Challenges of String Theory: Particle Physics and Cosmology

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Abstract. We review some of the mechanisms that string theory offers to address phenomenological questions, such as the origin of our observable four-dimensional space-time, the standard model gauge symmetries and matter, as well as the possible source of the inflationary epoch. These notes are not comprehensive but attempt to provide a general view of the current search for physics cases within the context of string model building.

1. Preliminaries
Two of the major contributions of the 20th century to science are doubtless the standard model of particle physics (SM) and the cosmological model ΛCDM. The former combines the revolutionary ideas of special relativity and quantum mechanics to describe the smallest components of matter while the latter is based on general relativity and describes not only the large scale universe, but also its evolution in time.

Despite its enormous experimental success, the SM leaves some open questions. Among them, there is the origin of its 19 free parameters, the number and structure of the families of quarks and leptons, the neutrino masses (which add 7 more free parameters), and the reason of the apparent unification of the coupling constants $\alpha_i$ at a large energy scale. One urgent question is also the existence of the conjectured Higgs boson, for which recent data from the LHC suggest a mass around 125 GeV [1, 2, 3].

On the other hand, the ΛCDM model provides the best fit to the observational data about the CMB, supernovae and the structure of the universe at large scales. Curiously, this model indicates that the SM can only describe about 4% of the universe. The remaining 96% corresponds to a mixture of dark matter and dark energy, whose nature is still to be unveiled. In the ΛCDM model, the existence of a period of exponential expansion known as inflation is conjectured. This possibility is the strongest candidate as explanation of the structure of our universe due to the observations by COBE [4] and WMAP [5]. Unfortunately, the cosmological model does not clarify which mechanisms trigger and terminate inflation. Nor does it explain the nature of any of the dark components of the universe.

Aiming at solving the aesthetical problems of the SM and completing our cosmological understanding of the universe, several ideas have been concocted. Among the most relevant ones, we find: supersymmetry [6, 7, 8, 9] and extra dimensions [10, 11] as plausible keys to the stability of the electroweak scale and the origin of dark matter; and grand unification theories (GUTs) [12, 13] as possible explanation of the origin of the three microscopic fundamental
forces and charge quantization. Although these proposals include some hurdles (such as the introduction of new free parameters and new exotic particles), they are considered to be promising solutions to the puzzles of contemporary physics. Nevertheless, they do not address why the major theories upon which our current understanding of physics rely are fully foreign to one another. Furthermore, none of these ideas deal with the quantization of gravity, which is fundamental to understand gravity at tiny scales and at the very early universe.

String theory is perhaps the best candidate to provide a quantized theory of gravity. Its only assumption is that all what surrounds us is composed of tiny vibrating strings instead of point-like particles. Surprisingly, this relatively simple assumption that originally attempted to address the physics of the strong interactions turned out to yield a quantum theory of gravity [14, 15, 16]. This breathtaking observation was further complemented by the fact that the emerging string theory also predicts the existence of extra dimensions and supersymmetry in our space-time (if the two-dimensional space-time swept by an evolving string –called worldsheet– is endowed with supersymmetry). Moreover, it leads to gauge symmetries as large as those proposed by GUTs, i.e. it has the potential of solving simultaneously several of our theoretical worries in both theories that are relevant to high energy physics [17, 18]. It is precisely this feature what encourages many physicists to expect string theory to be indeed the fundamental theory that Einstein longed for.

There exist 5 types of supersymmetric string theories in 10 dimensions: the heterotic (closed) strings which result from combining a supersymmetric string with a bosonic one and lead to $\mathcal{N} = 1$ SUGRA in 10D with gauge group $E_8 \times E_8$ or $SO(32)$; the (open) strings type I with gauge group $SO(32)$ which yields $\mathcal{N} = 1$ SUGRA in 10D; and the (closed) strings type II A and B which lead to $\mathcal{N} = 2$ supersymmetric fields without and with chirality, respectively. This means that, in contrast to field theory where an infinite number of actions are admissible, string theory is so restrictive that only 5 actions are allowed.

Since the “second revolution” in the 1990’s it is known that all 5 string theories are different expressions of a more fundamental theory in 11 dimensions called M-theory [19, 20]. It was found that the low-energy limit of M-theory is 11-dimensional supergravity and that, from this effective theory, it is possible to arrive at the effective limits of the heterotic string with gauge group $E_8 \times E_8$ and of the string theory type IIA. Further, compactifying the extra dimensions of the heterotic string with gauge group $SO(32)$ on a tiny (small volume) manifold renders the same result as compactifying the extra space of the heterotic string with gauge group $E_8 \times E_8$ on a huge (large volume) manifold. The same is true between the string theories type IIA and IIB. Further, the fifth theory, type I, studied at small coupling constant is equivalent to the heterotic string with gauge group $SO(32)$ at large coupling constant. This set of relations between the various string theories that arise in string theory is called a series of dualities. One could be thus tempted to conclude that, if one wishes to characterize the 4-dimensional solutions emerging from string theory, it might suffice to fully understand one type of strings. However, it becomes rather difficult to fully understand a theory as its coupling constant becomes large, QCD being the best known example of the difficulties that emerge in this case.

The name string theory is rather antidemocratic, since the theory contains other dynamic structures beside strings. In the 1990’s Joseph Polchinski discovered that the consistency of the theory (in presence of the cited dualities) requires the inclusion of objects that cover $p$ spatial dimensions called Dirichlet $p$-branes or simply D$p$-branes [21, 22], on which open strings with Dirichlet boundary conditions can end. These $p + 1$-dimensional objects\footnote{A time-like dimension is included.} can cover the whole 10-dimensional space-time predicted by string theory or only a subspace. Perhaps the most relevant aspects of D-banes are that they can yield gauge symmetries and that they appear in type I string theories, i.e. precisely where gauge groups cannot appear otherwise. The gauge
symmetries are born from considering all independent oriented open strings that can be attached to a D-brane or to a stack of superposed D-branes. These independent strings correspond to gauge fields in the low-energy effective theory confined to the space of the D-branes. For example, one single D-brane allows only one independent string and, therefore, one gauge boson in the emerging field theory, so that the corresponding gauge group is $U(1)$. The situation becomes even more interesting when more than one D-brane are involved. For example, for a stack of 2 superposed D-branes there are four independent open strings (see fig. 1), leading to the gauge group $U(2)$. In general, for stacks of $N$ D-branes there are $N^2$ open strings which generate the gauge groups $U(N) = SU(N) \times U(1)$.\(^2\)

It is important to mention that not every $p$ is admissible. Consistency of the theory demands that only $p = 0 \mod 2$ (0, 2, 4, 6, 8) and only $p = 1 \mod 2$ (−1, 1, 3, 5, 7, 9) appear in the type IIA string theory and type IIB string theory, respectively. A D($−1$) brane can be understood as an instanton whereas D0 branes are point-like particles. In type I string theory, also only odd values of $p$ are allowed, although these are restricted to the values $p = 1, 5, 9$. This is, as we shall see below, due to the fact that type I can be seen as the one half of the type IIB theory left untouched by a special “surgery” called orientifold.

With all these ingredients at hand, one of the tasks of today’s string theorists is to bring the enormous theoretical building that string theory represents in contact with the physics we know while giving answer to some (or all) of the persisting puzzles. This is the goal of a very active branch of string theory called string phenomenology, in which we are interested in these notes. A fast inspection of the aforementioned properties of string theory reveals that there are important differences between string theory and the successful theories we are trying to reproduce and improve. Those discrepancies are the fundamental challenges of string theory that we might enumerate as follows:

(i) 4D vs. 10D;
(ii) Matter states (quarks and leptons) from strings;
(iii) Exact (MS)SM matter spectrum and interactions (masses);
(iv) Stable cosmological constant.

The purpose of these notes is to provide some flavor of how string theory can give answer to these questions and what the prospects for observable physics are.

2. First challenge: 4D vs. 10D

The fact that our everyday life occurs in four dimensions is obvious to everybody. However, what would we expect if the additional dimensions were too small to be seen by the naked eye?

\(^2\) If the branes wrap orbifold-invariant planes, the resulting group is symplectic $Sp(N)$. Note that $SU(2) \simeq Sp(2)$.\(^2\)
A notable physical consequence of extra dimensions is that Newton’s law would change at very small scales. For instance, if an extra dimension is replaced by a circle of radius $R$ instead of an infinite line, the gravitational potential of a body with mass $M$ valid within the circle would be

$$V_5 = -G_5 \frac{M}{r_5^2}, \quad \text{with} \quad r_5^2 = r_4^2 + \zeta^2 = x^2 + y^2 + z^2 + \zeta^2,$$

instead of the usual $V_4 = -G_4 M/r_4$, where $G_4$ and $G_5$ denote the Newton constants in 4D and 5D respectively, and $\zeta$ represents the coordinate of the additional spatial dimension. Very sensitive experiments have searched for this discrepancy at scales as small as $\sim 1\text{mm}$ without success [24]. But the extra dimension(s), if it (they) existed, could well be beyond the experimental threshold.

