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Surveying standard model flux vacua on $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$

Jason Kumar and James D. Wells

Michigan Center for Theoretical Physics (MCTP), University of Michigan
Ann Arbor, MI 48109-1120, U.S.A.
E-mail: kumar.jason@gmail.com, jwells@umich.edu

Abstract: We consider the $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ Standard Model brane embedding in an orientifold of $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$. Within defined limits, we construct all such Standard Model brane embeddings and determine the relative number of flux vacua for each construction. Supersymmetry preserving brane recombination in the hidden sector enables us to identify many solutions with high flux. We discuss in detail the phenomenology of one model which is likely to dominate the counting of vacua. While Kähler moduli stabilization remains to be fully understood, we define the criteria necessary for generic constructions to have fixed moduli.

Keywords: D-branes, Superstring Vacua, Supersymmetry Breaking, Supersymmetry Phenomenology.

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Contents

1. Introduction

There is now substantial, though by no means conclusive, evidence that there are a very large number of vacua of string theory in which all [1 − 3] of the moduli [4 − 8, 2] can be fixed. Recently, there has been significant development in the techniques for describing and counting these vacua [9 − 12, 3]. Much of this work is statistical in nature, generating estimates of the number of vacua with specific gross properties [7, 13]. The actual construction of explicit models with all moduli vevs and their potentials identified is far more difficult [14 − 16], though constructions using Type IIA string compactifications seem more tractable than those of Type IIB [3].

Our interest in this work is to gain insight into the statistical distributions of vacua and their properties, and in particular how that can be connected to the problem of model-building [17 − 20, 22]. It seems clear that any concrete program for actually constructing stringy models of the real world (or even of toy models which resemble the world and have interesting phenomenology) can only be helped by a statistical study of the frequency with which low-energy properties occur on the landscape.

One direction is to characterize the distribution of vacua compatible with the standard model (SM) [17, 18, 21, 22]. This would be a difficult endeavor, even for the limited case of orientifolded Calabi-Yau three-fold compacitifications of Type IIB string theory. We will instead consider as an exercise the simple construction of Type IIB compactified on an orientifold of $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$. We hope the results of this survey can provide some intuition about the more general problems of constructing SM string embeddings on the landscape, and of determining the distribution of these embeddings.

In these constructions, some gauge dynamics will arise from open strings beginning and ending on branes. We would like this open string gauge theory to include the SM. As a result, we can roughly divide the open string gauge theory into two sectors: the visible sector containing the SM gauge group (with some extensions, such as to a Pati-Salam unification group or a left-right extension), and a hidden sector containing gauge groups not identified with the SM. The branes that are relevant for visible or hidden sector dynamics will wrap some holomorphic even-dimensional cycles of the Calabi-Yau, and will be extended in all of the non-compact directions.

One must first define what is meant by a string construction of the SM. We will consider a series of branes embedded in an orientifolded Calabi-Yau 3-fold compactification such that one set of branes yields the gauge group and chiral matter content of the SM. This set of branes will be called the visible sector. Various additional hidden sector branes will also be allowed, and there may be chiral exotics charged under either hidden sector branes, or both hidden and visible sector branes. The only demand we will make of the matter content is that there be no chiral exotics charged only under the visible sector. In particular, note that there is no restriction of any kind on vector-like matter, as this matter can receive a large mass and thus not conflict with experiment.

In section 2, we will review the general properties of our $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold. In section 3, we will codify the rules for constructing consistent brane embeddings. In section 4, we discuss actual SM embeddings, and the properties of associated flux vacua. In section 5, we discuss a model which dominates the counting of SM flux vacua. We close with a discussion of our results in section 6.

