Toward Realistic Intersecting D-Brane Models

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Abstract

We provide a pedagogical introduction to a recently studied class of phenomenologically interesting string models known as Intersecting D-Brane Models. The gauge fields of the Standard Model are localized on D-branes wrapping certain compact cycles on an underlying geometry, whose intersections can give rise to chiral fermions. We address the basic issues and also provide an overview of the recent activity in this field. This article is intended to serve non-experts with explanations of the fundamental aspects, and also to provide some orientation for both experts and non-experts in this active field of string phenomenology.

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1 INTRODUCTION

By now we have ample evidence that the Standard Model of particle physics extended by right-handed neutrinos or other mechanisms of neutrino masses describes nature with very high accuracy up to the energy range of the weak scale $E_W = 10^2$ GeV. The only missing ingredient is the Higgs particle itself, which is expected to be detected at the Tevatron or the LHC [[****AU: Please spell out LHC unless all readers will know this acronym.****]]. However, from a more formal point of view, the Standard Model is not completely satisfactory for ~~essentially~~ two reasons. First, it contains 26 free parameters (not counting arbitrary electric charges) like the masses and couplings of fermions and bosons, which have to be measured and among which no relation is apparent. Second, the Standard Model is formulated as a local four-dimensional quantum field theory and as such it does not include gravity. In fact, the Einstein theory of general relativity cannot simply be quantized according to the rules of local quantum field theory. Therefore, the physics we know of cannot describe our universe at very high energies where quantum effects of gravity become important.

Given these two formal shortcomings, in an ideal world we might hope that both problems actually have the same solution. Maybe there exists a fundamental quantum theory, which combines the Standard Model and general relativity into a unified framework and at the same time substantially reduces the number of independent parameters in the Standard Model. As this unified theory may be geometric in nature, one might envision that some of the structure of the Standard Model turns out to have a geometric origin as well.

~~We are currently in the situation that we do not know for sure~~ **Currently we do not know** what this final theory is, but at least we have a good candidate for it, which still has to reveal many of its secrets. This candidate theory of quantum gravity is called superstring theory and has been studied intensively during the last three decades. Since superstring theory is anomaly-free only in ten space-time dimensions, to make contact with the universe surrounding us we have to explain what happened to the other six dimensions without contradicting experiments. Compactifying à la Kaluza-Klein string theory on a compact six-dimensional space of very tiny dimensions, our visible world would be interpreted as an effective four-dimensional theory, where one only keeps the states of lowest mass. The question immediately arising is whether the formal string equations of motion allow for six-dimensional spaces such that the low-energy four-dimensional world resembles the Standard Model of particle physics. As a first approach, it would be too ambitious to require that all the couplings come out correctly. Instead, to begin with, one has to think about stringy mechanisms for generating gauge theories with chiral matter organized in replicated families.

The subfield of string theory concerned with such questions is called string phenomenology and has been pursued since the mid-1980s. ~~Already at that time~~ **By that time** is was clear that there exist two different types of ten-dimensional superstring theories containing gauge fields on the perturbative level. The heterotic string theories contain only closed-oriented strings and can support $SO(32)$ or $E_8 \times E_8$ gauge groups, whereas in the non-oriented Type I string theory, the gauge degrees of freedom arise from open strings, which can only support the gauge group $SO(32)$. In Type I string theory, the two possible orientations of the string are identified; in other words, one is gauging the word-sheet parity transformation.

From the mid-1980s to the mid-1990s, string theorists were mostly studying $E_8 \times E_8$ heterotic string compactifications, as it seemed to be more natural to embed the Standard Model gauge group $SU(3)_C \times SU(2)_W \times U(1)_Y$ into one of the E_8 factors and consider the second E_8 as a hidden gauge group, which might provide the infrared physics for supersymmetry breaking. In fact, it turned out that six-dimensional manifolds with $SU(3)$ holonomy, the so-called Calabi-Yau manifolds, also give rise to chiral fermions, which in the most simple scenario come in identical families where the multiplicity is given by one-half of the Euler number of the Calabi-Yau manifold. Many examples of such Calabi-Yau manifolds were constructed, including for instance toroidal orbifolds, hypersurfaces in weighted projective space, or toric varieties.

In the mid-1990s, string theory encountered an intellectual phase transition triggered by the realization that not only supersymmetric gauge theories but also string theories can be related by various dualities, some of which exchange weak and strong coupling. While ~~before people were~~ **previously scientists had been** merely studying perturbative aspects of string theory, now it was possible to move beyond the perturbative framework and to catch a glimpse of the non-perturbative physics of string theory. The conjectured web of string dualities relied on a speculative theory in eleven space-time dimensions, which was called M-theory. String theorists believe that this M-theory is actually the fundamental theory, of which the various string theories arise in certain perturbative limits.

In the process of establishing these dualities, it became clear that at the non-perturbative level, string theory is not only a theory of strings but also contains even higher dimensional objects called p-branes, which have p space-like and one time-like dimension. Surprisingly, the fluctuations of a certain subset of these p-branes, so-called D-branes, are again described by a string theory, which in this case is an open string theory with endpoints on the brane. ~~Since~~ **Because** at the massless level these D-branes support gauge fields, they are natural candidates for string phenomenology. The question is whether one can construct consistent string compactifications with D-branes in the background. The ~~easiest~~ **simplest** example is the aforementioned Type I string itself, which contains space-time filling D9-branes placed on top of the topological defect introduced by the gauging of the world-sheet parity. This already indicates that for getting models with D-branes, one should consider generalizations of the Type I string. Such models, now~~adays~~ called orientifolds, have been studied in the conformal field theory framework before [see (1) and references therein] and were ~~so to say~~ reinvented during the mid-1990s from a space-time point of view. The aforementioned defects were called orientifold planes.

Since their discovery, many orientifold models have been constructed, and there exists an extensive literature on this subject including some review articles, e.g., (2, 1). ~~The present~~ **This** article is not intended to be an additional review on general orientifold models, but instead focuses on a phenomenologically interesting class of orientifold models, which comes with its own intertwined history.

The class of models covered here has its origin in the observation that two generically intersecting D-branes can support chiral fermions on the intersection locus (3). Therefore, one is led to models ~~which~~ **that** not only contain D-branes on top of or parallel to orientifold planes, but also ~~one should~~ allow these D-branes to be placed such that there exist chiral intersections, as long as they do not violate the stringy consistency conditions. Historically, the first models of this kind were discussed in a T-dual formulation with magnetic fluxes in Type I string theory by ~~C.~~ Bachas (4). Providing the complete stringy picture of this early idea and showing its dual formulation in terms of intersecting D-branes, the first really intersecting D-brane models were constructed in (5, 6). Independently, supersymmetric compactifications of the Type I string to six-dimensions with magnetic fluxes were discussed in (7, 8). Intersecting branes and magnetized branes are equivalent descriptions (3). Therefore, without loss of generality, in this review article we stick to the more intuitive picture of geometrically intersecting D-branes. Non-chiral orientifold models with D-branes intersecting at angles had been considered even before [[**AU: Before what?**]] (9, 10, 11, 12, 13, 14).

~~To emphasize it again, i~~ **I**ntersecting branes provide a stringy mechanism for generating not only gauge symmetries but also chiral fermions, where family replication is achieved by multiple topological intersection numbers of various D-branes. Therefore, these models provide a beautiful geometric picture of some of the fundamental ingredients of the Standard Model.

After the introduction of these kinds of models, some generalizations and additional profound issues were discussed in (15, 16, 17). In the original models of (5), the closed string background was simply a flat torus, for which it could be shown that flat non-trivially intersecting D-branes always break supersymmetry explicitly at the string scale. Therefore, chiral models were necessarily non-supersymmetric. For a field theorist this is not a problem, as the Standard Model as we know is non-supersymmetric anyway. However, from the stringy point of view, supersymmetry is generally the mechanism which guarantees that string compactifications are stable. In order for a string vacuum to have a lifetime longer than the Planck (or better string) time, it seems desirable to start with a supersymmetric vacuum and then break supersymmetry softly in a controlled way. Though many papers in the literature deal with non-supersymmetric intersecting D-brane models, the reader should keep in mind that for these models, even though the open string sectors look amazingly similar to the Standard Model (18), one generically encounters stability problems in the closed string sector (19). Due to their popularity and some issues ~~which~~ **that** carry over to the supersymmetric models, we will also cover the non-supersymmetric models in this review, but our main focus will be on chiral supersymmetric models, first constructed in (20, 21). For them, Standard-like Models are much harder to construct and one has to consider more general than purely toroidal backgrounds, e.g., orbifolds.

In the original setting, one considered orientifolds of Type IIA string theory, which contain only orientifold six-planes, whose charges are canceled by introducing intersecting D6-branes. Such models can be defined on general six-dimensional manifolds, where the requirement of supersymmetry however implies this to be a Calabi-Yau manifold. Various generalizations with D-branes of other dimensions have been contemplated, but we think that the original models are the most natural class of intersecting D-brane models (~~as~~ for instance, **as** they are related to M-theory compactifications on G_2 manifolds). Therefore, throughout this article we will mainly work in this framework and only mention the possible generalizations.

Different aspects of these intersecting D-brane models have been discussed during the last four years in a large number of papers, which can be mainly categorized into three classes (we will provide the references in the appropriate sections of the main text). First, there are the stringy model building aspects, which in particular include the derivation of the stringy consistency conditions (R-R tadpole cancellation conditions) and the computation of the massless spectrum. Second, tools have been developed to compute for a given string model the four-dimensional low-energy effective action, which includes tree-level expressions for Yukawa couplings, higher point correlation functions, gauge couplings, Fayet-Iliopoulos terms, and Kähler potentials. Moreover, for the gauge couplings, ~~also~~ one-loop corrections have been computed. This program of determining the low-energy effective actionon-chiral states can typically obtain a string scale mass after deformations of the brane configuration. It is not complete yet and has mainly been applied to purely toroidal (orbifold) string backgrounds. Finally, using the results about the effective action, people discussed the phenomenological low-energy implications of intersecting D-brane models; some of them turn out to be rather model-independent, whereas others are not and might be used to discard certain models for phenomenological reasons. These are the three main aspects, but of course there exist relations of intersecting D-brane models to other branches of recent research such as M-theory compactification on G_2 manifolds or compactifications with non-trivial background fluxes. These latter developments will also be covered in this article, where, however, we do not provide a general introduction to G_2 manifolds or flux compactifications, as this would fill another review article. Another possible connection is to the phenomenological brane world ideas associated with possible large extra dimensions (22, 23, 24) that have been popular in recent years. ~~While~~ **Although** most intersecting D-brane constructions involve only small extra dimensions (within a few orders of magnitude of the inverse Planck scale), it is possible (and probably necessary for non-supersymmetric constructions) to consider internal spaces with large dimensions, providing a stringy realization of those ideas.

The aim of this article is twofold. First, it is intended to give a pedagogical introduction to the subject and to provide the main technical tools for the construction of intersecting D-brane models. It should allow non-experts to understand the main aspects of the subject and enable students to get started in this field. Second, we attempt to give as broad an overview as possible of developments in the field and to point out open questions. Of course, to be as complete as possible, we had to neglect many details, and we are aware that the topics we put special emphasis on reflect in some way our own preferences. We apologize to all those authors who feel that their work has not been covered to a degree they believe it deserves. Several articles of review type with slightly different emphases have appeared during the last **few** years (25, 26, 27, 28, 29, 30, 31).

2 ORIENTIFOLDS WITH INTERSECTING D-BRANES

Throughout this technical introduction into intersecting D-brane models, we assume that the reader is familiar at least at a textbook level (32, 33, 34, 35, 36, 37) with the basic notions of string theory including the concept of D-branes.

String compactifications from ten to four space-time dimensions have been studied throughout the history of string theory, but in the mid-1990s the second string theory revolution provided new insights into the constructions of four-dimensional vacua from M-theory. As with all the progress made during this exciting epoch, this had to do with the realization that string theory is not only a theory of either closed or open strings but also contains in its non-perturbative sector extended objects of higher dimensions, so called D-branes (see (38, 39, 36) for reviews on D-branes). These D-branes are charged under some of the massless fields appearing in the Ramond-Ramond (R-R) sector of the ten-dimensional Type IIA/B string theories. More concretely, a p-brane is an extended object with p space-like directions and one time-like direction and it couples to a $(p+1)$ form potential A_{p+1} as follows:

$$\mathcal{S}_p = Q_p \int_{D_p} A_{p+1}, \tag{1}$$

where the integral is over the $(p+1)$-dimensional world-volume of the D-brane and Q_p denotes its R-R charge. For BPS [[**AU: Please define BPS unless all readers will know this acronym.**]] D-branes in Type IIA string theory p is an even number and in Type IIB an odd one. Polchinski was the first to realize that the fluctuations of such D-branes can by themselves be described by a string theory (40), which in this case are open strings attached to the D-brane, i.e., with Dirichlet boundary conditions transversal to the D-brane and Neumann boundary conditions along the D-brane

$$\begin{aligned} \mu &= 0,\dots,p & \qquad \partial_\sigma X^\mu|_{\sigma=0,\pi} &= 0, \\ \mu &= p+1,\dots,9 & \qquad \partial_\tau X^\mu|_{\sigma=0,\pi} &= 0, \end{aligned} \tag{2}$$

where (σ, τ) denote the world-sheet space and time coordinates and X^μ the space-time coordinates. Their world-sheet superpartners are denoted as ψ^μ in the following. Upon quantization of an open string, the massless excitations $\psi^\mu_{-\frac{1}{2}}|0\rangle$ give rise to a $U(1)$ gauge field, which can only have momentum along the D-brane and is therefore confined to it. It is precisely the occurrence of these gauge fields which makes D-branes interesting objects for string model building. If one can construct string models with D-branes in the background, then one has a natural source of gauge fields, which are of fundamental importance in the Standard Model of particle physics. Placing N D-branes on top of each other the gauge fields on the branes transform in the adjoint representation of the gauge group $U(N)$.

In this section, we will discuss the general rules for constructing intersecting D-brane models. In subsection 2.1, employing the effective gauge and gravitational couplings, we discuss how the string scale of the intersecting D-brane models depends on the closed string moduli. Then in subsection 2.2, we discuss how chiral fermions arise at the intersection of D-branes. The fact that D-branes can intersect more than once in a compact space gives rise to the interesting feature of family replication, which will be discussed in subsection 2.3. In addition to R-R charges, D-branes also couple gravitationally which means that they have tension. To cancel the positive contribution to the vacuum energy from the tension of D-branes, we need to introduce negative tension objects known as orientifold planes. The notion of orientifolds will be discussed in subsection 2.4. The total R-R charge carried by the D-branes and orientifold planes has to vanish for consistency. Such tadpole cancellation conditions are derived in subsection 2.5. With the configuration of D-branes and orientifold planes that satisfy the tadpole conditions, one can derive the spectrum of massless open strings ending on the D-branes. The chiral part of the spectrum is summarized in subsection 2.6. In general, there are anomalous $U(1)$'s in intersecting D-brane models whose anomalies are canceled by the generalized Green-Schwarz mechanism as explained in subsection 2.7. In subsection 2.8, we discuss the conditions for the configuration of D6-branes to be supersymmetric. It turns out that they have to wrap around three-cycles known as special Lagrangian **(sLag)** cycles. Interestingly, the intersecting D6-brane models ~~which~~ **that** preserve $\mathcal{N}=1$ supersymmetry in four dimensions can be lifted to eleven-dimensional M-theory as compactifications on singular G_2 manifolds. The lift and the connection to how chiral fermions arise in the G_2 context are discussed in 2.9. As two warmup examples for later use, intersecting D-branes on T^6 and the $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold are presented in subsection 2.10.

2.1 The String Scale

The localization of gauge fields on D-branes provides a concrete stringy realization of the brane world scenario in which the Standard Model fields are confined on the branes, while gravity propagates in the bulk. As a result, the four-dimensional gauge couplings are determined by the volume of the cycles that the D-branes wrap around, whereas the gravitational coupling depends on the total internal volume. This opens up the possibility of lowering the string scale. More specifically, by dimensional reduction to four dimensions[1]:

$$\begin{aligned} \frac{1}{g_{YM}^2} &= \frac{M_s^{p-3} V_{p-3}}{(2\pi)^{p-2} g_s} \\ M_P^2 &= \frac{M_s^8 V_6}{(2\pi)^7 g_s^2}, \end{aligned} \tag{3}$$

[1]The factors of 2π were carefully worked out in (41).

where V_{p-3} is the volume of the $p-3$ cycle wrapped by a Dp-brane (which is in general different for different branes) and V_6 is the total internal volume. In this article, we will focus on models with intersecting D6-branes so that

$$g_{YM}^2 M_P = \sqrt{2\pi} M_s \frac{\sqrt{V_6}}{V_3}. \tag{4}$$

The experimental bounds on the masses of Kaluza-Klein replicas of the Standard Model gauge bosons imply that the volume of three-cycles cannot be larger than the inverse TeV scale generically. For a general internal space (such as a Calabi-Yau manifold), the volumes of the three-cycles are not directly constrained by the scale of the total internal volume, and can be much smaller than $\sqrt{V_6}$. In this case, a large Planck mass can be generated from a large total internal volume. This is precisely the idea of the large extra dimension scenario.

However, for intersecting D6-brane models in toroidal backgrounds, V_3 is of the same order as $\sqrt{V_6}$ (since for chiral models, there is no dimension transverse to all the branes) so the string scale is of the order of the Planck scale M_P. There is, however, more freedom than in theories with only closed strings (e.g., the heterotic string), and this could be used to lower the string scale to, e.g., 10^{16} GeV, a certainly desirable choice for Grand Unified Models.

2.2 Chirality

One of the main features of the Standard Model is that the light fermionic matter fields appear in chiral representations of the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge symmetry such that all gauge anomalies are canceled. Considering just parallel D-branes in flat space one does not get chiral matter on the branes, so that one has to invoke an additional mechanism to realize this phenomenologically very important feature. Essentially, ~~so far~~ two ways have been proposed to realize chirality for the D-brane matter spectrum. The first one is to place the D-brane not in flat space, but on so-called orbifold (or conifold) singularities; the second is to let the D-branes intersect at non-trivial angles (3). We will discuss here only the second mechanism in more detail and refer the interested reader for the first mechanism to the existing literature [see for instance (1, 42) and references therein].

To be more precise, consider two D6-branes sharing the four-dimensional Minkowskian space-time. This means that in the six-dimensional transversal space the branes are three-dimensional and wrap a three-dimensional cycle. In general position, two such branes do intersect in a point in the internal space. Consider the simple case of a flat six-dimensional internal space. Choosing light cone gauge, let us introduce complex coordinates $z^i = x^i + iy^i$ with $i = 0, \ldots, 3$. Then two D6-branes cover the z^0 plane and intersect in the other directions as shown in Figure 1.

Placing for convenience one D-brane along the x^i axes, an open string stretched between two intersecting D-branes has the following boundary conditions

$$\begin{aligned} \sigma = 0: &\qquad \partial_\sigma X^i = \partial_\tau Y^i = 0 \\ \sigma = \pi: &\qquad \partial_\sigma X^i + \tan(\Delta\Phi_i)\, \partial_\sigma Y^i = 0 \\ &\qquad -\tan(\Delta\Phi_i)\, \partial_\tau X^i + \partial_\tau Y^i = 0 \end{aligned} \tag{5}$$

which in complex coordinates read

$$\begin{aligned} \sigma = 0: &\qquad \partial_\sigma(Z^i + \overline{Z}^i) = \partial_\tau(Z^i - \overline{Z}^i) = 0 \\ \sigma = \pi: &\qquad \partial_\sigma Z^i + e^{2i\Delta\Phi_i}\partial_\sigma \overline{Z}^i = 0 \\ &\qquad \partial_\tau Z^i - e^{2i\Delta\Phi_i}\partial_\tau \overline{Z}^i = 0. \end{aligned} \tag{6}$$

Now, implementing these boundary conditions in the mode expansion of the fields

Z^i and $\overline{Z}^i$, one finds (3)

$$Z^i(\sigma,\tau) = \sum_{n\in} \frac{1}{(n+\epsilon_i)}\,\alpha^i_{n+\epsilon_i}\,e^{-i(n+\epsilon_i)(\tau+\sigma)} + \sum_{n\in} \frac{1}{(n-\epsilon_i)}\,\tilde{\alpha}^i_{n-\epsilon_i}\,e^{-i(n-\epsilon_i)(\tau-\sigma)} \quad (7)$$

with $\epsilon_i = \Delta\Phi_i/\pi$ for $i \in \{1,2,3\}$. Therefore the bosonic oscillator modes of the fields $Z^1,\ldots,Z^3$ are given by

$$\alpha^i_{n+\epsilon_i}, \quad \tilde{\alpha}^i_{n-\epsilon_i} \quad (8)$$

Similarly, for the world-sheet fermions the modes are $\psi^i_{n+\epsilon_i}$ and $\tilde{\psi}^i_{n-\epsilon_i}$ in the R-R sector and with an additional 1/2-shift in the Neveu-Schwarz Neveu-Schwarz (NS-NS) sector. Therefore, in analogy to the closed string sector, an open string between two intersecting D-branes can be considered as a twisted open string. As a consequence, for all ϵ_i non-vanishing, there are only two zero modes in the R-R sector, $\psi^1_0, \tilde{\psi}^1_0$, which give rise to a twofold degenerate R-R ground state. The GSO [[****AU: Define GSO unless all readers will know this acronym.****]]projection eliminates one half of these states, so that one is left with only one fermionic degree of freedom. Taking into account also the open string with the opposite orientation between the two D6-branes, one finally gets two fermionic degrees of freedom corresponding to one chiral Weyl-fermion from the four-dimensional space-time point of view. To summarize, we have found that two generically intersecting D6-branes give rise to one chiral fermion at the intersection point. If we now consider the intersection between a stack of M D6-branes with another stack of N D6-branes, it is clear that~~the for instance~~ **, for example, the** left-handed chiral fermion transforms in the bi-fundamental representation of the $U(M)\times U(N)$ gauge symmetry. We choose the convention that this is the $(\overline{M},N)$ representation of the gauge group. As such this result is not invariant under the exchange of the role of M and N. This can be remedied by giving an orientation to the branes and by assigning a sign to the intersection on each plane as shown in Figure 2. A negative intersection simply means that one gets a left-handed chiral fermion transforming in the conjugate representation of the gauge group. The intersection defined this way is anti-symmetric under exchange of the two branes.