The conjecture of very small compact dimensions was introduced by Theodor Kaluza and Oskar Klein [25, 26]. Although in their pioneering works they failed to unify gravity with electromagnetism, as they attempted, with the advent of string theory their ideas have gained vigor.

### 2.1. Compactifications

Following the line of thought of Kaluza and Klein, a consistent technique to achieve models from string theory that at least approach to our apparently four-dimensional universe is to allow the six additional dimensions predicted by string theory to form Ricci flat compact spaces. That is, one conjectures that the whole space-time of strings $X_{10}$ can be split into $X_6 \otimes \mathbb{M}^4$, where $X_6$ denotes a 6D compact manifold and $\mathbb{M}^4$ our observable Minkowski space-time. This method to “hide” the additional dimensions is called compactification. Since the consistent string theories are supersymmetric, one must impose further constraints to the shape of $X_6$ if one aims at a four-dimensional effective theory endowed with at most $N = 1$ supersymmetry. To arrive at 4D $N = 1$ supersymmetric field theories via compactification, the compact space must have SU(3) holonomy group (or a subgroup thereof larger than SU(2)).

These spaces can either be smooth like Calabi-Yau (CY) threefolds $CY_3$ [27] or contain a finite set of curvature singularities like an orbifold $\mathcal{O}_6$ [28, 29]. They are illustrated in fig. 2. Orbifolds are singular limits of CY manifolds, but the former are much simpler than the latter. The simplest orbifold is a circle $S^1$ divided by a $\mathbb{Z}_2$, i.e. where points on one side are identified with the points of the opposite side. By the identification, the two points at the poles become curvature singularities (also called orbifold singularities or fixed points - because they remain invariant under the $\mathbb{Z}_2$). On the other hand, one of the simplest $CY_3$ is the quintic, displayed in fig. 2a. In 6D, an Abelian toroidal orbifold is defined as the quotient of a $T^6$ by a discrete
isometry $\mathbb{Z}_N$ or $\mathbb{Z}_N \times \mathbb{Z}_M$ of the torus which acts on the whole compact space. Since the holonomy group of the orbifold corresponds to the symmetry that is moded out, the isometry group must be embeddable within $SU(3)$ but not in $SU(2)$ to retain $\mathcal{N} = 1$ supersymmetry in 4D. We shall be only interested in this kind of orbifolds.

Since the time of Kaluza and Klein, it is known that compactifying higher dimensional gravity leads to scalar fields in 4D called moduli that parametrize the size and shape of the compact space. These fields lack mass, becoming a source of extra massless matter that could be detected through their interactions. If there existed a mechanism to furnish them with potential (non-perturbatively or otherwise), these new fields would develop a mass and become important cosmologically. This can be considered both an advantage and a disadvantage: on the bright side, in the emerging effective field theory, these fields could account partly for the dark matter and for the conjectured inflaton. On the less appealing side, there might simply be too many moduli or they could acquire masses such that reheating or other cosmological events could be prevented. As we shall see, these worries, for which the term (cosmological) moduli problem was coined, are addressed with some success in different scenarios.

Independently of the personal taste, both CY and orbifold spaces can render some possibly useful 4D phenomenology. Originally, it was thought that this would only be true so long as the compact space be small enough to justify the experimental absence of evidence of deviations of Newton’s law. However, as we will shortly see, today we know that D-branes open up a way out of this constraint. There is a second reason to avoid very large extra dimensions though. Kaluza and Klein also taught us that the 4D matter particles appearing in the effective theory have some heavier ‘cousins’ dubbed Kaluza-Klein (KK) modes. The KK modes arise from the compactification of the torus which acts on the whole compact space. Since the holonomy group of the orbifold corresponds to the symmetry that is moded out, the isometry group must be only interested in this kind of orbifolds.

In the simplest string compactifications without D-branes, the theory itself dictates the size of the extra dimensions, Considering only closed strings, the heterotic strings are the best examples of theories that admit this compactification. In these scenarios, there are only two arbitrary constants: the string scale $M_s$ (inversely proportional to the size of the strings) corresponding to the energy scale at which the interactions between strings are important, and the string coupling $g_s$. In the effective theory, i.e. for energies below $M_s$, one can compare the outcome from compactifying the 10D gravity plus the gauge interactions in $X_6$ against the expected 4D gravity plus SM interactions. Assuming that $X_6$ has volume $V_6$, we obtain

$$S_{het} = \frac{M_s^8}{g_s^2} \int d^{10}x \, R_{10d} + \frac{M_s^6}{g_s^2} \int d^{10}x \, (\text{tr} \, F_{MN} F^{MN})_{10d}$$

$$\approx \frac{M_s^8 V_6}{g_s^2} \int d^4x \, R_{4d} + \frac{M_s^6 V_6}{g_s^2} \int d^4x \, (\text{tr} \, F_{\mu\nu} F^{\mu\nu})_{4d},$$

where $R_{n,d}$ is the Ricci scalar in $n$ dimensions, $M,N = 0,1,\ldots,9$, $M_{Pl} \sim 10^{18}$ GeV and
$g_{SM}^2 \sim 0.01$, assuming unification of all gauge interactions for simplicity. The second row is obtained by compactifying six dimensions and letting the moduli arising from the compact components of $R_{10d}$ and $F_{MN}$ not to yield observable interactions. Comparing the expressions before the integrals with those below the underbraces, we see that $M_s = M_{Pl} \cdot g_{SM} \sim 10^{17}$ GeV, which is very close to the proposed GUT scale $M_{GUT} \sim 10^{16}$ GeV. Further assuming that $g_s \approx g_{SM}$ close to the compactification scale, one finally finds that $V_6 \approx M_s^{-6}$ or, equivalently, that the size of the compact dimensions must be around $10^{-30}$ cm! This strong constraint (valid in heterotic constructions) can be considered a positive feature because it is a model independent prediction. However, it can also be a hindrance to get in touch with observations or to deal with precision coupling unification ($M_{GUT} \neq M_s$).

2.2. Brane-worlds

The might of D-branes is not only their capability to furnish type II string theories with gauge groups, but also the freedom they provide to design configurations in which the compact space can take the least beautiful shapes and still render 4D appealing physics. For one, with the aid of D-branes it is possible to escape the constraint on the size of the compact space. If the volume of the compact space is large, but our observable universe is confined to a D-brane that intersects only a region of the compact space with very small volume, the result is an acceptable 4D theory.

In string theory, brane-worlds are compactifications of type II (A or B) string theories, whose compact spaces contain some non-contractible loops wrapped by D$p$-branes, with $p \geq 3$. A natural consequence in these constructions is that certain kind of matter (particularly gauge bosons) is trapped within the intersection of the D-branes, isolated from some other physics occurring outside that intersection.

We can consider, for instance, a situation in which the gravity in 4D appears from gravity in 10D, i.e. from closed strings in a type II string theory, but the (unified) gauge forces are related to the open strings attached to a stack of D-branes. In this case, one arrives at

$$S_{II} = \frac{M^8_5}{g_s} \int d^{10} x R_{10d} + \frac{M^{p-3}_s}{g_s} \int d^{p+1} x \left( \text{tr} F_{\alpha \beta} F^{\alpha \beta} \right)_{(p+1)d}$$

$$\approx \frac{M^8_5 V_6}{g_s^2} \int d^4 x R_{4d} + \frac{M^{p-3}_s V_{p-3}}{g_s} \int d^4 x \left( \text{tr} F_{\mu \nu} F^{\mu \nu} \right)_{4d}.$$ 

The differences between this expression and eq. (3) appear in the second integral. First, the different power in the factor $1/g_s$ is because open strings couple with only one power of $g_s$. Second, the indices $\alpha, \beta$ run only over the $p+1$ dimensions of the D-branes. Let us assume that 3 spatial dimensions of the D-branes cover our $M^d$ space-time; this implies that $V_6 = V_{p-3} \cdot V_\perp$, where $V_{p-3}$ is the compact part of the volume of the D-branes and $V_\perp$, the compact volume of $X_6$ perpendicular to the dimensions of the D-branes. Comparing the coefficients of the integrals to the expected theory in 4D, we find that $M_s$ is a function of the compact volume perpendicular to the D-brane $V_\perp$ and the total volume $V_6$ of $X_6$: $M_{p-3}^{l-p} \approx 10^{16}$ GeV · $\sqrt{V_6/V_\perp}$. Whereas observations constrain the magnitude of $V_6$ (or equivalently $V_{p-3}$), one can adjust $V_\perp$ at will to get arbitrarily small values of $M_s$. This freedom allows one to set $M_s$ as low as few TeV without affecting any experimental bounds.