2. $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold

We will focus on brane constructions of the SM on an orientifold of $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$. [17, 18, 23, 24]. The $\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold group is thus supplemented by the orientifold element ΩR; the orbifold group is generated by the elements

$$\begin{aligned} \alpha &: (z_1, z_2, z_3) \to (-z_1, -z_2, z_3) \\ \beta &: (z_1, z_2, z_3) \to (z_1, -z_2, -z_3) \end{aligned} \tag{2.1}$$

where the involution is given by

$$R : (z_1, z_2, z_3) \to (-z_1, -z_2, -z_3) \tag{2.2}$$

and Ω is the worldsheet parity operator.

This orientifold group generates 64 O3-planes and 12 O7-planes. The O3-planes stretch in the 4 non-compact directions. The O7-planes appear in three sets which all stretch in the non-compact directions, but which wrap different 4-cycles in the compact directions. In particular, the O7-planes respectively wrap the tori $T_1^2 T_2^2$, $T_1^2 T_3^2$ and $T_2^2 T_3^2$. The $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$ orbifold itself breaks $N = 8$ supersymmetry down to $N = 2$. The orientifold action further breaks supersymmetry down to $N = 1$.

For our choice of discrete torsion, this model has 3 Kähler moduli (which determine the size of each T^2 factor) and 51 complex structure moduli. 48 of these complex structure moduli arise at the fixed points of elements of the orbifold group. For a different choice of discrete torsion, the 48 moduli arising at the fixed points would instead be Kähler moduli.

In this compactification, branes may wrap either the cycles of the T^6, or the shrunken cycles arising at the orbifold fixed points. We will only be interested in the cycles of the T^6. We may describe the relevant branes of Type IIB string theory (D3-,D5-,D7- and D9-branes) in the so-called magnetized D-brane formalism [25]. This formalism arises from the realization that lower dimensional branes can be described by magnetic fluxes on the worldvolume of a D9-brane. Essentially, we factorize the six-torus into $(T^2)^3$, and we assign to each brane an ordered pair (n_i, m_i) for each T_i^2, where m_i is the number of times that the brane wraps this T^2, and n_i gives the amount of constant magnetic flux on this cycle normalized as

$$\frac{m_i}{2\pi}\int_{T_i^2} F^i = n_i\,. \tag{2.3}$$

We see that a brane in which 3 of the m's are zero must be a D3-brane extended in the non-compact dimensions, but with no wrapping on the torus. Similarly, a D5-brane wrapping a 2-cycle of the compact space will have only one non-zero m, while a D7-brane wrapping a 4-cycle of the compact space will have two non-zero m's. For a D9-brane, all m's will be non-zero.

It is often easier to picture this from the T-dual prescription, which is Type IIA string theory on a $T^6/\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold. In this picture, the involution R' which appears in the orientifold element $\Omega R'$ is given by

$$R' : (z_1, z_2, z_3) \rightarrow (z_1^*, z_2^*, z_3^*)\,. \tag{2.4}$$

The IIB branes are dualized to D6-branes of Type IIA which wrap a one-cycle of each torus. The (n_i, m_i) ordered pairs just give the winding numbers on each cycle of the torus T_i^2.

If Δ_a is any particular brane, its orientifold image Δ'_a will have wrapping numbers given by $n_i^a \rightarrow n_i^a$, $m_i^a \rightarrow -m_i^a$. It is easiest to see this by examining the T-dual IIA picture, in which the m's are the wrapping numbers along the three real directions which are inverted by the involution R.

3. Rules for brane constructions

The rules for constructing brane models in type IIB theories are detailed in the literature. We wish here to distill these discussions into a list of rules that must be followed to construct D-brane models in our setup. When all rules are satisfied, a mathematically consistent theory results.

Wrapping numbers: In the previous section we discussed the salient properties of the $T^6/\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold, and showed that a brane Δ_a could be written as a set of its wrapping numbers on the three two-tori cycles

$$\Delta_a = (n_1^a, m_1^a)(n_2^a, m_2^a)(n_3^a, m_3^a)\,. \tag{3.1}$$

The n^i and m^i numbers must be co-prime integers.

