Chiral 4d String Vacua with D-branes and Moduli Stabilization

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We discuss type IIB orientifolds with D-branes, and NSNS and RR field strength fluxes, with D-brane sectors leading to open string spectra with non-abelian gauge symmetry and charged chiral fermions. The closed string field strengths generate a scalar potential stabilizing most moduli. Hence the models combine the advantages of leading to phenomenologically interesting (and even semirealistic) chiral open string spectra, and of stabilizing the dilaton and most geometric moduli. We describe the explicit construction of two classes of non-supersymmetric models on $\text{T}^6$ and orbifolds/orientifolds thereof, with chiral gauge sector arising from configurations of D3-branes at singularities, and from D9-branes with non-trivial world-volume magnetic fields. The latter examples yield the chiral spectrum of just the Standard Model.

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1 Introduction

One of the most remarkable features of string theory is that, despite its complexity, it admits vacua with low-energy physics surprisingly close to the structure of observed particles and interactions. In particular there exist by now several classes of constructions (e.g. heterotic compactifications, type II models with D-branes at singularities, intersecting D-branes, compactifications of Horava-Witten theory, etc) leading to four-dimensional gravitational and non-abelian gauge interactions, with charged chiral fermions. Within each class, particular explicit models with spectrum very close to that of the (Minimal Supersymmetric) Standard Model have been constructed. On the other hand, a generic feature of all these constructions, is the existence of a (very often large) number of moduli, which remain massless in the construction, unless some supersymmetry breaking mechanism is proposed. Even in the latter case, one often encounters runaway potentials for some moduli, with no minima at finite distance in moduli space. From this viewpoint these models are relatively far from describing physics similar to the observed world.

Recently, it has been shown that, in the setup of Calabi-Yau compactification of type II string theory (or also M-theory), there exists a natural mechanism which stabilizes most moduli of the compactification. This is achieved by considering compactifications with non-trivial field strength fluxes for closed string $p$-form fields. This proposal has already been explored in different setups, leading to large classes of models with very few unstabilized moduli. Hence this mechanism is one of the most interesting recent insights to address the long-standing problems of moduli in phenomenological string models. Unfortunately, compactifications with field strength fluxes have centered on simple models, which lead to uninteresting gauge sectors, from the phenomenological viewpoint. In particular, the class of models studied are naturally non-chiral, since the corresponding gauge sectors arise from too simple stacks of parallel D-branes.

Our purpose in the present paper is to construct models which combine the interesting features from the above two approaches. In fact, we construct string compactifications with interesting 4d chiral gauge sectors and flux stabilization of (most) moduli. The models are based on introducing NSNS and RR 3-form fluxes in compactifications of type IIB theory with D-branes. We discuss two classes of models, exploiting different mechanisms to lead to chirality in the D-brane open string sector. In the first class, we construct models with D3-branes located at orbifold singularities, and chirality arises due to the orbifold projection. In the second class, we construct models with magnetised D9-branes, namely spacetime filling D9-branes with non-trivial gauge bundles on their world-volume, and chirality arises from a non-zero index of the Dirac operator in the Kaluza-Klein compactification of the charged 10d fermions (this mechanism is related, in the absence of fluxes, to the appearance of chirality at D-brane intersections via mirror symmetry). Although the complete models turn out to
be non-supersymmetric, we discuss the stability properties of the resulting models.

For simplicity we center on constructions on $T^6$ and orbifolds/orientifolds thereof. The rules we describe are however quite general and we expect that the techniques and our new observations are useful in constructing other models, supersymmetric or not, in these and other geometries.

The paper is organized as follows. In Section 2 we review the construction of toroidal orientifold compactifications with NSNS and RR 3-form fluxes. In section 3 we discuss diverse classes of D-brane configurations leading to chiral gauge sectors, and issues arising in the possible introduction of 3-form fluxes for them. In Section 4 we describe the construction of models with D3-branes at singularities and NSNS and RR 3-form fluxes, and provide an explicit example with an open string spectrum yielding a 3-family $SU(5)$ GUT. In Section 5, we describe the construction of models with D-branes with world-volume magnetic fields. We construct explicit models with 3-form fluxes and configurations of magnetised D-branes, with an open string spectrum yielding just the Standard Model, with three fermions families.

These constructions provide the first examples in a presumably rich and interesting class of realistic models. In Section 6 we make our final comments.

This work elaborates over techniques and examples in [1], see also [2] for models closely related to those of section 5.

## 2 Review of fluxes

Compactifications of type II theories (or orientifolds thereof) with NSNS and RR field strength fluxes have been considered, among others, in [3, 4, 5, 6, 7, 8, 9, 10, 11]. In this section we review properties of type IIB compactifications with 3-form fluxes.

### 2.1 Consistency conditions and moduli stabilization

Type IIB compactification on a Calabi-Yau threefold $X_6$ with non-trivial NSNS and RR 3-form field strength backgrounds $H_3, F_3$ have been extensively studied. In particular the analysis in [4, 6] provided, in a quite general setup, the consistency conditions such fluxes should satisfy. They must obey the Bianchi identities

$$dF_3 = 0 \quad dH_3 = 0$$

and they should be properly quantized, namely for any 3-cycle $\Sigma \subset X_6$

$$\frac{1}{(2\pi)^2 \alpha'} \int_\Sigma F_3 \in \mathbb{Z} \quad ; \quad \frac{1}{(2\pi)^2 \alpha'} \int_\Sigma H_3 \in \mathbb{Z} \quad .$$

We should clarify that due to the flux backreaction, the metric is not the Ricci-flat metric in the Calabi-Yau, but rather conformal to it due to a non-trivial warp factor (sourced by the fluxes and objects in the background) [4, 6].
The fluxes hence define integer 3-cohomology classes in $H^3(X_6, \mathbb{Z})$.

A subtlety in flux quantization in toroidal orientifolds was noticed in [7, 8]. Namely, if flux integrals along some 3-cycle are integer but odd, consistency requires the corresponding 3-cycle to pass through an odd number of exotic O3-planes. For simplicity we restrict to the case where all flux integrals are even integers.

An important observation [4, 6] is that, to allow for non-trivial fluxes in compactifications to 4d Minkowski space, it is crucial to include orientifold 3-planes in the compactification, so we consider type IIB orientifolds with these objects. The simplest way to understand the need of these objects is to notice the type IIB supergravity Chern-Simons coupling

$$\int_{M_4 \times X_6} H_3 \wedge F_3 \wedge C_4$$

where $C_4$ is the IIB self-dual 4-form gauge potential. This coupling implies that upon compactification the flux background contributes to a tadpole for $C_4$, with positive coefficient $N_{\text{flux}}$ (in D3-brane charge units). Moreover, fluxes contribute positively to the energy of the configuration, due to the 2-form kinetic terms. The only way to cancel these tadpoles is to introduce O3-planes, objects with negative RR $C_4$-charge and negative tension, to cancel both the RR tadpole and also to compensate the vacuum energy of the configuration.

Having O3-planes in the configuration, it is natural to consider the possibility of adding $N_{Q_3}$ explicit D3-branes as well. The RR tadpole cancellation constraint hence reads

$$N_{Q_3} + N_{\text{flux}} + Q_{O3} = 0$$

(2.4)

We normalize charge such that a D3-brane in covering space has charge +1. With this convention an O3-plane has charge $-1/2$, and

$$N_{\text{flux}} = \frac{1}{(4\pi^2 \alpha')^2} \int_{X_6} H_3 \wedge F_3 = \frac{1}{(4\pi^2 \alpha')^2} \frac{i}{2\phi_I} \int_{X_6} G_3 \wedge \overline{G}_3$$

(2.5)

where $\phi_I$ is the imaginary part of the IIB complex coupling $\phi = a + ig_s$, and

$$G_3 = F_3 - \phi H_3$$

(2.6)

Finally, in order to satisfy the equations of motion, the flux combination $G_3$ must be imaginary self-dual with respect to the Hodge operation defined in terms of the Calabi-Yau metric in $X_6$

$$s_6 G_3 = i G_3$$

(2.7)

These conditions guarantee the existence of a consistent supergravity solution for the different relevant fields in the configuration, metric, and 4-form, which have the form of a warped compactification (similar to a black 3-brane solution, since the same fields are sourced) [4, 6].

