at some intermediate scale around $10^{10} - 10^{11}$ GeV. It turns out to be surprisingly difficult to conceive a mechanism which induces such a hierarchical supersymmetry breaking (second problem).

Even though in superstring theory the cosmological constant is zero before supersymmetry breaking one needs a mechanism which keeps it at zero after the breaking (third problem).

As I already indicated, it is believed that the solution of these generic problems lies in the non-perturbative regime of string theory. This can only be analysed under the assumption that the dominant non-perturbative effects are field theoretic in nature. Those are certainly contained in any ‘stringy’ non-perturbative effect. The real assumption is that they are dominant.

Any non-perturbative effective potential arising in field theory displays a dependence on the gauge coupling constant which is of the form $V_{\text{eff}} \sim e^{-1/g^2}$. A popular example for such a $g$-dependence is generated by gaugino condensation. This was already investigated in the context of supergravity in the early 80’s [21] and first applied to string theory in ref. [22]. One assumes a so-called ‘hidden sector’ (no renormalizable couplings with the observable sector), † which consists of an asymptotically free non-Abelian gauge theory, weakly coupled at $M_{\text{Pl}}$ but which becomes strongly coupled at some lower scale $\Lambda_c$. In such a theory the gauginos (supersymmetric partners of

‡ There are other scalar fields in the spectrum which share the property of being a flat direction of $V_{\text{eff}}$. These fields are called moduli and their VEVs parametrize continuously connected families of string vacua. However, they are not related to the tree-level gauge couplings.

† The non-perturbative effect encountered in some toy string theories are of order $e^{1/g}$ [21]. It is important to further investigate those effects.

Diamond They appear in many string vacua, the $E_8$ in $E_8 \times E_8$ is only one example.
the gauge bosons) condense and induce an effective potential of the form (at leading order)

\[ V_{\text{eff}} \sim \sum_a c_a \exp \left( -\frac{24\pi^2}{b_a g_a^2} \right), \tag{3} \]

where the sum runs over the different factors of the hidden gauge group, \( b_a \) are the 1-loop coefficients of the \( \beta \) function, and \( c_a \) are some constants. The field dependence of \( 1/g_a^2 \) [eq. (2)] enters \( V_{\text{eff}} \) and generates a potential for \( \phi \). Unfortunately, the minimum of this potential occurs at \( \langle \phi \rangle = \infty (g^2 = 0) \) corresponding to a free-field theory. This unacceptable state of affairs can change, once loop corrections to gauge couplings are taken into account. In general, below \( M_{\text{Pl}} \), the running gauge couplings are given by

\[ \frac{1}{g_a^2(p^2)} = \phi + b_a \ln \frac{M_{\text{Pl}}}{p^2} + \Delta_a(T), \tag{4} \]

where \( \Delta_a \) are the (infra-red finite) threshold corrections \[23\]. They arise from integrating out the heavy fields and are generically gauge-group-dependent. Thus they destroy the universality of the tree-level gauge coupling. If the masses of the heavy fields depend on the VEV of a light scalar field \( T \) (Higgs or modulus), these threshold corrections can induce further field dependence (in addition to the dilaton) into the gauge couplings and via (3) also into the effective potential.

It was observed in ref. \[24\] and further investigated in ref. \[25\] that a hidden sector of two gauge groups with very closely matched \( \beta \) functions (e.g. \( SU(N) \times SU(N+1) \)) and a difference \( \delta (\delta = \frac{\Delta_{N+1}}{N+1} - \frac{\Delta_N}{N}) \) in the field-independent (constant) threshold corrections might stabilize the dilaton VEV. Indeed, minimization of eq. (3) generates a VEV whose size is controlled by large \( N \) and to leading order in \( N \) is given by \( \text{Re } S = \frac{N^2 \delta}{16\pi^2} \) (for example \( SU(9) \times SU(10) \) with \( \delta = 4 \) leads to \( \text{Re } S = 2.1 \) or \( \alpha_{\text{GUT}} \sim 1/25 \)). Thus, by taking into account the constant threshold corrections, it is possible to fix the dilaton VEV at a phenomenologically acceptable value. Unfortunately, without any \( T \) dependence of \( \Delta_a \), supersymmetry is unbroken at this minimum.

The computation of \( \Delta_a \) in string theory can be performed explicitly when the complete mass spectrum is known \[26\]. In ref. \[27\] the field dependence of \( \Delta_a(T) \) was computed for a particular class of string vacua (orbifolds) and the dependence on a specific set of massless modes \( T \) (untwisted moduli). It was found that

\[ \Delta_a(T) = A_a \left( \ln(|\eta(T)|^4) + \ln(1 + T) \right), \tag{5} \]
where $A_a$ are some constants and $\eta$ is the Dedekind $\eta$ function. The functional form of $\Delta_a(T)$ strongly depends on the vacuum chosen. However, the computation of $\Delta_a(T)$ can be somewhat simplified by observing that the two terms in eq. (5) correspond to the contribution of the heavy and light states running in the loop [28]. The contribution of the light states can be calculated rather easily from the tree-level couplings alone. By employing symmetry arguments (for example the duality mentioned above) this might be sufficient to infer the full structure of $\Delta_a$. Clearly it is important to understand the field dependence of $\Delta_a$ for a larger class of vacua and also the dependence on other (non-moduli) fields.

It was shown in ref. [18] that once the dilaton is fixed supersymmetry can be broken by including field-dependent $\Delta_a(T)$. Now the potential has to be minimized in a multidimensional field space, which can have non-supersymmetric minima. This computation was performed for the particular class of vacua where $\Delta_a$ is given by eq. (5). Even though this is a small subset of all vacua, it is encouraging that supersymmetry breaking occurs, suggesting that it might happen under much more general conditions. As I already indicated, supersymmetry breaking alone is not enough: it has to occur at the right scale. By considering hidden gauge groups containing some ‘hidden matter’, the generated hierarchy can be improved [30]. The second important aspect is that the same mechanism also determines the VEV of the modulus $T$. This leads to a (partial) lifting of the vacuum degeneracy.

However, we are left with a number of shortcomings. First of all, in almost all models a cosmological constant is induced after supersymmetry breaking. This prohibits a further analysis of soft supersymmetry-breaking terms and more detailed low-energy phenomenology. Secondly, the two mechanisms proposed above, the fixing of $\langle \phi \rangle$ and supersymmetry breaking, have not been put together in a realistic string vacuum. Thus we are still at the stage of suggesting possible scenarios of how the string might choose to break supersymmetry and fix the gauge coupling constant. The recent developments are very encouraging, though.

Finally, let me speculate about another interesting physical effect related to $\Delta_a$. It is likely that $\Delta_a$ also depends on scalar fields charged under the observable gauge

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- For the example of the Calabi–Yau manifold referred to earlier, this procedure has been carried out in ref. [29].
group. When minimizing $V_{\text{eff}}$, this might be another mechanism to induce gauge symmetry breaking.

To summarize, I indicated some recent technical progress in computing the low-energy effective Lagrangian. This is an important step in comparing the Lagrangian with the SM. Furthermore, by including threshold corrections of the gauge couplings into the non-perturbative effective potential, one is able to fix the dilaton VEV at a phenomenologically acceptable value. The field dependence of these corrections can also induce supersymmetry breaking. Unfortunately, we have made little progress on the problem of the cosmological constant.

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