2.3 Family Replication

In the last section we have seen that intersecting D-branes can be a source for chiral fermions, which makes them very interesting candidates for model building. However, chiral fermions in the Standard Model come in three families differing only by their mass scale. Therefore, it is important to search for a mechanism for family replication. As we will see by considering intersecting branes on compact backgrounds, such a mechanism automatically arises.

In the non-compact flat background depicted in Figure 1, it is clear that the intersection number can only be ± 1. However, in the compact case ~~like for instance~~ **such as** a torus, it can be easily seen that the intersection number can be larger than one. Assuming for simplicity that the background is a six-dimensional torus with complex structure chosen such that it can be written as $T^6 = T^2 \times T^2 \times T^2$, a large class of D6-branes cover only a one-dimensional cycle on each factor T^2. Such D6-branes have been called factorizable in the literature and are described by three pairs of wrapping numbers (n^i, m^i) along the fundamental 1-cycles of three T^2s. In Figure 3, we have shown two such wrapped D6-branes with wrapping numbers $(1,0)(1,1)(2,1)$ for the first D-brane and $(0,1)(1,-1)(1,-1)$ for the second one.

~~From the picture one reads off that~~ **[[**AU: Or recast the sentence for**

clarity.]] ~~t~~ **T**he intersection number between the two D6-branes is $I_{ab} = 6$, which is just one simple example of the general expression for the intersection number

$$I_{ab} = \prod_{i=1}^{3} (n_a^i\, m_b^i - m_a^i\, n_b^i). \tag{9}$$

By deforming the D6-branes ~~it is clear that~~ **,** one can easily generate additional intersections, but they always come in pairs with positive and negative sign, so that the net number of chiral fermions remains constant. Therefore, what really counts the net number of chiral fermions is the topological intersection number, which only depends on the homology classes of the two branes. In our case, the homological three-cycles are simply products of three one-cycles, where the homological one-cycles on each T^2 are characterized just by the wrapping numbers (n_a, m_a).

Generalizing the set-up we have introduced so far, we consider compactifications of Type IIA string theory on a six-dimensional manifold $\mathcal{M}$. To preserve $\mathcal{N} = 1$ supersymmetry, the manifold $\mathcal{M}$ is a Calabi-Yau manifold. As a topological space, $\mathcal{M}$ has homological three-cycles π_a, $a \in \{1, \ldots, K\}$, on which we can wrap N_a D6-branes. From the effective four-dimensional point of view, we obtain gauge fields of $\prod_{a=1}^{K} U(N_a)$ localized on the seven-dimensional world-volume of the D6-branes. Additionally, one gets chiral fermions localized on the four-dimensional intersection locus of two branes which come with multiplicity given by the topological intersection number $\pi_a \circ \pi_b$ and transform in the $(\overline{N}_a, N_b)$ representation of the gauge group.

As we will discuss in the following sections, tadpole cancellation and supersymmetry impose certain constraints on the three-cycles **that** the D6-branes are wrapped around.

2.4 Orientifolds

As has been pointed out in (19), non-supersymmetric models, though easy to handle, are unstable in the sense that their perturbative scalar potential, which is due to the so-called NS-NS tadpoles, gives rise to runaway behavior for many of the closed string moduli fields including the dilaton. In addition, for constructions based on toroidal-type compactifications, the volume of three cycles is of the same order of magnitude as the square root of the volume of the internal space. Thus, for the case of intersecting D6-branes there is no direction in the internal space that can be taken large compared to the Planck radius while keeping the correct values of gauge couplings and the Planck scale in four-dimensions (see eqs.(3) in section 2.1.). Therefore, in this case the string scale cannot be much below the Planck scale and thus both the NS-NS tadpole contributions to the potential as well as radiative corrections in the effective theory are large, i.e., of the order of the Planck scale. There is a chance that NS-NS tadpoles might be stabilized by non-perturbative effects or by turning on fluxes, but nevertheless, in order to stay on firm ground from the string theory perspective, we prefer to mainly consider supersymmetric models.

The set-up introduced so far with intersecting D6-branes in Type IIA compactifications always breaks supersymmetry. This can be seen as follows. For a globally supersymmetric background the vacuum energy has to vanish. However, all D6-branes have a positive contribution to the vacuum energy, as their tension is always positive and therefore they break supersymmetry. The only way to finally find non-trivial supersymmetric models is by introducing objects of negative tension into the theory. It is well known that such objects exist in string theory and that they naturally occur in so-called orientifold models.

An orientifold is the quotient of Type II string theory by a discrete symmetry group G including the world-sheet parity transformation $\Omega:(\sigma,\tau)\to(-\sigma,\tau)$. As a consequence, the resulting string models contain non-oriented strings and their perturbative expansion also involves non-oriented surfaces like the Klein-bottle. Dividing out by such a symmetry, new objects called orientifold planes arise, whose presence can be detected for instance by computing the Klein-bottle amplitude

$$K=\int_0^\infty \frac{dt}{t}\mathrm{Tr}\left(\frac{\Omega}{2}\, e^{-2\pi t(L_0+\overline{L}_0)}\right). \tag{10}$$

These objects, though non-dynamical, do couple to the closed string modes and in particular they carry tension and charge under some of the R-R fields. In other words, there exist non-vanishing tadpoles of the closed modes on the orientifold planes which, as it turns out, can have opposite sign than the corresponding terms for D-branes. Since the overall charge one puts on a compact space has to vanish by Gauss's law, the contribution from the orientifold planes and the D-branes have to cancel. We would like to emphasize that for orientifolds, the presence of D-branes in the background is in most cases not an option but a necessity.

Which are the appropriate orientifolds to consider so that intersecting D6-branes might cancel the tadpoles? Clearly we need O6-planes, meaning that the world-sheet parity has to be dressed with an involution, locally reflecting three out of the six internal coordinates, and being a symmetry of the internal space. Let us assume that $\mathcal{M}$ admits a complex structure so that we locally can introduce complex coordinates z^i. Now, we consider Type IIA string theory divided out by $\Omega\overline{\sigma}(-1)^{F_L}$, where F_L denotes the left-moving space-time fermion number [2] and

[2] Note that the $(-1)^{F_L}$ factor was not explicitly written down in many of the papers on intersecting D-brane models.

$\overline{\sigma}$ an isometric anti-holomorphic involution of $\mathcal{M}$. This acts on the Kähler class J and the holomorphic covariantly constant three-form Ω_3 as

$$\overline{\sigma}\, J = -J, \qquad \overline{\sigma}\, \Omega_3 = e^{2i\varphi}\overline{\Omega}_3 \tag{11}$$

with $\varphi \in \mathbb{R}$. For $\varphi = 0$ in local coordinates this can be thought of as complex conjugation. As a result we get an orientifold O6-plane localized at the fixed point locus of $\overline{\sigma}$, which topologically is a three-cycle π_{O6} in $H_3(\mathcal{M}, \mathbb{Z})$. To cancel the resulting massless tadpoles, we introduce appropriate configurations of intersecting D6-branes wrapping homological three-cycles π_a. For $\overline{\sigma}$ to be a symmetry of the brane configuration, one also needs to wrap D6-branes on the $\overline{\sigma}$ image three-cycles π'_a. As a new feature, in orientifold models it is also possible to get orthogonal and symplectic gauge symmetries. The rule is very simple. If a three-cycle is invariant under the anti-holomorphic involution one gets either $SO(2N_a)$ or $SP(2N_a)$ gauge symmetry; if the cycle is not-invariant, one gets $U(N_a)$. Figure 4 depicts in a simplified way the set-up discussed in this section.

2.5 R-R Tadpole Cancellation

As we have mentioned already, tadpole cancellation provides some constraints on the positions of the O6-planes and D6-branes, which we now summarize. Historically, for deriving the tadpole cancellation conditions, one used an indirect method by first computing, using conformal field theory techniques, the one-loop Klein-bottle, annulus and Möbius strip diagrams, and extracting from the corresponding tree-channel amplitudes the infrared divergences due to massless tadpoles. Employing a direct method using the Dirac Born Infeld action, here we essentially follow (43, 44), where ~~also~~ more details of the derivation can be found.

Consider the part of the supergravity Lagrangian where the R-R field C_7 appears

$$\mathcal{S} = -\frac{1}{4\kappa^2}\int_{\mathbb{R}^{3,1}\times\mathcal{M}} dC_7\wedge\star dC_7 + \mu_6\sum_a N_a\int_{\mathbb{R}^{3,1}\times\pi_a} C_7 \tag{12}$$

$$+ \mu_6\sum_a N_a\int_{\mathbb{R}^{3,1}\times\pi'_a} C_7 - 4\mu_6\int_{\mathbb{R}^{3,1}\times\pi_{\mathrm{O6}}} C_7, \tag{13}$$

where the ten-dimensional gravitational coupling is $\kappa^2 = \frac{1}{2}(2\pi)^7(\alpha')^4$ and the R-R charge of a D6-brane reads $\mu_6 = (\alpha')^{-\frac{7}{2}}/(2\pi)^6$. Note that here we have assumed that the orientifold planes are of type $O^{(-,-)}$, i.e., they carry negative tension and R-R charge. Recall that D-branes in this convention carry positive tension and R-R charge. Such models have also been called orientifolds without vector structure. The resulting equation of motion for the R-R field strength $G_8 = dC_7$ is

$$\frac{1}{\kappa^2}\, d\star G_8 = \mu_6\sum_a N_a\,\delta(\pi_a) + \mu_6\sum_a N_a\,\delta(\pi'_a) - 4\mu_6\,\delta(\pi_{\mathrm{O6}}), \tag{14}$$

where $\delta(\pi_a)$ denotes the Poincaré dual three-form of π_a. Since the left-handed side in Equation (14) is exact, the R-R tadpole cancellation condition boils down to just a simple condition on the homology classes

$$\sum_a N_a\,(\pi_a + \pi'_a) - 4\pi_{O6} = 0. \tag{15}$$

The above condition [[**AU: Please recast this sentence for clarity.**]]implies that the overall three-cycle all the D-branes and orientifold planes wrap is trivial in homology. This is a restrictive condition but it is moderate enough to admit non-trivial solutions with branes are not simply placed right on top of the orientifold plane. Note that so far we have not assumed supersymmetry and that therefore Equation (15) does not automatically guarantee the NS-NS tadpoles to be canceled as well.

However, it is important to note that the above method, using the Dirac Born Infeld action together with supergravity, does not take into account all of the R-R charges carried by the D-branes. The reason is that D-brane charges are classifed by K-theory groups rather than homology groups, and the R-R fields in general are not simply p-forms ~~like above~~ (see, e.g., (45) for a more detailed discussion). Indeed, as pointed out in (46), the cancellation of homological R-R charges [i.e., the conditions (15) above] are not sufficient to ensure that all the R-R tadpoles vanish and hence the consistency of the models. The inconsistencies due to uncanceled K-theory charges would show up as discrete global anomalies (47) either in the low-energy spectrum or on the world-volume of a probe D-brane (46). One way to heuristically derive these constraints is to introduce probe D-branes with an $Sp(2n)$ gauge group, and require ~~that~~ the total number of fundamental representations in their world-volume theory to be even. These K-theory constraints are widely unnoticed in the model building literature because for simple models, they are automatically satisfied. However, these consistency constraints are far from trivial. For example, such K-theory constraints for the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold were derived in (48, 49) and have **been** shown to play an important role in the construction of more realistic models. We will discuss such constraints in more detail in subsection 2.10.

2.6 The Massless Spectrum

For model building purposes, it is very important to have control over the massless spectrum arising from any kind of string compactification. For the orientifold models with intersecting D6-branes, the chiral spectrum arising from the various open string sectors can be determined just from the intersection numbers of the three-cycles **that** the D6-branes are wrapped around. For simplicity let us assume that all D6-branes wrap three-cycles not invariant under the anti-holomorphic involution, so that the gauge symmetry is $\prod_a U(N_a)$. For this case the general rule for determining the massless left-handed chiral spectrum is presented in Table 1. Open strings stretched between a D-brane and its $\o\sigma$ image are the only ones left invariant under the combined operation $\Omega\overline{\sigma}(-1)^{F_L}$. Therefore, they transform in the antisymmetric or symmetric representation of the gauge group, indicating that the price **that** we have to pay by considering intersecting D-branes in an orientifold background is that more general representations are possible for the chiral fermions. Sometimes this is an advantage, like for constructing $SU(5)$ Grand Unified Models, but sometimes the absence of such fermions imposes new conditions on the possible D-brane set-ups.

The rule for the chiral spectrum in Table 1 is completely general and, as was demonstrated in (43), the chiral massless spectra from many orientifold models discussed using conformal field theory methods in the existing literature can be understood in this framework.

Moreover, one can easily check that the R-R tadpole cancellation condition (15) together with [[**AU: Please clarify–the table guarantees? Maybe the data from the table?****]]Table 1 guarantees the absence of non-Abelian gauge anomalies. Naively, there exist Abelian and mixed Abelian-non-Abelian anomalies, as well as gravitational anomalies. However, we shall see in the subsequent section that all of these are canceled by a generalized Green-Schwarz mechanism.

To apply Table 1 to concrete models, one has to compute the intersection numbers of three-cycles, which by itself is ~~in general~~ not an easy task. However, there exist backgrounds for which generic rules can be presented. ~~Besides~~ **In addition to** the simplest case of just a torus T^6, toroidal orbifolds, such as $T^6/\mathbb{Z}_N$ or $T^6/\mathbb{Z}_N \times \mathbb{Z}_M$, are natural candidates for string backgrounds. Therfore, let us discuss the application of Table 1 to such orbifolds in some more detail.

~~Recall that t~~ **T**he spectrum in Table 1 is meant to be computed using the intersection numbers on the resolved orbifold and not on the ambient torus. There are some three-cycles π_a on the orbifold space ~~which~~ **that** are inherited from the torus. In the Kaluza-Klein reduction on the orbifold, they correspond to massless modes in the untwisted closed string sector of the theory. In general, three-cycles π_a^t on the torus are arranged in orbits of length N under a $\mathbb{Z}_N$ orbifold group, i.e.,

$$\pi_a^o = \sum_{j=0}^{N-1} \Theta^j \, \pi_a^t, \tag{16}$$

where Θ denotes the generator of $\mathbb{Z}_N$. Such an orbit can then be considered as a three-cycle of the orbifold, where the intersection number is given by

$$\pi_a^o \circ \pi_b^o = \frac{1}{N} \left(\sum_{j=0}^{N-1} \Theta^j \, \pi_a^t \right) \circ \left(\sum_{k=0}^{N-1} \Theta^k \, \pi_b^t \right). \tag{17}$$

Beside [[****AU: Beside–do you mean in addition to or next to?****]]these untwisted three-cycles, certain twisted sectors of the orbifold action can give rise to additional so-called twisted three-cycles, which correspond to massless fields in the twisted sectors of the orbifold. Since these twisted three-cycles are not explicitly needed in this article, we refer the reader to the existing literature (43, 50, 51, 52) to see how these twisted cycles can be appropriately dealt with.

Table 1 only gives the chiral spectrum of an intersecting D6-brane model. To compute the generally moduli-dependent non-chiral spectrum, one has to employ the usual techniques of conformal field theory. Therefore, the Higgs sector of a given model is under less analytic control than the chiral matter sector.

2.7 Generalized Green-Schwarz Mechanism

Given the chiral spectrum of Table 1, we have stated that the non-Abelian gauge anomalies of all $SU(N_a)$ factors in the gauge group vanish. On the other hand, the Abelian, the mixed Abelian-non-Abelian, and the mixed Abelian-gravitational anomalies naively do not. However, as string theory is a consistent theory, it provides another mechanism to cancel these anomalies. This is the so-called Green-Schwarz mechanism (53), which can be generalized to the intersecting D-brane case (16). Here let us discuss in some more detail the mixed Abelian-non-Abelian anomalies.

Computing the $U(1)_a - SU(N_b)^2$ anomalies in the effective four-dimensional gauge theory one finds

$$A_{ab} = \frac{N_a}{2} \left(-\pi_a + \pi'_a\right) \circ \pi_b. \tag{18}$$

for each pair of stacks of D-branes. On each stack of D6-branes there exist Chern-Simons couplings of the form

$$\int_{\mathbb{R}^{1,3}\times\pi_a} C_3 \wedge \mathrm{Tr}\left(F_a \wedge F_a\right), \qquad \int_{\mathbb{R}^{1,3}\times\pi_a} C_5 \wedge \mathrm{Tr}\left(F_a\right) \tag{19}$$

where F_a denotes the gauge field on the D6$_a$-brane. Now we expand every three-cycle π_a and π'_a into an integral basis (α^I, β_J) of $H_3(M, \mathbb{Z})$ with $I, J = 0, \ldots, h_{21}$.

$$\pi_a = e^a_I\, \alpha^I + m^J_a\, \beta_J, \qquad \pi'_a = (e^a_I)'\, \alpha^I + (m^J_a)'\, \beta_J. \tag{20}$$

This allows us to define the four-dimensional axions Φ_I and 2-forms B^I as

$$\begin{aligned} \Phi_I &= \int_{\alpha^I} C_3, \qquad \Phi^{I+h^{(2,1)}+1} = \int_{\beta_I} C_3, \\ B^I &= \int_{\beta_I} C_5, \qquad B_{I+h^{(2,1)}+1} = \int_{\alpha^I} C_5. \end{aligned} \tag{21}$$

In four dimensions $(d\Phi_I, dB^I)$ and $(d\Phi^{I+h^{(2,1)}+1}, dB_{I+h^{(2,1)}+1})$ are Hodge dual to each other. The general couplings (19) can now be dimensionally reduced to four dimensions and yield axionic couplings of the form

$$\int_{\mathbb{R}^{1,3}} \Phi_I \wedge \mathrm{Tr}\,(F_a \wedge F_a)\,, \qquad \int_{\mathbb{R}^{1,3}} B^I \wedge F_a. \tag{22}$$

The tree-level contribution to the mixed gauge anomaly described by these couplings takes the form depicted in Figure 5, and, adding up all these terms taking the R-R tadpole conditions into account, one can show that the result has precisely the form (18) and cancels the field theoretic anomaly (43). By the same mechanism~~also~~ , the Abelian and mixed gravitational–gauge anomalies are canceled, where the latter ~~ones~~ arise from the $U(1)_a - G - G$ triangle diagram and are given by

$$A_a^{(G)} = 3\, N_a\, \pi_{O6} \circ \pi_a. \tag{23}$$

This anomaly is canceled by the Chern-Simons coupling

$$\int_{\mathbb{R}^{1,3}\times\pi_a} C_3 \wedge \mathrm{Tr}\,(R \wedge R)\,. \tag{24}$$

A second important effect of these couplings is that some of the $U(1)$ gauge fields pair up with the axions to become a massive gauge field. The axionic couplings have the detailed form

$$\int_{\mathbb{R}^{1,3}} N_a(\pi_{a,I} - \pi'_{a,I})\, B^I \wedge F_a \tag{25}$$

with $\pi_{a,I} \in \{e_I^a, m_a^I\}$ depending on the index I. It has been pointed out in (18, 54) that in general, not only the anomalous $U(1)$s receive a mass, but also some of the anomaly-free ones, which are given by the kernel of the matrix

$$M_a^I = N_a(\pi_{a,I} - \pi'_{a,I}). \tag{26}$$

Therefore, to determine the low-energy spectrum, one has to carefully analyze these quadratic couplings. The massive $U(1)$s still give rise to perturbative global

$U(1)$ symmetries of the low-energy effective action (55), which severely constrain the allowed couplings.

2.8 Supersymmetric D-branes

So far we were not assuming anything more about the D6-branes than that they are wrapping some homological three-cycles in the background geometry. If one is interested in supersymmetric models, further constraints on the bulk geometry and the cycles on which the D6-branes wrap have to be imposed. Throughout this section we assume that $\mathcal{M}$ is a Calabi-Yau manifold so that the closed string bulk sector of the Type IIA orientifold preserves $\mathcal{N}=1$ supersymmetry. First of all, one has to require that each D-brane by itself preserves supersymmetry, i.e., it has to be a BPS brane. As was shown in (56), this implies that the three-cycles the D6-branes are allowed to wrap have to be ~~so-called~~ special Lagrangian (sLag) cycles, which are defined as follows.

On a Calabi-Yau manifold, there exist a covariantly constant holomorphic three-form, Ω_3, and a Kähler 2-form J. A three-cycle π_a is called Lagrangian if the restriction of the Kähler form on the cycle vanishes

$$J|_{\pi_a} = 0. \tag{27}$$

If the three-cycle in addition is volume minimizing, which can be expressed as the property that the imaginary part of the three-form Ω_3 vanishes when restricted to the cycle,

$$\Im(e^{i\varphi_a}\,\Omega_3)|_{\pi_a} = 0, \tag{28}$$

then the three-cycle is called a sLag cycle. The parameter φ_a determines which $\mathcal{N}=1$ supersymmetry is preserved by the brane. Thus, different branes with different values for φ_a preserve different $\mathcal{N}=1$ supersymmetries. One can show that (28) implies that the volume of the three-cycle is given by

$$\mathrm{Vol}(\pi_a) = \left| \int_{\pi_a} \Re(e^{i\varphi_a}\,\Omega_3). \right| \tag{29}$$

A shift of $\varphi_a \to \varphi_a + \pi$ corresponds to exchanging a D-brane by its anti-D-brane, where the D-brane really satisfies (29) without taking the absolute value. Therefore a supersymmetric cycle π_a is calibrated with respect to $\Re(e^{i\varphi_a}\Omega_3)$.