In a rather bottom-up approach, brane-worlds refer to higher dimensional field theories endowed with certain hypersurfaces (not precisely D-branes). The work by Lisa Randall and Raman Sundrum [10, 11] (see also [30]) is a good example of this case. We will not focus on this work.
This way it turns out possible that some stringy effects due to the compactification process be already observable at LHC. In particular, the KK modes have masses that go as multiples of $M_s$ in general compactifications. If $M_s$ is accessible to current experiments, heavier states with the properties of the photon, neutrinos, electrons, etc. would display signatures in the accelerators.

3. Second challenge: matter fields from string theory

We call matter to all the fields that enter the SM, i.e. gauge bosons and chiral fermions. However, in a first step, we are interested in obtaining matter representations that transform as adjoints and as fundamentals of an arbitrary gauge group. We pursue here also a mechanism that yields small gauge groups, such as those of the SM or the smallest GUTs. The way how all kinds of matter appear depends on the particular scenario.

3.1. Compactifications

In the heterotic strings as well as in the string type I, the compactification of the additional dimensions must also affect their original gauge groups in order to retain the (modular) symmetries of the original stringy actions. The effect is an explicit breakdown of the original gauge groups ($E_8 \times E_8$ or $SO(32)$) to a subgroup thereof. Compactifying on the simplest CY threefold, the so-called standard embedding, the gauge groups are broken according to

$$E_8 \times E_8 \rightarrow E_6 \times E_8,$$

$$SO(32) \rightarrow SO(26).$$

Based on this breaking, the branching rules of the original representations read

$$\begin{align*}
\begin{pmatrix} 248, 248 \end{pmatrix} &\rightarrow \begin{pmatrix} 78, 1 \end{pmatrix} \oplus \begin{pmatrix} 248, 1 \end{pmatrix} \oplus 3 \begin{pmatrix} 27, 1 \end{pmatrix} \oplus 3 \begin{pmatrix} 27, 1 \end{pmatrix} \oplus 8 \begin{pmatrix} 1, 1 \end{pmatrix}, \\
496 &\rightarrow 325 + 6 \times 26 + 16 \times 1.
\end{align*}$$

Not all of these states survive the compactification due to further consistency conditions on the compactifications. However, the relevant feature to notice here is that, the splitting of the adjoint representations yields fundamental representations that can give rise to states of the (supersymmetric) SM in addition to the gauge bosons of the new unbroken gauge group. For example, a $27$-plet of $E_6$ contains all the gauge representations of a family of quarks and leptons of the SM.

Similarly, orbifold compactifications lead to the breakdown of the original gauge symmetry like in eq. (5). The difference is that, unlike in CY compactifications, the rank is never broken by the compactification. So, e.g. in the simplest case (orbifold standard embedding), $E_8$ is broken to $E_6 \times SU(3)$ or $E_6 \times SU(2) \times U(1)$ or $E_6 \times U(1)^2$. Further, invariance of the resulting states (what we call here the it spectrum of the model) under the action of the orbifold impose further constraints, yielding frequently less states than the branching rules might suggest. For example, compactifying the $E_8 \times E_8$ heterotic string on a $T^6/\mathbb{Z}_2$ orbifold leads to the unbroken gauge group $E_6 \times U(1)^2 \times E_8$ and the non-adjoint matter stemming from the 10D vector multiplets is just $3 \times \begin{pmatrix} 27, 1 \end{pmatrix} \oplus \begin{pmatrix} 27, 1 \end{pmatrix} \oplus 3 \times \begin{pmatrix} 1, 1 \end{pmatrix}$, where $U(1)$ charges have been omitted.

In an orbifold, there is an additional source of matter fields. At the singularities, some strings that would be open in the string theory compactified on a smooth manifold are actually closed around the singularities due to the orbifold. To picture this let us take a $T^2/\mathbb{Z}_2$ orbifold. In this case, there are four singularities and the compact space of the torus is reduced to only half of it. Taking into account these ingredients, one arrives at the picture displayed in fig. 3. The original smooth space of the torus is deformed to a tetrahedron. As depicted in the figure, there are some closed strings around the vertices. If one unfolds the tetrahedron, it is clear...
Figure 3: Orbifold $T^2/\mathbb{Z}_2$. The bulk states arise from the original 10D super Yang Mills and 10D gravity. The local or twisted states arise from the compactification itself and are attached to the singularities. These can be vector multiplets that form a local gauge group as well as chiral multiplets that form complete representations under the local gauge group.

that some edges have to be cut, rendering evident that the closed strings on the tetrahedron are open on a smooth (plane) manifold. The resulting states from those special strings are called twisted or simply local, as they attached to the singularities. In contrast, The states coming from the original vector and gravity multiplets in 10D (i.e. from the original string theory) are free to move everywhere. Retaking our previous $T^6/\mathbb{Z}_6$–II example, those local states yield the representations (omitting again the U(1) charges)

$$32 \times (27, 1) \oplus 10 \times (27, 1) \oplus 217 \times (1, 1).$$

Counting a 27 as a generation of quarks and leptons (as is done in E6 GUTs), including the bulk or untwisted states, one obtains in this example a net number (generations minus anti-generations) of 24 generations.

3.2. Brane-worlds
In brane-worlds, we have learnt that vectorial multiplets accounting for the existence of gauge bosons arise from the open strings attached to a stack of $N$ D$p$ branes. The dimensionality of the branes make no difference on the resulting gauge group (which can be SU($N$) or Sp($N$)). In order to obtain chiral supermultiplets, it is bit more complicated.

First, recall that the type II theories (where D-branes appear) yield $N = 2$ supersymmetry in 10D. Thus, if one simply compactifies the additional dimensions onto a CY manifold or an admissible $O_6$, the emerging 4D theory has $N = 2$ supersymmetries too. To reduce the amount of remaining supersymmetry, it is needed to perform an orientifold (see e.g. [31] for an introduction to orientifolds). An orientifold is a generalization of an orbifold where the discrete isometry group $G_I$ is accompanied by a discrete operation of orientation-reversal in the worldsheet $\Omega$, so that the final result is a theory of unoriented strings. In fact, the type I string theory can be obtained by letting the orientifold $G_I = 1$ and $\Omega = \mathbb{Z}_2$ act on the type IIB string theory. A very important feature is that the invocation of an orientifold very often implies the introduction of open strings (absent in type IIB but present in type I strings). These new strings may be considered an analog of the twisted strings of the orbifold. Consistency of the theory (tadpole cancellation) restricts the amount of open-string states appearing after the orientifold and, in fact, this is the way how the gauge group SO(32) (and not any other) arises from type IIB.
The second source of chiral matter fields are intersecting branes (for a useful review of these constructions, see [33]). The starting point are D-branes that cover the 4D space that should correspond to ours and that wrap some of the non-contractible cycles of the compact space. For example, considering $X_6 = T^2 \times T^2 \times T^2$, there are 6 non-contractible cycles corresponding to the independent loops of the three 2-tori. If one introduces D6 branes in the type IIA string theory, three spatial brane dimensions can wrap three of those cycles. It is then clear that, if there are more than one D6 branes, one can accommodate the extra dimensions of the branes in different cycles, yielding points of intersection of the D-branes in the compact space.

As depicted in fig. 4a, at the intersection of two D-brane stacks there are open strings whose ends connect both stacks. It turns out that these strings yield physical states living in the intersection space (in this case, the 4D space we assume they both share) that are charged under the gauge groups associated to the stacks. In fact, the resulting states form (anti-)fundamental representations under those groups. For example, taking the upper left intersection of fig. 4a, one finds that the open strings would lead to the matter representations $(3, 2)$ or $(\bar{3}, 2)$ with some U(1) charges associated to the gauge groups of the parent D-branes $SU(3) \times SU(2) \times U(1)^2$.

Figure 4b depicts a situation in which the intersections proposed before occur in some region of a CY$_3$, where non-contractible cycles exist.

4. Third challenge: the (MS)SM spectrum & interactions

One of the final goals of string phenomenology is to reproduce the gauge forces and particle content of the SM. As we shall see shortly, this goal has been achieved in different scenarios at least at the supersymmetric level (the MSSM). The next step is then to address more serious questions like reproducing the matter interactions contained in the Lagrangian of the SM, providing thereby a plausible solution to the extant issues, such as the origin of baryon stability, the strong CP problem, the flavor problems, etc. Despite all the effort invested on the latter task, it is perhaps fair to say that this challenge is still open.