RR tadpoles: Gauss's law imposes the constraint that there may be no net charge in a compact space. As a result, RR charges which stretch over all non-compact directions must cancel. This constraint can be rephrased as the statement that, for a consistent brane embedding, all RR tadpoles must cancel. It is easy to implement these tadpole conditions by constructing a brane-charge vector $\vec{Q}$ such that Q_0 is the $D3$ brane charge, and Q_i are the $D7_i$ brane charges:[1]

$$\begin{aligned} Q_0 &= n_1 n_2 n_3 \quad (D3 \;\; \text{charge}) \\ Q_1 &= -n_1 m_2 m_3 \quad (D7_1 \;\; \text{charge}) \\ Q_2 &= -m_1 n_2 m_3 \quad (D7_2 \;\; \text{charge}) \\ Q_3 &= -m_1 m_2 n_3 \quad (D7_3 \;\; \text{charge}) \end{aligned} \tag{3.2}$$

This charge vector, summing over all branes Δ_a and their images Δ'_a (where $m_i^a \to -m_i^a$), must equal the sum of all orientifold plane charges:

$$\sum_a N_a(\vec{Q}(\Delta_a) + \vec{Q}(\Delta'_a)) = 32\vec{Q}(\mathcal{O})\,, \tag{3.3}$$

where N_a is the number of branes for each stack a. Since $\vec{Q}(\mathcal{O}) = (1,1,1,1)$ and $\vec{Q}(\Delta_a) = \vec{Q}(\Delta'_a)$, we have as the final condition

$$\sum_a N_a \vec{Q}(\Delta_a) = (16,16,16,16)\,. \tag{3.4}$$

Three-form fluxes will contribute to the $D3$ charge but not the $D7_i$ charge. One unit of flux contributes 32 units of $D3$ charge, thereby changing the RR tadpole conditions to

$$\sum_a N_a \vec{Q}(\Delta_a) + (32N_{\text{flux}}, 0, 0, 0) = (16,16,16,16)\,, \tag{3.5}$$

where N_{flux} is a non-negative integer, and again, this final sum is only over the branes and not their images. The factor of 32 arises from the reduction in size of the 3-cycle volume due to the orbifold action [31] and from the condition that there by no exotic branes [32].

It is interesting to note that the cancellation of RR tadpoles implies the cancellation of anomalies in the worldvolume gauge theories of the embedded branes [33, 28]. Indeed, the only consistency conditions, from the string theory point of view, will amount to tadpole cancellation conditions. Consistency of the overall string compactification implies consistency of the low-energy description, although this may appear through an anomaly inflow which cancels local anomalies. This fact will have important implications later for the number of chiral exotics in the embeddings we consider.

[1] The minus signs in the definitions of $Q_{1,2,3}$ serve the purpose of assigning $+1$ charge to pure $D7_i$ branes, just as there is a charge of $+1$ to a pure $D3$ brane. Note that $D7_i$ refers to a D7-brane which is not wrapped on the torus T_i^2, but is wrapped on the other two.

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K-theory constraints: the RR tadpole conditions are requirements on the total $D3$ and $D7$ charges of the brane stacks. There are analagous constraints on the total $D5$ brane and $D9$ brane charges [25]. We can form a new vector $\vec{Y}$ of these charges, such that

$$\begin{aligned} Y_0 &= m_1 m_2 m_3 \quad (D9 \text{ charge}) \\ Y_1 &= m_1 n_2 n_3 \quad (D5_1 \text{ charge}) \\ Y_2 &= n_1 m_2 n_3 \quad (D5_2 \text{ charge}) \\ Y_3 &= n_1 n_2 m_3 \quad (D5_3 \text{ charge}) \,. \end{aligned} \tag{3.6}$$