We remark that the condition (2.7) should not be regarded as an additional constraint on the fluxes. Rather, for a set of fluxes in a fixed topological sector (i.e. in a fixed cohomology
class), it is a condition on the scalar moduli which determine the internal metric. The scalar potential is minimized at points in moduli space where (2.7) is satisfied, while fluxes induce a positive scalar potential at other points. Hence introduction of fluxes leads to a natural mechanism to stabilize moduli. Explicit expressions will be discussed later on, for the moment let us state that generically the dilaton and all complex structure moduli are stabilized by this mechanism. On the other hand, Kahler moduli are not stabilized [4, 6, 7, 8].

2.2 Supersymmetry

The conditions for a configuration with 3-form fluxes to preserve some supersymmetry have been studied in [12], and applied in explicit constructions in [7, 8, 11].

The 10d $N = 2$ type IIB real supersymmetry transformation parameters $\epsilon_L, \epsilon_R$, can be gathered into a complex one $\epsilon = \epsilon_L + i \epsilon_R$. It is chiral in 10d, satisfying $\Gamma_{10d}\epsilon = -\epsilon$, with $\Gamma_{10d} = \gamma^0 \ldots \gamma^9$. Compactification on $X_6$ splits this spinor with respect to $SO(6) \times SO(4)$ e.g. as

$$\epsilon_L = \xi \otimes u + \xi^* \otimes u^*$$  \hspace{1cm} (2.8)

where $\xi$ is a 6d chiral spinor $\Gamma_{6d}\xi = -\xi$, and $u$ is a 4d chiral spinor $\Gamma_{4d}u = u$. For $X_6$ of generic $SU(3)$ holonomy only one component of $\xi$ is covariantly constant and provides susy transformations in 4d.

On the other hand, the presence of the O3-planes and D3-branes in the background preserves only those $\epsilon$ satisfying

$$\epsilon_R = -\gamma^4 \ldots \gamma^9 \epsilon_L$$  \hspace{1cm} (2.9)

Such spinors are of the form $\epsilon = 2\xi \otimes u$.

The conditions for a flux to preserve a supersymmetry associated to a particular spinor component of $\xi$ are [12]

$$G\xi = 0 \ ; \ G\xi^* = 0 \ ; \ G\gamma^m\xi^* = 0$$  \hspace{1cm} (2.10)

where $G = \frac{1}{6} G_{mnl}[\gamma^m \gamma^n \gamma^l]$.

Let us introduce complex coordinates $z^i, \bar{z}^i$ and define the highest weight state $\xi_0$ satisfying $\gamma^i\xi_0 = 0$. Then the O3-planes preserve $\xi_0$ and $\gamma^{ij}\xi_0$. Of these, a general Calabi-Yau (on which $z_i$ are complex coordinates) preserves only $\xi_0$, since it is $SU(3)$ invariant.

The conditions that a given flux preserves $\xi_0$, can be described geometrically [12, 7, 8] as

a) $G_3$ is of type (2, 1) in the corresponding complex structure.

b) $G_3$ is primitive, i.e. $G_3 \wedge J = 0$ where $J$ is the Kahler form.

For explicit discussion of these conditions see below. Notice that a $G_3$ flux which is not (2, 1) in a complex structure, may still be supersymmetric if it preserves other spinor $\xi'_0$. 

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(although it does not preserve $\xi_0$). In such case, $G_3$ would be of type $(2,1)$ in a different complex structure where $\xi_0'$ is the spinor annihilated by the new $\gamma^7$. In a general Calabi-Yau, however, there is a prefered complex structure, so that a supersymmetric flux should obey the above conditions with respect to it.

Since the techniques to find consistent (possibly supersymmetric) fluxes at particular values of the stabilized moduli (and vice versa) have been discussed in the literature, we will not delve into their discussion.

3 Chiral D-brane configurations

The models considered in [8, 7] succeed in leading to $N = 1$ or non-supersymmetric low-energy theories, with stabilization of most moduli. However, they are unrealistic in that they are automatically non-chiral since the only gauge sectors live on parallel D3-branes whose low-energy spectra are non-chiral.\footnote{The D3-brane world-volume spectra are at best $N = 1, 0$ deformations of $N = 4$ theories, by flux-induced operators breaking partially or totally the world-volume supersymmetry}

We are interested in constructing models containing gauge sectors with charged chiral fermions, and a bulk with flux-induced moduli stabilization. There are several possibilities to do this, corresponding to the different ways to build configurations of D-branes containing chiral fermions. Equivalently, the different ways of enriching the simple configuration of parallel D3-branes, to lead to chiral open string spectra.

D3-branes at singularities

One possibility is to use compactification varieties containing singular points, e.g. orbifold singularities. Locating D3-branes at the singularity leads to chiral gauge sectors, with low energy spectrum given by a quiver diagram [13].\footnote{For phenomenological model building in this setup with no field strength fluxes, see [14].} A simple set of models can be constructed starting with type IIB theory on $T^6/\mathbb{Z}_3$ modded out by the $\Omega R$ orientifold projection, where $\Omega$ is worldsheet parity and $R : z_i \to -z_i$. This is particularly promising, since it is the simplest orbifold which can lead to three families in the sector of D3-branes at singularities. However, it is not possible to obtain $N = 1$ supersymmetric models in this setup, for the following reason. In the complex structure where the spinor invariant under $\mathbb{Z}_3$ satisfies $\gamma^7 \xi_0 = 0$, the $\mathbb{Z}_3$ orbifold action reads

$$\theta : (z_1, z_2, z_3) \to (e^{2\pi i/3} z_1, e^{2\pi i/3} z_2, e^{-4\pi i/3} z_3)$$

Fluxes preserving the same spinor $\xi_0$ should be of type $(2,1)$ in this complex structure, namely linear combinations of $dz_1^1 dz_2^2 dz_3^3$, $dz_1^1 dz_2^2 dz_3^3$ and $dz_1^1 dz_2^2 dz_3^3$. Such fluxes are not invariant under the orbifold action, and cannot be turned on. In other words, the only possible
fluxes are not supersymmetric. The same problem arises for other promising orientifolds, like $Z_3 \times Z_2 \times Z_2$.

So we will not pursue the construction of $N=1$ susy models in the setup of D3-branes at singularities, and center on non-supersymmetric models. An amusing possibility is to build models where the D-brane sector preserves some supersymmetry, while the closed string sector is non-supersymmetric due to the combination of fluxes and orbifold action. We provide an example of this kind in section 4.

**Magnetised D-branes**

Although it is not often explicitly stated, it is also possible to obtain chiral fermions from wrapped parallel D-branes, if the geometry of the wrapped manifold or the topology of the internal world-volume gauge bundle are non-trivial. The chiral fermions arise from the Kaluza-Klein reduction of the higher-dimensional worldvolume fermions if the index of the corresponding Dirac operator is non-zero.