In this section, we shall discuss some of the methods employed in string theory to achieve models that approach at the (MS)SM and some qualities of the emerging models. We shall also mention how these constructions might lead to answers to persisting puzzles.
4.1. Compactifications

**Heterotic CY compactifications.** Both kinds of compactifications, on smooth spaces (with and without fluxes) and on orbifolds, have been actively investigated since the 80's, although they were partly left aside when D-branes were discovered. During the last decade, perhaps due to the invention of more powerful computers, there have been intensive scans of the landscape of possible CY spaces leading to promising models, which have produced interesting results [34, 35, 36, 37, 38, 39, 40] (GUT models were derived from U(N) line bundles [41, 42]). Around 200 models with the exact spectrum of the MSSM have been found. The gauge group of these models is just $SU(3) \times SU(2) \times U(1)_Y \times U(1)_{B-L}$. The matter spectrum includes just three generations of quarks and leptons, some pairs of Higgs supermultiplets and no exotics of any kind.

Despite these appealing features, their advantages play also as challenges. One problem is that in the SM the symmetry $U(1)_{B-L}$ appears only globally whereas in these constructions it is a gauge symmetry. The fact that there are no exotics particles of any kind implies that there are no additional particles that may be charged only under that symmetry, such that could trigger a spontaneous breakdown of $U(1)_{B-L}$. Getting rid of this symmetry requires desperate measures. One method to solve it is by assuming that the sneutrinos (the bosonic partners of the neutrinos) develop somehow an expectation value [43].

Another complicated problem is related to the nature of CY manifolds. CY manifolds are smooth spaces, however even writing down the corresponding metric is a challenge. Therefore, this compactification scheme must abandon the calculability of (CFT in) string theory and apply low-energy geometric methods to arrive at e.g. the size of the interactions between the emerging matter fields. One requires then numerical methods combined with cohomology techniques that allow one to deal with the interactions of fields on a very complex geometric background. This effort has revealed some progress [44, 45], but further refinement must be achieved, to figure out e.g. what the explicit dependence of the coupling constants on the moduli is.

The qualitative picture of Yukawa couplings in CY compactifications is as follows. The Yukawa coupling is either order one or zero, however, it appears difficult to obtain hierarchical couplings at a generic point in moduli space. This does not mean that the CY is incapable of reproducing phenomenology, but it might mean that one would have to move close to a special point in moduli space. Whether this point can, at least to some extent, be identified with a configuration close to an orbifold point (or a partial blow-up of an orbifold) represents an interesting question. In fact, in the CY model [38, 45] there is no light Higgs pair at a generic point in moduli space: the mu-term only vanishes if the size of a certain cycle shrinks to zero. A $\mu$ term of order TeV forces us to move to a rather special point in moduli space.

**Heterotic orbifold compactifications.** Let us now focus on orbifold compactifications of the heterotic string, which have also gained much momentum in recent years. In fact, since the birth of string theory, successful models reproducing several features of the MSSM have been found [47, 48, 49, 50, 51, 52, 53]. The new ingredient that has contributed to the appearance of many phenomenologically appealing heterotic orbifold models is a feature which we would like to refer to as **gauge group topography**: different gauge groups can be realized at different points in internal space. In the interesting orbifolds that have been constructed, the 10D gauge group $E_8 \times E_8$ gets broken to different subgroups at different orbifold singularities. Figure 5 illustrates this situation in a two-dimensional $Z_2$ orbifold. To understand fig. 5, we must only understand how the local gauge groups emerge. They are comprised out of the gauge bosons corresponding to generators that fulfill certain (local) orbifold invariance conditions, which can be recast schematically as

\[ \text{generator} \cdot \text{local shift} = 0 \mod 1, \]  

(7)
where “generator” is any of the generators of the 10D gauge group of the heterotic string, and the local shift [54, 51] introduced here depends on the fixed point [55]. Since these invariance conditions can be different at different fixed points, different gauge groups can live at the various fixed points.

In the field-theoretic description, one would say that the local gauge groups are made up from the generators for which the corresponding gauge boson has a non-vanishing profile at the singularity. This also means that the coset, \((E_8 \times E_8)/G_{\text{local}}\) with \(G_{\text{local}}\) denoting the local gauge group, comprises generators where the gauge bosons’ profile has to vanish at the singularity. From this point of view it is also clear what the massless gauge interactions in 4D are: they are mediated by the bosons with flat profile, i.e. which fulfill the invariance conditions (7) everywhere. In other words, the 4D gauge group emerges as the intersection of local gauge groups in \(E_8 \times E_8\).

Figure 5 displays not only the gauge groups at the singularities of the orbifold, but also the intersections between the different local groups. This is useful for two reasons. First, one notices that, even if two singularities are endowed with the same gauge group, their internal structure might be different and therefore lead to non-trivial intersections (see the vertices on the right). Second, focusing on the bottom left vertices, \(SO(10)\) and \(SU(6) \times SU(2)\), one finds that local GUTs (like \(SO(10)\)) can lead to GUTs at the intersection. This is a hint that shall be used below.

Since the 4D gauge group is the intersection of the local gauge groups at the singularities, this implies that, when the gauge group topography is non-trivial, the local gauge groups at the orbifold singularities have to be larger than the SM, \(G_{\text{local}} \supset G_{\text{SM}}\). This then leads to the picture of local grand unification [56].

From the requirement \(G_{\text{local}} \supset G_{\text{SM}}\) we see immediately that the local gauge groups \(G_{\text{local}}\) can be groups that have been discussed in the context of GUTs, such as \(SU(5)\) and \(SO(10)\). It is important to note that fields confined to a region with certain gauge symmetry have to furnish complete representations of this symmetry. Thus, matter fields localized at a singularity with a GUT symmetry would appear as GUT multiplets in the low-energy theory although the gauge symmetry of the low-energy theory was smaller. This statement applies, in particular, to matter...
living in regions with SO(10) symmetry: it will furnish complete SO(10) representations, i.e. a 16-plet will give rise to a complete family of quarks and leptons.

On the other hand, matter fields can also come from the bulk. Such fields do, like the gauge bosons, feel the symmetry breaking at every fixed point, and thus appear in split multiplets of the various local gauge groups. Together with what we have discussed in the previous paragraph, this provides a very simple scheme allowing to understand the simultaneous existence of complete and split GUT multiplets in nature (solving thereby the well-known doublet-triplet splitting problem of GUTs).

Based on these observations, a strategy to arrive at semi-realistic heterotic orbifold models (i.e. models with three chiral generations plus one pair of Higgs below the compactification scale) has been developed [57, 58]. The idea is to search for models containing 16-plets living at singularities with SO(10) symmetry plus a pair of Higgs multiplets coming from the bulk. As a first shot, one would aim at all three generation arising from 16-plets in order to explain the origin of the structure of the SM families. However, this situation is disfavored in heterotic orbifolds. The next simplest possibility is to demand the third generation to be different from the other two, i.e. two generations can rise from 16-plets whereas the third one may come in split multiplets from the bulk or singularities without a GUT symmetry.

In fig. 6, we present the structure of one of the heterotic orbifolds constructed in [57, 58] and further studied in [52]. Focusing our attention on two of the six orbifolded dimensions, the resulting space is again a tetrahedron which is topologically equivalent to the pillow of fig. 6. The local symmetry of two (degenerate) singularities is SO(10), which contains SU(5). If we let these two dimensions be larger than the other four, the emerging theory is a 6D \( \mathcal{N} = 1 \) supersymmetric theory where the (now macroscopic) singularities of interest enjoy an SU(5) group with the representations \( 10 \oplus \overline{5} \oplus 1 \), i.e. with the complete 16-plet split in SU(5) representations. If now the pillow is shrunk to submillimetric sizes, the emerging 4D theory has the gauge group of the SM and the two localized 16-plets appear with the correct quantum numbers of the quarks and leptons. Furthermore, since the gauge group stems from a local SO(10), the quarks and fields are charged under a gauged U(1)\(_{B-L}\) symmetry, pretty much as in the CY case discussed above. One useful remark at this point is that the two localized generations arising from the SO(10) local GUTs have exactly the same quantum numbers. In fact, they are identical. This characteristic originates the appearance of a discrete \( D_4 \) family symmetry [52, 59], which has several phenomenological applications, including the suppression of flavor changing neutral currents in gravity mediated supersymmetry breaking [60].

The origin of the third generation of quarks and leptons and the Higgs multiplets in these constructions is not so simple. We have mentioned that states living in the bulk form split
There is also a set of exotic matter. The exotics are vectorlike because for each multiplets. These are distinguished from the lepton doublets and their antiparticles.