The requirement on these charges is merely that they be even under the summation

$$\sum_a R_a \vec{Y}^a = (0,0,0,0) \bmod 2\,, \tag{3.7}$$

where R_a is the rank of the gauge group for brane stack a. In fact, this is equivalent to the demand that if one inserts a probe D-brane with SU(2) gauge group, then there should be no global anomaly (i.e., the number of Weyl fermions in the probe brane worldvolume theory must be even [28, 29]). For brane stacks at orbifold fixed points, $R_a = N_a/2$ and the constraint becomes

$$\sum_a N_a \vec{Y}^a = (0,0,0,0) \bmod 4\,. \tag{3.8}$$

This is often called the K-theory constraint, since it is a restriction on the theory not captured by pure homology [30]. It might appear rather innocuous, and in many cases in the early literature it had no strong impact on the results. But in making theories with magnetized $D9$ branes, this constraint can be quite restrictive, as we will encounter in the next section.

It turns out that adding discrete B-field(s) along the tori [26] will enable us to more easily find solutions to these K-theory constraints. However, these B-fields will also change the RR-tadpole constraints, making them more difficult to solve. In any case, they can potentially introduce obstructions to the vector bundles which would be necessary for us to obtain the open string gauge groups which we will need for our SM construction.[2] As a result, we will not consider turning on any of these discrete B-fields.

NSNS tadpoles: NSNS tadpole cancellation (i.e., no uncompensated brane tensions) is achieved if the $N = 1$ supersymmetry that remains from the orientifold is preserved in all D-brane sectors. Operationally, this leads to the turning on of a Fayet-Iliopoulos (FI) term if the tadpole is not cancelled. Each brane has a FI-term which is determined in terms of the set of three Kähler moduli parameters $\mathcal{A}_1$, $\mathcal{A}_2$ and $\mathcal{A}_3$ by the equation[3]

$$\sum_{i=1}^{3} \tan^{-1}(m_i^a \mathcal{A}_i, n_i^a) \bmod 2\pi = \theta \sim \xi_{FI}\,, \tag{3.9}$$

[2]We are indebted to S. Sethi and G. Shiu for discussions of these points.

[3]This derives from the fact that two branes preserve common supersymmetries if they can be related by an SU(d) rotation in d complex dimensions [33].

where $\alpha = \tan^{-1}(y, x)$ is defined from 0 to 2π (or equivalently, $-\pi$ to π) in such a way that $\sin\alpha = \frac{y}{\sqrt{x^2+y^2}}$ and $\cos\alpha = \frac{x}{\sqrt{x^2+y^2}}$, and where "mod" is defined such that $\theta \in [-\pi, \pi]$.

It is important to note that if NSNS tadpoles are not cancelled [39], supersymmetry will **not** necessarily be broken. Indeed, the FI-term will appear in the D-term potential in the form

$$V_{D_a} = \frac{1}{2g^2}(\sum q_i|\phi_i|^2 + \xi)^2 , \tag{3.10}$$

where the ϕ are scalars charged under the diagonal U(1) of the gauge group associated with the stack of branes Δ_a. This implies that even if $\xi \neq 0$, an appropriate scalar can possibly get a vev to cancel the D-term potential. This corresponds to brane recombination, in which the two branes that bind are those under which the veved scalar is charged. Thus, we should think of this FI-term as simply contributing to a real constraint equation involving both open string moduli and the Kähler moduli. We will indeed find that our brane embeddings will necessarily be very rich, yielding many scalars charged under each U(1) with both positive and negative sign. As such, it is most likely that a non-vanishing FI-term will in fact lead to a deformation of the brane system which preserves supersymmetry.