The simplest such setup is type IIB compactified on $T^6$ (with an additional orientifold, and possibly orbifold projections), with configurations of D9-branes spanning all of spacetime (namely wrapped on $T^6$) with non-trivial gauge bundles on $T^6$. The simplest bundles are given by constant magnetic fields on each of the $T^2$, so we call them magnetised D-brane configurations. They have been considered in [15] in the absence of closed string field strength fluxes. Magnetised D9-branes are related to D3-branes in that, due to world-volume Chern-Simons couplings, the non-trivial magnetic fields induce non-zero D3-brane charge on them (indeed, in the limit of very large magnetic fields, their charges reduce to those of a set of parallel D3-branes).

In the absence of fluxes, these configurations are related by T-duality to configurations of intersecting D6-branes, so that any model of the latter kind can be easily translated [17, 18] to a magnetised D9-brane setup. The advantage of using the magnetised D9-brane picture with O3-planes is that it is now straightforward to include NSNS and RR fluxes in the configuration, by applying the tools reviewed above (for the situation without D-branes). Notice that in the T-dual version of intersecting branes this corresponds to turning on a complicated set of NSNS, RR and metric fluxes, see below. For the class of fluxes we consider, the picture of magnetised D9-branes is more useful.

In section 5 we review magnetised D9-brane configurations, first without NSNS and RR fluxes, and then discuss the introduction of the latter, and provide some explicit examples with interesting chiral gauge sectors.

**Intersecting D6-branes**

Much progress has been made in D-brane model building using type IIA D6-branes wrapped on intersecting 3-cycles in an internal space (e.g. [16, 17, 19, 20, 21, 22, 23, 24])

\[^{4}\text{See [25] for early work leading to non-chiral models, and [26] for reviews.}\]
However it is difficult to introduce NSNS and RR field strength fluxes in those setups; the difficulties can be seen from different perspectives. The models usually contain O6-planes, hence we need combinations of fluxes which source the RR 7-form. A possible combination is the RR 0-form field strength of type IIA (i.e. a cosmological constant of massive type IIA) and the NSNS 3-form field strength. This combination of fluxes has not been studied in the literature, so it is not a convenient starting point (see [27] for preliminary non-chiral results in this direction).

One may think that T-dualizing three times a model with O3-planes and RR and NSNS fluxes would yield the desired configurations with O6-planes. However, T-duality acts in a very non-trivial way on $H_3$, transforming some of its components into non-trivial components of the T-dual metric, which is no longer Calabi-Yau [28, 29]. The final configuration indeed would contain fluxes (RR and ‘metric fluxes’) which source the RR 7-form, and would lead to Poincare invariant 4d models (consistently with T-duality). However, a full analysis of the T-dual geometry, and its properties (e.g. possible 3-cycles on which to wrap D6-branes) is lacking.

Hence we will not attempt to discuss intersecting D-brane models with NSNS and RR fluxes.

4 NSNS and RR fluxes in a type IIB orientifold with D3-branes at singularities.

4.1 D3-branes at singularities

We would like to consider compactifications of type IIB theory, of the form $M_4 \times X_6$, with D3-branes located at singular points in the transverse space. Namely, their world-volume is $M_4 \times P$, with $P$ a singular point in $X_6$. Since we are interested in the massless states in the open string spectrum, these are only sensitive to the local geometry around the singular point $P$, and are insensitive to the global structure of the compactification. Hence, it is enough to consider a local model for the geometry around $P$.

A simple class of singularities, on which string theory can be quantized exactly in $\alpha'$ (and hence, on which we know the open string spectrum on D3-branes at the singularity) are orbifold singularities. For instance, an orbifold singularity $C^3/\mathbb{Z}_N$ is obtained from flat 6d space $\mathbb{R}^6 = C^3$ by identifying points related by the order $N$ action

$$\theta : (z_1, z_2, z_3) \rightarrow (e^{2\pi i a_1/N} z_1, e^{2\pi i a_2/N} z_2, e^{2\pi i a_3/N} z_3)$$

which generates $\mathbb{Z}_N$. Here $a_i \in \mathbb{Z}$, and $\sum a_i = \text{even}$, so that the space is spin. We will consider supersymmetry preserving geometries, which requires $\sum a_i \in 2N\mathbb{Z}$.

The origin in $C^3$ descends in the quotient to a conical singular point, at which we locate the D3-branes. The spectrum on a stack of $K$ D3-branes at the $C^3/\mathbb{Z}_N$ singularity is obtained
from the spectrum of $K$ D3-branes in flat space (which is 4d $N=4$ supersymmetric $U(K)$ Yang-Mills), by keeping states invariant under the $\mathbb{Z}_N$ action. In the latter, one should take into account the (Chan-Paton) action of $\mathbb{Z}_N$ on the gauge degrees of freedom, given by conjugation by a $K \times K$ matrix

$$\gamma_\theta = \text{diag} \left( 1_{n_0}, e^{2\pi i/N}1_{n_1}, \ldots, e^{2\pi i(N-1)/N}1_{n_{N-1}} \right)$$  \hspace{1cm} (4.2)$$

where $\sum_a n_a = K$. The $n_a$ are moreover constrained by certain consistency conditions, the twisted RR tadpole cancellation conditions. A simple (but not always unique) choice satisfying them is $n_r = k$ and hence $K = Nk$.

The computation of the spectrum has been discussed and gives the following $N=1$ supermultiplet content [13, 14]

\begin{align*}
N = 1 \text{ Vector Multiplet} & \quad \prod_a U(n_a) \\
N = 1 \text{ Chiral Multiplet} & \quad \sum_{a=1}^N \left[ (\Box, \Box_{n+a_1}) + (\Box, \Box_{n+a_2}) + (\Box, \Box_{n+a_3}) \right] \\
\end{align*}

A concrete nice example is the $\mathbb{C}^3/\mathbb{Z}_3$ singularity, with $(a_1, a_2, a_3) = (1, 1, -2)$, and which leads to a spectrum with a triplicated set of chiral multiplets (and hence of chiral fermions).

\begin{align*}
N = 1 \text{ Vector Multiplet} & \quad U(k) \times U(k) \times U(k) \\
N = 1 \text{ Chiral Multiplet} & \quad 3 \left[ (\Box, \Box) + (\Box, \Box) + (\Box, \Box) \right] \\
\end{align*}

For future convenience, we simply point out that for D3-branes at a singularity, mapped to itself under an orientifold action, an additional projection should be imposed. For such an orientifold of the $\mathbb{C}^3/\mathbb{Z}_3$ singularity, the spectrum is

\begin{align*}
N = 1 \text{ Vector Multiplet} & \quad SO(k-4) \times U(k) \\
N = 1 \text{ Chiral Multiplet} & \quad 3 \left[ (\Box, \Box) + (1, \Box) \right] \\
\end{align*}

The analysis in this section applies both to models without and with fluxes. In the presence of the latter, the fluxes do not change the chiral structure of the theory, but may lead to additional operators in the D3-brane world-volume field theory (like e.g. gaugino masses, etc). These are computable in the regime of dilute fluxes, namely in the large volume limit [30].

\section*{4.2 Models with fluxes: A 3-family $SU(5)$ GUT example}  

In this section we discuss the construction of a compactification with fluxes, and D3-branes at an orbifold singularity. Consider type IIB on $\mathbb{T}^6$ modded by the orientifold action $\Omega R$, and quotient by the above $\mathbb{Z}_3$ action

$$\theta : (z_1, z_2, z_3) \rightarrow (e^{2\pi i/3}z_1, e^{2\pi i/3}z_2, e^{-4\pi i/3}z_3)$$  \hspace{1cm} (4.3)$$

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This is a T-dual version of the model in [31].