We see that there is a net number of three SM generations plus a pair of supersymmetric Higgs fields. In the model illustrated in fig. 6, one finds that the interactions of these fields is of order unity. (This observation has been found to be relevant to explain the supersymmetric scenario known as natural supersymmetry in which the third generation fermions and especially the stop are much lighter than all other superpartners [61].) The other quarks and leptons of the third generation are found at other singularities of the orbifold. As we shall see, interactions among fields located at singularities are exponentially suppressed by the distance that separate them.

In table 1 we show a summary of all matter representations of a semi-realistic orbifold model. We see that there is a net number of three SM generations plus a pair of supersymmetric Higgs multiplets. These are distinguished from the lepton doublets \( \ell_i \) thanks to the gauged \( U(1)_{B-L} \) symmetry. There is also a set of exotic matter. The exotics are vectorlike because for each exotic \( X \) there is its antiparticle \( \bar{X} \) too. A crucial point here is the existence of the SM singlet \( X \) and some or all of them develop VEVs in a generic supersymmetric vacuum, \( U(1)_{B-L} \) is spontaneously broken down to a \( \mathbb{Z}_2 \) matter parity (useful to prevent rapid proton decay [62, 52]). Notice that this is an important advantage against the CY compactifications discussed earlier.

The second effect of the VEVs is that all exotics acquire masses very close to the compactification scale provided that superpotential couplings of the type \( S^n X \bar{X}, n \in \mathbb{Z} \), are null.
allowed by all the symmetries of the orbifold, i.e. provided that these couplings satisfy gauge invariance and certain stringy selection rules [63, 64, 65, 66]. It has been explicitly verified that those couplings exist and that indeed all vectorlike exotics attain masses through this Higgs-like mechanism.

Another important feature of the model under consideration is that the gauge group includes, besides $G_{SM}$ and $U(1)_{B-L}$, a hidden sector. The hidden sector is understood here as a group under which the particles of the (MS)SM are uncharged. In this case, the hidden gauge group is $SU(4) \times SU(2)$, but $SU(2)$ is spontaneously broken by the VEVs of some $N_i$, leaving only $SU(4)$ as hidden group. This is relevant because a hidden sector with a pure Yang-Mills theory can trigger spontaneous breaking due to gaugino condensation. Roughly speaking, the hidden gauginos $\lambda$ form bound states $\langle \lambda \lambda \rangle$ at the scale at which the interactions become strong,

$$\Lambda \sim M_{GUT} e^{-8\pi^2/b_0 g^2(M_{GUT})}, \quad b_0 = \beta \text{-function coefficient of hidden group}, \quad (8)$$

(which follows from the corresponding renormalization group flow,) inducing thereby a supersymmetry-breaking $F$-term which is transmitted to the observable sector via gravitational interactions. The result is that the scale of supersymmetry breaking in the observable sector is given by [67]

$$m_{3/2} \sim \Lambda^3/M_{Pl}^2.$$

Assuming that there is unification of all the SM interactions at $M_{GUT} \sim 10^{16}$ GeV, implying $g^{-2}(M_{GUT}) \sim 2$, in the current case $m_{3/2} \sim \text{few MeV}$. This scale is rather low given the current constraints on supersymmetric particles. However, the important message we learn from this is that these models induce naturally low-scale supersymmetry breaking [68].

With these ingredients in hand, one can investigate e.g. the structure of the Yukawa matrices. There are also superpotential couplings of the type $Y^{\alpha}_{ij} q_h \bar{u}_j$, where $Y^u$ is a $3 \times 3$ matrix whose entries are of the form $S^0/M_{Pl}^{n-3}$. What is interesting is that, as announced before, the entry $Y^u_{33}$ is order unity due to the fact that there is (almost) no obstruction that prevents interactions among bulk fields. However, since the other quarks are localized, the stringy selection rules [63, 64, 65, 66] do not allow trilinear couplings such as $q_{1/2} h_u u_{1/2}$, so that powers of fields $S$ are required. Since $(S) \sim 0.1 M_{Pl}$, each power of $S$ adds a suppression factor $1/10$. So, the generic picture is that all quarks (up and down type) acquire masses once the Higgs mechanism takes place (i.e. once $h_u$ and $h_d$ acquire VEVs); however, there is a natural hierarchy between the top mass and the masses of all other quarks which depends on the specifics of each model. Similarly, also the Yukawa couplings of the leptons are non-trivial, implying that, as in the SM, all leptons and quarks acquire masses.

Nonetheless, to arrive at specific predictions from these constructions, such as the exact texture of the CKM matrix, the precise value of the quark and lepton masses, etc., one needs additionally perform a precise computation of the coefficients $\lambda$ of couplings like $\lambda S^0 q_h \bar{u}_j$ directly from string theory. The advantage here is that, in contrast to CY compactification, we have the power of CFT tools that allow us to perform an analytic computations. In fact, couplings between three and four fields have already been computed [69, 70, 71, 72].

The starting point of the computation are the vertex operators $V$ associated with the strings that yield the fields of interest. Vertex operators are operators in the 2D CFT of the worldsheet (the surface that the strings sweep as they evolve) that carry the whole information of the fields in the emerging 10D QFT (the so-called target space). The form of the vertex operators is known and strongly depends on two features: 1) whether the field is a boson or a fermion; and 2) whether the coupling is computed at tree or loop level in the worldsheet. For the first feature, there are vertex operators for bosons $V^b$ and for fermions $V^f$. The second feature has to do with the fact that interactions between strings form different surfaces. For example, three closed strings interacting at tree level form the surface of a pair of trousers, which is topologically
equivalent to a sphere when the ‘wholes’ of the trousers are pulled to infinity. Similarly, at one loop, the surface subtended by the string interaction is a 2D torus; at two loops, the surface has two handles, and so on. Focusing on tree-level interactions, the picture associated to the scattering amplitude of six strings is displayed in fig. 7.

Once the specifics of the vertex operators are known, we can compute the scattering amplitudes for the effective fields and deduce the coupling terms in the 4D effective field theory. A term $\Phi^L$ in the superpotential, with $\Phi$ a chiral superfield with components $(\phi, \psi)$, can be inferred most straightforwardly from an interaction of the form $\psi \psi \phi^L - 2$. The computation required in this case is the following correlation function

$$\lambda = \langle V_1^f(z_1, \bar{z}_1)V_2^f(z_2, \bar{z}_2)V_3^b(z_3, \bar{z}_3)V_4^b(z_4, \bar{z}_4)\ldots V_L^b(z_L, \bar{z}_L) \rangle,$$  \hspace{1cm} (9)

where the vertex operators have to be adjusted to the form needed for interactions on a sphere. The points $z_i$ on the complex plane are points of insertion on the sphere of interaction. The correlation function eq. (9) is a path integral, where the action is the string action of the heterotic string compactified on the orbifold, which is well known. (Further details can be found in ref. [66] or directly in the original literature [63, 64].)

In the simplest case of trilinear interactions, it has been found that the coupling strength goes like $\lambda \sim e^{-\alpha |\mathbf{u}(T_i)|^2}$, where $\alpha$ is a numerical factor and $\mathbf{u}(T_i)$ is the vector that parametrizes the distance between the fixed points at which the fields on interest are attached. This non-trivial interaction among distant strings is possible due to the existence of worldsheet instantons. Notice that $\mathbf{u}$ depends on the Kahler (also called volume) moduli $T_i$, whose values must be fixed (see sec. 5.1) prior to computing the actual value of $\lambda$. There is a second challenge. At non-renormalizable level, the integral expressions of eq. (9) become increasingly complicated. Therefore, it is necessary to implement computer algorithms that can deal with complicated integral in complex variables.

Despite all the complications, there is some progress in the computations and, some particular cases (where the value of the moduli has been fixed to order unity values and the VEVs of the SM singlets $S$ has been tuned), one can arrive at the following eigenvalues of the Yukawa matrices.

$$|Y_u^{\text{diag}}| \sim \text{diag}(0.0003, 0.06, 1.1), \quad |Y_d^{\text{diag}}| \sim \text{diag}(0.0002, 0.001, 0.04).$$

These values, although not quite realistic, are pretty close to the values of the Yukawa couplings that one expects in a scenario where $\tan \beta \sim 1$.