However, we should distinguish between branes that satisfy the NSNS tadpole conditions somewhere in Kähler moduli space, and those that cannot satisfy them anywhere in the moduli space (such as $\overline{D3}$ branes, for example). Branes that cannot satisfy the tadpole conditions anywhere in moduli space have the potential to destabilize the solution. An example of such an instability was shown in [40], where it was found that the presence of sufficiently many anti-D3 branes would result in a classical instability through which $\overline{D3}$-branes and fluxes would annihilate. More generally, $\overline{D3}$-branes will also contribute a term to the potential arising from their vacuum energy, which can destabilize the Kähler moduli of the solution unless it is tuned to be small. What distinguishes an $\overline{D3}$-brane from a supersymmetric brane (such as a D3-brane) is its orientation. In our previous language, the $\overline{D3}$-brane maximally violates the NSNS tadpole constraint, and cannot be made to satisfy it anywhere in Kähler moduli space.

On the other hand, a brane that can satisfy the NSNS tadpole conditions somewhere will indeed be supersymmetric and stable (with unbroken gauge symmetry) when the FI-term vanishes. Therefore, the only potentially destabilizing contribution which they make to the potential is from the FI-term, and in fact this term merely provides another constraint which can be satisfied (restoring supersymmetry).

Although we do not have a general argument that branes that cannot satisfy the NSNS tadpole conditions anywhere in Kähler moduli space will lead to instability, our experience with $\overline{D3}$-branes leads us to strongly suspect that such branes have the potential to destabilize the solution. As such, we will demand that all branes be able to satisfy the NSNS tadpole constraints somewhere in Kähler moduli space, but not necessarily at the same point. This will ensure that our solutions are likely to be truly stable, though our limitation may exclude some stable solutions. Indeed, one may nevertheless be able to add a small number of $\overline{D3}$ branes to such solutions for the purpose of breaking supersymmetry. However, one must then be careful to tune their contribution to avoid destabilizing the solution. We will not make this choice rigid, in order to preserve the potentially more attractive possibility of a different mechanism for supersymmetry breaking (such as IASD fluxes). We will discuss the issue of supersymmetry breaking in more detail in section 5.

Gauge groups of brane stacks: we will choose brane stacks to either lie on orientifold planes or on orbifold fixed points away from orientifold planes. This will ensure that we have an odd number of generations [21]. The resulting gauge groups are different in these two cases.

For a stack with N_a branes on an orientifold plane,[4] the gauge group[5] is $\mathrm{USp}(N_a)$. This gauge group only makes sense if N_a is even. ($\mathrm{USp}(2)$ is isomorphic to $\mathrm{SU}(2)$, which we will use below.) Only pure $D7$ branes and $D3$ branes can be located on orientifold planes and give rise to these $\mathrm{USp}(N_a)$ gauge groups. For a stack with N_a branes at an orbifold fixed point, but not on an orientifold plane, the gauge group is $\mathrm{U}(N_a/2)$. Again, N_a must be even.

Intersection numbers and chiral matter content: the chiral matter content is obtained by computing intersection numbers of one brane stack with another:

$$I_{ab} = \prod_{i=1}^{3} \det \begin{pmatrix} n_i^a & m_i^a \\ n_i^b & m_i^b \end{pmatrix} = \left(n_1^a m_1^b - m_1^a n_1^b\right)\left(n_2^a m_2^b - m_2^a n_2^b\right)\left(n_3^a m_3^b - m_3^a n_3^b\right). \tag{3.11}$$

Let us label a as a stack of a magnetized branes (either magnetized $D7$ or $D9$). The matter content arising from a string beginning and ending on a will be

$$aa \text{ matter}: \quad \mathrm{U}\left(\frac{N_a}{2}\right) \text{ vector multiplet plus 3 adjoint chirals}. \tag{3.12}$$