Now introduce NSNS and RR 3-form field strength fluxes. As discussed in previous section, the models turn out be non-supersymmetric in the closed string sector, since the orbifold and the $\mathbb{Z}_3$ invariant fluxes necessarily preserve different supersymmetries. Introduce a 3-form flux $^5$

$$G_3 = 2dz_1dz_2dz_3$$ (4.4)

Such fluxes have been considered in $T^6/\Omega R$ in [7, 8]. In this geometry all geometric moduli are Kahler, hence they are not stabilized. The above flux however, stabilizes the dilaton, which is fixed at a value $\phi = e^{2\pi i/3}$. The flux is imaginary self-dual with respect to the underlying metric. For the configuration to make sense it is crucial that the flux (4.4) is invariant under the action of $\theta$. It is also important to notice that the $\mathbb{Z}_3$ quotient of $T^6/(\Omega R)$ does not contain closed 3-cycles which are not closed in $T^6/\Omega R$, hence proper quantization is not spoilt. This is because the collapsed cycles at $\mathbb{Z}_3$ singularities are 2- and 4-cycles, hence do not impose additional quantization constraints.

Notice that this flux preserves some of the supersymmetries of the underlying $T^6/(\Omega R)$ geometry, namely the spinors $\gamma^i \xi_0$. These are however broken by the orbifold projection. It is possible that this kind of breaking of supersymmetry has some particularly nice features, since the interactions between untwisted modes is sensitive to supersymmetry breaking only via effects involving twisted modes. It would be interesting to analyze the impact of this property on the violations of the no-scale structure of the low energy supergravity effective theory for these models (i.e. the degree of protection against $\alpha'$ or $g_s$ corrections).

For the above flux we have $N_{\text{flux}} = 12$, hence cancellation of RR tadpoles requires the introduction of 20 D3-branes. To define the model completely, we need to specify the configuration of these 20 D3-branes. In the $\Omega R$ orientifold of $T^6$ there is one point, the origin $(0,0,0)$, fixed under $\Omega R$ and $\theta$. At this point, cancellation of RR twisted tadpoles requires the presence of D3-branes, with a Chan-Paton matrix satisfying

$$\text{Tr} \gamma_\theta = -4$$ (4.5)

In addition, there are other 26 points fixed under $\theta$ (and gathered in 13 pairs under $\Omega R$), where there is no twisted RR tadpole. If D3-branes are present, they should have traceless Chan-Paton matrix. Finally there are 63 points fixed under $\Omega R$ (gathered in 21 trios under $\mathbb{Z}_3$) at which we may locate any number (even or odd) of D3-branes.

A simple solution would be to locate the 20 D3-branes at the origin, with

$$\gamma_\theta = \text{diag} (1_4, e^{2\pi i/3}1_8, e^{4\pi i/3}1_8)$$ (4.6)

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$^5$In what follows we absorb the normalization factor $1/(4\pi^2\alpha')^2$ appearing in 2.5 in the definition of the fluxes.
leading to an $N = 1$ supersymmetric sector (to leading approximation, since interactions with the closed sector would transmit supersymmetry breaking), with spectrum (4.1)

\[
\begin{align*}
N &= 1 \text{ Vector Multiplet} & SO(4) \times U(8) \\
N &= 1 \text{ Chiral Multiplet} & 3 \left[(4, \tilde{8}) + (1, 28)\right]
\end{align*}
\]

A more interesting possibility, which we adapt from [32], is to locate 11 D3-branes at the origin, with

\[
\gamma_\theta = \text{diag}(1, e^{2\pi i/3}1_5, e^{4\pi i/3}1_5)
\]

This leads to a gauge sector with

\[
\begin{align*}
N &= 1 \text{ Vector Multiplet} & U(5) \\
N &= 1 \text{ Chiral Multiplet} & 3[\tilde{5} + 10]
\end{align*}
\]

One should now be careful in locating the additional D3-branes. We have introduced an odd number of D3-branes on top of the O3-plane at the origin, and this implies that it is an $\tilde{O}3^-$-plane in notation of [33], i.e. there exists a $\mathbb{Z}_2 B_{RR}$ background on an $\mathbb{RP}_2$ around the O3-plane. In order to be consistent with the fact that our flux has even integral over the different 3-cycles implies that there should exist other $\tilde{O}3^-$-planes in the configuration. The conditions in [7, 8] state that for any 3-plane over which the integrated flux of $H_3$ is even (any 3-plane in our case), the number of $\tilde{O}3^-$-planes must be even. Here we consider a configuration, appeared in a model without fluxes [32], which satisfies this constraint.

Let us denote $A$, $B$ or $C$ the coordinate of an O3-plane in a complex plane, according to whether $z_i = 1/2$, $z_i = e^{2\pi i/3}/2$ or $z_i = (1 + e^{2\pi i/3})/2$. The remaining 9 D3-branes in the model are located on top of O3-planes at the points

\[
\begin{align*}
(A, A, A) & \quad (B, B, C) & \quad (C, C, B) \\
(A, 0, 0) & \quad (B, 0, 0) & \quad (C, 0, 0) \\
(0, A, 0) & \quad (0, B, 0) & \quad (0, C, 0)
\end{align*}
\]

(4.8)

This set is invariant under exchange of fixed points by $\mathbb{Z}_3$, and introduces the right number of $\tilde{O}3^-$-planes at the right places. The additional D3-branes do not lead to additional gauge symmetries.

Thus the final model contains a 3-family $SU(5)$ GUT gauge sector (although without adjoint chiral multiplets to break it down to the Standard Model), as the only gauge sector of the theory. In addition, its closed string sector is non-supersymmetric, but the cosmological constant vanishes at leading order. The impact of the supersymmetry breaking on the gauge sector is computable in the large volume limit [30]. It would also be interesting to construct other models based on the $\mathbb{Z}_3$ orbifold, or other orbifold models. We leave these interesting question for future work.
5 Magnetised D-branes

In this section we describe configurations of magnetised D9-branes, first before the introduction of fluxes, and then introducing them. We first consider the case of toroidal compactifications and orientifolds thereof. These models, in the absence of bulk fluxes, are T-dual to the models of intersecting D6-branes in toroidal compactifications [19], toroidal orientifolds [17] and orbifolds [23]. Subsequently we describe the introduction of fluxes and provide two explicit examples with the gauge symmetry and chiral fermion content of just the Standard Model.

5.1 Magnetised D-branes in toroidal compactifications

We start with the simple case of toroidal compactification, with no orientifold projection. Consider the compactification of type IIB theory on $\mathbf{T}^6_6$, assumed factorizable.

We consider sets of $N_a$ D9-branes, labelled $D9_a$-branes, wrapped $m^i_a$ times on the $i^{th}$ 2-torus ($\mathbf{T}^2_i$) in $\mathbf{T}^6$, and with $n^i_a$ units of magnetic flux on ($\mathbf{T}^2_i$). Namely, we turn on a world-volume magnetic field $F_a$ for the $U(1)_a$ gauge factor in $U(N_a)$, such that

$$m^i_a \frac{1}{2\pi} \int_{\mathbf{T}^2_i} F^i_a = n^i_a$$

Hence the topological information about the D-branes is encoded in the numbers $N_a$ and the pairs $(m^i_a, n^i_a)$.

This description automatically includes other kinds of lower dimensional D-branes. For instance, a D7-brane (denoted $D7_{(i)}$) sitting at a point in $\mathbf{T}^2_i$ and wrapped on the two remaining two-tori (with generic wrapping and magnetic flux quanta) is described by $(m^i, n^i) = (0, 1)$ (and arbitrary $(m^j, n^j)$ for $j \neq i$); similarly, a D5-brane (denoted $D5_{(i)}$) wrapped on $\mathbf{T}^4_i$ (with generic wrapping and magnetic flux quanta) and at a point in the remaining two 2-tori is described by $(m^j, n^j) = (0, 1)$ for $j \neq i$; finally, a D3-brane sitting at a point in $\mathbf{T}^6$ is described by $(m^i, n^i) = (0, 1)$ for $i = 1, 2, 3$. This is easily derived by noticing that the boundary conditions for an open string ending on a D-brane wrapped on a two-torus with magnetic flux become Dirichlet for (formally) infinite magnetic field.