Let us define locally the holomorphic 3-form Ω_3 and the Kähler form J by

$$\Omega_3 = dz_1 \wedge dz_2 \wedge dz_3, \quad J = i \sum_{i=1}^{3} dz_i \wedge d\bar{z}_i. \tag{30}$$

Let us choose $\overline{\sigma}$ to be just complex conjugation in local coordinates. Then from $\overline{\sigma}(\Omega_3) = \overline{\Omega}_3$ and $\overline{\sigma}(J) = -J$ it follows that the fixed three-cycle of the anti-holomorphic involution is a sLag cycle with $\varphi_a = 0$. Therefore, to finally obtain a globally $\mathcal{N} = 1$ supersymmetric intersecting D-brane model, all D6-branes have to wrap sLag three-cycles, which are calibrated with respect to the same three-form $\Re(\Omega_3)$. It has been checked in (57, 43) that for such globally supersymmetric configurations, ~~indeed~~ the NS-NS tadpoles cancel precisely if the R-R tadpoles are canceled. One can ~~indeed show that precisely~~ **show that** if two branes are relatively supersymmetric, one of the four complex world-sheet bosons becomes massless and extends the massless chiral fermion at the intersection point to a complete $\mathcal{N} = 1$ chiral supermultiplet.

In section 2.10.1 we shall see that for the toroidal (orbifold) compactifications, this condition becomes a simple geometric condition on the intersection angles of each D-brane with respect to the orientifold plane.

2.9 Lift to G_2 Compactifications of M Theory

In this section we briefly discuss the relation between intersecting D6-brane models and M-theory compactifications on G_2 manifolds. This needs some more advanced mathematical notions and is not really relevant for understanding the rest of the review. For completeness, however, we summarize some key ideas here.

Globally $\mathcal{N} = 1$ supersymmetric intersecting D6-brane models have also shed light on how chiral fermions arise in G_2 compactifications of M theory. D6-branes and O6-planes are special because they correspond to pure geometry at strong coupling (unlike other branes which carry additional sources, i.e., M-branes or G-fluxes). Therefore, from the number of supercharges **that** the background preserves, the globally $\mathcal{N} = 1$ supersymmetric intersecting D6-brane models are expected to lift up to eleven-dimensional M-theory compactification on singular G_2 manifolds (58, 21, 59, 60, 61). In the Type IIA picture, chiral fermions are localized at the intersection of D6-branes. Away from the intersections of IIA D6-branes and/or O6-planes, the IIA configuration corresponds to D6-branes and O6-planes wrapped on (disjoint) smooth supersymmetric three-cycles, which we denote generically by Q. The corresponding G_2 holonomy space hence corresponds to fibering a suitable Hyperkähler four-manifold over each component of Q. That is an A-type ALE [[****AU: Please spell out ALE unless all readers will know this acronym.****]]singularity for N overlapping D6-branes, and a D-type ALE space for D6-branes on top of O6-planes [with the Atiyah-Hitchin manifold for no D6-brane, and its double covering for two D6-branes etc., as follows from (62, 63)]. Intersections of objects in type IIA therefore lift to co-dimension 7-singularities, which are isolated up to orbifold singularities. It is evident from the IIA picture that the chiral fermions are localized at these singularities.

The structure of these singularities has been studied directly in the G_2 context in (58). One starts by considering the (possibly partial) smoothing of a Hyperkähler ADE [[****AU: Please spell out ADE unless all readers will know this acronym.****]]singularity to a milder singular space, parameterized by a triplet of resolution parameters (D-terms or moment maps in the Hyperkähler construction of the space). The kind of 7-dimensional singularities of interest are obtained by considering a three-dimensional base parameterizing the resolution parameters, on which one fibers the corresponding resolved Hyperkähler space. The geometry is said to be the unfolding of the higher singularity into the lower one. This construction guarantees that the total geometry admits a G_2 holonomy metric. To determine the matter content arising from the singularity, one decomposes the adjoint representation of the A-D-E [[****AU: Is this different, A-D-E vs. ADE above?****]]group associated with the higher singularity with respect to that of the lower. One obtains chiral fermions with quantum numbers in the corresponding coset, and multiplicity given by an index which for an isolated singularity is one. This construction arises in the M-theory lift of the intersecting D6-brane models. For example, at points where two stacks of N D6-branes and M D6-branes intersect, the M-theory lift corresponds to a singularity of the G_2 holonomy space that represents the unfolding of an A_{M+N-1} singularity into a 4-manifold with an A_{M-1} and an A_{N-1} singularity. By the decomposition of the adjoint representation of A_{M+N-1}, we expect the charged matter to be in the bi-fundamental representation of the $SU(N) \times SU(M)$ gauge group, in agreement with the IIA picture. A different kind of intersection arises when N D6-branes intersect with an O6-plane, and consequently with the N D6-brane images. The M-theory lift corresponds to the unfolding of a D_N type singularity into an A_{N-1} singularity. The decomposition of the adjoint representation predicts the appearance of chiral fermions in the antisymmetric representation of $SU(N)$, in agreement with the IIA picture.

2.10 Examples

So far we have presented the main conceptual ingredients for constructing intersecting D6-brane models in a fairly general way. In order to see how this formalism works, let us work out two simple examples in more detail.

2.10.1 Intersecting D6-branes on the torus

As in section 2.2, we assume that the six-dimensional torus factorizes as $T^6 = T^2 \times T^2 \times T^2$. Introducing complex coordinates $z^i = x^i + iy^i$ on the three T^2 factors, the anti-holomorphic involution $\overline{\sigma}$ is chosen to be just complex conjugation $z^i \to \overline{z}^i$. Then as shown in Figure 6, on each T^2 there exist two different choices of the complex structure, which are consistent with the anti-holomorphic involution. Next we introduce factorizable D6-branes, which are specified by wrapping numbers (n^i, m^i) along the fundamental cycles $[a^i]$ and $[b^i]$ respectively $[a'^i]$ and $[b^i]$ on each T^2. It is useful to express also the branes for the tilted tori in terms of the untilted 1-cycles $[a^i]$ and $[b^i]$ by writing $[a'^i] = [a^i] + \frac{1}{2}[b^i]$. Then a three-cycle can be written as a product of three 1-cycles

$$\pi_a = \prod_{i=1}^{3} \left(n_a^i \; [a^i] + \widetilde{m}_a^i \; [b^i] \right). \tag{31}$$

with $\widetilde{m}_a^i = m_a^i$ for untilted tori and $\widetilde{m}_a^i = m_a^i + \frac{1}{2} n_a^i$ for tilted ones. Using the fundamental intersection number $[a^i] \circ [b^i] = -1$ with all the remaining ones vanishing, the intersection number between two three-cycles can be computed as

$$I_{ab} = \prod_{i=1}^{3}(n_a^i\,\widetilde{m}_b^i - \widetilde{m}_a^i\,n_b^i) = \prod_{i=1}^{3}(n_a^i\,m_b^i - m_a^i\,n_b^i). \tag{32}$$

To work out the tadpole cancellation conditions, one has to determine the three-cycle of the O6-plane and the action of the anti-holomorphic involution on the D6-branes.

Independent of the tilt on each T^2, the O6-plane is wrapping the cycle $2[a^i]$, so that the entire three-cycle reads $\pi_{O6} = 8\prod_i [a^i]$. The action of $\overline{\sigma}$ on a general three-cycle is simply $(n^i, \widetilde{m}^i) \to (n^i, -\widetilde{m}^i)$. Expanding the general tadpole cancellation condition for the homological R-R charges (15), one obtains the four independent equations

$$\begin{aligned}
[a^1][a^2][a^3] \quad &: \quad \sum_{a=1}^{K} N_a \prod_i n_a^i = 16 \\
[a^i][b^j][b^k] \quad &: \quad \sum_{a=1}^{K} N_a\, n_a^i\, \widetilde{m}_a^j\, \widetilde{m}_a^k = 0, \ \text{with}\ i \neq j \neq k \neq i.
\end{aligned} \tag{33}$$

These formulas were derived initially in (5) using a conformal field theory approach. As discussed in subsection 2.5, these tadpole conditions should be supplemented with some additional K-theory constraints. The K-theory constraints for the toroidal orientifold were derived in (26), and together with the tadpole conditions (33) above provide the main constraints for building semi-realistic Standard-like Models in toroidal orientifolds.

In evaluating the supersymmetry conditions we first consider the non-compact situation and a factorizable D-brane which intersects the x^i-axes on each T^2 at an angle φ_a^i. With J and Ω_3 chosen as in (30), we notice that a factorizable D-brane always satisfies $J|_{\pi_a} = 0$. Expanding the second condition $\Im(\Omega_3)|_{\pi_a} = 0$ leads to

$$0 = (dy^1\,dy^2\,dy^3 - dy^1\,dx^2\,dx^3 - dx^1\,dy^2\,dx^3 - dx^1\,dx^2\,dy^3)|_{\pi_a}. \tag{34}$$

Using $\frac{dy^i}{dx^i}|_{\pi_a} = \frac{m^i_a}{n^i_a} u^i$ with $u^i = \frac{R^i_2}{R^i_1}$, this can be brought to the form

$$\prod_{i=1}^{3} \widetilde{m}^i_a - \sum_{i\neq j\neq k\neq i} \widetilde{m}^i_a \, n^j_a \, n^k_a (u^j \, u^k)^{-1} = 0. \tag{35}$$

A further constraint arises from the condition $\Re(\Omega_3)|_{\pi_a} > 0$, which takes the form

$$\prod_{i=1}^{3} n^i_a - \sum_{i\neq j\neq k\neq i} n^i_a \, \widetilde{m}^j_a \, \widetilde{m}^k_a (u^j \, u^k) > 0. \tag{36}$$

These two conditions are equivalent to the maybe more familiar supersymmetry condition

$$\phi^1_a + \phi^2_a + \phi^3_a = 0 \bmod 2\pi. \tag{37}$$

We conclude that for a given D-brane with definite wrapping numbers, the supersymmetry condition (37) puts a constraint on the complex structure moduli $u^i = \frac{R^i_2}{R^i_1}$. If all $\phi^i_a \neq 0$ then the D6-branes preserve $\mathcal{N} = 1$ supersymmetry, and if some angles are vanishing either $\mathcal{N} = 2$ or the maximal $\mathcal{N} = 4$ supersymmetry is preserved.

~~It has been shown that u~~ **U**sing factorizable branes on T^6, no non-trivial, globally supersymmetric, tadpole-cancelling configuration of intersecting D6-branes exists. The physical reason for this is that the moment the D6-branes do not lie entirely on the x-axes, the tension of the branes in the perpendicular y-directions cannot be compensated, as there are no orientifold planes with negative tension along these directions. In the T-dual picture with magnetic fluxes, more general non-factorizable configurations of branes were investigated (64, 65).

2.10.2 Intersecting D6-branes on the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold

~~As pointed out at the end of the last section, t~~ **T**o obtain non-trivial supersymmetric models, one needs more orientifold planes extending also along y-directions.

The easiest way to obtain these is by considering not just tori but toroidal orbifolds, of which the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold is the simplest (13, 20, 21). The orbifold action of the two $\mathbb{Z}_2$ symmetries is defined as

$$\Theta : \begin{cases} z_1 \to -z_1 \\ z_2 \to -z_2 \\ z_3 \to z_3 \end{cases} \qquad \Theta' : \begin{cases} z_1 \to z_1 \\ z_2 \to -z_2 \\ z_3 \to -z_3 \end{cases} . \tag{38}$$

As it stands, this model is not completely defined~~, as~~ **;** there are two possible choices for the signs of the action of Θ' in the Θ twisted sector and vice versa. This freedom is called discrete torsion, and here we consider the model in which one keeps the $(1,1)$ forms in the twisted sectors and kills the $(2,1)$ forms. Therefore, this model has the Hodge numbers $(h_{21}, h_{11}) = (3, 51)$, which means that there are precisely eight three-cycles in the untwisted sector. These are

$$[a^1][a^2][a^3]^t,\ [a^i][a^j][b^k]^t,\ [a^i][b^j][b^k]^t,\ [b^1][b^2][b^3]^t \tag{39}$$

with $i \neq j \neq k \neq i$, and where the upper index indicates that so far these three-cycles are defined on the ambient T^6.

Given the rules from section 2.6 ~~of~~ **about** how to deal with three-cycles in the orbifold case, we have to carefully distinguish between three-cycles in the ambient T^6 and three-cycles on the orbifold space. Under the action of $\mathbb{Z}_2 \times \mathbb{Z}_2$, a three-cycle on T^6 has 3 images, which homologically are identical to the original cycles. Therefore a three-cycle in the bulk of the orbifold space can be identified with $\pi^B = 4\pi^t$. Applying the rule for the intersection number we get $\pi_a^B \circ \pi_b^B = 4\pi_a^t \circ \pi_b^t$. Therefore, the cycles π_a^B do not span the integral homology lattice $H_3(M, \mathbb{Z})$, which suggests that there exist smaller three-cycles in the orbifold space. This is indeed the case. By choosing the three-cycles to run through the origin, we obtain three-cycles which are given by $\pi_a^o = \frac{1}{2}\pi_a^B$, which have intersections on the orbifold $\pi_a^o \circ \pi_b^o = \pi_a^t \circ \pi_b^t$. Therefore, the untwisted three-cycles on the orbifold space have the same form as in (31) and the same intersection form (32), with the only difference that the basis of three-cycles is now defined on the orbifold space instead of the torus.

Working out the fixed point locus of the four orientifold projections $\Omega\overline{\sigma}(-1)^{F_L}$, $\Omega\overline{\sigma}\Theta(-1)^{F_L}$, $\Omega\overline{\sigma}\Theta'(-1)^{F_L}$, $\Omega\overline{\sigma}\Theta\Theta'(-1)^{F_L}$ and expressing everything in terms of three-cycles in the orbifold, we obtain

$$\pi_{O6} = 4\prod_i [a^i]^o - \sum_{i\neq j\neq k\neq i} 4^{1-\beta_j-\beta_k}[a^i][b^j][b^k]^o, \tag{40}$$

where $\beta_j = 0$ for an untilted T^2 factor and $\beta_j = 1/2$ for a tilted one. Therefore, the four tadpole cancellation conditions for the homological R-R charges now read [3]

$$\begin{aligned} [a^1][a^2][a^3]^o \quad &: \quad \sum_{a=1}^K N_a \prod_i n_a^i = 8, \\ [a^i][b^j][b^k]^o \quad &: \quad \sum_{a=1}^K N_a\, n_a^i\, \widetilde{m}_a^j\, \widetilde{m}_a^k = -2^{3-2\beta_j-2\beta_k}, \ \text{with } i\neq j\neq k\neq i. \end{aligned} \tag{41}$$

Note the changes on the right-hand side of (41) as compared to the purely toroidal case (33). To ensure consistency of the models, these tadpole conditions should be supplemented with additional K-theory constraints. For untilted tori, the K-theory constraints for this $\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold (49) read:

$$\begin{aligned} \sum_{a=1}^K N_a \prod_i m_a^i \quad &\in \quad 2\mathbb{Z}, \\ \sum_{a=1}^K N_a\, n_a^i\, n_a^j\, m_a^k \quad &\in \quad 2\mathbb{Z}, \ \text{with } i\neq j\neq k\neq i. \end{aligned} \tag{42}$$

[3] In (61) a different convention was used such that there appeared an overall factor of two in all four R-R tadpole cancellation in conditions (41). This is consistent with the rank of the gauge group, as the rule in (61) was that a stack of N_a branes carries a gauge group $U(N_a/2)$.

It is straightforward to generalize these conditions to cases where some or all of the tori are tilted.

Finally, for the intersection number between a D-brane and the orientifold plane one obtains

$$\pi_{O6} \circ \pi_a^o = 4 \prod_i \widetilde{m}_a^i - \sum_{i \neq j \neq k \neq i} 4^{1-\beta_j-\beta_k}\, \widetilde{m}_a^i\, n_a^j\, n_a^k. \tag{43}$$

The supersymmetry conditions are the same as for the toroidal case.

The equations developed in the last two subsections provide the main tools for constructing quasi-realistic intersecting D-brane models in these two most simple backgrounds.

3 SEMI-REALISTIC INTERSECTING D-BRANE MODELS

In this section we give an overview of the different intersecting D-brane world models explicitly constructed so far. Essentially, there are two philosophical attitudes toward approaching this problem, which differ in their assumptions about the size of the string scale, i.e., the energy scale where stringy effects become relevant. In particular, because of stability and phenomenological considerations it is usually assumed that M_s is low (e.g., $1 - 100$ TeV) for non-supersymmetric constructions, ~~while~~ **although** supersymmetric studies usually assumed that M_s is much closer to the Planck scale.

3.1 Non-supersymmetric Standard-like Models

Here we review different approaches to construct semi-realistic non-supersymmetric Standard-like Models and highlight some fairly general phenomenological features of such models. The first explicit chiral intersecting D-brane models were constructed in (5, 7, 15, 16, 17, 18), where, except in (7, 15, 16), the background space was simply chosen to be an $\Omega\overline{\sigma}(-1)^{F_L}$ orientifold of a factorizable torus T^6. These articles triggered a lot of subsequent work using essentially the same framework and ideas but generalizing the D-brane set-ups in certain ways. Before we list all these different constructions, we would like to present a prototype model, which shows that the particle content of intersecting D-brane models can come quite close to **that of** the Standard Model.

3.1.1 A simple semi-realistic model

In (18) the authors were considering the simple toroidal orientifold set-up mentioned above. Using a bottom-up approach, they introduced four stacks of D6-branes with the wrapping numbers chosen as shown in Table 2. The intersection numbers between these four stacks of D6-branes give rise to the chiral fermions listed in Table 3, which transform in the various bi-fundamental representations of the (naive) gauge group $SU(3)_C \times SU(2)_W \times U(1)_a \times U(1)_b \times U(1)_c \times U(1)_d$. The hypercharge is given by the linear combination $Q_Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c + \frac{1}{2}Q_d$. For more details and the phenomenological implications, we refer the reader to (18). Here we would like to simply list some features typical for such intersecting D-brane models:

- There are many (infinite) non-supersymmetric intersecting D-brane constructions with the Standard Model particle spectrum, where in most cases one needs additional "hidden" branes to satisfy tadpole cancellation.
- Models with only bi-fundamental matter necessarily contain right-handed neutrinos.
- One obtains additional $U(1)$ factors, which partly receive a mass via the generalized Green-Schwarz mechanism (see section 2.7); the condition that

$U(1)_Y$ remains massless imposes further conditions on the parameters in Table 2.

- All the massive former $U(1)$ gauge symmetries survive as perturbative global symmetries and can be identified with baryon number Q_a and lepton number Q_d, stabilizing the proton and preventing Majorana neutrino masses.

- One can show that in the (bc) sector there are additional non-chiral (tachyonic) fields, which might have an interpretation as Higgs particles; condensation of these fields corresponds to D-brane recombination in string theory (see for instance (66, 16, 67, 68, 69, 70)).

~~Since~~ **Because** the models are not supersymmetric, there are typically uncanceled NS-NS tadpoles, contributing to the (dilaton-dependent) cosmological constant which is of the order of M_s^4. In addition, in the effective theory below the string scale, there are large radiative corrections of the order of M_s. Therefore, typically these models require M_s of the order of the TeV scale. However, as emphasized in the subsection 2.1, for the toroidal constructions with intersecting D6-branes, the internal space cannot be much larger than the Planck volume, and the string scale M_s is restricted to be of the order of the Planck scale.[4]

3.1.2 Generalizations

Many generalizations of the above construction have been considered in the literature. Here we only list, in non-chronological order, the ones for which only

[4] This is not, however, a fundamental problem of the intersecting brane world scenario, ~~since~~ **because** the chiral spectrum of Table 3 can be achieved in models where the string scale can be lowered to a TeV (71).

non-supersymmetric models are possible or have been considered. For more details on the various constructions we refer the reader to the original literature.

The straightforward generalization of the above set-up is to introduce more than 4 four stacks [[**AU: "4 four stacks" OK?**]]of D6-branes (72, 73, 74) to realize directly the Standard Model gauge group. Similarly, one can try to find Grand Unified–like Models in this toroidal set-up (75, 76).

If one is giving up supersymmetry, then of course there is no need to introduce orientifold planes in the first place, and one can simply start with intersecting D6-branes in Type IIA (15, 16).

Another approach is not to work with D6-branes but instead with D4- respectively D5-branes [[**AU: Do you mean D4- and D5-branes, respectively?**]], where in order to achieve chirality one has to perform an additional orbifold in the transverse space (15, 16). Therefore, the models constructed in (15, 16, 77, 78, 79, 80, 71, 81, 82, 83, 84, 85) can be regarded as a hybrid of the two ways to obtain chiral fermions, ~~namely~~ as intersecting branes at singularities.

Giving up supersymmetry one can also start with orientifolds of Type O string theory (86).

A peculiarity about intersecting D-branes has been pointed out in (57, 67, 87), namely that one can build models in which at each intersection between two branes an $\mathcal{N} = 1$ supersymmetry is preserved, even though it is not preserved globally. In such models the absence of one-loop corrections to the Higgs mass weakens the gauge hierarchy problem and allows one to enhance the string scale up to 10 TeV. Such so-called quasi-supersymmetric models have also been studied in (88, 89) from a field theory perspective, and additional models have been constructed in (90).

In (19, 91) Type IIA orientifolds on the $\mathbb{Z}_3$ orbifold were considered and a non-supersymmetric three-generation flipped $SU(5)$ Grand Unified Model (92) was constructed explicitly.

3.2 Supersymmetric Models on the $\mathbb{Z}_2 \times \mathbb{Z}_2$ Orientifold

In order to show how semi-realistic supersymmetric intersecting D-brane models can arise and what their salient features are, we now describe in ~~some~~ more detail the four-dimensional chiral $\mathcal{N} = 1$ supersymmetric intersecting D-brane models constructed in (20, 21). Unlike the non-supersymmetric models discussed so far, these supersymmetric intersecting D-brane models are stable because both the NS-NS and the R-R tadpoles are canceled. As discussed in section 2.9, these models also have the additional interesting feature that when lifted to M theory, they correspond to chiral G_2 compactifications (21, 61).

The background geometry of this class of models is the $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold as described in subsection 2.10.2. As explained, there exist two choices of complex structure of T^2 that are compatible with the orientifold symmetry: rectangular or tilted (see Figure 6). If the Standard Model sector D-branes are not on top of the orientifold planes and the T^2 are rectangular, as in the toroidal models discussed in (5), the number of chiral families is even.[5] Hence, we consider models with one tilted T^2. This ~~mildly~~ **slightly** modifies the closed string sector but has an important impact on the open string sector ~~since~~ **because** the number of chiral families can now be odd. ~~Due to~~ **Owing to** the smaller number of O6-planes in tilted configurations, the R-R tadpole conditions (41) are very stringent for more

[5] The weak sector can come from D-branes on top of an orientifold plane ~~since~~ **because** $Sp(2) \simeq SU(2)$, in which case **an** odd number of families can be obtained without tilted tori (48, 49).

than one tilted T^2, so we focus on models with only one tilted T^2.