How generic are these results? Evidently, these results are quite encouraging. One might be tempted to claim that one of the goals of string phenomenology is about to become true. So it is
natural to wonder whether these results are common to all or a large subset of compactifications. To answer this question, we have performed a scan based on the local GUT strategy depicted before. That is, we have started a search of models endowed with local GUTs and selected those among them that fulfill a series of priors such as couplings unification, absence of exotics and sizable Yukawa couplings for the top quark. The results were called the mini-landscape [57, 58] because they correspond to a set of about 300 models with properties similar to those explained before identified within a larger set of about 10 million (not so appealing) models. Due to the amount of data that has been studied, the huge task of finding those models would not have been possible without the orbifolder [73]. The orbifolder is a program design in C++ specifically for the computation of all the properties of Abelian heterotic orbifolds, which is now made public through the site \url{http://projects.hepforge.org/orbifolder/}. It can also be used on-line via our web interface: \url{http://stringpheno.fisica.unam.mx/orbifolder}. The idea of that publication is to allow one to overcome the technicalities of the heterotic compactifications and focus on a more phenomenological quest towards more appealing constructions and a better understanding of the models that are already available in the literature.

The message that one should draw from that fruitful search is that there exist patches in the landscape of string vacua which lead to phenomenologically promising models. It seems also sensible to say that the main property of those patches is that they are endowed with local GUTs at a certain level. The task to accomplish in the future is then to analyze thoroughly the models and to figure out whether they lead to any kind of predictions.

Unfortunately, as we mentioned earlier, contact between these orbifold constructions and particle phenomenology requires a full understanding of the details of the vacua that each model offers. These details are associated with the dynamics of the moduli fields that describe all geometrical features of the particular compactification. For instance, the masses of quarks and leptons as well as the dynamics of supersymmetry breakdown depend on the vacuum expectation values (VEVs) of the moduli, which are arbitrary at tree level in many string constructions. Therefore, in order to achieve a prediction from string theory and simultaneously to avoid severe cosmological constraints, a mechanism that fixes the VEVs of the moduli and gives them large masses is needed. This mechanism is usually called moduli fixing or stabilization and will be discussed below.

Before discussing some results on brane-worlds, a final comment is in order. Both CY and orbifold compactifications have been applied in some other string constructions, such as M-theory compactifications [74, 75, 76, 77] and type I phenomenology [78, 79]. Although some phenomenological prospects have been identified in these scenarios, it is fair to say that they do not compete with the results obtained in CY and orbifold compactifications.

4.2. Brane-worlds

Just as in the case of the promising constructions obtained by compactifying the heterotic strings, in the type II theories there has been a lot of progress in recent years. The main idea is encoded in the Madrid model [80, 81], which is displayed in fig. 4. In that picture, type II string theories (A or B) appear compactified on a CY and include an orientifold background, so that one ends up with a 4D $N = 1$ supersymmetric theory. Without paying much attention on the cancellation of tadpoles, the minimal model that allows for one generation of quarks and leptons of the (MS)SM contains four stacks of D6 or D5 branes, depending on whether one starts with type IIA or type IIB. There is a stack of 3 D-branes called baryonic, a stack of 2 D-branes called left (for left-chiral quarks and leptons are supposed to arise from this stack), and two “stacks” of one brane, one called leptonic (only leptons couple to this one) and the other called right (right-chiral particles are attached to this brane). The emerging gauge group in 4D (where all branes share the whole space) is, as expected, $U(3)_u \times U(2)_b \times U(1)_c \times U(1)_d$.

As we have learnt before, in these scenarios the gauge bosons are associated with those
strings whose both ends are attached to the same stack. So e.g. we find the vector multiplets corresponding to the gluons in the baryonic stack and the weak vector bosons, in the left stack. Chiral matter states arise at the intersections of different stacks of branes. The direct consequence of the intersections proposed in fig. 4 is that the matter representations obtained at the four intersections are

\[(3, 2)_{a_q, b_q, c_q, d_q} \oplus (1, 2)_{a_\ell, b_\ell, c_\ell, d_\ell} \oplus (1, 1)_{a_e, b_e, c_e, d_e} \oplus (3, 1)_{a_d, b_d, c_d, d_d},\]

where the subindices denote the charges under the four different U(1) symmetries. Two observations must be made here. First, a linear combination of the U(1)s can yield the hypercharge, so that the resulting fields from this configuration are \(q_L, \ell_L, e_R\) and \(\bar{d}_R\). The second observation is that this is not a complete generation of quarks and leptons.

In order to rescue this model, one must rely on one unavoidable effect of the orientifold that we have not mentioned until here. Another way to interpret the worldsheet orientation reflection is by adding a D-brane parallel (not superposed) to the right stack, whose charges of the associated states be opposite to the \(c_i\) charges. We shall dub this D-brane right*. This way, two new matter representations are induced:

\[(3, 1)_{a_u, b_u, c^{*u}, d_u} \oplus (1, 1)_{a_\nu, b_\nu, c^{*\nu}, d_\nu} \rightarrow \bar{u}_R \text{ and } \bar{\nu}_R\]

that complete an (MS)SM matter generation.

Once the full matter content of one generation of quarks and leptons has been achieved, the next goal is to get the Higgs multiplets. This is achieved by proposing that, at some point, the right and right* brane stacks intersect the left stack. Each of these intersections will induce a matter representation \((1, 2)\) with some U(1) charges that shall combine to provide the correct hypercharges of \(h_u\) and \(h_d\).

This relatively simple setup provides all what one needs for the stable matter of nature, namely the first generation of quarks and leptons. For arriving at the second and third generations, one needs an additional ingredient. The key point is once more that chiral matter appear at the intersections of D-brane stacks. Thus, if e.g. there was a way to let the baryonic and the right (right*) stacks intersect more than once, one would obtain more than one copy of \(q_R\). (From now on, we shall refer to \(\bar{u}_R\) and \(\bar{d}_R\) as \(q_R\).)

In the original Madrid model, the six extra dimensions of the type II strings were not compactified on a CY, but on a factorizable 6D torus \(T^6 = T^2 \times T^2 \times T^2\). Since D5 and D6 branes do not cover the whole compact space, it is reasonable not to expect that the stacks intersect in every single \(T^2\). Another important element here is that they must not intersect.
exactly perpendicularly; the can intersect at angles. Supposing e.g. that the baryonic and the right/right* stacks intersect on one of the tori at an angle $\sim 3\pi/10$, one achieves the picture displayed in fig. 8. This leads to three copies of $q_R$. In fig. 8, the horizontal/green line represents the baryonic stack where as the other line represents the right/right* stack. Notice that a slightly different angle leads to a different number of quark copies. The crucial point here is that branes have to wrap non-contractible compact cycles for this mechanism to operate. Therefore, this result is also valid for much more complicated spaces than a torus, like $CY_3$, so long as they exhibit such cycles and the D-brane stacks wrap them. This kind of arrangement has been successfully done for all brane stacks [80, 81], leading to exactly three (MS)SM generations plus a pair of Higgs multiplets.

Despite its elegant properties, the original Madrid has a very important problem. Tadpole cancellation ensures that all equations of motion of the theory have solutions. If this is not satisfied, the theory is not consistent globally, since tadpoles are a result of considering the whole compactification scheme. In fact, the original Madrid model is challenged by this constraint. The way to fix this problem is by adding further brane stacks and, therefore, further chiral and vector multiplets to the effective theory, so that, besides the matter content of the MSSM, one ends up with an extra bunch of exotic chiral matter. Another more subtle issue is anomaly cancellation. Generically, the U(1)s that appear from D-brane stacks are anomalous, i.e. are broken explicitly by quantum effects. It is then difficult (but not impossible) to come up with a linear combination of all the U(1)s involved that yields the correct hypercharge for the SM states and simultaneously does not suffer from anomalies.

As in heterotic compactifications, there is another more phenomenological question one must pose to a promising construction. What are the interactions between the emerging fields and how large are they? The interactions between strings originating at different brane intersections have been computed via worldsheet instantons that can connect strings that are disconnected in the compact space [32]. This will usually involve the presence of three different D-brane stacks, which determine the boundary conditions of the worldsheet instanton contributing to this Yukawa coupling, see figure 9.

Roughly speaking, the instanton contribution to the Yukawa coupling will be given by evaluating the classical action $e^{-S_{cl}}$ on the surface of minimal area connecting the three intersections. As a result, as before, Yukawa couplings will depend on several moduli of the

![Figure 9: Yukawa couplings arise from instantonic interactions among states living at closely located intersections.](image)
theory, such as D-brane positions (open string moduli) and the compact manifold metric (closed string moduli).

In the Madrid model, since intersections such as the one exhibited in fig. 9 only appear for one generation of quarks \(^4\), the result is that only the third generation of quarks and leptons can acquire a mass once the Higgs mechanism occurs. This situation can be alleviated by the introduction of non-perturbative effects. In particular, it was found that D-brane instantons (i.e. instantons in 10D space-time) can generate admissible couplings \(^82\).