Furthermore, a and its image a' ($m_i^b \to -m_i^b$) can intersect with any other brane c and its orientifold image c' (applicable only if c is magnetized $D7$ or $D9$-brane stack):

$$\begin{aligned}
ac \text{ matter}: &\quad I_{ac} \text{ copies of } (\Box_a, \overline{\Box}_c) \text{ chirals} \\
aa' \text{ matter}: &\quad \left(\frac{I_{aa'}}{2} - 2I_{a,\mathcal{O}}\right) \text{ copies of } \Box\!\Box \text{ chirals} \\
&\quad \left(\frac{I_{aa'}}{2} + 2I_{a,\mathcal{O}}\right) \text{ copies of } \begin{smallmatrix}\Box\\ \Box\end{smallmatrix} \text{ chirals} \\
ac' \text{ matter}: &\quad I_{ac'} \text{ copies of } (\Box_a, \Box_c) \text{ chirals}
\end{aligned} \tag{3.13}$$

where $I_{a,\mathcal{O}}$ is the intersection number summed over the each orientifold plane:

$$I_{a,\mathcal{O}} = m_1^a n_2^a n_3^a + n_1^a m_2^a n_3^a + n_1^a n_2^a m_3^a - m_1^a m_2^a m_3^a\,. \tag{3.14}$$

[4] Sometimes, N_a branes on an orientifold brane are identified with their images, leading some authors to say by convention that there are $2N_a$ branes for this case.

[5] There is some freedom in the choice of action of the orientifold group on the Chan-Paton indices. Indeed, Denef et al. [16] use this freedom to construct (non-standard model) embeddings with $SO(N)$ gauge groups. One can consider whether these discrete choices allow one to find SM constructions, and how these constructions relate to the ones we consider.

Let us label b as a stack of pure $D7$ or $D3$ branes lying on an orientifold plane. The matter arising from strings beginning and ending on b are

$$bb \text{ matter}: \quad \mathrm{USp}(N_b) \text{ vector multiplet plus 3 copies of } \begin{smallmatrix}\square\\[-2pt]\square\end{smallmatrix} \text{ chirals} \tag{3.15}$$

b can also intersect with any other brane c to yield

$$bc \text{ matter}: I_{bc} \text{ copies of } (\square_b, \overline{\square}_c) \text{ chirals}\,. \tag{3.16}$$

There is no intersection of b with its image or with the image of any other brane that contributes more to the total matter content.

4. Standard model embeddings

One goal of a string model building exercise is to construct a visible sector that has a chance of reducing to the SM at low energies. For us, this will mean a visible sector that contains the SM gauge group with the correct chiral matter. We need four different brane stacks in order to achieve a SM embedding: a stack for each of the three gauge groups of the SM, plus another stack to enable $\mathrm{SU}(3)\times\mathrm{SU}(2)_L$ singlets to intersect with hypercharge. Not only would we like these stacks to yield a gauge group containing the SM with the right chiral matter content, but we would also like this visible sector to preserve the same $N=1$ supersymmetry as the orientifold.

Requiring the visible sector to contain the SM with 3 generations forces us to include visible sector branes with large charges. This will (in known constructions) mean that the visible sector includes branes with $D3$-brane charge larger than that carried by the $O3$-planes. To compensate for this, one must include a source of negative $D3$-brane charge. One approach is to allow anti-$D3$ branes into the spectrum, thereby breaking supersymmetry but allowing large fluxes that stabilize the complex structure moduli. A second approach is to construct a visible sector entirely out of $D7$ branes and then using magnetized $D9$ branes with induced negative $D3$ brane charge in the hidden sector to cancel the RR tadpole. A third approach is to make the magnetized $D9$ branes part of the visible sector, and use only pure $D7$ and $D3$ branes in the hidden sector to cancel the RR tadpoles.

All of these approaches to finding a SM embedding for type IIB are laudable. However, we wish to focus on the second approach since it has the advantage of an economical gauge group structure of the visible sector (i.e., USp(2) groups rather than U(2) groups can generate the weak SU(2)) and allows a wide variety of supersymmetry breaking mechanisms and scales without destabilizing the solution.