To simplify the supersymmetry conditions within our search for realistic models, we consider a particular Ansatz for the intersection angles of the branes with the x-axes: $(\phi_1, \phi_2, 0)$, $(\phi_1, 0, \phi_3)$ or $(0, \phi_2, \phi_3)$ with $\sum_i \phi_i = 0$ for each brane. Focusing on tilting just the third torus, the search for theories with $U(3)$ and $U(2)$ gauge factors carried by branes at angles and three left-handed quarks turns out to be very constraining, at least within our Ansatz. A D6-brane configuration with wrapping numbers $(n_a^i, \widetilde{m}_a^i)$ which gives rise to a three-family supersymmetric Standard-like Model is presented in Table 4.

The four D6-branes labeled C_1 are split into two parallel but not overlapping stacks of three and one branes leading to an adjoint breaking of $U(4)$ into $U(3) \times U(1)$. Consequently, a linear combination of the two $U(1)$'s is actually a generator within the non-Abelian $SU(4)$ arising for coincident branes. This ensures that this $U(1)$ is automatically non-anomalous and massless (free of linear couplings to untwisted moduli) (16, 15, 18), which turns out to be crucial for the appearance of the Standard Model hypercharge.

For convenience we consider the four D6-branes labelled A_1 to be away from the O6-planes in all three complex planes. ~~This implies~~ This leads to two D6-branes that can move independently, giving rise to a gauge group $U(1)^2$, plus their Θ, Θ' and $\Omega\overline{\sigma}(-1)^{F_L}$ images. These $U(1)$'s are also automatically non-anomalous and massless. In the effective theory, this corresponds to Higgsing of $USp(8)$ down to $U(1)^2$.

The surviving non-Abelian gauge group is $SU(3)_C \times SU(2)_W \times Sp(2) \times Sp(2) \times Sp(4)$. The $SU(3)_C \times SU(2)_W$ corresponds to the ~~MSSM, while~~ **minimal supersymmetric Standard Model (MSSM), whereas** the last three factors form a quasi-hidden sector, i.e., most states are charged under one sector or the other, but there are a few which couple to both. In addition, there are three non-anomalous $U(1)$ factors and two anomalous ones. The generators Q_3, Q_1 and Q_2 refer to the $U(1)$ factor within the corresponding $U(n)$, while Q_8, Q'_8 are the $U(1)$'s arising from the higgsed $USp(8)$. $Q_3/3$ and Q_1 are essentially baryon (B) and lepton (L) number, respectively, while $(Q_8+Q'_8)/2$ is analogous to the generator T_{3R} occurring in left-right symmetric extensions of the Standard Model. The hypercharge is defined as:

$$Q_Y = \frac{1}{6}Q_3 - \frac{1}{2}Q_1 + \frac{1}{2}(Q_8+Q'_8). \tag{44}$$

From the above comments, Q_Y is non-anomalous guaranteeing that $U(1)_Y$ remains massless. There are two additional surviving non-anomalous $U(1)$'s, i.e., $B-L = Q_3/3 - Q_1$ and $Q_8 - Q'_8$. The gauge bosons corresponding to the anomalous $U(1)$ generators $B+L$ and Q_2 acquire string-scale masses, so those generators act like perturbative global symmetries on the effective four-dimensional theory.

The spectrum of chiral multiplets in the open string sector is tabulated in Table 5. There are also vector-like multiplets in the model, but they are generically massive so we do not tabulate them here [they can be found in (93)]. The theory contains three Standard Model families, multiple Higgs candidates, a number of exotic chiral (but anomaly-free) fields, and multiplets which transform in the adjoint or singlet representation of the Standard Model gauge group.

For more details and phenomenological features, please consult the original literature (93, 94). Here, we would like to highlight some of the special features of this supersymmetric model:

- The model involves an extended gauge structure, including two additional $U(1)'$ factors, one of which has family non-universal and therefore flavor changing couplings. Extended gauge structure is quite generic among string models, and more so for intersecting D-brane models.

- There are additional Higgs doublets, suggesting such effects as a rich spectrum of Higgs particles, neutralinos, and charginos, perhaps with nonstandard couplings due to mixing and flavor-changing effects.
- In addition to the three chiral families of the Standard Model, there are chiral exotic states, i.e., chiral states with unconventional Standard Model quantum numbers. It was argued in (93) that these states may decouple from the low energy spectrum due to hidden sector charge confinement.
- There exists a quasi-hidden non-Abelian sector, which becomes strongly coupled above the electroweak scale. The dynamics of the strongly coupled hidden sector leads to dynamical supersymmetry breaking with dilaton and untwisted complex structure moduli stabilization, as studied in detail in (95). Charge confinement modifies the low energy spectrum by causing some exotics to disappear, while anomaly considerations imply that new composite states may emerge (93) (see also (96, 97)).

~~Just as in~~ **Like** the non-supersymmetric constructions of Standard-like Models, the model does not have the conventional form of gauge unification, as each gauge factor is associated with a different set of branes. However, the string-scale couplings are predicted in terms of the ratio of the Planck and string scales and a geometric factor as discussed in section 2.1. The explicit dependence of the tree-level holomorphic gauge kinetic function on the dilaton and the complex structure moduli will be discussed in section 4.2. ~~As common to~~ **Like** [[****AU: Change OK? Or how should this read?****]] all intersecting D6-brane constructions, the Yukawa couplings among chiral matter are due to world-sheet instantons associated with the string world-sheet stretching among the intersections where the corresponding chiral matter fields are localized (15). The details of the Yukawa coupling interactions will be discussed in Subsection 4.1.1.

3.2.1 Supersymmetric grand unified models

The set-up with intersecting D6-branes on orientifolds also allows for the construction of Grand Unified Models, based on the Georgi-Glashow $SU(5)$ gauge group (98). Such non-supersymmetric Grand Unified Models were constructed in (19, 92, 99, 100) and supersymmetric ones in (21, 101, 102). (For additional work on non-supersymmetric Grand Unified Models see also (76, 103).)

The supersymmetric constructions of such models have a lift on a circle to M-theory and provide examples of Grand Unified Models of strongly-coupled M-theory compactified on singular seven-dimensional manifolds with G_2 holonomy (58, 21, 59, 60, 61). (See also section 2.9.)

The key point in these constructions is the appearance of anti-symmetric representations, i.e., **10** of $SU(5)$, which can emerge at the intersection of the D-brane with its orientifold image (see Table 1). Thus, **10**-plets, along with the bi-fundamental representations $(\overline{\mathbf{5}}, \mathbf{N_b})$ at the intersections of $U(5)$ branes with $U(N_b)$ branes, form the chiral particle content of the quark and lepton families. It turns out that the gauge boson for the diagonal $U(1)$ factor of $U(5)$ is massive, and the anomalies associated with this $U(1)$ are canceled via the generalized Green-Schwarz mechanism, as explained in section 2.7.

For toroidal and $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold compactifications, there are three copies of the adjoint representations on the world-volume of the branes. They are moduli associated with the splitting and Wilson lines of the branes that wrap the same three-cycles which in these two cases are not rigid. Turning on appropriate vacuum expectation values (VEVs) of these adjoint representations can spontaneously break $SU(5)$ down to the Standard Model gauge group. ~~As mentioned s~~ **S**uch VEVs have a geometric interpretation in terms of the appropriate parallel splitting of the $U(5)$ branes.

~~Since~~ **Because** all the Standard Model gauge group factors arise from branes wrapped on parallel, but otherwise identical, cycles, this construction provides a natural framework for gauge coupling unification and thus a natural embedding of the traditional grand unification (98) into intersecting D-brane models.

The explicit supersymmetric constructions of Grand Unified Models were given for the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold models. The first such example (21) was a four-family model. Further systematic analysis (101) revealed that within $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold models with factorizable three-cycles, all the three-family models necessarily also contain three copies of **15**-plets, the symmetric representations of $SU(5)$. (Analogous observations have been reported in (104) for supersymmetric $SU(5)$ models in the $\mathbb{Z}_4 \times \mathbb{Z}_2$ orbifold background.) There are approximately twenty such models (101), which are not fully realistic:

- The additional **15**-plets decompose under $SU(3)_C \times SU(2)_Y \times U(1)_Y$ as as $(\mathbf{6},\mathbf{1})(-\frac{2}{3}) + (\mathbf{1},\mathbf{3})(+1) + (\mathbf{3},\mathbf{2})(+\frac{1}{6})$ and thus contain additional exotic Standard Model particles.
- ~~Since~~ **Because** the chiral states are charged under the $U(1)$ factor of $U(5)$ and the only candidates for the Higgs fields are in the adjoint **24** and fundamental **5** representations, the fermion masses can arise only from the Yukawa couplings of the type: $\overline{\mathbf{5}}\ \mathbf{10}\ \overline{\mathbf{5}}_H$ (subscript H refers to the Higgs fields), while the couplings of the type $\mathbf{10}\ \mathbf{10}\ \mathbf{5}_H$ are absent due to the $U(1)$

charge conservation (19). The absence of perturbative Yukawa couplings to the up-quark families is generic for these constructions (supersymmetric or not).

- Within this framework one can address the long-standing problem of doublet-triplet splitting, i.e., ensuring that after the breaking of $SU(5)$, the doublet of $\mathbf{5}_H$, responsible for the electroweak symmetry breaking, remains light, while the triplet becomes heavy. The mechanism, suggested within M-theory on G_2 holonomy manifolds (105), allows for the $SU(5)$ breaking via Wilson lines with different discrete quantum numbers for the doublet and the triplet, which in turn forbids the mass term for the doublet. However, the Wilson lines in the present context are continuous rather than discrete, ~~due to~~ **owing to** the generic non-rigid nature of three cycles on orbifold compactifications. Although the current constructions of the models are not fully realistic, generalizations to examples with rigid three-cycles may provide an avenue to address the appearance of genuinely discrete Wilson lines.

3.2.2 Systematic search for supersymmetric Standard-like Models

In the previous two subsections, we have seen that supersymmetric Standard-like and Grand Unified Models can be constructed within the intersecting D6-brane framework. Subsequently, systematic searches for supersymmetric three-family Standard-like Models have been carried out within $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold constructions with factorizable three-cycles.

~~As mentioned above, w~~ **W**ithin this framework, the systematic search for three-family $SU(5)$ Grand Unified Models (101) produced three-family models which

~~however necessarily also~~ contain three copies of **15**-plets. However, this feature is specific to this specific orbifold and it remains to be seen whether it persists for more general models.

As for the Standard-like Model constructions with gauge group factors arising from different intersecting D6-branes, sets of models with fewer Higgs doublets (106) were obtained. However, all these three-family models still possess additional exotics. Subsequently, a systematic search for supersymmetric Pati-Salam models based on the left-right symmetric gauge symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R$ was presented in (107). The gauge symmetry can be broken down to the Standard Model one via D6-brane splitting and a further D- and F-flatness preserving Higgs mechanism from massless open string states in an $\mathcal{N} = 2$ subsector. Among the models that also possess at least two confining hidden gauge sectors, where gaugino condensation can in turn trigger supersymmetry breaking and (some) moduli stabilization, the search revealed eleven models. Two models realize gauge coupling unification of $SU(2)_L$ and $SU(2)_R$ at the string scale. However, all these models still possess additional exotic matter.

In another related work (108), the study of splitting of D6-branes parallel to orientifold planes, within $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifolds, led to the examples of four-family Standard-like orientifold models without chiral exotics. The starting point is a one-family $U(4) \times Sp(2f)_L \times Sp(2f)_R$, $(f = 4)$ model, which is broken down to a four-family $U(4) \times U(2)_L \times U(2)_R$ model by parallel splitting of the D-branes, originally positioned on the O-planes. The chirality of the model is changed ~~due to the fact that~~ **because** the original branes were positioned on top of an orientifold singularity. Both the string theory and field theory aspects of these specific D-brane splittings are discussed in detail in (108).

These systematic searches for realistic models seem to suggest that an extended Higgs sector is ubiquitious in intersecting D-brane models. It is therefore quite remarkable that a simple D-brane configuration, introduced in (109), yields just the MSSM chiral spectrum and its minimal Higgs content. As the authors of (109) pointed out, in toroidal compactifications this model must be seen as a local construction, where extra R-R sources such as hidden sector branes and/or background fluxes should be added. Several attempts have been made to embed this local construction into a global model. First, it was shown in (110) that this local model can be embedded into an abstract conformal field theory construction known as Gepner orientifold (see subsection 3.3). These Gepner constructions are located at special points in the Calabi-Yau moduli space where the geometric intuition is lost, [6] and so it is therefore desirable to find an embedding into a geometrical construction. However, it proves difficult to do so for a toroidal (orbifold) background without introducing anti-branes because the cancellation of R-R tadpoles requires some R-R charges of a D-brane to have the same sign as that of an O6-plane. Peculiar as it might seem, ~~it was pointed out in an earlier work~~ **D-branes with this property do exist** (20, 21)~~ that D-branes with this property do exist~~ **.** Armed with this observation, two independent attempts, (108) and (48, 49), were made to embed the local model of (109) into a $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, and indeed a consistent global realization of (109) was found in (48, 49). Unfortunately, in the original version of (108), only the homological R-R charges are canceled, but the K-theory constraints (49) are not satisfied, resulting in the massless spectrum with discrete global anomalies (47). In the revised version of (108), employing the K-theory constraints derived in (49), a

[6]For instance, it is not straightforward how to introduce background fluxes.

consistent model was obtained with minor modifications. The Higgs sector of the model is no longer minimal, unlike the construction in (48, 49). It should be noted that the hidden sector D-branes introduced in these global models (48, 49, 108) have non-trivial intersections with the Standard Model sector D-branes and so there are chiral exotics.

3.3 Supersymmetric Models on More General Backgrounds

In ~~the~~ recent years, other supersymmetric intersecting D-brane models have been constructed with the aim ~~to find~~ **of finding** realizations of the MSSM. Essentially, two different classes of string backgrounds were considered. First, using the methods reviewed in section 2, more complicated orbifold backgrounds like a $\mathbb{Z}_4$, $\mathbb{Z}_4 \times \mathbb{Z}_2$ or $\mathbb{Z}_6$ orbifold have been studied. Second, to move beyond toroidal orbifolds and to consider intersecting branes on more general Calabi-Yau spaces, methods to treat Gepner model orientifolds were developed. Let us briefly review these activities in the following two sections.

3.3.1 Other toroidal orbifolds

One way to generalize the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifolds studied above is to include additional shift symmetries in the $\mathbb{Z}_2$ actions (111, 112, 113). These have the effect of eliminating some of the orientifold planes present in the original models, which would make it much harder to find interesting supersymmetric models. On the other hand, it also gives rise to twisted sector three-cycles, which allows for more general fractional D6-branes. Some of these models were constructed in (111, 112, 113).

Employing the topological methods introduced in section 2 (43), chiral supersymmetric intersecting D-brane models have been studied so far on the $\mathbb{Z}_4$ (50), $\mathbb{Z}_4 \times \mathbb{Z}_2$ (51, 114, 104) and $\mathbb{Z}_6$ (52) toroidal orbifolds. In the first two cases, ~~it turned out that~~ semi-realistic MSSM-like models could only be achieved after certain D-brane recombination processes were taken into account (see the original papers for more details). For the $\mathbb{Z}_6$ model (52), the authors were performing an exhaustive search for MSSM-like models and found a class of interesting D-brane configurations, which gave rise to the MSSM spectrum without the complication of brane recombinations.

Moreover, there are both four- and six-dimensional toroidal backgrounds, where so far only the non-chiral solutions to the tadpole cancellation condition, with D6-branes placed parallel to the orientifold planes, have been considered (9, 10, 11, 12, 115, 13, 116, 14).

3.3.2 Gepner Model orientifolds

One of the unattractive phenomenological features of all the toroidal orbifold models discussed above is that they give rise to too many adjoint scalars. Geometrically this means that the three-cycles π_a one is considering have too many deformations, which are counted by $b^1(\pi_a)$. ~~Not only f~~ **F**or this reason, **among others,** it is desirable to have many more backgrounds available. However, for more general algebraic Calabi-Yau spaces, not very much is known about sLag three-cycles, which prevents a direct geometric approach to the problem as pursued for instance in (43, 117, 118, 44).

One ~~way out~~ **solution** is to use Gepner models, which are exactly solvable conformal field theories known to describe certain symmetric points in the moduli space of distinguished Calabi-Yau manifolds. ~~Since~~ **T**he description of D-branes and orientifold planes in this context is a subject of its own, **but** here we ~~would like to only mention~~ **mention only** that after some first attempts (119, 120), ~~during the last few years~~ methods have **recently** been developed to treat Gepner model orientifolds very efficiently (121, 122, 123, 124, 110, 125, 126), allowing a systematic computer search for MSSM like models. Specifically, the impressive results of (127) provide large classes of three-family Standard-like Models with no chiral exotics. It remains to be seen whether some of these models also satisfy the more refined Standard Model constraints. Note, however, that these exact conformal field theory models are located at very special points in the Kähler and complex structure moduli space, where the geometric intuition is lost. ~~Since~~ **Because** one expects that all radii are of string scale size, couplings, such as Yukawa couplings, are not expected to possess hierarchies associated with the size of the internal spaces, ~~such~~ as in the case of the toroidal orbifolds with D-branes. In addition, the introduction of supergravity fluxes is not straightforward to perform (see section 5).

4 LOW ENERGY EFFECTIVE ACTIONS

The models presented in the previous section 3 provide a starting point for the study of couplings in the effective low energy theory whose massless spectrum was determined by techniques presented in section 2. For the orientifold models with intersecting D6-branes compactified on orbifolds, the calculation of such couplings can be done by employing the conformal field theory techniques on orbifolds (128). The tree-level calculations can in principle be performed both for the supersymmetric and the non-supersymmetric constructions; for the one-loop calculations, supersymmetry is a necessary ingredient to obtain an unambiguous finite answer. The summary of the explicit results will therefore focus primarily on the supersymmetric constructions (Alternatively, part of the low energy effective action can also be determined by a dimensional reduction of the ten-dimensional supergravity theory (129, 130, 131)).

The calculation of couplings for chiral superfields at the intersection of D6-branes and, in particular, the Yukawa couplings for such states are clearly of phenomenological interest. As discussed in section 2.2, the states at such intersections correspond to the open string excitations stretched between the two intersecting D6-branes. As a consequence, the bosonic string oscillator modes are like given in eq.(8). Therefore, physical string excitations at the two D6-brane intersections are associated with the twisted open-string sectors, which are analogous to the closed string twisted sectors on orbifolds (128).

The string amplitudes for these excitations can in turn be calculated employing conformal field theory techniques (128). The correlation functions of the fermionic string excitations can be obtained in a straightforward manner by employing a world-sheet bosonization procedure. On the other hand, the correlation functions for the bosonic excitations involve the calculation of the correlation functions for the so-called bosonic twist fields, i.e., $\sigma_{\epsilon_i}(x)$, evaluated at the world-sheet location x on the disc. The bosonic twist field ensures that the bosonic open string fields $X^i(z)$ (in the i-th toroidal direction) have the correct twisted boundary conditions. Here z is the world-sheet coordinate. These boundary conditions are encoded in the following operator product expansion (128, 132):

$$
\begin{aligned}
\partial X^i(z)\sigma_{\epsilon_i}(x) &\sim (z-x)^{\epsilon_i - 1}\tau_{\epsilon_i}(x) + \dots \\
\partial \tilde{X}^i(z)\sigma_{\epsilon_i}(x) &\sim (z-x)^{-\epsilon_i}\tau'_{\epsilon_i}(x) + \dots \\
\bar{\partial} X^i(\bar{z})\sigma_{\epsilon_i}(x) &\sim -(\bar{z}-x)^{-\epsilon_i}\tau'_{\epsilon_i}(x) + \dots
\end{aligned}
$$

$$\bar{\partial}\tilde{X}^i(\bar{z})\sigma_{\epsilon_i}(x) \quad \sim \quad -(\bar{z}-x)^{\epsilon_i-1}\tau_{\epsilon_i}(x)+\ldots \tag{45}$$

and similarly for $\sigma_{-\epsilon_i}(x)$. Here τ_{ϵ_i} and τ'_{ϵ_i} correspond to the excited bosonic twist fields. Employing the so-called stress-energy method (128), which allows one to determine the correlation functions of bosonic twist fields by employing the properties of the operator product expansion of the conformal field theory stress-energy tensor with the twist fields, along with the above operator product expansions (45), enables one (132) to determine the bosonic twisted sector string amplitudes. The application of these calculations to the four-point couplings and Yukawa couplings will be discussed in subsection 4.1.1.

Another set of couplings involves the calculation of the Kähler potential for the states at the intersection, in particular the explicit dependence of the leading term which is bi-linear in powers of chiral superfields at the intersection. The corresponding string amplitudes involve the correlation functions containing both states at the intersection (open string states) and toroidal moduli fields (closed string states) (133). Such couplings will be discussed in subsection 4.1.2.

Another important topic is the calculation of the gauge couplings. In particular, determination of the holomorphic gauge kinetic function in terms of the dilaton and the toroidal moduli both at the tree-level and one-loop–level is an important task and will be discussed in section 4.2.

4.1 Correlation Functions for States at D-brane Intersections

In this section we summarize the results for the tree-level calculations for the chiral matter appearing at the D6-brane intersections (see section 2.2). In the supersymmetric constructions, the states at intersections correspond to the full massless chiral supermultiplet. The couplings of most interest are the tri-linear superpotential couplings, such as the coupling of quarks and leptons to the Higgs fields. On the other hand, the four-point couplings are also of interest, ~~since~~ **because** they indicate the appearance of higher order terms in the effective Lagrangian; for example, certain four-fermion couplings could contribute to the flavor-changing neutral currents in the Standard-like Model constructions (see (134, 135)) and in the Grand Unified Models, triggering proton decay (see (136)).

4.1.1 The four-point and three-point functions: Yukawa couplings

The explicit calculations of the three-level four-point and three-point correlation functions for the states appearing at the D-brane intersections were done in (137, 132, 134, 136). Generalizations to n-point functions were addressed in (138).