Unfortunately, this situation is in general much worse. Even though this kind of couplings can render phenomenologically useful interactions, in globally consistent models (i.e. those with tadpole cancellation) Yukawa couplings for all generations are forbidden by some U(1)s of the construction. Arguably, non-perturbative effects such as D-brane instantons can induce couplings of order unity (required for the top quark) while producing tiny couplings for the neutrinos, leading to semi-realistic Yukawa textures \(^82, 83\) in type IIB brane-worlds with D7 branes.

There have been several models since the Madrid model getting closer to the properties of the SM. Recently, a thorough scan of type IIA models with intersecting branes has been performed. The main result is that there seem to be many (perhaps equivalent)\(^5\) consistent models (i.e. with all tadpoles canceled) with the following features \(^84, 85, 86\):

- gauge group either \(U(3) \times Sp(2) \times U(1)^n\) or \(U(3) \times U(2) \times U(1)^n\), \(n \in \mathbb{Z}\), times a hidden sector gauge group;
- some models with and without \(U(1)_{B-L}\);
- hidden gauge group of rank smaller than four or no hidden gauge group;
- three (MS)SM matter generations;
- sizable Yukawa couplings for a generation of quarks and leptons;
- supersymmetric \(\mu\)-term and proton-decay operators forbidden by gauge symmetries;
- large amount of vectorlike exotics; and
- large number of Higgs multiplet pairs.

All these properties are very promising. However, it still unclear how to generate non-perturbative couplings in a consistent manner, such that masses for the first generations can arise. In this case, it is not known whether the exotics can acquire large masses. Research along these lines is on course.

Another more subtle issue of this problem in these constructions is that there is no explanation to why the coupling constants of the known microscopic interactions seem to unify at large energy scales. The idea of GUTs is simply not consistent with this kind of setups.

5. Fourth challenge: towards cosmology

5.1. Moduli stabilization

As we have seen, many of the phenomenological properties of the string constructions we have discussed depend on some free parameters encoded in scalar fields called moduli. The moduli parametrize the metric of the compact space (size and shape) as well as the position and angles (or equivalently magnetization) of the D-branes. Since the moduli are truly free at perturbative level, i.e. they exhibit no perturbative potential that binds them, a ‘good’ vacuum with phenomenologically appealing properties is equivalent to any other arbitrary vacuum. In fact, if the moduli remained free, transitions between the different universes predicted by the

\(^4\) To figure this out, one must place the left stack vertically and close to the left axis in fig. 8.

\(^5\) It is still under debate whether the models in the classification are different nor whether all exotics are truly vectorlike in the most promising constructions.
different vacua would be unavoidable. That definitively does not describe a universe like ours.
Furthermore, even if those vacuum transitions could be avoided, free scalar fields can also be
interpreted as mediators of as yet unknown forces and/or contribute in excess to the energy
density of the universe, spoiling thereby cosmological processes like reheating, recombination
and big bang nucleosynthesis. In short: moduli are not welcome!
There is only one way to solve this issue: **moduli stabilization**. This process consists in
searching for all possible sources of vacuum energy related to the moduli at perturbative and non-
perturbative level. There are some typical complications here. First, even after including non-
perturbative elements—which are not as well understood as the perturbative ones—, frequently
the resulting scalar potentials do not have a minimum at a finite value of the moduli; this is
regularly the problem of the dilaton in heterotic string theory. Second, a normal problem of
supergravity theories (like the effective limits of any string theory) is that, when there is a stable
vacuum, it is usually at values for which $V(\text{min}) < 0$. This *anti-de Sitter* vacuum is bad news
because the vacuum in which we live has $V(\text{min}) \sim 10^{-120}$, i.e. we live in a *de Sitter* vacuum
which is fairly flat, almost Minkowski.
In string phenomenology, besides constructions that potentially explain particle physics,
moduli stabilization has been independently studied in several scenarios. In type IIB string
theory, it has been shown that conjugating fluxes, non-perturbative effects (gaugino
condensation) and/or $\alpha'$ corrections could yield stable vacua [87, 88, 89]. There has
also been some progress in fixing Kähler and complex structure moduli in heterotic
compactifications on Calabi-Yau [90] and (generalized) half-flat [91] manifolds. In heterotic
orbifold compactifications [28, 29], it is frequently argued that the absence of fluxes (other than
the NS flux) can render much more difficult—if not impossible—the search for a stable vacuum.
However, it is known that the inherent target space modular symmetries [92, 93, 94] together
with nonperturbative effects such as gaugino condensation and string worldsheet instantons can
provide a scalar potential with metastable minima, at least in toy models [95, 96, 97, 98, 99].
Actually, this question has been recently addressed in explicit heterotic orbifolds [100, 101].

Some of the essential features of moduli stabilization are captured in the so-called KKLT
model [89]. Let us then describe how KKLT arrive to stabilized moduli in a de Sitter vacuum.
The starting point is the type IIB theory compactified on a CY with strong warping on a
particular region and that allows for fluxes to stabilize a large subset of moduli without breaking
supersymmetry (all the complex structure moduli) [87]. After fixing these moduli, the fluxes
render a constant contribution $W_0$ to the superpotential. In this ideal setup, the dilaton has
acquired a mass too and all tadpoles have been canceled. The next assumption is that *only* the
overall volume modulus, a Kähler modulus $T$, is left untouched and thus the string coupling
depends only on this modulus, $g_s^{-2} = \text{Re } T \equiv t$. Allowing for a strong interacting hidden sector
whose gauginos condensate inducing a term proportional to $\Lambda$ (see eq. (8)) in the superpotential.
The result is the superpotential

$$W = W_0 + Ae^{\alpha T}, \quad |W_0| \ll 1,$$

which, together with the tree-level Kähler potential $K = -3 \log \text{Re } T = -3 \log t$ yields the scalar potential

$$V(t) = -a^2 A^2 e^{-2at} / 6t.$$  \hfill (10)

This scalar potential is always negative and, therefore, its minimum must render an anti-de Sitter vacuum. This, by itself, is not so beautiful, although it must be pointed out that the modulus $t$ has been stabilized.

The next ingredient one must add in order to achieve a de Sitter vacuum is anti-D3 branes
in the region of strong warping. These branes have the effect of breaking supersymmetry while
introducing a new term in the scalar potential of the form $D/t^3$ (or $D/t^2$ if one introduces
D7 branes. In order to arrive at a de Sitter vacuum, $D$ must be of order $W_0^2$, very small. Once all these effects are included, the scalar potential of the KKLT proposal takes the shape for $V(t)$ displayed in fig. 10. In this figure, the values taken for the parameters are $A = 1, a = 0.1, W_0 = -10^{-4}$ and $D = 3 \times 10^{-9}$.

One of the unpleasant features of the KKLT proposal is that it requires an uplifting sector (the anti-D3 branes) which have no purpose other than pushing the minimum of $V$ to positive values. It has been shown that this might actually not be needed at all, as long as $T$ is not the only unfixed modulus [102]. In fact, if matter fields are available, it is very simple to arrive at a stabilized vacuum. Another unpleasant feature is the amount of fine-tuning needed to arrive at a de Sitter vacuum: in particular, the numbers $W_0$ and $D$ must be adjusted very carefully not to lose the minimum of the potential. A very simple example of moduli stabilization without uplifting nor fine-tuning, based partly on the proposal of [102], is presented in Appendix A. Finally, the sector in which the SM matter should be located is not addressed. Hence, a realization of this model within string theory seems very challenging.

Despite its challenging features, the KKLT scenario has served as inspiration for many stringy attempts to arrive at realistic vacua and in type II theories (and M-theory), this model has become the standard method to stabilize moduli. The key generic ingredients are: 1) most moduli stabilized in a supersymmetric fashion, 2) non-perturbative effects in order to fix the remaining modulus ($T$ in the case of type II theories, and the dilaton $S$ for the heterotic strings), and 3) a uplifting sector which, as we have seen, can contain a set of branes or some matter fields. It seems thus clear that using this or a similar method, one can obtain a stable compactification scheme in which some cosmological problems are avoided.

5.2. Inflation

A successful string model should be able not only to reproduce the particle physics of our universe, but also to explain the structure and dynamics of the cosmos. One persisting problem is to explain the origin and dynamics of the rapid period of expansion of the observable dimensions, known as inflation (see [103] for a useful review). In field theory, it is assumed that there is some field, the inflaton, whose potential is almost flat and has a minimum for some value of the inflaton. Many single and multi-field inflationary models have been concocted and lead to acceptable $e$-folds, primordial fluctuations and non-gaussianities.

In the context of string theory, there has been an effort mostly in the type IIB string [104, 105, 106, 107] to derive inflation from first principles. Remarkably, it has been found that, in certain scenarios, admissible fluctuations and nongaussianities are more natural in a universe dominated by multi-field dynamics [108], as is always the case in string constructions.