The simplest structure for such a four-stack embedding is the left-right model, as advocated by [17, 18]. In this model, the $\mathrm{SU}(2)_{L,R}$ groups arise from USp(2) gauge theories living on stacks of branes, rather than from U(2). This feature can be quite attractive, as otherwise there will be additional $\mathrm{U}(1)_{L,R}$ anomalies which must be cancelled. Of course it is possible to cancel such anomalies, either in a strictly field theoretic context through the use of U(2) anti-doublets as part of the chiral matter [35], or in a string theory context through the Green-Schwarz mechanism. However, the use of anti-doublets may not be desirable from a phenomenological standpoint. Although these difficulties can be solved, they are avoided altogether in models which contain USp(2) groups in the visible sector.

We will consider only Pati-Salam left-right constructions in which the $\mathrm{SU}(2)_{L,R}$ groups arise from USp(2) groups. As a result, two of the four brane stacks (those generating $\mathrm{U}(3), \mathrm{U}(1) \subset \mathrm{U}(4)$) have the same wrapping numbers, while the other two stacks are either pure D3-branes or pure D7-branes. The two SU(2) branes must have different intersections with respect to the U(3) brane (to account for the chiral bifundamental matter). As a result, we must pick the SU(2) stacks from two distinct choices out of the four pure D3/D7-branes

$$\begin{aligned}
&D7_1(1,0)(0,1)(0,-1) \\
&D7_2(0,1)(1,0)(0,-1) \\
&D7_3(0,1)(0,-1)(1,0) \\
&D3(1,0)(1,0)(1,0)\,.
\end{aligned} \tag{4.1}$$

There are thus six choices we can make, but without loss of generality, we will choose $D7_2$ and $D7_3$ [17, 18]. Similar statements will apply for the other cases.

Since the U(3) and U(1) brane stacks have the same wrapping numbers, the only choice we have left is with these six wrapping numbers. These are chosen subject to the constraint that the brane be supersymmetric somewhere in Kähler moduli space, and that it have the right intersection numbers with the two SU(2) branes. These intersection conditions can be rephrased as

$$\begin{aligned}
n_1 m_2 n_3 &= 3 \\
n_1 n_2 m_3 &= -3\,.
\end{aligned} \tag{4.2}$$

If we wish to have no magnetized D9-branes in the visible sector, then must have $n_1 = 1$,$m_1 = 0$. We see then our only choices for the wrapping numbers of the U(4) stack are

$$\begin{aligned}
&\text{choice 1 } b_{1,2,3} = 0 \;\; (1,0)(3,1)(3,-1) \\
&\text{choice 2 } b_{1,2,3} = 0 \;\; (1,0)(1,3)(1,-3)\,.
\end{aligned} \tag{4.3}$$

Thus, going with choice 1, we can identify the full visible sector of a three-generation model of this type as four stacks of $D7$ branes

$$\begin{aligned}
N_a &= 6 \;\; (1,0)(3,1)(3,-1) \\
N_b &= 2 \;\; (0,1)(1,0)(0,-1) \\
N_c &= 2 \;\; (0,1)(0,-1)(1,0) \\
N_d &= 2 \;\; (1,0)(3,1)(3,-1)\,.
\end{aligned} \tag{4.4}$$

Gauge groups for this are $\mathrm{SU}(3) \times \mathrm{SU}(2)_1 \times \mathrm{SU}(2)_2 \times \mathrm{U}(1)_a \times \mathrm{U}(1)_d$, which yields the matter content given by table 1. The H_u and H_d fields are vector complements that do not contribute to the net chirality and are thus not accounted for by the intersection numbers between branes [18]. They are charged states in the (bc) sector despite $I_{bc} = 0$, and we