As discussed in the introduction of this section, the non-trivial part in the calculation involves the evaluation of the correlation functions of four (three) bosonic twist fields, which signify the fact that the states at the intersection arise from the sector with twisted boundary conditions on the bosonic and fermionic string fields. The conformal field theory techniques employed are related to the study of bosonic twist fields of the closed string theory on orbifolds (128). For technical details of the specific calculation of the four- and three-point functions, employing conformal field theory techniques, we refer the reader to references (132, 133), and for a detailed calculation of the classical part of Yukawa couplings to reference (137).

The calculations have been done in the case of intersecting D6-branes wrapping factorizable three-cycles of a six-torus $T^6 = T^2 \times T^2 \times T^2$. Thus, in each T^2, the D6-branes wrap one-cycles, and the problem reduces to a calculation of correlation functions of bosonic twist fields associated with the twisted sectors at intersections of D6-branes wrapping the one-cycles of a T^2. The final answer is therefore a product of contributions from correlation functions on each of the three T^2 (137).

In particular, the following four-point correlation functions of bosonic twist fields are of interest:

$$\langle \sigma_\nu(x_1)\sigma_{-\nu}(x_2)\sigma_\nu(x_3)\sigma_{-\nu}(x_4)\rangle \tag{46}$$

and

$$\langle \sigma_\nu(x_1)\sigma_{-\nu}(x_2)\sigma_{-\lambda}(x_3)\sigma_\lambda(x_4)\rangle. \tag{47}$$

The first one corresponds to the bosonic twist field correlation function of states appearing at the intersection of two pairwise parallel branes with intersection angle $\pi\,\nu$ (see Figure 7). This correlation function is a key ingredient in the calculation of the four-fermion couplings, **an ingredient** that contributes to the flavor-changing neutral currents in the Standard-like models (135) and to the proton decay amplitudes in the Grand Unified Models (136).

The second amplitude (47) corresponds to the bosonic twist field correlation function of states appearing at the intersection of two branes intersecting at respective angles $\pi\nu$ and $\pi\lambda$ with the third set of parallel branes (See Figure 8). This correlation function is specifically suited for taking the limit of the world-sheet coordinate $x_2 \to x_3$ which factorizes to a three point function associated with the intersection of three branes. This latter result is particularly interesting ~~since~~ **because** it provides a key element in the calculation of the Yukawa coupling.

By employing the stress-energy conformal field theory techniques and the properties of the operator product expansions of the bosonic twist fields (45) one can determine (132) both the classical part and the quantum part of such amplitudes and thus the exact tree-level answer for the corresponding couplings. The calculation of the quantum part depends only on the intersection angles and is thus insensitive to the scales of the internal space and relative position of the branes. On the other hand the classical part carries information on the actual separation among the branes and the overall volume of T^2 as well. An important part in the calculation is the determination of an overall normalization of the four-point amplitude, which can be done by factorizing the amplitude in the limits $x \to 0$ or $x \to 1$, where it reduces to a product of the two three-point amplitudes. Namely, in these limits, the four-point amplitude contains a dominant contribution from the exchanges of the intermediate open string winding states around the compact directions. The dominant contribution can be interpreted as the s (or t)-channel exchange of the massless gauge bosons living on the world-volume of the D6-branes. Thus, in this limit the amplitude is completely determined in terms of the gauge-coupling which in turn determines the normalization of the full four-point amplitude. For details see (132, 136).

We shall skip the technical details and in the following only quote the result for the exact (string) tree-level expression for the Yukawa coupling for two massless fermionic states and one massless bosonic state, appearing at the intersections of three D6-branes (132):

$$Y = \sqrt{2} g_0 2\pi \prod_{j=1}^{3} \left[\frac{16\pi^2 \Gamma(1-\nu_j)\Gamma(1-\lambda_j)\Gamma(\nu_j+\lambda_j)}{\Gamma(\nu_j)\Gamma(\lambda_j)\Gamma(1-\nu_j-\lambda_j)} \right]^{\frac{1}{4}} \sum_{m \in \{\mathrm{ws-inst.}\}} \exp\left(-\frac{A_j(m)}{2\pi\alpha'} \right) , \tag{48}$$

where $A_j(m)$ is the area of the m-th triangle (world-sheet instanton) formed by the three intersecting branes on the j-th two-torus and $g_0 = e^{\Phi/2}$, with Φ corresponding to the Type IIA dilaton. In order to derive (48) it was assumed that all three fields have canonically normalized kinetic energies. For a discussion of phenomenological implications the above results, see subsection 6.4.

In earlier works (94, 109) the leading classical contribution to such Yukawa couplings was calculated. Also previously, a comprehensive analysis and computation of the full classical contribution (which contains all the open string moduli dependence) was performed in (137). Intriguingly, in (139) the three-point function was calculated in the mirror-dual Type IIB theory. Here a purely classical, leading order in α', computation gave already the full world-sheet instanton corrected Type IIA superpotential contribution to the Yukawa couplings [see eq.(48)]. This can be considered as a nice confirmation of mirror symmetry. As expected, the Kähler potential contributions to the Yukawa couplings only agreed to leading order in α'.

The prefactor in (48) corresponds to the quantum part of the bosonic twist correlator; it has a suggestive factorizable form associated with the angles of states appearing at each intersection, resulting in the following complex structure dependence of the Kähler potential for chiral superfields at D6-brane intersections:

$$K = \frac{1}{4\pi} \sum_{\nu} \prod_{j=1}^{3} \sqrt{\frac{\Gamma(\nu_j)}{\Gamma(1-\nu_j)}}\, \Phi_\nu\, \Phi^*_\nu, \, . \tag{49}$$

(For simplicity in the above expression the Planck scale was set to $M_P = 1$.) On the other hand, the classical part of the Yukawa coupling is proportional to the area of the intersection triangles and thus depends on the toroidal Kähler structure moduli. It includes a contribution from the part of the Kähler potential that depends on toroidal Kähler moduli (139). Thus the full Kähler potential takes the form displayed in the next subsection [see Equation (50)]. After the inclusion of toroidal two-form field potentials and the asscoiated Wilson lines (137), the remaining part of the Yukawa coupling describes the superpotential trilinear coupling as a holomorphic function of toroidal Kähler moduli; this coupling typically takes a form of modular theta functions [for further details, see (137, 139)]. This splitting of the three point function (48) into the leading Kähler potential contribution and the superpotential contribution has been confirmed as a part of the calculations described in the following subsection.

Higher tree-level n-point correlation functions for chiral superfields at D-brane intersections have been studied in (134, 138). The one-loop calculation of the three point functions for such states was done in (140) and it leads to new results for the one-loop corrections to the Kähler potential for the corresponding chiral superfields at D-brane intersections.

4.1.2 The closed-open string amplitudes - Kähler potential

A direct calculation of the tree-level leading order Kähler potential for chiral superfields at D-brane intersections and their dependence on the closed string sector moduli involves the determination of the string amplitudes for the two open string sector vertices and an arbitrary number of closed sector moduli vertices. For the toroidal (orbifold) backgrounds, these calculations have been carried out explicitly for any number of toroidal complex and Kähler structure moduli in (133). First explicit results for the four-point string amplitudes were derived and then by further employing the symmetry structure of higher n-point functions, sets of differential equations were obtained for the n-point functions with an arbitrary number of vertices for the toroidal moduli. These could be explicitly solved, thus resulting in the explicit string amplitudes with any number of toroidal closed sector moduli. As a consequence, the leading order tree-level Kähler potential for chiral superfields at D-brane intersections, and its explicit dependence on both the toroidal Kähler and complex structure moduli, could be derived. Specifically, the Kähler potential for the open string sector chiral superfields $\Phi_{\nu_{ab}}$, appearing at the intersection of the stack a and stack b of D6-branes, takes the form (133):

$$K = \frac{1}{4\pi}\left[\prod_{i=1}^{3}(T_i + T_i^*)^{-\nu_{ab}^i}\sqrt{\frac{\Gamma(\nu_{ab}^i)}{\Gamma(1-\nu_{ab}^i)}}\right]\Phi_{\nu_{ab}}\Phi_{\nu_{ab}}^*\,, \tag{50}$$

where again $\pi\,\nu_{ab}^i$ denotes the angle of the a- and b- D6-brane intersection in the i-th two-torus and T_i is the Kähler modulus of the i-th two-torus. (For simplicity, again, the Planck scale is set to $M_{pl} = 1$.) Note that the dependence of the above Kähler potential on the angles and thus implicitly on the toroidal complex structure moduli is the same as the one obtained from the Yukawa coupling calculation (48). In addition, (50) also contains ~~the~~ information on the toroidal Kähler moduli.

4.2 Gauge Couplings

The last function that specifies the effective four-dimensional $\mathcal{N} = 1$ supersymmetric theory is the gauge kinetic function. The tree-level gauge kinetic function for each stack of D6-branes can be determined in a straightforward way by reducing the D6-brane world-volume kinetic energy action along the three-cycle wrapped by the stack of D6-branes in the internal space. For a supersymmetric three-cycle, π_a, the tree-level gauge kinetic function is a holomorphic function of complex structure moduli fields and it is of the form (57, 43, 141)

$$f_a = \frac{M_s^3}{(2\pi)^4}\left[e^{-\varphi}\int_{\pi_a}\Re(\Omega_3) + 2i\int_{\pi_a}C_3\right], \tag{51}$$

where C_3 denotes the R-R three-form. For supersymmetric three-cycles on toroidal orbifolds, this holomorphic gauge kinetic function takes an explicit form in terms of the toroidal complex structure moduli U^i and the dilaton field S

$$\begin{aligned} S &= \frac{M_s}{2\pi}e^{-\varphi}\prod_i R_1^i + \frac{i}{4\pi}C^0\,, \\ U^i &= \frac{M_s}{2\pi}e^{-\varphi}\,R_1^i\,R_2^j\,R_2^k + \frac{i}{4\pi}C^i\,, \end{aligned} \tag{52}$$

with $i \neq j \neq k \neq i$. For example in the case of the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, the gauge coupling function takes the form [see, e.g., (57, 95))]:

$$f_a(U^i, S) = \frac{1}{4}\left[n_a^1\, n_a^2\, n_a^3\, S \,-\, n_a^1\, \widetilde{m}_a^2\, \widetilde{m}_a^3\, U^1 \,-\, \widetilde{m}_a^1\, n_a^2\, \widetilde{m}_a^3\, U^2 \,-\, \widetilde{m}_a^1\, \widetilde{m}_a^2\, n_a^3\, U^3\right] , \tag{53}$$

where as usual n_a^i and $\widetilde{m}_a^i$ are the wrapping numbers of the three-cycle π_a and the pre-factor $\frac{1}{4}$ is the dimension of the orbifold/orientifold group.

~~We emphasize that since~~ **Because** the gauge coupling for each gauge group factor depends on the volume of the corresponding three-cycle π_a, in general the intersecting D-brane constructions do not have gauge coupling unification in the sense of Grand Unified Models. However, for each gauge group factor the tree-level gauge coupling is calculable in terms of the toroidal complex structure moduli, the dilaton, and the wrapping numbers of the three-cycle π_a. The results for the tree-level gauge kinetic functions were employed as the starting point to address the renormalization group running of gauge couplings from the string scale to the electroweak regime for the semi-realistic constructions (93, 141, 142) as well as in the study of the moduli stabilization and supersymmetry breaking due to gaugino condensation in the hidden sector of supersymmetric semi-realistic constructions (95) (for the respective phenomenological implications see subsections 6.2 and 6.6.1).

The tree-level couplings receive corrections at the one-loop level due to the so-called threshold corrections of the heavy string modes. These explicit calculations are involved ~~since~~ **because** the complete massive string spectrum is needed. For the perturbative heterotic string theory on orbifolds, these threshold corrections were first calculated in (143). The calculation of the gauge coupling threshold corrections for the intersecting D6-branes on a toroidal orbifold backgrounds amounts to similar complexity and the explicit results have been computed in (144). While the one-loop corrections for open string sectors preserving $\mathcal{N}=4$ supersymmetry vanish, the $\mathcal{N}=2$ sectors depend on both the toroidal complex and Kähler moduli. For explicit expressions, please consult (144). These corrections bear similarities with the heterotic orbifold corrections (143). For details and specific explicit calculation for the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, see again (144). In addition, there are $\mathcal{N}=1$ sector corrections; they can be cast in a compact expression which for the $SU(N_a)$ gauge couplings takes the form:

$$\Delta_{ab} = -b_{ab} \ln \frac{\Gamma(1-\nu^1_{ba})\Gamma(1-\nu^2_{ba})\Gamma(1+\nu^1_{ba}+\nu^2_{ba})}{\Gamma(1+\nu^1_{ba})\Gamma(1+\nu^2_{ba})\Gamma(1-\nu^1_{ba}-\nu^2_{ba})}, \tag{54}$$

where $b_{ab} = N_b I_{ab} \mathrm{Tr}(Q_a^2)$ and $\pi\,\nu^i_{ba}$ denote the intersection angles of a and b D6-branes in the i-th two-torus. These angles can be expressed in terms of the wrapping numbers of the π_a and π_b three-cycles and the toroidal complex structure moduli U^i. [For the explicit formula, see, e.g., (132).] The total correction from the $\mathcal{N}=1$ sector is obtained as a summation over all bs, i.e., stacks of all the other D6-branes wrapping the three-cycles π_b.

These threshold corrections could play an important role in the study of the renormalization group running of gauge couplings; in particular, they can modify the effective string scale. In addition, ~~since~~ **because** the threshold corrections depend on both the Kähler and complex structure moduli, they could play an important role in the strong infrared "hidden sector" dynamics and the possibility of stabilizing all toroidal moduli.

5 FLUX VACUA WITH MAGNETIZED D-BRANES

As we have seen in the previous sections, intersecting D-brane worlds provide a simple geometric framework within which many semi-realistic particle physics models can be constructed. However, just like other supersymmetric string constructions, these intersecting D-brane models suffer from the usual moduli problem. Typically, these models contain a lot of moduli (from both closed and open string sectors) which remain massless before supersymmetry is broken and hence, if not stabilized, would lead to serious phenomenological problems as well as loss of predictivity. Deeply related to the moduli problem is the question of how supersymmetry is broken. In studying the phenomenological consequences of string theory, one traditionally starts with an $\mathcal{N} = 1$ supersymmetric string vacuum whose low energy spectrum contains the Standard Model. The hope is that the same mechanism that breaks supersymmetry and gives masses to the superpartners of the Standard Model could also lift all the moduli. Strong D-brane gauge dynamics, for example, can result in gaugino and matter condensations, generating a non-perturbative Veneziano-Yankielowicz-type potential (145, 146) that can in principle stabilize certain closed string moduli. However, such non-perturbative dynamics can be analyzed only at the level of an effective super Yang-Mills theory with the leading instanton contribution. Moreover, without fine-tuning, the vacua stabilized by non-perturbative effects typically have a large non-vanishing cosmological constant, thus rendering these models unrealistic for further phenomenological studies (see section 6.6). Therefore, in practice, one often treats the supersymmetry breaking sector as a black box and simply parameterizes our ignorance of supersymmetry breaking with the VEVs of the auxiliary fields of some moduli without specifying how they acquire a VEV [see, e.g., (147)].

Recently there have been some interesting attempts to understand simultaneously these two central problems in string phenomenology – moduli stabilization and supersymmetry breaking – by considering compactification with background flux (148, 149, 150, 151, 152, 153). The idea, which is most conveniently expressed in the framework of Type IIB string theory, is that a superpotential can be generated by the NS-NS and R-R three-form flux background (148, 150). The superpotential thus generated depends on the dilaton and the complex structure moduli, and so these moduli are generically lifted. Interestingly, depending on how the gauge and chiral sectors are embedded, supersymmetry can be softly broken by the flux. More importantly, the resulting soft SUSY breaking terms can be calculated in a systematic way, perturbatively (154, 155, 156, 157, 158). We shall see toward the end of this section how such soft terms are generated from the local D-brane physics point of view. From the effective four-dimensional supergravity ~~point of view~~ **perspective**, the effect of the background flux is to introduce a microscopic source for the auxiliary fields of the moduli which [[****AU: What, exactly, breaks supersymmetry? Possible to recast for clarity?****]] as a result breaks supersymmetry.

Thus, flux compactification provides a rather attractive framework for string phenomenology. However, in order to explore quantitatively its phenomenological features, it is important to construct some concrete examples in which realistic features of the Standard Model, such as chirality, can be incorporated. The general techniques of constructing chiral flux vacua have been developed in (159, 160, 161), although no chiral models which are free of tadpole instability have been found. More recently, it was ~~realized in~~ **recognized by** (48, 49) that a crucial ~~ingredient~~ **step** in constructing stable chiral flux vacua is to introduce additional pairs of $D9 - \overline{D9}$-branes, which nonetheless are BPS because of the magnetic flux on their world-volumes. Chiral flux vacua (both supersymmetric and non-supersymmetric) that are free of tadpoles have been constructed in (48, 49). Furthermore, the low energy spectrum of these models is remarkably close to the MSSM, and hence they provide a proof of concept that realistic particle physics features can be embedded in flux compactification[7]. Interestingly, in cases where supersymmetry is broken softly by the background flux, the vacuum remains Minkowski [[**AU: OK?**]]after supersymmetry breaking (at least to lowest order) ~~since~~ **because** the NS-NS tadpoles are absent. Subsequently, there have been further interesting attempts ~~in constructing~~ **to construct** realistic models within this framework, and more examples of three- and four-family chiral flux vacua, including supersymmetic ones, have been found in (163, 164).

Although this review has so far ~~been~~ focused on Type IIA string theory with intersecting D6-branes, the intersecting D-brane models discussed here are related (in the absence of flux) by a simple duality to Type IIB orientifolds with magnetized D9-branes [see, e.g., (160) for the details of such map]. Hence, there is an alternative, albeit less geometrical, description of the same models in Type IIB string theory. In fact, the techniques of building intersecting D-brane models that we have discussed can be readily adapted to construct chiral flux vacua, which we will review below.

5.1 Three-form Fluxes in Type IIB String Theory

Various aspects of flux compactifications have been discussed in the literature. Instead of providing a comprehensive overview of flux compactifications (which is not the main purpose of this review), we will only sketch here some of the

[7]Local models of flux compactification with realistic particle physics features have been considered in (162).

basic results (149, 153, 165, 166) relevant to string model building. Consider Type IIB string theory in the presence of a non-trivial three form background flux $G_3 = F_3 - \tau H_3$. Here F_3 denotes the R-R and H_3 the NS-NS three form flux, and τ is the complex dilaton-axion field. The background fluxes must obey the Bianchi identity and be properly quantized, i.e., they take values in $H^3(\mathcal{M}, \mathbb{Z})$. In toroidal (orbifold) backgrounds, such quantization conditions are particularly simple:

$$\frac{1}{(2\pi)^2\alpha'}\int_\Sigma H_3 \in N_{min} \times \mathbb{Z}, \qquad \frac{1}{(2\pi)^2\alpha'}\int_\Sigma F_3 \in N_{min} \times \mathbb{Z}, \tag{55}$$

where Σ is a three-cycle in the *covering space*, N_{min} is a positive integer which reflects the fact that in an orientifold (orbifold), there can exist three-cycles which are smaller that in the covering space. For example, for the $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, taking into account also the orientifold projection, $N_{min} = 8$, and N_{flux}, defined in (57), is a multiple of 64.

Turning on F_3 and H_3 fluxes has two important effects. First, the Chern-Simons terms in the Type IIB effective supergravity action

$$S_{CS} = \frac{1}{2\kappa_{10}^2}\int d^{10}x \frac{C_4 \wedge G_3 \wedge \overline{G}_3}{4i\,\mathrm{Im}\tau}\,, \tag{56}$$

when integrated over the six-dimensional manifold $\mathcal{M}$ induces a tadpole for the R-R four-form gauge potential C_4. In particular, the D3 charge contribution to the tadpoles is of the form

$$N_{flux} = \frac{1}{(4\pi^2\alpha')^2}\int_\mathcal{M} H_3 \wedge F_3\,. \tag{57}$$

The second effect is that the kinetic term for G_3 (suppressing the warp factor)

$$V = \frac{1}{4\kappa_{10}^2 \mathrm{Im}\tau}\int_\mathcal{M} d^6y\, G_3 \wedge \star_6 \overline{G}_3\,, \tag{58}$$

induces a scalar potential, which can be written as

$$V = \frac{1}{2\kappa_{10}^2 \mathrm{Im}\tau} \int_{\mathcal{M}} d^6y\, G_3^- \wedge \star_6 \overline{G}_3^- - \frac{i}{4\kappa_{10}^2 \mathrm{Im}\tau} \int_{\mathcal{M}} d^6y\, G_3 \wedge \overline{G}_3\,, \tag{59}$$

where again we suppress the warp factor. Here, $G_3^{\pm}$ is the imaginary self-dual/anti-self-dual (ISD/IASD) part of G_3, i.e., it satisfies $\star_6 G_3^{\pm} = \pm i G_3^{\pm}$. The second term in (59) is a topological term equal in magnitude to N_{flux} and gives rise to the NS-NS tadpole of the flux. Contrarily, the first term in (59) is a positive semi-definite F-term potential (148, 150), which precisely vanishes if the flux is imaginary self dual, i.e., $G_3^- = 0$.

It has been shown in (150) that this F-term potential V_F can be derived from the Gukov-Vafa-Witten superpotential (148)

$$W = \int_{\mathcal{M}} \Omega_3 \wedge G_3\,, \tag{60}$$

which apparently depends only on the complex structure moduli and the dilaton and vanishes if these moduli are chosen such that G_3 is imaginary self-dual. Self-duality implies that the three-form flux has a $(2,1)$ and a $(0,3)$ component with respect to the complex structure of the underlying Calabi-Yau, i.e., $G_3 = G_3^{(2,1)} + G_3^{(0,3)}$. For supersymmetric minima, one gets the additional conditions that (i) $G_3^{(0,3)} = 0$, and (ii) primitivity of G_3, *i.e.*, $G_3 \wedge J = 0$; the latter condition is automatically satisfied on a Calabi-Yau manifold. Taking the back-reaction of the fluxes on the geometry into account, one finds it is quite moderate in the (topological) sense that one gets a warped Calabi-Yau metric [although the metric can be strongly warped as in (167, 165)]. To summarize, the flux-induced scalar potential allows one to freeze the complex structure moduli and the dilaton at its minima.