D9-branes with world-volume magnetic fluxes are sources for the RR even-degree forms, due to their worldvolume couplings

$$\int_{D9_a} C_{10} \ ; \ \int_{D9_a} C_8 \wedge \text{tr} F_a \ ; \ \int_{D9_a} C_6 \wedge \text{tr} F_a^2 \ ; \ \int_{D9_a} C_4 \wedge \text{tr} F_a^3$$

Consistency of the configuration requires RR tadpoles to cancel. Following the discussion in [19], leads to the conditions

$$\sum_a N_a m^1_a m^2_a m^3_a = 0$$

\[\text{Notice the change of roles of } n \text{ and } m \text{ as compared with other references. This however facilitates the translation of models in the literature to our language.}\]
\[ \sum_a N_a m_a^1 n_a^2 n_a^3 = 0 \quad \text{and permutations of 1, 2, 3} \]
\[ \sum_a N_a m_a^1 n_a^2 n_a^3 = 0 \quad \text{and permutations of 1, 2, 3} \]
\[ \sum_a N_a n_a^1 n_a^2 n_a^3 = 0 \]  \hspace{1cm} (5.3)

Which amounts to cancelling the D9-brane charge as well as the induced D7-, D5- and D3-brane charges.

Introducing for the \( i \)th 2-torus the even homology classes \([0]_i\) and \([T^2]_i\) of the point and the two-torus, the vector of RR charges of one D9-brane in the \( a \)th stack is
\[ [Q_a] = \prod_{i=1}^{3} (m^i_a[T^2]_i + n^i_a[0]_i) \]  \hspace{1cm} (5.4)

The RR tadpole cancellation conditions (5.3) read
\[ \sum_a N_a [Q_a] = 0 \]  \hspace{1cm} (5.5)

The conditions that two sets of D9-branes with worldvolume magnetic fields \( F^i_a \), \( F^i_b \) preserve some common supersymmetry can be derived from [16]. Indeed, it is possible to compute the spectrum of open strings stretched between them and verify that it is supersymmetric if
\[ \Delta^1_{ab} \pm \Delta^2_{ab} \pm \Delta^3_{ab} = 0 \]  \hspace{1cm} (5.6)
for some choice of signs (in order to preserve \( \xi_0 \), the choice should be all positive signs). Here
\[ \Delta_i = \arctan [(F^i_a)^{-1}] - \arctan [(F^i_b)^{-1}] \]  \hspace{1cm} (5.7)
and
\[ F^i_a = \frac{n^i_a}{m^i_a R_{x_i} R_{y_i}} \]  \hspace{1cm} (5.8)
which follows from (5.1).

The spectrum of massless states is easy to obtain. The sector of open strings in the \( a a \) sector leads to \( U(N_a) \) gauge bosons and superpartners with respect to the 16 supersymmetries unbroken by the D-branes. In the \( a b + b a \) sector, the spectrum is given by \( I_{ab} \) chiral fermions in the representation \((N_a, \overline{N}_b)\), where
\[ I_{ab} = [Q_a] \cdot [Q_b] = \prod_{i=1}^{3} (n^i_a m^i_b - m^i_a n^i_b) \]  \hspace{1cm} (5.9)
is the intersection product of the charge classes, which on the basic classes \([0]_i\) and \([T^2]_i\) is given by the bilinear antisymmetric form
\[
\begin{pmatrix}
0 & -1 \\
1 & 0 \\
\end{pmatrix}
\]  \hspace{1cm} (5.10)
The above multiplicity can be computed using the $\alpha'$-exact boundary states for these D-branes [17], or from T-duality with configurations of intersecting D6-branes. We now provide an alternative derivation which remains valid in more complicated situations where the worldsheet theory is not exactly solvable. Consider for simplicity a single two-torus. We consider two stacks of $N_a$ and $N_b$ branes wrapped $m_a$ and $m_b$ times, and with $n_a$, $n_b$ monopole quanta. Consider the regime where the two-torus is large, so that the magnetic fields are diluted and can be considered a small perturbation around the vacuum configuration. In the vacuum configuration, open strings within each stack lead to a gauge group $\text{U}(N_a m_a)$ and $\text{U}(N_b m_b)$ respectively, which is subsequently broken down to $\text{U}(N_a) \times \text{U}(N_b)$ by the monopole background, via the branching

$$U(N_a m_a) \times U(N_b m_b) \rightarrow U(N_a)^{m_a} \times U(N_b)^{m_b} \rightarrow U(N_a) \times U(N_b)$$

(5.11)

Open $ab$ strings lead to a chiral 10d fermion transforming in the bifundamental $(\mathbf{a}, \mathbf{b})$ of the original $U(N_a m_a) \times U(N_b m_b)$ group. Under the decomposition (5.11) the representation splits as

$$(\mathbf{a}, \mathbf{b}) \rightarrow (\mathbf{a}, \ldots; \mathbf{b}, \ldots) \rightarrow m_a m_b (\mathbf{a}, \mathbf{b})$$

(5.12)

The 8d theory contains chiral fermions arising from these, because of the existence of a nonzero index for the internal Dirac operator (coupled to the magnetic field background). The index is given by the first Chern class of the gauge bundle to which the corresponding fermions couples. Since it has charges $(+1, -1)$ under the $a^{th}$ and $b^{th}$ $U(1)$'s, the index is

$$\text{ind} \, \hat{D}_{ab} = \int_{T^2} (F_a - F_b) = \frac{n_a}{m_a} - \frac{n_b}{m_b}$$

(5.13)

Because of the branching (5.12), a single zero mode of the Dirac operator gives rise to $m_a m_b$ 8d chiral fermions in the $(\mathbf{a}, \mathbf{b})$ of $U(N_a) \times U(N_b)$. The number of chiral fermions in the 8d theory in the representation $(\mathbf{a}, \mathbf{b})$ of the final group is given by $m_a m_b$ times the index, namely

$$I_{ab} = m_a m_b \int_{T^2} (F_a - F_b) = n_a m_b - n_a n_b$$

(5.14)

The result (5.9) is a simple generalization for the case of compactification on three two-tori.

Notice that the field theory argument to obtain the spectrum is valid only in the large volume limit. However, the chirality of the resulting multiplets protects the result, which can therefore be extended to arbitrarily small volumes. This kind of argument will be quite useful in the more involved situation with closed string field strength fluxes, where we do not have a stringy derivation of the results.

### 5.2 Magnetised D-branes in toroidal orientifolds

We are interested in adding orientifold planes into this picture, since they are required to obtain supersymmetric fluxes. Consider type IIB on $\mathbf{T}^6$ (with zero NSNS B-field) modded
out by $\Omega R$, with $R : z_i \rightarrow -z_i$. This introduces 64 O3-planes, which we take to be all $O3^-$. It also requires the D9-brane configuration to be $\mathbb{Z}_2$ invariant. Namely, for the $N_a$ D9$_a$-brane with topological numbers $(m^i_a, n^i_a)$ we need to introduce their $N_a \Omega R$ images D9$_a'$ with numbers $(-m^i_a, n^i_a)$.