In general, the inflaton must be an (MS)SM singlet scalar whose potential energy satisfies a
Figure 11: Modified hybrid inflation. The field $\psi$ develops quickly a VEV inducing an inflationary potential for $\phi$. We have taken $V = (\psi - 3)^2/2 + \phi^2/16$.

series of constraints in order to produce inflation during an appropriate period of time. In the simplest case of slow-roll inflation, the scalar potential of the inflaton $V(\phi)$ must be very flat and the number of e-folds $-\int d\phi V/V'$ during the inflationary epoch must be something between 50 and 60. Since the moduli are scalar fields and, in some cases, their potential is very flat, it is natural to expect that some modulus play the role of the inflaton.

Interestingly, in the heterotic orbifolds with semi-realistic matter content, it has been found that some scalar fields whose squared masses lie between $10^{-100}$ and $10^{-30}$ in Planck units [100]. Since these quantities correspond to the curvature of the potential on the directions of those fields in moduli space, there are some almost flat directions which, nevertheless, lead to a minimum. The left panel of fig. 11 is a cartoon of this situation involving only two fields, where we have denoted them by $\phi$ and $\psi$.

Under these conditions, we can consider two different situations: when the coupling between $\phi$ and $\psi$ cannot be neglected, and when $V(\phi, \psi) \approx V(\phi) + V(\psi)$. In the former case, we might arrive at the scenario described by hybrid inflation models. If the field $\psi$ rapidly develops a VEV, the resulting scalar potential $V(\phi) = V(\phi, \psi_{\text{min}})$ is just the ideal slow-roll potential, as displayed in the right panel of fig. 11. In the latter case, the fields with quasi-flat potential might be suitable to accommodate the multi-field curvaton-inflaton scenario investigated in [108], yielding more accessible values for cosmological observables, such as the non-linearity parameter $f_{NL}$. Of course, whether or not we get such behaviors in actual orbifold models is still a question that must be studied.

On the other hand, given that most of the moduli stabilization mechanism are based on the proposal of KKLT, it is natural to investigate what prospects there are for inflationary scenarios in that case. In KKLT-based models, the Hubble constant during the last stages of a string theory inflation model should be below the expected gravitino mass, $H \lesssim 1$ TeV [104, 109]. Therefore, one should find a realistic particle physics model where the non-perturbative string theory dynamics occurs at the LHC scale or even lower, and inflation occurs at a density at least 30 orders of magnitude below the Planck energy density [104, 109]. This represents a problem since the inflationary energy scales are close to the compactification scales, which are much higher than $m_{3/2}$.

This is not the first time that string theory and supergravity have encountered cosmological problems associated with the small value of the gravitino mass and of the moduli fields. The famous gravitino problem and the cosmological moduli problem are haunting us for more than two decades [110, 111]. Now we see that the smallness of the gravitino mass leads to an additional problem in the context of string cosmology [104, 109, 112].

There is an interesting idea introduced in [107] to solve this issue, which does not involve giving up the solution of the naturalness problem through supersymmetry (i.e. the value of...
Inflation
Low-scale vacuum
Volume) ~ Log(Φ

\[V\]

10^{-18} M_P
10^{-47} M_P
1.3 M_P
19 M_P

Figure 12: An illustration of the scenario proposed in [107]. At relatively small volume, high-scale inflation occurs due to fine-tuned quantum corrections. After inflation the volume modulus evolves over a long range of many Planck scales, eventually settling in the large volume minimum with TeV gravitino mass. Although the barrier protecting from decompactification is very small compared to the initial energies, an attractor solution guides the fields to the minimum and prevents overshooting.

\[m_{3/2}\] is not increased radically) nor introducing new non-standard contributions to the Kähler potential plus huge fine-tuning [113]. The idea is that inflation should end with a runaway in field space, with the true minimum lying a very long distance in field space from the location where inflation occurs (e.g. about twenty Planckian distances, like in fig. 12). The problem to solve is that characteristic inflationary energy scales are much larger than those appropriate for supersymmetry breaking. The advantage of a runaway epoch is that evolution along runaway potentials like \(e^{-\lambda \phi}\) is one of the few efficient ways of naturally dissipating large quantities of energy and reducing the scale of the potential by many orders of magnitude. If the true minimum of the scalar potential lies a long way along the runaway direction, it naturally has much lower characteristic energy scales than apply during inflation.

In this case, supergravity models of high-scale inflation consistent with low-scale supersymmetry breaking, \(m_{3/2} \lesssim 1\) TeV, should have three stages. In the first, inflation occurs at high energy scales with \(m_{3/2} \gg 1\) TeV during inflation. In the second, inflation ends with the fields fast-rolling towards a runaway direction. In the third stage, the presence of small initial quantities of radiation drives the fields to a minimum. During the evolution towards the minimum large amounts of energy are dissipated. Due to this dissipation of energy, overshooting (i.e. \(\phi \rightarrow \infty\)) is avoided and the field falls into the global minimum of the potential in which \(m_{3/2} \sim 1\) TeV. This scenario is illustrated in fig. 12.

The proposed scenarios are promising. However, a crucial step is missing. One must compute all details in explicit constructions and pursue actual predictions, which do not involve black boxes. In this sense, current efforts in heterotic constructions as those cited above are needed and hopefully shall render useful fruits.

6. Final remarks
During the last decades, string phenomenologists have worked in many different scopes in order to bring string theory down to Earth. Models that reproduce many properties of the MSSM exist in all different scenarios, but they all are challenged by the computation of the magnitude of the
couplings, which eventually might lead to predictions. The situation is even more complicated when cosmological features are addressed. It seems nevertheless that the current tools developed in string theories could serve to overcome those obstacles; one must just perform the crucial computations which, in some cases, will require new algorithms and mathematical methods.

Before concluding, it is worth discussing one recurring worry: the so-called landscape problem. The issue is that there is a plethora of CY, orbifold and other similar spaces with mathematically consistent properties to compactify the extra dimensions of string theories, and only very few or one (if at all) can reproduce our universe. How can one be wise enough to find out which one is the correct one? This formulation of the issue can lead to believe that it is impossible to circumvent it. However, it is useful to compare it with finding out both the correct metric describing the gravitation dynamics of our universe and the action that describes particle physics. The number of mathematically admissible metrics and actions is actually infinite. And one must make use of one’s intuition to arrive at the correct results. In string theory, the situation is somewhat better. The mathematical consistency of the theory implies several constraints that yield a finite (but very large) set of possible compactification spaces and vacua. The task of a string phenomenologist is then to use (once again) her intuition and observable data to arrive at the compactification that describes our physics.

Acknowledgements
I am thankful to the organizers of the XIII Mexican School on Particles and Fields for their invitation. I am also indebted to O. Loaiza-Brito for the hospitality during the workshop and to G. Contreras for his support. This work was partially supported by CONACyT grants 82291 and 151234, and DGAPA-IACOD grant IA101811.

Appendix A. A simple exercise: de Sitter vacua with two moduli
Consider a hypothetical (type II) string compactification where all moduli excepting the overall volume $T$ and a (bulk) matter field $\Phi$ have been fixed in a supersymmetric fashion. Including the Kähler potential at tree level and supposing a generic superpotential with gaugino condensation, the resulting model in the supergravity limit is defined by

\[ W = \Phi^3 + A \Phi e^{-aT}, \quad K = -3 \log(T + \bar{T}) + \Phi \bar{\Phi}. \]  

(A.1)

Let us take $T = \alpha + i \tau$ and $\Phi = re^{i\varphi}$ and suppose that $\tau = \varphi = 0$. The corresponding scalar potential with $A = -1$ and $a = 15$ is depicted in fig. A1. We find that the minimum is located...
at \( r = 4 \times 10^{-3} M_{Pl} \) and \( \alpha = 0.7 M_{Pl} \), which leads to the vacuum energy (cosmological constant) \( V(\min) \approx 3 \times 10^{-12} M_{Pl}^4 \). In this vacuum, supersymmetry is broken because the gravitino mass is non-vanishing, \( m_{3/2} = e^{K/2}|W| \approx 10^{-7} M_{Pl} \).

The main property of this model is the coupling between the gaugino condensate \( (\lambda\lambda) \) and the matter field giving rise to the superpotential given in eq. (A.1). The value chosen for \( a \) is quite natural as it can arise from a simple SU(3) Yang-Mills theory in the hidden sector. Surprisingly, the gravitino mass is quite suppressed and, although is still unrealistic, there might be variations of the current model that provide an acceptable \( m_{3/2} \) as well as a proper cosmological constant. Despite its simplicity, this model seems much than the proposal by KKLT, since it avoids all naturalness issues present in previous attempts of stabilization.

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