Although the discussion here is in the context of Type IIB string theory, there should be an alternative description in Type IIA string theory where intersecting D6-branes (the subject of this review) can be introduced. However, under duality, the three-form fluxes that we consider here become metric fluxes on the Type IIA side and the underlying geometry becomes non-Kähler (166, 168, 169). The types of three-cycles that the D6-branes can wrap around in such non-Kähler geometries are not well understood. Alternatively, one can study directly Type IIA orientifolds with background fluxes [see (170, 171, 172) and references therein]. For recent efforts in obtaining examples of flux compactifications in massive Type IIA supergravity with intersecting D6-branes, see (173, 174, 175).

5.2 Semi-realistic Flux Vacua

In the following, we summarize the techniques for constructing consistent Type IIB orientifolds with magnetized D9-branes in toroidal (orbifold) backgrounds developed in (159, 160).

A stack of N_a magnetized D9-branes on toroidal orbifolds is characterized by three pairs of integers (n_a^i, m_a^i) which satisfy

$$\frac{m_a^i}{2\pi}\int_{T_i^2} F_a^i = n_a^i, \tag{61}$$

where the m_a^i denote the wrapping numbers of the D9-brane around the i-th two-torus T_i^2, n_a^i is the magnetic flux, and the F_a^i is the corresponding $U(1)$ magnetic field-strength on the D9-brane. The orientifold projection on these quantum numbers acts as: $\Omega\overline{\sigma}(-1)^{F_L} : (n_a^i, m_a^i) \to (n_a^i, -m_a^i)$. In the T-dual intersecting D6-branes picture, these quantum numbers correspond to the wrapping numbers (n_a^i, m_a^i) of the homology one-cycles $([a_i], [b_i])$ of the i-th two-torus T_i^2 that the D6-branes wrap around. [For a detailed dictionary between these two T-dual descriptions, see, e.g., (160). More general aspects of the T-dual picture were discussed in (176).]

Because of the orientifold projection, the magnetized D9-branes set-up as a whole does not carry any net D5- and D9-brane charges. However, there are additional discrete K-theory charges that needed to be taken into account (49). Other than this subtlety, the tadpole cancellation conditions for the magnetized D9-brane sector simply amount to the cancellation of the D3- and three types of D7-brane charges. Such conditions can be deduced from the conditions (41) in the T-dual intersecting D6-brane picture in section 2. Note, however, that the background fluxes introduce an additional contribution to the D3 tadpole (57). The corresponding tadpole conditions, here specifically written for the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold, read (159, 160):

$$\begin{aligned} \mathrm{D3-charge} \quad & \sum_a N_a\, n_a^1\, n_a^2\, n_a^3 &=& \; 8 - \tfrac{N_{flux}}{4}, \\ \mathrm{D7_i-charge} \quad & \sum_a N_a\, n_a^i\, m_a^j\, m_a^k &=& \; -8 \quad \text{for } i \neq j \neq k \neq i. \end{aligned} \tag{62}$$

[For consistency conditions as applied to other orbifolds, see (177).] The large positive D3-charge contribution from N_{flux} makes it hard to satisfy these tadpole conditions without introducing anti-D3 branes which give rise to instabilities (159, 160).

Fortunately, it was ~~realized in~~ **recognized by** (48, 49) that the negative D3-brane charge needed for the cancellation of R-R tadpoles can be accounted for by introducing additional pairs of $D9-\overline{D9}$-branes which are nonetheless BPS because of the magnetic flux supported on their world-volumes. This observation led to the first example (48, 49) of a three-family supersymmetric Standard Model in flux compactification. In this construction, the Standard Model sector is based on the local MSSM module introduced in (109). Therefore, as in (109), there is only a pair of Higgs doublets in the low energy spectrum and thus precisely the minimal Higgs content of the MSSM. However, the $D9 - \overline{D9}$ pairs which carry the needed negative D3-charge have a non-vanishing "intersection product" with the Standard Model building block. Hence in addition to the gauge and matter content of the MSSM, there are also some additional chiral exotics.

Subsequently, additional semi-realistic flux vacua have been constructed in (163, 164). In particular, by considering such BPS $D9 - \overline{D9}$ pairs as part of the observable sector, a broader class of Standard Model-like vacua with three and four families of chiral matter, and larger units of flux, including supersymmetric ones, were constructed in (164). These models have typically chiral exotics and more than one-pair of Higgs doublets. It is fair to say that a fully realistic model of flux compactification has yet to be found.

Let us now briefly comment on two related issues pertinent to these chiral flux vacua: generation of soft supersymmetry breaking terms and stabilization of open string moduli. First, we can understand heuristically how soft supersymmetry breaking terms are generated by the background flux. As discussed above, the background three-form flux carries R-R charge and tension. [[****AU: Your use of "respectively" in the following: Is this standard in the field? I am unfamiliar with this construction. Possible to clarify? Perhaps just omit the word?****]] More precisely, ISD (respectively IASD) flux carries the same type of R-R charge and tension as that of a D3-brane (respectively $\overline{D3}$-brane). Therefore, a D3-brane (respectively $\overline{D3}$-brane) will be attracted to a region where the IASD (respectively ISD) flux is maximum. Recall that the positions of D3-branes correspond to world-volume scalars, and so the energy needed to move the D3-branes away from the maximum flux region would reflect as soft masses on the world-volume gauge theory. The analysis for the D7-brane sector is more involved, but one can again understand how soft terms are generated by studying the induced D3-brane charge (due to the background flux) carried by the D7-branes.

For the same reason that soft terms are generated, the background flux can also induce a mass to some of the open string moduli and thus provide a way to stabilize them (178, 179, 158). Finally, although the Kähler moduli do not enter the flux-induced superpotential, a linear combination of some toroidal Kähler moduli and open string moduli enters the Fayet-Iliopoulos D-term, and so we expect such linear combination of closed and open string moduli to be fixed.

Much work needs to be done before a fully realistic model of flux compactification (i.e., a model not only with a realistic low energy spectrum and couplings, but also with all its moduli stabilized) can be found. However, the developments described here have undoubtedly pointed to an interesting direction in string phenomenology.

6 PHENOMENOLOGICAL ISSUES

No fully realistic intersecting D-brane model has been constructed yet. Furthermore, many of the phenomenological features, such as the gauge and Yukawa couplings, or the masses and other properties of the low energy particles, are model dependent or depend on details of supersymmetry breaking. Nevertheless, it is useful to survey here some of the phenomenological features that have emerged in various constructions, with an emphasis on the difficulties, possibilities for new physics, and things to watch for in the future. Many of the technical aspects or detailed consequences of specific models were discussed in earlier sections. Here we focus on general issues.

Let us start with two general comments. The first concerns the fundamental string scale M_s. As discussed in subsection 2.1, most toroidal (orbifold) constructions have assumed either that M_s and the inverse sizes of the extra dimensions are very large, i.e., within a few orders of magnitude of the Planck scale, or else that M_s is much lower, e.g., in the 1- 1000 TeV range. The latter can occur for either supersymmetric or non-supersymmetric constructions, when the overall volume of the extra dimensions is much larger than the volumes of the three-cycles wrapped by the D6-branes, and provides a stringy implementation of the phenomenological brane world models with large extra dimensions (23) [for recent reviews, see e.g., (180, 181, 182)].

Most non-supersymmetric constructions, including many toroidal ones, have assumed a low M_s, of the order of the TeV scale, to avoid large radiative corrections governed by M_s which aggravate the Higgs-hierarchy problem and to avoid large contributions of $\mathcal{O}(M_s^4)$ to the cosmological constant resulting from NS-NS tadpoles. However, as discussed in subsection 2.1, for the purely toroidal constructions with intersecting D6-branes, there is typically no direction in the internal space that would be transverse to all the D6-branes, and thus the size of the internal space is constrained to be of the order of the Planck 6-volume, and M_s is restricted to be within a few orders of magnitude to the Planck scale. The consistent implementation of the non-supersymmetric constructions with a low M_s would therefore be possible only for more general Calabi-Yau spaces (like fractional D6-branes on toroidal orbifolds). On the other hand, for the supersymmetric constructions , the NS-NS tadpoles cancel and the radiative corrections below M_s are at most logarithmic. Therefore such constructions with large M_s, as dictated by toroidal (orbifold) internal spaces, are stable, resulting in a calculable spectrum and effective Lagrangians at M_s.

We should also point out that models (57, 67) with locally supersymmetric spectrum of the MSSM, which however do not cancel R-R tadpoles for toroidal orientifolds, are also of interest ~~since~~ **because** they may provide a prototype D-brane configuration of the MSSM-sector. [R-R tadpole free implementation of such construction was realized, at the expense of Standard Model chiral exotics, within $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifolds in (48, 108); see subsection 3.2.2.]

Another issue to keep in mind is that there may be new physics at the TeV scale beyond the Standard Model or MSSM. Most of the explicit constructions lead, e.g., to additional Higgs or exotic matter or additional $U(1)$ gauge symmetries at the TeV scale [as do most heterotic constructions; see, e.g., (183, 184, 185)]. It is of course possible that these are defects of the models and that there is nothing beyond the MSSM at the TeV scale. On the other hand, one should keep open the possibility of a rich spectrum of "top-down" motivated new physics, especially of the kinds that appear so commonly in constructions.

6.1 The Spectrum

Most explicit intersecting D-brane constructions that contain the MSSM spectrum also involve additional matter states. As described in section 3, it is straightforward to construct non-supersymmetric or locally supersymmetric intersecting D-brane models with only the Standard Model spectrum, but existing fully supersymmetric constructions are highly constrained and always contain additional matter. This is true of most heterotic constructions as well. It is often considered the goal of model building to come as close to the MSSM as possible, but it is also possible that additional matter really does exist at the TeV scale.

Such states may either be chiral or non-chiral, although non-chiral states can typically obtain a string scale mass after deformations of D-brane configurations. For toroidal (orbifold) models, these states involve two stacks of D-branes that wrap the same one-cycle in one two-torus; parallel splitting of the two stacks of D-branes in this two-torus renders the non-chiral matter massive (21). Chiral matter occurs in (non-Abelian) anomaly-free combinations, and is frequently necessary to cancel the anomalies associated with additional gauge factors that would not be present for the MSSM spectrum alone. There are stringent constraints from precision electroweak physics on new matter that is chiral with respect to $SU(2)_W \times U(1)_Y$ (186) that essentially exclude the possibility of a fourth ordinary or mirror family or other representations that are chiral under the Standard Model gauge group, except possibly for rather tuned ranges of masses or other compensations. However, states that are singlets or vector-pairs under the Standard Model group but chiral under additional gauge factors are allowed in those cases in which the construction allows a mechanism for generating fermion masses in the hundreds of GeV range (scalar masses may be due to soft supersymmetry breaking). Additional matter may also have important implications for gauge coupling unification, as discussed in subsection 6.3.

6.1.1 Extended Higgs Sector

One ubiquitous possibility is an extended Higgs sector, involving more than one pair of Higgs doublets (often many pairs) and often Standard Model singlets whose VEVs could break additional gauge factors. Additional doublets would lead to a rich Higgs spectrum detectable at colliders and could mediate flavor-changing neutral currents. Higgs singlets S could couple to doublets with superpotential couplings $W = hSH_uH_d$, so that a TeV-scale VEV could lead to an effective μ parameter $\mu_{eff} = h\langle S\rangle$, elegantly solving the μ problem. This would be an implementation of some form of the next to minimal supersymmetric Standard Model (NMSSM) [see, for example, (187) and references therein] or its $U(1)'$ extension (188). Such models differ dramatically from the MSSM, e.g., by the expanded spectrum; possible large doublet-singlet mixing with implications for Higgs masses, production, and decays (189); a different allowed range for $\tan\beta$; expanded possibilities for electroweak baryogenesis because of a strong first order phase transition and new sources of CP violation [see (190, 191) and references therein]; and an enlarged and modified neutralino sector, extending the possibilities for cold dark matter (192, 191, 193). It should be stressed that no existing construction has all of these features or all of the couplings needed for a fully realistic model. For example, the supersymmetric construction (20) has no chiral singlet S to generate a μ_{eff}, though its role could be played by a field in the $\mathcal{N} = 2$ sector if it did not acquire a large mass (93). [Another possibility for a μ term would be a D-brane splitting in models in which the Higgs doublets are non-chiral, as in the locally supersymmetric model in (57, 67), although it is not clear why the splitting would be sufficiently small.]

6.1.2 Exotic quarks and leptons

Heavy exotic (i.e., with non-standard Standard Model representations) quarks and leptons, presumably vector-like pairs with respect to $SU(2)_W \times U(1)_Y$, are also possible, with exotic quarks especially producing distinctive effects at a hadron collider [see, e.g., (194)]. These are familiar in E_6 grand unification [see, e.g., (195)], and often emerge in string constructions as well, e.g., associated with the remnants of a would-be fourth family (20).

6.1.3 Chiral exotics

Still more exotic (and probably unwanted) possibilities exist. These include the open string sector moduli, typically in the adjoint (anti-symmetric) representation for the $U(N)$ ($Sp(2N)$) gauge symmetry, as well as fields in the symmetric and anti-symmetric representation for the $U(N)$ gauge factors, associated with the intersection of D6-branes with its orientifold image (see subsection 2.6). It is expected that the inclusion of fluxes would induce a back-reaction that would give a mass to some of the open string moduli (160). Another possibility is to introduce rigid three-cycles like fractional branes in toroidal orbifolds. D6-branes wrapping such three-cycles do not have massless moduli in the adjoint representation.

Many constructions also involve intersections between the ordinary and hidden sector branes, leading to states that are charged under the non-Abelian factors of both, i.e., the hidden sector is really only quasi-hidden. ($U(1)$ gauge bosons also typically couple to both sectors.) These mixed states often carry exotic electric charges (such as $1/2$). The laboratory and astrophysical constraints on fractional charges are severe (196). Fortunately, the quasi-hidden groups are often strongly coupled, so that such states may be confined, and may even lead to observable composite states with more conventional quantum numbers (93).

6.1.4 Grand Unification exotics

As discussed in subsection 3.2.1, it is possible to construct Grand Unified gauge groups such as $SU(5)$ from intersecting branes, and both fundamental and adjoint

Higgs, and fundamental and antisymmetric matter representations, appear. The supersymmetric three-family $SU(5)$ models that have been constructed always also include symmetric **15**-plet representations (101), which contain highly exotic states such as color sextets and weak triplets. Other phenomenological aspect of Grand Unification, such as doublet-triplet splitting (105) due to discrete Wilson lines, are touched on in subsection 3.2.1.

6.2 The Gauge Group

Physics beyond the Standard Model may involve extended gauge groups. In particular, intersecting D6-brane models respectively D-brane models in general [[****AU: Construction with respectively OK? Please recast for clarity if possible.****]]involve $U(N)$ ($Sp(2N)$) groups for planes not parallel (parallel) to orientifold planes. These may break to the Standard Model gauge group $G_{SM} = SU(3)_C \times SU(2)_W \times U(1)_Y$ and additional $U(1)$ and non-Abelian factors (the latter most commonly involves a quasi-hidden sector non-Abelian group, as discussed in subsection 6.6). At intermediate stages in all explicit $\mathbb{Z}_2 \times \mathbb{Z}_2$ examples, $SU(3)_C$ is embedded into a Pati-Salam $SU(4)$ in which lepton number is the fourth color (197). In some cases, there is also an embedding of $SU(2)_W \times U(1)_Y$ into $SU(2)_W \times SU(2)_R$. Here we focus on additional $U(1)$ factors, which frequently occur in intersecting D-brane constructions, as well as other types of string constructions (198, 199) and alternative approaches to move beyond the Standard Model, such as dynamical symmetry breaking (200) and Little Higgs models (201, 202). Because of their generality, extra $U(1)$ symmetries and their associated heavy Z' bosons are probably the best motivated extension of the Standard Model after supersymmetry. Experimental limits on an extra Z' are very model dependent, depending on the Z' mass, gauge coupling, and couplings to the left- and right-handed quarks and leptons. However, typical limits from the combination of Z-pole and other precision experiments and direct searches for $p\bar{p} \to e^+e^-,\ \mu^+\mu^-$ are typically $M_{Z'} > 500 - 800$ GeV and the $Z - Z'$ mixing less than a few $\times 10^{-3}$ (see, e.g., (203)).

There are several sources of extra $U(1)$ symmetries in intersecting D-brane models. One is that stacks of branes not parallel to O6-planes yield $U(N) \sim SU(N) \times U(1)$ groups, where the $U(1)$s are typically anomalous. Recall that the $U(1)$ anomalies are canceled via a generalized Green-Schwarz mechanism. As described in subsections 2.7 and 3.1, the Z' gauge bosons associated with these extra $U(1)$ factors will typically acquire string-scale masses by Chern-Simons terms, even if they are not anomalous (18). [Field theoretic analogs have been studied recently in (204).] However, the $U(1)$'s will survive as perturbative global symmetries of the theory, often restricting possible Yukawa couplings and/or leading to conserved baryon and lepton numbers, stabilizing the proton and preventing Majorana neutrino masses (18). For non-supersymmetric models with a TeV string scale, this implies new Z' gauge bosons with masses generated without a Higgs mechanism. Experimental constraints from their mixing with the Z have been examined in (205, 206, 207), where lower bounds on $M_{Z'}$ and therefore on the string scale in the TeV range were obtained, somewhat more stringent than typical bounds on E_6-motivated Z' bosons.

Non-anomalous additional $U(1)$'s may arise from the breaking of $SU(N)$ factors by parallel splitting of $U(N)$ branes, such as the extra $U(1)_{B-L}$ emerging from $SU(4)$ in (20, 21) and other typical Standard-like Model constructions; or from the splitting of $Sp(2N)$ branes parallel to O6-planes, such as the $Q_8 - Q_{8'}$ in

(20, 21). [See (108) for a general discussion.] As discussed in (93), such $U(1)$ factors need to be broken by the VEVs of Standard Model singlets charged under the $U(1)$. The breaking could be at the TeV scale if it is driven by the same type of terms which drive electroweak breaking (208), or at a scale intermediate between the TeV and Planck scale if it is along a D and (tree-level) F-flat direction (209), provided there are appropriate Standard Model singlet fields. In some cases, the only candidates are the bosonic partners of right-handed neutrinos (93), and the needed couplings are not always present.

Experimental implications of a TeV-scale Z' are significant. These include the effects at colliders of the Z' and associated exotics needed for anomaly cancellation (see, e.g., (203)), and the effects of the extended Higgs and neutralino sectors for colliders and cosmology, commented on in subsection 6.1. The Z' couplings are often family-nonuniversal in both intersecting brane and heterotic constructions, implying flavor-changing neutral currents after fermion mixing is turned on [see, e.g., (210, 211, 212)]. These could be significant, e.g., for rare B, K, and μ decays.

6.3 Gauge Coupling Unification

It is well known that the observed (properly normalized) low energy gauge couplings $\alpha_1^{-1} \equiv \frac{3}{5}\alpha_Y^{-1}$, α_g^{-1}, and α_s^{-1} associated respectively with $U(1)_Y$, $SU(2)_W$, and $SU(3)_C$ are roughly consistent with gauge unification at a scale $M_U \sim 3\times10^{16}$ GeV when the β functions are calculated assuming the MSSM particle content [see, e.g., (213)]. The value of $\alpha_s \sim 0.13$ predicted from α and $\sin^2\theta_W$ is slightly larger than the observed value (~ 0.12), even accounting for uncertainties in the sparticle spectrum, but could be due to high scale threshold effects in traditional

Grand Unified theories. In heterotic string constructions one expects to maintain gauge unification at the string scale, which is typically an order of magnitude larger than the apparent GUT [[**AU: Please spell out GUT unless all readers will know this acronym.**]]scale M_U. However, the normalization of gauge couplings at M_s is modified for higher Kač-Moody embeddings, which are common for $U(1)_Y$ but not for $SU(3)_C \times SU(2)_W$. Also, most constructions involve additional matter which can modify the β functions. These two effects can each modify the predicted α_s and $\log M_U$ by $O(1)$, compared to the 10% corrections that are needed, i.e., traditional gauge unification is lost unless these two effects are absent or somehow compensate.

As discussed in subsection 4.2, traditional gauge unification is lost in most intersecting D-brane constructions because the gauge coupling at the string scale for each stack of D-branes depends on stack-dependent moduli, i.e., on the volume of the three-cycle wrapped by the stack. [One exception are supersymmetric Grand Unified Models (21, 101), in which the Standard Model gauge factors all are derived from a single stack. A local construction (90), based on a three-family $U(3)^3$ sector also provides a gauge coupling unification of the three gauge factors at the string scale.] Furthermore, as discussed in subsection 6.1 the constructions (including the Grand Unified ones) typically involve exotic states that will modify the running. One must therefore hope that these effects will somehow compensate to yield the observed couplings.

One approach is to predict the low energy couplings (including those for any additional gauge factors) for a given construction in terms of the spectrum. This was done as a function of M_s and the volume moduli in (93) for the supersymmetric model (20) described in subsection 3.2. It was found that the predicted couplings were typically smaller than the observed ones due to the extra chiral matter. The analysis was refined in (95), in which it was shown that due to gaugino condensation, the toroidal complex structure moduli and dilaton are fixed (see also subsection 6.6). One could then predict $\alpha_s^{-1} \sim 52.2$ and $\alpha(M_Z)^{-1} \sim 525$, much larger than the observed values ~ 8.5 and 128, due to the extra chiral matter. The weak angle, which is a ratio, came out better, with the predicted $\sin^2\theta_W \sim 0.29$ not too far from the observed 0.23. ~~While~~ **Although** not successful, this illustrates the possibility that a more realistic construction might lead to the observed couplings.