The RR tadpole cancellation conditions read

$$\sum_a N_a [Q_a] + \sum_a N_a [Q_{a'}] - 32 [Q_{O3}] = 0$$

(5.15)

with $[Q_{O3}] = [0]_1 \times [0]_2 \times [0]_3$. More explicitly

$$\sum_a N_a m^1_a m^2_a n^3_a = 0 \quad \text{and permutations of } 1, 2, 3$$

$$\sum_a N_a n^1_a n^2_a n^3_a = 16$$

(5.16)

Namely, cancellation of induced D7- and D3-brane charge. Notice that there is no net D9- or D5-brane charge, in agreement with the fact that the orientifold projection eliminates the corresponding RR fields.

The rules to obtain the spectrum are similar to the above ones, with the additional requirement of imposing the $\Omega R$ projections. This requires a precise knowledge of the $\Omega R$ action of the different zero mode sectors. The analysis is simplest in terms of the T-dual description, where it amounts to the geometric action of the orientifold on the intersection points of the D-branes. The result, which is in any case can be derived in the magnetised brane picture, is as follows [17].

The $aa$ sector is mapped to the $a'a'$ sector, hence suffers no projection. We obtain a 4d $U(N_a)$ gauge group, and superpartners with respect to the $N = 4$ supersymmetry unbroken by the brane.

The $ab + ba$ sector is mapped to the $b'a' + a'b'$ sector, hence does not suffer a projection. We obtain $I_{ab}$ 4d chiral fermions in the representation $(\Box_a, \Box_b)$). Plus additional scalars which are massless in the susy case, and tachyonic or massive otherwise.

The $ab' + b'a$ sector is mapped to the $ba' + a'b$. It leads to $I_{ab'}$ 4d chiral fermions in the representation $(\Box_b, \Box_a)$ (plus additional scalars).

The $aa' + a'a$ sector is invariant under $\Omega R$, so suffers a projection. The result is $n_{\Box_a}$ and $n_{\Box_b}$ 4d chiral fermions in the $\Box_a, \Box_b$ representations, resp, with

$$n_{\Box_a} = \frac{1}{2} (I_{aa'} + 8 I_{a,O3}) = -4 m^1_a m^2_a m^3_a (n^1_a n^2_a n^3_a + 1)$$

$$n_{\Box_b} = \frac{1}{2} (I_{a'a} - 8 I_{a,O3}) = -4 m^1_a m^2_a m^3_a (n^1_a n^2_a n^3_a - 1)$$

(5.17)

where $I_{a,O3} = [Q_a] \cdot [Q_{O3}]$.

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7There is also an additional discrete constraint, which we skip for simplicity, see footnote 10 in [1].

8We do not consider branes for which $a = a'$ here.
5.3 A model with fluxes and D-branes, with just the Standard Model spectrum

In this section we consider the introduction of fluxes, constructing two explicit models with the chiral content of the Standard Model.

Consider type IIB theory on $T^6$ and mod out by the orientifold action $\Omega R$. Let us introduce the flux.

$$G_3 = \frac{2}{\sqrt{3}} e^{-\pi i/6} (dz_1 d\bar{z}_2 dz_3 + dz_1 d\bar{z}_2 dz_3 + dz_1 dz_2 d\bar{z}_3)$$

(5.18)

This flux [8] stabilizes the dilaton at $\phi = e^{2\pi i/3}$, and the $T^6$ geometry at a factorized product of three $T^2$, all with complex structure parameter $\tau = e^{2\pi i/3}$. The flux is manifestly $(2,1)$, and is supersymmetric for the subset of Kahler moduli satisfying the primitivity condition. The above choice of factorized geometry corresponds to $J = \sum A_i dz_i d\bar{z}_i$ and makes the flux supersymmetric.

The flux contributes to the 4-form tadpoles with $N_{\text{flux}} = 12$ units. There is also the contribution of $-32$ arising from the 64 O3-planes.

In order to cancel the remaining tadpole, we introduce a set of branes, which will contain the chiral gauge sector. We introduce $N_a$ D9$_a$-branes wrapping $m^i_a$ times on the $(T^2)_i$, and with $n^i_a$ units of world-volume magnetic flux on them.

In order to build interesting examples, we would like to choose a set of branes with the chiral spectrum of just the Standard Model $^9$. Sets of branes of this kind were described in [21], mainly using the T-dual picture of D6-branes at angles. We may however use a suitable translation of their result. We take a configuration contributing 20 units to the 4-form tadpole (and zero to others), for instance the choice

$$\rho = 1, \beta^1 = \beta^2 = 1, \epsilon = 1, n^a_2 = n^2_d = 2, n^1_b = 1, n^1_c = 4$$

(5.19)

in table 2 in [21] $^{10}$. This leads to

| $N_a$ | $(m^1_a, n^1_a)$ | $(m^2_a, n^2_a)$ | $(m^3_a, n^3_a)$ |
|-------|-----------------|-----------------|-----------------|
| 3     | (0, 1)          | (1, 2)          | (1/2, 1)        |
| 2     | (-1, 1)         | (0, 1)          | (3/2, 1)        |
| 1     | (3, 4)          | (0, 1)          | (1, 0)          |
| 1     | (0, 1)          | (-1, 2)         | (3/2, 1)        |

$^9$There are two subtleties at this point. The first is that the set of consistent D-branes in a configuration with flux is classified by a modified K-theory group. The consistency of the configurations considered is discussed in appendix B in [1]. The second is that our models contain D7-branes, whose moduli are also expected to be stabilized by the 3-form fluxes [9]. Hence, there may be additional restrictions, not taken into account, on the structure of the D7-brane stacks in our examples below.

$^{10}$This requires introducing a non-zero $B_{NSNS}$ field in the third two-torus. This has been discussed in section 6 of [1].
This choice cancels all tadpoles. It leads to gauge group $U(3) \times U(2) \times U(1)^2$. It furthermore reduces to $SU(3) \times SU(2) \times U(1)$ after taking into account the $B \wedge F$ couplings; our choice of parameters is such that the $U(1)$ remaining massless precisely corresponds to hypercharge\footnote{Coupling of $B$ fields to closed string sector $U(1)$'s\cite{35} could have modified the structure of the surviving $U(1)$, as compared with the situation in\cite{21} (see also\cite{34}). However, we have checked that the fields coupling to world-volume $U(1)$'s do not couple to closed string $U(1)$'s, so the condition for a massless hypercharge in\cite{21} remains valid.}. Finally, it leads to chiral fermions multiplicites given by

$$I_{ab} = 1, I_{ab'} = 2, I_{ac} = -3, I_{ac'} = -3,$$

$$I_{bd} = 0, I_{bd'} = -3, I_{cd} = -3, I_{cd'} = 3 \quad (5.20)$$

(others are zero), which leads to exactly the chiral spectrum of the Standard Model. It is important to point out that since $\prod_i m_a^i = 0$, there are really no D9-branes, and the Wess-Zumino terms of \cite{36} are not present. The dynamical fermion content is non-anomalous.

The main features of this model are: The closed string sector is supersymmetric, while the open string sector is not. Open string tachyons are however avoided by choosing particular regions in the Kahler moduli space. Nevertheless, the model may lead to runaway potentials for Kahler moduli when the brane tension is taken into account, unless some Kahler moduli stabilization mechanism is included in the construction.

There is another interesting kind of model that we would like to describe, and which is very similar in some respects to the models leading to deSitter vacua in string theory\cite{37}. The idea is to introduce too much flux, so that one overshoots the O3-plane RR tadpole, and to introduce a set of D-branes carrying net anti-D3-brane charge to cancel it. The extra contribution from the antibrane tension leads to a positive vacuum energy, which may turn the non-compact geometry into a deSitter space. This requires in addition a mechanism to stabilize the Kahler moduli, and to avoid a runaway behaviour instead of a deSitter vacuum, on which we will not enter. Hence our model is only reminiscent of\cite{37}, and not an explicit realization of their proposal. It is however interesting to describe configurations of this kind.