In a general D-brane construction, there ~~will not be a~~ **is no** simple relation between the three gauge couplings at M_s. However, it was observed in (141) that under certain circumstances, there ~~would be~~ **is** a tree-level relation

$$\alpha_1^{-1} = \frac{2}{5}\alpha_s^{-1} + \frac{3}{5}\alpha_g^{-1}, \tag{63}$$

a special case of the canonical GUT relation in which all three are equal. This could come about in models in which the weak hypercharge satisfies the left-right symmetry relation $Q_Y = \frac{1}{2}(B-L) + Q_{3R}$, with the additional assumptions that $U(1)_{B-L}$ derives from the same stack as $SU(3)_C$ (Pati-Salam embedding) and that there is a left-right symmetry that ensures the same coupling for Q_{3R} and $SU(2)_W$. It was shown in (141) that if (63) holds, then from the observed low energy couplings and from the contributions to the β functions from exotic matter, one can predict the value of M_s and the volume moduli. In fact it turned out that an effective β function coefficient was always an even integer leading to a discrete set of possible values for the string scale. For example, for no exotics one finds $M_s \sim 2 \times 10^{16}$ GeV with volume radii $R_s = 2.6/M_s$ and $R_W = 3.3/M_s$ for the $SU(3)_C$ and $SU(2)_W$ branes, respectively. The addition of exotic matter can lead to very different M_s and radii. Of course, there is no guarantee that after stabilization M_s and the moduli would actually take these values. ~~While~~ **Although** the first assumption (on $U(1)_{B-L}$) is satisfied by existing supersymmetric constructions because $SU(3)_C \times U(1)_{B-L} \subset SU(4)$, the second (on Q_{3R}) is not satisfied in most known constructions such as (20, 21), which are not left-right symmetric. [Examples in which it does hold were given for the locally supersymmetric construction (109), two models among supersymmetric constructions in (107) and a four-family model in (108).] In (127) the frequency of this relation was statistically investigated in the ensemble of MSSM like Gepner model orientifolds. It was found that for approximately 10% of these models, this relation was satisfied [which could clearly be seen in the overall plot in (127)].

6.4 Yukawa Couplings

Yukawa couplings and the pattern of fermion masses and mixings are one of the least understood aspects of nature. In the context of the Standard Model or MSSM, or in simple Grand Unification extensions, it is often assumed that some sort of additional family symmetry might lead to textures (hierarchies of elements including zeroes) in the fermion Yukawa matrices to explain the observed patterns. The ratio $\tan\beta$ of the VEVs of the neutral Higgs fields from the two Higgs doublets H_u and H_d of the MSSM may also play a role. The recent observation of neutrino oscillations further complicates the situation because of the possibility of Majorana masses.

In existing string constructions (including heterotic), the possible Yukawa and other superpotential interactions are typically very much restricted by additional symmetries (e.g., the perturbative global symmetries that remain after anomalous

$U(1)$'s are broken by the Green-Schwarz mechanism) or by stringy selection rules such as orbifold and orientifold projections. Such restrictions may be weakened in more general constructions, but are an important feature of existing examples, and they may lead, e.g, to texture zeros. For example, one of the families of quark and lepton doublets in the supersymmetric model (20) in subsection 3.2 has no Yukawa couplings due to the Q_2 symmetry and remains massless, or the $SU(5)$ models described in subsection 3.2.1 have no $\mathbf{10}\ \mathbf{10}\ \mathbf{5}_H$ couplings ~~due to~~ **owing to** the $U(1)$ of $U(5)$. Many models [e.g., (18, 20)] have conserved B and L ~~due to~~ **owing to** global and local $U(1)$'s, stabilizing the proton and preventing Majorana neutrino masses. Similarly, in existing intersecting D-brane models (unlike some heterotic constructions) H_d and lepton doublets are clearly distinguished even though they have the same Standard Model quantum numbers because two of the relevant branes are distinct.

String constructions also allow natural mechanisms for hierarchies of Yukawa couplings. For example, free fermionic models can lead to small effective couplings from higher dimensional operators. As described in subsection 4.1.1 intersecting D-brane constructions allow for a geometrical origin of hierarchies, because allowed Yukawa couplings are ~~due to~~ **caused by** world-sheet instantons and are proportional to $\exp(-A)$, where A is the area of the triangle connecting the three intersecting branes.

~~Another aspect is that e~~ **E**xisting supersymmetric intersecting D-brane constructions contain more than a single $H_{u,d}$ pair, as described in subsection 6.1 (this is true for many heterotic constructions, as well), with each having different Yukawa matrices. Thus, hierarchies of their VEVs could be an additional mechanism for achieving hierarchies of masses and nontrivial mixings, and in generating otherwise vanishing masses. Of course, the actual VEVs would depend on the details of how supersymmetry is broken. In particular, in schemes of radiative electroweak breaking (in which negative Higgs mass squares are generated from positive ones at a higher scale by renormalization group running) there will be a strong tendency for only those Higgs fields with large Yukawa couplings to actually acquire VEVs. There has been relatively little phenomenological work on these sorts of extended Higgs sectors.

In specific intersecting D-brane models on toroidal (orbifold) compactifications, the Yukawa couplings often factorize in terms of the family indices for the left and right-handed fermions, e.g., the couplings $h^k_{i,j}$ between H^k_u, Q_i and $\bar{U}_j$ are proportional to products $a^k_i b^k_j$. This can occur, for example, if the non-trivial intersections for Q_i and $\bar{U}_j$ occur in different two-tori (137) or if the orientifold and orbifold projections associate each $\bar{U}_j$ with a distinct H^j_u (94). The factorization does not hold in more general examples [e.g., (107)]. Factorization could actually pose a problem for a construction with only a single pair of $H_{u,d}$ doublets, because it allows only one massive state of each fermion type (u-type, d-type, e-type). Some means must therefore be found to populate other terms in the mass matrices. Possibilities include accepting additional Higgs pairs, modifying the D-brane geometry, invoking (non-aligned) four-point interactions in non-supersymmetric models with low M_s (214), or allowing for (non-aligned) supersymmetry breaking A terms (if allowed by the supersymmetry breaking mechanism) (214). There are also potential problems with the minimal two-doublet structure if the electroweak symmetry is promoted to $SU(2)_L \times SU(2)_R$, because the $SU(2)_R$ symmetry would ensure equal Yukawa matrices for the u and d, preventing a nontrivial CKM quark mixing matrix (67, 109, 137). [This is also one reason $SO(10)$ models require more than a single Higgs multiplet coupling to fermions (215).]

A more detailed analysis of the Yukawa couplings for the supersymmetric multi-Higgs model (20) described in subsection 3.2 was made in (94). It was shown that for appropriate values of some (unknown) volume moduli, one could obtain nontrivial masses and mixing for two families. Near the symmetric points (small splitting between stacks of branes) one obtains the GUT-like result of similar d and charged lepton mass matrices, as well as similar u and Dirac neutrino masses. The Dirac neutrino masses are problematic because the model has no non-perturbative mechanism to generate Majorana masses for a seesaw mechanism. The Yukawa structure for non-supersymmetric models with one or two pairs of Higgs fields were studied in (18, 214). A locally supersymmetric model with a single Higgs pair (whose global embedding was realized in (110, 48, 49)) was considered in (109, 137), where it was emphasized that having only one massive family is actually an excellent first approximation, ~~since~~ **because** m_t, m_b, and m_τ are much larger than the other generations.

6.5 Flavor-changing Effects and Proton Decay

In the Standard Model there are no flavor-changing neutral currents (FCNC) mediated by the Z, γ, or Higgs at tree-level, and FCNC at loop-level are suppressed (the GIM mechanism). However, there are enhanced FCNC effects in most extensions of the Standard Model, including new loop effects in supersymmetry and new interactions in dynamical symmetry breaking. Similarly, the only sources of CP violation are the phases in the quark (and lepton) mixings, possible neutrino Majorana phases, and a possible strong CP parameter θ_{QCD}. For small quark mixings, all but θ_{QCD} lead to extremely small neutron, atomic, and electric dipole moments (EDM), ~~while~~ **whereas** most extensions of the Standard Model lead to enhanced effects. Therefore, experimental studies of rare decays and suppressed mixings, such as $\mu \to 3e$, $K_L - K_S$ mixing and rare B decays, as well as refined ~~electric dipole moment~~ **EDM** experiments, are an excellent way to search for new physics.

6.5.1 FCNC

There are a number of sources of FCNC in string constructions (in addition to the standard sparticle loops in supersymmetric constructions). The tree-level calculation of the string four-point amplitudes (134) (see subsection 4.1) produces flavor-changing four-point operators in the effective action. For non-supersymmetric constructions with a low M_s, the analysis of such operators was carried out in (135, 214), where it was shown that there could be significant effects from both Kaluza-Klein modes and stretched heavy string modes. For example, Kaluza-Klein excitations couple non-universally to states located at different positions, and therefore to FCNC. The authors of (135, 214) studied the constraints on these operators from experimental founds on FCNCs, EDMs, and supernova cooling by neutrino emission induced by four-fermi operators, and they showed that the FCNC ~~especially~~ severely restrict the string scale to be higher than $\sim 10^4$ TeV. This suggests that such non-supersymmetric constructions have a severe fine-tuning problem and makes it unlikely that other effects, such as the $U(1)$ gauge bosons which acquire a string-scale mass by the Chern-Simons terms (205, 206, 207) described above, will be observable.

A number of other (field theoretic) sources of FCNC may be expected from intersecting D-brane (and other) string constructions and may be observable in future experiments. The most promising are additional TeV-scale $U(1)$'s with family-nonuniversal couplings, as described in subsection 6.2; multiple Higgs doublets, for which the neutral components can mediate FCNC; or extended non-Abelian groups that can survive down to low energies, such as the embedding of $SU(2)_W$ into $Sp(6)$ at ~ 100 TeV, leading to $K_L \to \mu^{\pm} e^{\mp}$ (108).

6.5.2 CP violating phases

The CP violating phases for supersymmetric constructions can appear in the Yukawa couplings which depend on the VEVs of the (complex) Kähler moduli [for a detailed discussion of this moduli dependence, see (137)]. Another source of the CP violating phases can be complex soft supersymmetry breaking masses and the μ parameters; in intersecting D-brane constructions the complex soft supersymmetry breaking mass parameters are due to the complex VEVs of closed sector moduli, as discussed briefly in subsection 6.6.

6.5.3 Strong CP problem

In (216) a mechanism to solve the strong CP problem was proposed, which could have a realization within intersecting D6-brane models. This mechanism is reminiscent of the (chiral) anomaly inflow mechanism. Specifically, the proposal employs an additional bulk $U(1)_X$ gauge factor under which quarks are not charged, and the flux associated with the NS-NS three-form field strength H_3. The anomaly cancellation takes place ~~due to~~ **owing to** a Chern-Simons term of the Type IIA supergravity and terms in the expansion of the D6-brane world-volume Chern-Simons action. A specific non-supersymmetric intersecting

D6-brane model that explicitly realizes this mechanism was constructed in (216). It remains an open problem to implement this mechanism for the supersymmetric intersecting D6-brane constructions with supersymmetric H_3 fluxes.

6.5.4 Proton Decay

Supersymmetric Grand Unified theories (186) allow proton decay by dimension 5 or dimension 6 operators (we assume that dimension 4 R parity-violating terms that could lead to unacceptable rates are absent). The dimension 5 operators (via heavy colored fermion exchange) lead to too rapid proton decay unless they are somehow forbidden, ~~while~~ **whereas** the dimension 6 operators from heavy gauge boson exchange typically lead to a lifetime of the order of 10^{36} yr, too long to observe in planned experiments (the current limit of $\sim 4 \times 10^{33}$ yr for $p \to e^+\pi^0$, which may be improved to $\sim 10^{35}$ yr).

The expectations for proton decay in supersymmetric intersecting D-brane constructions have been studied recently in (136). In many intersecting D-brane constructions, baryon number is conserved perturbatively and the proton is stable. However, the proton can decay in the Grand Unified constructions described in 3.2.1. The four-fermion contact operator for $\mathbf{10}^2\overline{\mathbf{10}}^2$ in intersecting D-brane $SU(5)$ models for the four states located at the same intersection (where there is no suppression from area factors) was calculated in (136) (see also subsection 4.1). This operator has an enhancement, relative to the standard Grand Unified Models, ~~due to~~ **caused by** the exchange of Kaluza-Klein excitations of the color triplet gauge bosons, which leads to the decay amplitude $\propto \alpha_{GUT}^{-1/3}$. In order to further increase the decay amplitude, the string coupling was taken to be $\mathcal{O}(1)$, thus leading to the M-theory on G_2 holonomy space (see subsection 2.9), and the gauge coupling threshold corrections (217) were included. However the final result did not have additional large enhancement factors, suggesting a lifetime of around 10^{36} yr, comparable to ordinary supersymmetric Grand Unification.

6.6 Moduli Stabilization and Supersymmetry Breaking

Intersecting D-brane constructions on toroidal (orbifold) backgrounds possess a large number of closed and open string sector moduli, thus leading to a large vacuum degeneracy. In fact, the vacuum degeneracy problem is generic for supersymmetric string constructions. As mentioned before, this problem has been addressed via two mechanisms: (1) implementation of the strong D-brane gauge dynamics that can lead to gaugino and matter condensations and generates a non-perturbative superpotential for the closed string sector moduli fields; (2) introduction of supergravity fluxes whose back-reaction introduces a moduli dependent potential. It is expected that in a realistic framework, a combination of both mechanisms will play a role in obtaining string vacua with (all) moduli stabilized, broken supersymmetry and potentially realistic cosmological constant. In the following we shall summarize the phenomenological implications, studied for these two mechanisms.

6.6.1 Strong D-brane gauge dynamics

Explicit supersymmetric intersecting D6-brane constructions typically possess a quasi-hidden gauge sector that has a number of non-Abelian confining gauge group factors, typically with $Sp(2N)$ gauge symmetries. The non-perturbative superpotential of the Veneziano-Yankielowicz type (145, 146) is a sum of exponential factors (associated with each confining gauge factor):

$$W_a(U^i, S) = \frac{\beta_a}{32\pi^2}\frac{\Lambda^3}{e}\exp(\frac{8\pi^2}{\beta_a}f_a(U^i, S))\,, \tag{64}$$

where the dynamically generated scale Λ is roughly of the order of the string scale M_s, β_a is the beta function of the specific gauge group factor and $f_a(U^i, S)$ denotes the corresponding gauge kinetic function, which for intersecting D6-branes depends on the complex structure moduli U^i and the dilaton field S. Eq.(64) accounts only for the leading instanton contribution. One should also point out that for a specific number of "flavor" (matter) N_f and "color" (gauge) N_c degrees of freedom, there are subtleties; e.g., ~~for~~ $Sp(2N_c)$ gauge factors~~,~~ can lead to the quantum lift of the moduli space ($N_f = N_c + 1$) or absence of the non-perturbatively generated global superpotential ($N_f > N_c + 2$). [For a review see, e.g., (218) and references therein; for the implementation of strong gauge dynamics in the effective actions from heterotic strings, see (219).] Classes of semi-realistic supersymmetric intersecting D6-branes constructions, e.g., (106, 107), have the property that the hidden sector gauge group factors satisfy $N_f < N_c + 1$, resulting in confining infrared dynamics and the non-perturbative superpotential of the type (64).

For toroidal (orbifold) compactifications, as discussed in subsection 4.2, the tree-level gauge kinetic function $f_a(U^i, S)$ (53) depends on the dilaton S and three toroidal complex structure moduli U^i (some of the toroidal complex structure moduli are fixed by the supersymmetry constraints in the D6-brane sector) and the specific wrapping numbers $(n_a^i, \widetilde{m}_a^i)$ of the three-cycle π_a, wrapped by a stack of N_a D6-branes. For a specific supersymmetric semi-realistic construction (20, 21), the non-perturbative superpotential (64), associated with the confining $Sp(2) \times Sp(2) \times Sp(4)$ sector, resulted (95) in the minimum of the potential that stabilized the remaining toroidal complex structure modulus U and the dilaton S, and broke supersymmetry. It would also be interesting to implement the threshold corrections to the gauge kinetic function (144) as discussed in subsection 4.2. For $\mathcal{N}=2$ sectors these corrections depend also on toroidal Kähler moduli, and thus the non-perturbative superpotential (64) could in principle allow for the stabilization of the toroidal Kähler moduli as well.

When supersymmetry is broken ~~due to~~ **by** such a non-perturbative superpotential, the gaugino masses m_{λ_a} can be determined in terms of F-breaking terms associated with S and U^i moduli directions:

$$m_{\lambda_a} = (\partial_{\phi^i} f_a(\Phi^i))K^{\Phi^i\,\bar{\Phi}^j}\bar{F}_{\bar{\Phi}^j}\,. \tag{65}$$

Here $K^{\Phi^i\,\bar{\Phi}^j}$ is the inverse of the Kähler metric of the moduli Φ^i, and F_{Φ^j} are the F-breaking-terms for the moduli $\Phi^j = \{S,\ U^i\}$. Unlike the heterotic constructions and simple Grand Unified theories, the gaugino masses and gauge couplings at the string scale depend on more than one modulus, i.e., S and U_i, which in general have complex VEVs and thus lead to non-universal and complex (indicating significant CP-violating phases) gaugino masses. Unfortunately for the specific model studied in (95), these masses were too heavy, i.e., $\mathcal{O}(10^8)$ GeV.

The study of soft supersymmetry breaking parameters of the charged matter sector requires detailed information on the moduli dependence of the leading term in the Kähler potential for the charged matter; this Kähler potential was recently determined in (133, 132) and discussed in subsection 4.1.2 (specifically, see Equation 50). Unfortunately, for the specific example studied in (95), the minimum of the non-perturbative superpotential produced a large negative cosmological constant, and thus these vacua do not provide realistic backgrounds for a detailed study of the soft supersymmetry breaking parameters of the charged matter sector. However, one can assume that the non-perturbative mechanism for supersymmetry breaking does not introduce a large cosmological constant, and then one can parameterize such soft masses via F_{Φ_i}-breaking terms associated with moduli Φ_i by employing the standard supergravity techniques. Such a study was recently performed in (220). [For an earlier work see (221).] In the regime where the F_{U_i}-breaking terms~~,~~ **that are** associated with the U_i moduli~~,~~ are dominant, the mass parameters do not depend on the Yukawa couplings and have a pattern different from the heterotic string.

In principle, the strong gauge dynamics can also lead to composite (baryon-type) states whose constituents include states that are chiral exotics, i.e., states charged both under the Standard Model gauge factors and the hidden strong gauge sectors. This scenario could provide another mechanism to remove chiral exotics from the light spectrum (see (93)).

6.6.2 Supergravity fluxes

Supergravity fluxes provide another mechanism to stabilize the compactification moduli. The supersymmetric flux compactifications are better understood on the Type IIB side (for details see section 5). Semi-realistic constructions of Type IIB vacua consist of the magnetized D-brane sector (T-dual to the intersecting D6-branes) and the G_3 fluxes stabilizing the toroidal complex structure moduli (in the T-dual picture Kähler moduli) and the dilaton-axion field.

Typical semi-realistic examples have fluxes that break supersymmetry via a $(0, 3)$ component of G_3, and recent phenomenological studies focused on the implied generation of soft supersymmetry breaking terms in the low energy effective action (154, 155, 156, 157, 158, 222, 223, 224, 220). These terms have been derived by employing two complementary approaches:

- The soft supersymmetry breaking mass terms due to fluxes were obtained by expanding the resulting Dirac-Born-Infeld action for the D3 and D7 branes (154, 155, 158) to the lowest order in the coordinates transverse to the D-brane world-volume.
- Employing the standard supergravity formalism, one can parameterize the soft supersymmetry breaking terms via the supersymmetry breaking VEVs of the auxiliary F and D components of the chiral and vector supermultiplets (156, 157).

Both approaches are expected to be equivalent. A third approach using the F-theory description of a certain orientifold has been pursued in (65).

We have given in section 5 a heuristic argument why such soft terms are generated. In the following we summarize the specific results. For the imaginary self-dual G_3, the soft supersymmetry breaking mass terms are absent for the matter associated with the open string states on the D3-branes. However, for the anti-D3-branes these masses are non-vanishing, and specifically they stabilize the open-string modulus associated with the position of the anti-D3-brane (154, 156). This point also plays a very important role in getting de-Sitter vacua via the KKLT construction (165).

On the other hand, such mass terms for the open string states on D7-branes are due to the non-supersymmetric $(0,3)$-components of G_3 fluxes (158, 157). Interestingly, the supersymmetric $(2,1)$-components of G_3 fluxes can induce superpotential mass terms for the D7-brane moduli, including those associated with D7-brane "intersections" (158), thus providing a stabilization mechanism for them. Assuming a homogeneous flux, the scale of such mass terms is of the order $\frac{M_s^2}{M_{pl}}$. For the supersymmetry breaking masses to be of the TeV scale, this implies that the string scale is in the intermediate regime. This is reminiscent of the gravity mediated supersymmetry breaking mechanism. More generally, such mass terms measure the local flux density and so M_s that appears in the above estimate should be the local string scale which can in principle be much smaller because of the non-trivial warp factor.

6.7 Cosmological Aspects

The main focus of this review has been the particle physics aspects of intersecting D-brane models. For completeness, however, let us briefly mention some cosmological aspects of this scenario as well. A comprehensive overview of string/brane cosmology is beyond the scope of this review. Here, we shall only sketch some highlights of this subject that are particularly relevant to intersecting D-brane worlds. For details and references, we refer the readers to some excellent reviews (225, 226, 227).

There has been widespread hope that string theory may provide a microscopic origin for inflation. The discovery of D-branes has opened up several new possibilities. In this review, we have focused on D-brane models that preserve $\mathcal{N} = 1$ supersymmetry, for otherwise the D-brane configurations are generically unstable with too short a life-time [[****AU: What, exactly, would spell disaster for particle physics? Please recast this sentence for clarity.****]]that would spell disaster for particle physics today. However, in the early universe, the initial configuration of D-branes is not necessarily perfectly stable. Instead the D-branes could intersect at non-supersymmetric angles, or there could be additional pairs of branes and anti-branes separated in the compact dimensions. These instabilities drive the system of D-branes to a neighboring stable configuration, so we can think of the supersymmetric models that we have discussed at length in this review as the endpoints of such dynamical processes. In fact, a natural candidate for the inflaton field in this scenario is the open string mode whose VEV describes the inter-brane separation (228). The dynamics of inflation is therefore governed by the interaction between D-branes. This idea of brane inflation (228) has been applied to construct inflationary models arising from the collision of branes and anti-branes (228, 229, 230, 231, 232), as well as branes intersecting at angles (233, 234, 235). In particular, (232), which is by far the most detailed model of inflation from string theory, demonstrated that the brane inflation proposal can be implemented in a string model where the geometric moduli are stabilized by the background fluxes (a concrete mechanism that we discussed in section 5).

Interestingly, the cosmic string network produced at the end of D-brane inflation offers an exciting opportunity to test stringy physics from cosmological observations [see, e.g.,(227, 236) for some reviews and references]. Toward the end of brane inflation, the inflaton potential becomes tachyonic. The condensation of this complex tachyon mode results in the formation of cosmic strings, rather than other cosmological defects such as monopoles or domain walls. Finally, in addition to cosmic D-strings, there could in general be stable D-branes carrying K-theory charges in the intersecting D-brane models discussed here. They could be interesting candidates for superheavy dark matter (237).

7 CONCLUSIONS AND OUTLOOK

In this ~~article~~ **review** we have provided a pedagogical introduction to string theoretic intersecting D-brane models, which we hope suits the need of students to have a comprehensive though not too technical guideline for this topic. ~~In the second part w~~ **W**e have **also** tried to briefly review much of the work on intersecting D-brane models carried out so far, including an overview ~~on~~ **of** model building attempts, ~~including~~ **such as** recent flux compactifications, as well as ~~on~~ the structure of the low-energy effective action. The latter of course is very important for concrete phenomenological applications of these models.

During the short history of intersecting D-brane constructions, ~~it has happened several times that~~ new momentum was brought into the field from other branches of string theory research like M-theory compactifications on G_2 manifolds or flux compactifications. Clearly, all these model building schemes are intimately related. After more than four years of intense research, it has become clear that intersecting D-brane models provide a general phenomenologically appealing class of string compactifications, [[****AU: Possible to break this sentence into two and/or recast for clarity?****]] which can also be considered as honest string theory realizations of some of the ideas concerned with more phenomenologically motivated brane world models.

Even though we have a nice geometric framework, we are still lacking a completely convincing model realizing the MSSM. One can find isolated mechanisms for realizing most of the features of the Standard Model, like family replication, hierarchical Yukawa couplings, absence of extra gauge symmetries and vector-like matter, etc., but all concrete models studied so far do not realize all Standard Model properties at the same time: They either have extra chiral exotic matter as is typical for supersymmetric constructions, or the models fail at the level of couplings, such as gauge and Yukawa couplings. Of course only a very few classes of models, primarily based on toroidal orbifolds, were constructed, and even fewer were studied in detail. In addition, the techniques are not yet available to study more general intersecting D-brane models on ~~say~~ **, for example,** generic smooth Calabi-Yau spaces. ~~In fact it seems to be the case that, for instance,~~ **It seems that** the notorious appearance of extra vector-like matter is related to the fact that we are only considering models at very special, highly symmetric points in moduli space, like orbifold or Gepner points. ~~In view of the impression~~ **Considering** that the finer details of the Standard Model are far from being very natural, there is no guarantee that nature has finally stabilized in a string vacuum ~~which~~ **that** is highly symmetric and treatable with the simple methods developed so far. Therefore, it would be interesting to develop the tools to study more generic intersecting D-brane models.

Alternatively, it is entirely possible that physics at the TeV scale is richer than the MSSM, and that some of the features found in existing constructions, such as extended gauge symmetries, extra chiral matter, and flavor changing neutral currents, really exist. The LHC and future experimental probes are eagerly awaited to refine the target of our theoretical investigations.

~~As reviewed in this article, c~~ **C**oncerning the low energy effective field theory, considerable progress has been made in computing, for example, Yukawa couplings, the Kähler potential, or the resulting soft supersymmetry breaking terms for very simple toroidal backgrounds. However, much more work is needed to derive similar results for more general backgrounds.

~~As should be clear~~ **Clearly**, for each Calabi-Yau manifold there does exist a plethora of consistent intersecting D-brane models. In view of these, one might ask whether there does exist any chance to find the/a realistic string vacuum. This picture becomes even more severe when one also takes **into account** the so-called landscape of flux compactifications ~~into account~~ . It was proposed that complementary to a model-by-model search, one could study the statistical distribution of string theory vacua (238) [see (239) for a statistical analysis of intersecting D-branes] to obtain an estimate of the chances of finding a realistic model ~~and maybe in which region of the parameter space one should look~~ **, as well as to identify a region of the parameter space in which to look**.

For the moment we can only hope that continuous work on both approaches—the model-by-model search and the statistical analysis—will eventually lead us to a realistic string model from which, once the background is fixed, all features of the low-energy effective theory can be derived. However, whether such a model is in any sense unique is not guaranteed, as we will always measure the physical parameters with some finite accuracy. Having one string model which describes our world within the accuracy of our measurements would nevertheless be considered a milestone in our understanding of nature.

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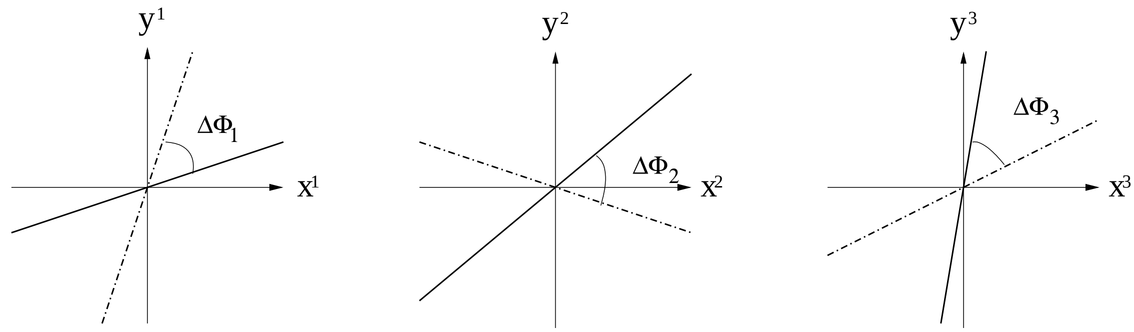

Figure 1: Intersecting D6-branes.

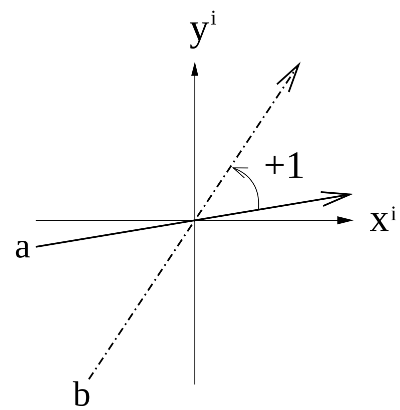

Figure 2: Oriented intersection.

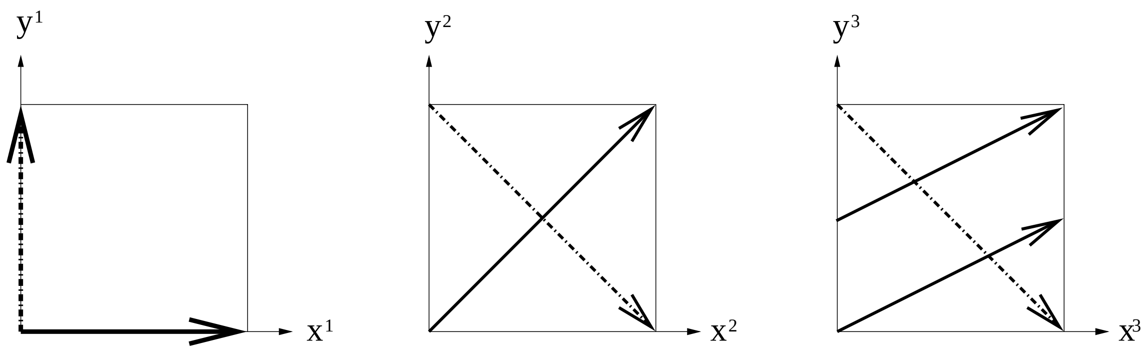

Figure 3: Intersecting D6-branes on a torus

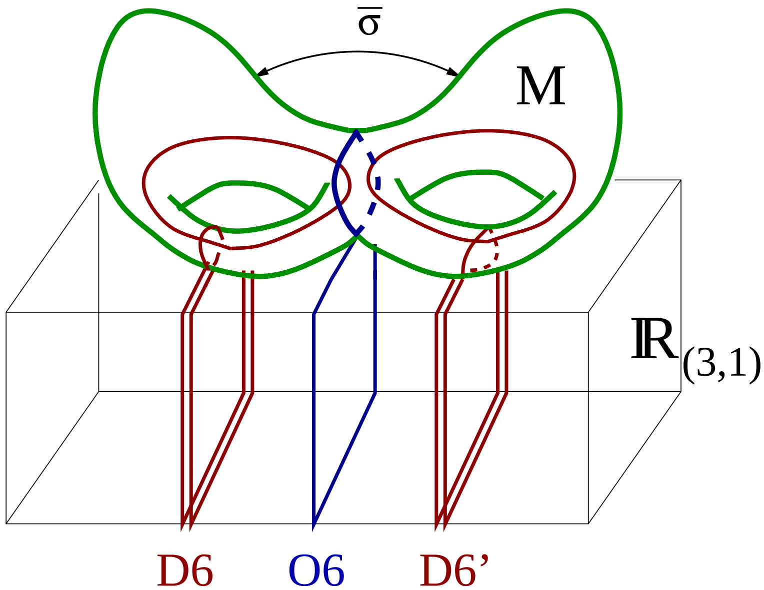

Figure 4: Schematic image of an $\Omega\overline{\sigma}(-1)^{F_L}$ orientifold with O6-planes and intersecting D6-branes. In reality the O6-plane and the D6-branes would cover the entire flat Minkowski space.

Table 1: Chiral spectrum for intersecting D6-branes

Representation	Multiplicity
${\begin{smallmatrix}\square\\\square\end{smallmatrix}}_a$	$\frac{1}{2}\left(\pi'_a \circ \pi_a + \pi_{\mathrm{O6}} \circ \pi_a\right)$
$\square\!\square_a$	$\frac{1}{2}\left(\pi'_a \circ \pi_a - \pi_{\mathrm{O6}} \circ \pi_a\right)$
$(\overline{\square}_a, \square_b)$	$\pi_a \circ \pi_b$
$(\square_a, \square_b)$	$\pi'_a \circ \pi_b$

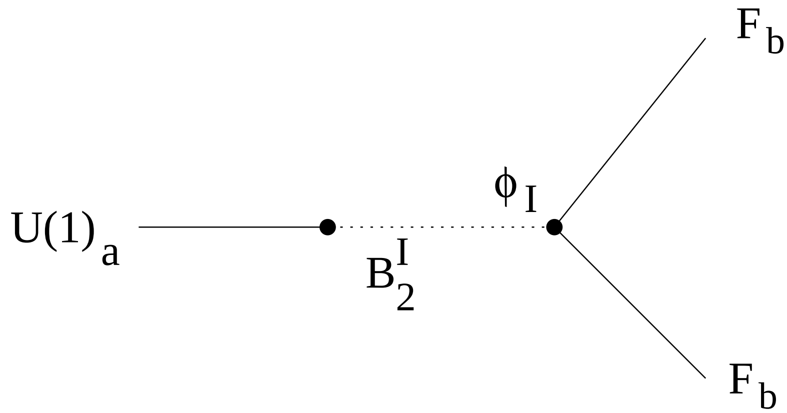

Figure 5: Green-Schwarz mechanism.

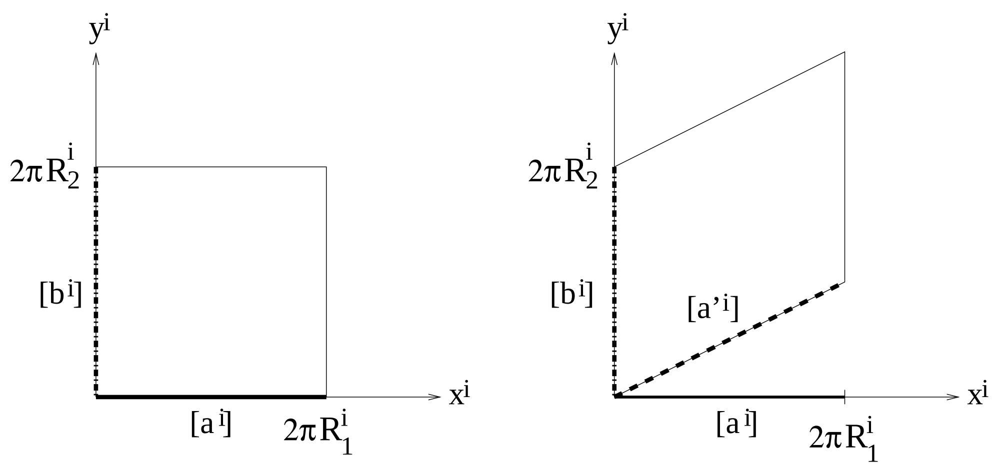

Figure 6: Choices of T^2s

Table 2: Wrapping numbers for a semi-realistic non-supersymmetric model. The parameters are defined as $\overline{\beta}^{1,2} = 1 - \beta^{1,2}$, $\beta^3 = 1/2$, $\rho = 1, 1/3$, $\epsilon = \pm 1$ and $n_a^2, n_b^1, n_c^1, n_d^2 \in \mathbb{Z}$.

N_a	$(n^1, \widetilde{m}^1)$	$(n^2, \widetilde{m}^2)$	$(n^3, \widetilde{m}^3)$
$N_a = 3$	$(1/\overline{\beta}^1, 0)$	$(n_a^2, -\epsilon\overline{\beta}^2)$	$(1/\rho, -1/2)$
$N_b = 2$	$(n_b^1, \epsilon\overline{\beta}^1)$	$(1/\overline{\beta}^2, 0)$	$(1, -3\rho/2)$
$N_c = 1$	$(n_c^1, -3\rho\epsilon\overline{\beta}^1)$	$(1/\overline{\beta}^2, 0)$	$(0, -1)$
$N_d = 1$	$(1/\overline{\beta}^1, 0)$	$(n_d^2, \overline{\beta}^2\epsilon/\rho)$	$(1, -3\rho/2)$

Table 3: Chiral massless spectrum of the semi-realistic four stack model. $(.)^c$ denotes the charge conjugated field.

Intersection	Matter	Rep.	Y
(a,b)	Q_L	$(3,2)_{(1,-1,0,0)}$	$1/6$
(a',b)	q_L	$2\times(3,2)_{(1,1,0,0)}$	$1/6$
(a,c)	$(U_R)^c$	$3\times(\overline{3},1)_{(-1,0,1,0)}$	$-2/3$
(a',c)	$(D_R)^c$	$3\times(\overline{3},1)_{(-1,0,-1,0)}$	$1/3$
(b',d)	L_L	$3\times(1,2)_{(0,-1,0,-1)}$	$-1/2$
(c,d)	$(E_R)^c$	$3\times(1,1)_{(0,0,-1,1)}$	1
(c',d)	$(N_R)^c$	$3\times(1,1)_{(0,0,1,1)}$	0

Table 4: D6-brane configuration for the three-family $\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold model.

Type	N_a	$(n_a^1,m_a^1)\times(n_a^2,m_a^2)\times(n_a^3,m_a^3)$	Gauge Group
A_1	4	$(0,1)\times(0,-1)\times(2,0)$	$Q_8, Q_{8'}$
A_2	1	$(1,0)\times(1,0)\times(2,0)$	$Sp(2)_A$
B_1	2	$(1,0)\times(1,-1)\times(1,3/2)$	$SU(2), Q_2$
B_2	1	$(1,0)\times(0,1)\times(0,-1)$	$Sp(2)_B$
C_1	3+1	$(1,-1)\times(1,0)\times(1,1/2)$	$SU(3), Q_3, Q_1$
C_2	2	$(0,1)\times(1,0)\times(0,-1)$	$Sp(4)$

Sector	$SU(3)_C \times SU(2)_Y \times Sp(2)_B \times Sp(2)_A \times Sp(4)$	$(Q_3, Q_1, Q_2, Q_8, Q_8')$	Q_Y	$Q_8 - Q_8'$	F
A_1B_1	$3 \times 2 \times (1, \overline{2}, 1, 1, 1)$	$(0, 0, -1, \pm 1, 0)$	$\pm\frac{1}{2}$	± 1	H_U
	$3 \times 2 \times (1, \overline{2}, 1, 1, 1)$	$(0, 0, -1, 0, \pm 1)$	$\pm\frac{1}{2}$	∓ 1	H_U
A_1C_1	$2 \times (\overline{3}, 1, 1, 1, 1)$	$(-1, 0, 0, \pm 1, 0)$	$\frac{1}{3}, -\frac{2}{3}$	$1, -1$	[illegible]
	$2 \times (\overline{3}, 1, 1, 1, 1)$	$(-1, 0, 0, 0, \pm 1)$	$\frac{1}{3}, -\frac{2}{3}$	$-1, 1$	[illegible]
	$2 \times (1, 1, 1, 1, 1)$	$(0, -1, 0, \pm 1, 0)$	$1, 0$	$1, -1$	[illegible]
	$2 \times (1, 1, 1, 1, 1)$	$(0, -1, 0, 0, \pm 1)$	$1, 0$	$-1, 1$	[illegible]
B_1C_1	$(3, \overline{2}, 1, 1, 1)$	$(1, 0, -1, 0, 0)$	$\frac{1}{6}$	0	
	$(1, \overline{2}, 1, 1, 1)$	$(0, 1, -1, 0, 0)$	$-\frac{1}{2}$	0	
B_1C_2	$(1, 2, 1, 1, 4)$	$(0, 0, 1, 0, 0)$	0	0	
B_2C_1	$(3, 1, 2, 1, 1)$	$(1, 0, 0, 0, 0)$	$\frac{1}{6}$	0	
	$(1, 1, 2, 1, 1)$	$(0, 1, 0, 0, 0)$	$-\frac{1}{2}$	0	
B_1C_1'	$2 \times (3, 2, 1, 1, 1)$	$(1, 0, 1, 0, 0)$	$\frac{1}{6}$	0	
	$2 \times (1, 2, 1, 1, 1)$	$(0, 1, 1, 0, 0)$	$-\frac{1}{2}$	0	
B_1B_1'	$2 \times (1, 1, 1, 1, 1)$	$(0, 0, -2, 0, 0)$	0	0	
	$2 \times (1, 3, 1, 1, 1)$	$(0, 0, 2, 0, 0)$	0	0	
A_1A_1	$3 \times 8 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	
	$3 \times 4 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, \pm 1, \pm 1)$	± 1	0	
	$3 \times 4 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, \pm 1, \mp 1)$	0	± 2	
	$3 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, \pm 2, 0)$	± 1	± 2	
	$3 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, 0, \pm 2)$	± 1	∓ 2	
A_2A_2	$3 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	
B_1B_1	$3 \times (1, 3, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	
	$3 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	
B_2B_2	$3 \times (1, 1, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	
C_1C_1	$3 \times (8, 1, 1, 1, 1)$	$(0, 0, 0, 0, 0)$	0	0	

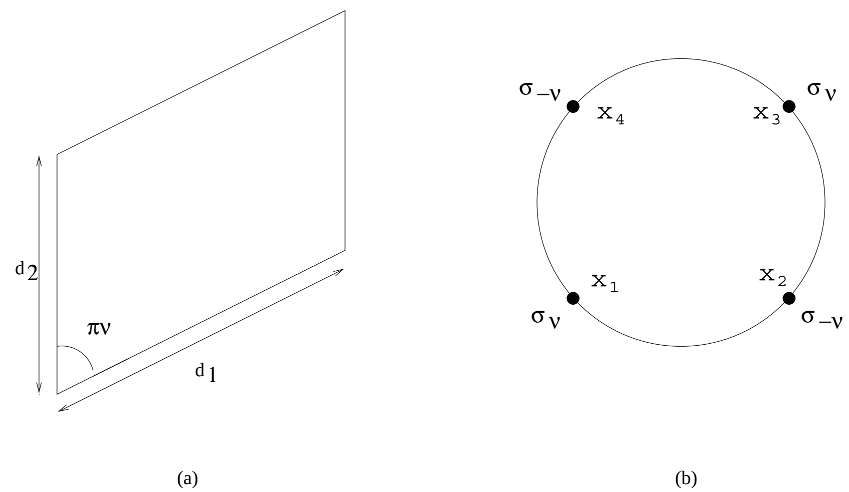

Figure 7: Target space: the intersection of two parallel branes separated by respective distances d_1 and d_2 and intersecting at angles $\pi\nu$ (Figure a). World-sheet: a disk diagram of the four twist fields located at $x_{1,2,3,4}$ (Figure b). The calculation involves a map from the world-sheet to target space.

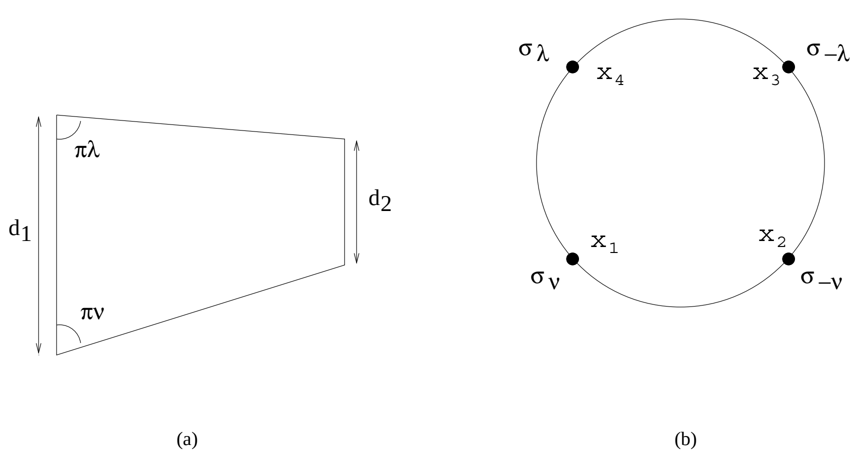

Figure 8: Target space: the intersection of two branes ~~intersecting respectively~~ with the two parallel branes at angles $\pi\nu$ and $\pi\lambda$, respectively (a). World-sheet: a disk diagram of the four twist fields located at $x_{1,2,3,4}$ (b). The calculation involves a map from the world-sheet to target space, allowing for a factorization to a three-point function.

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