Let us consider a flux similar to the above, but with larger quanta

$$G_3 = 2 \times \frac{2}{\sqrt{3}} e^{-\pi i/6} (d\bar{z}_1 dz_2 dz_3 + dz_1 d\bar{z}_2 dz_3 + dz_1 dz_2 d\bar{z}_3) \quad (5.21)$$

This flux stabilizes moduli at the same values as the above one, but leads to a larger contribution to the RR 4-form tadpole, $N_{\text{flux}} = 48$ units.

In order to cancel the remaining tadpole, we introduce a set of branes, which will contain the chiral gauge sector, contributing with $-16$ units of RR 4-form tadpole. We choose another example from the general class in \cite{21}, with

$$\rho = 1, \beta^1 = -\beta^2 = 1, \epsilon = 1, n_a^2 = n_a^2 = -1, n_b^1 = n_c^1 = 2 \quad (5.22)$$

in table 2 in \cite{21}. This leads to...
which cancels all RR tadpoles, and leads to a chiral sector with just the Standard Model spectrum, exactly as above.

Very interestingly, our configuration corresponds to D7-branes with world-volume magnetic fields mimicking anti-D3-brane charge. This is a very explicit configuration realizing the proposal in [38] to replace antiD3-brane of [37] by world-volume anti-instantons. In particular, allows to address the absence of scalar vevs: there are regions in Kahler moduli space where no scalar tachyons are present, i.e. it is not possible to restabilize the supersymmetric vacuum since there are no scalars with the correct charges to cancel the FI term.

Clearly, many other scenarios can be devised. For instance one may break supersymmetry in the closed string sector, and consider a supersymmetric D-brane configuration, as in our example in section 4. Hopefully our examples have provided a good illustration of the model building possibilities in this setup.

5.4 Supersymmetric models?

It is a natural question to wonder about the construction of supersymmetric chiral models with NSNS and RR fluxes. This has been attempted in [2, 1] using orientifolds of $T^6/(Z_2 \times Z_2)$ without success. The difficulty arises as follows: Chirality requires the introduction of non-trivial magnetic fields in the three two-tori, and this introduces several D-brane charges, which require several kinds of orientifold planes to cancel their charge in a supersymmetry-preserving way. In order to obtain several kinds of orientifold planes one needs orbifold quotients (since due to the group law the product of two orientifold actions is an orbifold action), with the above mentioned $Z_2 \times Z_2$ being one of the simplest. The existence of orbifold projections usually modifies the flux quantization conditions (due to the requirement of proper quantization over cycles collapsed at the singularity, or similar subtleties), and requires flux quanta larger than in toroidal models. Such large fluxes generate a tadpole $N_{\text{flux}}$ exceeding the negative value from the orientifold planes, so that tadpole cancellation requires the introduction of antibranes, which render the model non-supersymmetric.

Although there is no general theorem in this direction, it seems difficult to construct supersymmetric chiral models with NSNS and RR fluxes using toroidal orbifolds. We expect that they however exist for orientifolds of more general Calabi-Yau manifolds, although such models would be more difficult to construct explicitly.
6 Final comments

Compactifications with field strength fluxes are a most promising avenue for model building, in that they provide a canonical mechanism to stabilize most moduli of the compactification. In this paper we have provided the basic model building rules for compactifications with D-branes, leading to chiral gauge sectors, and moduli stabilization by fluxes. We have described different approaches to achieve this aim, and provided explicit examples with D3-branes at singularities (with supersymmetry broken in the closed string sector) and D-branes with world-volume magnetic fluxes (with a supersymmetric closed string sector and a non-supersymmetric open string sector).

Many further directions remain open. In the context of moduli stabilization, it would be interesting to explore the interplay between the flux induced scalar potential with other sources of potential for moduli, like non-supersymmetric sets of D-branes, or non-perturbative corrections. This step is crucial in order to understand the fate of the moduli which are not stabilized by the fluxes. It would also be interesting to determine the set of values at which moduli stabilize, in order to understand for instance what properties the underlying model must have in order to lead to e.g. small 4d gauge couplings, or large radii.

Finally, our models contain several of the ingredients involved in the construction of deSitter vacua in string theory. It would be interesting to improve the kind of techniques discussed in the present paper, aiming towards building explicit models of this kind of constructions. This is essential in order to flesh out recent discussions on the discretuum of ‘realistic’ vacua in string theory, and their implications for particle physics and cosmology.

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References

[1] J. F. G. Cascales, A. M. Uranga, ‘Chiral 4d N = 1 string vacua with D branes and NSNS and RR fluxes’, JHEP 0305 (2003) 011, hep-th/0303024.

[2] R. Blumenhagen, D. Lust, T. R. Taylor, ‘Moduli Stabilization in Chiral Type IIB Orientifold Models with Fluxes’, hep-th/0303016.

[3] J. Polchinski, A. Strominger, ‘New vacua for type II string theory’, Phys. Lett. B388 (1996) 736, hep-th/9510227;
K. Becker, M. Becker, ‘M theory on eight manifolds’, Nucl. Phys. B477 (1996) 155, hep-th/9605053;
J. Michelson, ‘Compactifications of type IIB strings to four-dimensions with nontrivial classical potential’, Nucl. Phys. B495 (1997) 127, hep-th/9610151;
S. Gukov, C. Vafa, E. Witten, ‘CFT’s from Calabi-Yau four folds’, Nucl. Phys. B584 (2000) 69, Erratum-ibid. B608 (2001) 477, hep-th/9906070;

[4] K. Dasgupta, G. Rajesh, S. Sethi, ‘M theory, orientifolds and G-flux’, JHEP 9908 (1999) 023, hep-th/9908088;

[5] S. Gukov, ‘Solitons, superpotentials and calibrations’, Nucl. Phys. B574 (2000) 169, hep-th/9911011;
T. R. Taylor, C. Vafa, ‘R R flux on Calabi-Yau and partial supersymmetry breaking’, Phys. Lett. B474 (2000) 130, hep-th/9912152;
K. Behrndt, S. Gukov, ‘Domain walls and superpotentials from M theory on Calabi-Yau three folds’, Nucl. Phys. B580 (2000) 225, hep-th/0001082;
B. R. Greene, K. Schalm, G. Shiu, ‘Warped compactifications in M and F theory’, Nucl. Phys. B584 (2000) 480, hep-th/0004103;
G. Curio, A. Klemm, D. Lust, S. Theisen, ‘On the vacuum structure of type II string compactifications on Calabi-Yau spaces with H fluxes’, Nucl. Phys. B609 (2001) 3, hep-th/0012213;
M. Haack, J. Louis, ‘M theory compactified on Calabi-Yau fourfolds with background flux’, Phys. Lett. B507 (2001) 296, hep-th/0103068;
J. Louis, A. Micu, ‘Type 2 theories compactified on Calabi-Yau threefolds in the presence of background fluxes’, hep-th/0202168;

[6] S. B. Giddings, S. Kachru, J. Polchinski, ‘Hierarchies from fluxes in string compactifications’, hep-th/0105097.
[7] A. R. Frey, J. Polchinski, ‘N=3 warped compactifications’, hep-th/0201029.

[8] S. Kachru, M. Schulz, S. Trivedi, ‘Moduli stabilization from fluxes in a simple iib orientifold’, hep-th/0201028.

[9] P. K. Tripathy, S. P. Trivedi, ‘Compactification with flux on K3 and tori’, hep-th/0301139.

[10] R. D’Auria, Sergio Ferrara, S. Vaula, ‘N=4 gauged supergravity and a IIB orientifold with fluxes’, New J.Phys. 4 (2002) 71, hep-th/0206241; S. Ferrara, M. Porrati, ‘N=1 no-scale supergravity from IIB orientifolds’, Phys. Lett. B545 (2002) 411, hep-th/0207135.

[11] M. Berg, M. Haack, B. Kors, ‘An Orientifold with fluxes and branes via T duality’, Nucl. Phys. B669 (2003) 3, hep-th/0305183.

[12] M. Graña, J. Polchinski, ‘Supersymmetric three form flux perturbations on AdS5’, Phys. Rev. D63 (2001) 026001, hep-th/0009211.

[13] M. R. Douglas, G. W. Moore, ‘D-branes, quivers, and ALE instantons’, hep-th/9603167; M. R. Douglas, B. R. Greene, D. R. Morrison, ‘Orbifold resolution by D-branes’, Nucl. Phys. B506 (1997) 84, hep-th/9704151.

[14] G. Aldazabal, L. E. Ibáñez, F. Quevedo, A. M. Uranga, ‘D-branes at singularities: A Bottom up approach to the string embedding of the standard model’, JHEP 0008 (2000) 002, hep-th/0005067; D. Berenstein, V. Jejjala, R. G. Leigh, ‘The Standard model on a D-brane’, Phys. Rev. Lett. 88 (2002) 071602, hep-ph/0105042; L. F. Alday, G. Aldazabal, ‘In quest of just the standard model on D-branes at a singularity’, JHEP 0205 (2002) 022, hep-th/0203129.

[15] C. Bachas, ‘A Way to break supersymmetry’, hep-th/9503030; C. Angelantonj, I. Antoniadis, E. Dudas, A. Sagnotti, ‘Type I strings on magnetized orbifolds and brane transmutation’, Phys. Lett. B489 (2000) 223, hep-th/0007090; G. Pradisi, ‘Magnetized (shift)orientifolds’, hep-th/0210088.

[16] M. Berkooz, M. R. Douglas, R. G. Leigh, ‘Branes intersecting at angles’, Nucl. Phys. B480 (1996) 265, hep-th/9606139.

[17] R. Blumenhagen, L. Goerlich, B. Kors, D. Lust, ‘Noncommutative compactifications of type I strings on tori with magnetic background flux’, JHEP 0010 (2000) 006, hep-th/0007024.

[18] R. Rabadán, ‘Branes at angles, torons, stability and supersymmetry’, Nucl. Phys. B620 (2002) 152, hep-th/0107036.
[19] G. Aldazabal, S. Franco, L. E. Ibanez, R. Rabadan, A. M. Uranga, ‘D=4 chiral string compactifications from intersecting branes’, J. Math. Phys. 42 (2001) 3103, hep-th/0011073; ‘Intersecting brane worlds’, JHEP 0102 (2001) 047, hep-ph/0011132.

[20] R. Blumenhagen, B. Kors, D. Lust, ‘Type I strings with F flux and B flux’, JHEP 0102 (2001) 030, hep-th/0012156;

[21] L. E. Ibanez, F. Marchesano, R. Rabadan, ‘Getting just the standard model at intersecting branes’, JHEP 0111 (2001) 002, hep-th/0105155.

[22] S. Forste, G. Honecker, R. Schreyer, ‘Orientifolds with branes at angles’, JHEP 0106 (2001) 004, hep-th/0105208;
R. Blumenhagen, B. Kors, D. Lust, T. Ott, ‘The standard model from stable intersecting brane world orbifolds’, Nucl. Phys. B616 (2001) 3, hep-th/0107138;
D. Cremades, L. E. Ibanez, F. Marchesano, ‘SUSY Quivers, Intersecting Branes and the Modest Hierarchy Problem’, hep-th/0201205; ‘Intersecting brane models of particle physics and the Higgs mechanism’ hep-th/0203160; ‘Standard Model at Intersecting D5-branes: lowering the string scale’, hep-th/0205074.
C. Kokorelis, ‘GUT model hierarchies from intersecting branes’, hep-th/0203187; ‘New standard model vacua from intersecting branes’, hep-th/0205147.

[23] M. Cvetic, G. Shiu, A. M. Uranga, ‘Chiral four-dimensional N=1 supersymmetric type 2A orientifolds from intersecting D6 branes’, Nucl. Phys. B615 (2001) 3, hep-th/0107166;
‘Three family supersymmetric standard - like models from intersecting brane worlds’, Phys. Rev. Lett. 87 (2001) 201801, hep-th/0107143.

[24] R. Blumenhagen, L. Gorlich, T. Ott, ‘Supersymmetric intersecting branes on the type 2A T6 / Z(4) orientifold’, hep-th/0211059;
M. Cvetic, I. Papadimitriou, G. Shiu, ‘Supersymmetric three family SU(5) grand unified models from type IIA orientifolds with intersecting D6-branes’, hep-th/0212177;
G. Honecker, ‘Chiral supersymmetric models on an orientifold of Z4 × Z2 with intersecting D6-branes’, hep-th/0303015.

[25] R. Blumenhagen, L. Gorlich, B. Kors, ‘Supersymmetric 4-D orientifolds of type IIA with D6-branes at angles’, JHEP 0001 (2000) 040, hep-th/9912204;
S. Forste, G. Honecker, R. Schreyer, ‘Supersymmetric Z(N) x Z(M) orientifolds in 4-D with D branes at angles’, Nucl. Phys. B593 (2001) 127, hep-th/0008250.

[26] A. M. Uranga, ‘Chiral four-dimensional string compactifications with intersecting D-branes’, hep-th/0301032;
R. Blumenhagen, V. Braun, B. Kors, D. Lust, ‘The standard model on the quintic’, hep-th/0210083.

[27] K. Behrndt, M. Cvetic, ‘Supersymmetric intersecting D6-branes and fluxes in massive type IIA string theory’, hep-th/0308045.

[28] S. Gurrieri, J. Louis, A. Micu, D. Waldram, ‘Mirror symmetry in generalized Calabi-Yau compactifications’, hep-th/0211102.

[29] S. Kachru, M. B. Schulz, P. K. Tripathy, S. P. Trivedi, ‘New supersymmetric string compactifications’, hep-th/0211182.

[30] M. Grana, ‘MSSM parameters from supergravity backgrounds’, hep-th/0209200; P. G. Cámara, L. E. Ibáñez, A. M. Uranga, ‘Flux-induced SUSY-breaking soft terms’, to appear.

[31] C. Angelantonj, M. Bianchi, G. Pradisi, A. Sagnotti, Ya.S. Stanev, ‘Chiral asymmetry in four-dimensional open string vacua’, Phys. Lett. B385 (1996) 96, hep-th/9606169.

[32] J. Lykken, E. Poppitz, S. P. Trivedi, ‘Branes with GUTs and supersymmetry breaking’, Nucl. Phys. B543 (1999) 105, hep-th/9806080.

[33] E. Witten, ‘Baryons and branes in anti-de Sitter space’, JHEP 9807 (1998) 006, hep-th/9805112.

[34] I. Antoniadis, E. Kiritsis, J. Rizos, ‘Anomalous U(1)s in type 1 superstring vacua’, Nucl. Phys. B637 (2002) 92, hep-th/0204153.

[35] G. Aldazabal, L. E. Ibáñez, A.M. Uranga, ‘Gauging away the strong CP problem’, hep-ph/0205250.

[36] A. M. Uranga, ‘D-brane, fluxes and chirality’, JHEP 0204 (2002) 016, hep-th/0201221.

[37] S. Kachru, R. Kallosh, A. Linde, S. P. Trivedi, ‘De sitter vacua in string theory’, hep-th/0301240.

[38] C. P. Burgess, R. Kallosh, F. Quevedo, ‘De Sitter string vacua from supersymmetric D terms’, hep-th/0309187.