Sector	N_{copies}	$\text{SU}(3) \times \text{SU}(2)_L \times \text{SU}(2)_R$	$\text{U}(1)_a$	$\text{U}(1)_d$	$\text{U}(1)_{B-L}$
(ab)	3	$(d_L, u_L) \sim (3{,}2{,}1)$	1	0	1/3
(ac)	3	$(d_R^c, u_R^c) \sim (\bar{3}{,}1{,}2)$	-1	0	$-1/3$
(db)	3	$(l_L, \nu_L) \sim (1{,}2{,}1)$	0	1	-1
(dc)	3	$(l_R^c, \nu_R^c) \sim (1{,}1{,}2)$	0	-1	1
(bc)	1	$(H_u, H_d) \sim (1{,}2{,}2)$	0	0	0

Table 1: Minimal spectrum of MSSM in left-right model visible sector brane construction used in this paper.

include them in the table to provide a complete picture of what is the minimum visible sector allowed by this framework consistent with SM needs.

The visible sector gauge groups can be broken down to the SM as illustrated in [18]. As pointed out in [17, 18], the visible sector is automatically supersymmetric as long as the Kähler moduli parameters satisfy $\mathcal{A}_2 = \mathcal{A}_3$. We can think of this as an entire plane in Kähler moduli space that is supersymmetric for the visible sector. It will actually be necessary to impose this constraint to forbid giving vev's to scalars charged under $\text{SU}(3)_{qcd}$, as we will see shortly.

The visible sector identified above cannot stand alone as it does not cancel the RR tadpole conditions. We need to introduce a hidden sector for that purpose. We will classify our choice of consistent hidden sector into two categories: those that introduce one new NSNS tadpole constraint and those that introduce two. Of course, it is consistent to consider hidden sectors that introduce more constraints. Such constraints will generically not be satisfied, but deformation of the brane embedding due to the veving of open string fields will usually allow such solutions to preserve supersymmetry. But the addition of larger numbers of branes will make it more difficult for the hidden sector to satisfy the RR-tadpole conditions. For this reason (as well as the computational difficulty in searching for solutions with many branes), we will content ourselves in this work with finding hidden sectors that impose at most two more constraints. We will not attempt to study the non-perturbative structure of the superpotential; the question of whether or not the Kähler moduli are actually fixed [44] in a supersymmetric[6] solution thus remains open.

4.1 Solutions with one NSNS tadpole constraint

As noted above, the visible sector alone does not satisfy the RR tadpole conditions. The contributions from the branes of the visible sector to the $\vec{Q}$ vector are

$$\vec{Q}_{\text{vis}} \equiv \sum_{k=a,b,c,d} N_k \vec{Q}(\Delta_k) = (72, 8, 2, 2)\,. \tag{4.5}$$

We need the visible sector plus hidden sector of branes to cancel the RR tadpole conditions

$$\vec{Q}_{\text{vis}} + \vec{Q}_{\text{hid}} + (32 N_{\text{flux}}, 0, 0, 0) = (16, 16, 16, 16) \tag{4.6}$$

[6]The reference to "supersymmetric" solutions perhaps requires some clarification. It is true that in a no-scale model, generic ISD fluxes will in fact break supersymmetry with F-terms for the Kähler moduli. However, generic non-perturbative corrections to both the Kähler potential and superpotential will break no-scale structure [9, 8]. This permits solutions where the F-terms are zero.

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which leads to the requirement that we find a hidden sector such that

$$\vec{Q}_{\rm hid} = (-56 - 32N_{\rm flux}, 8, 14, 14)\,. \tag{4.7}$$

In addition, we demand that the hidden sector preserve the same $N = 1$ supersymmetry as the visible sector.

The very large negative $Q_0^{\rm vis}$ charge needed is problematic. Pure supersymmetric $D3$ branes or fluxes only add positively to this charge and pure $D7$ branes do not add to Q_0 at all. What we need is a brane configuration with large negative Q_0 and small or negative contributions to $Q_{1,2,3}$. Furthermore, one can demonstrate from the NSNS tadpole conditions that a brane can preserve supersymmetry only if at most one of its Q charges is negative [20]. The only choice of brane that can preserve the same supersymmetry as the orientifold and carry a negative charge is a magnetized $D9$ brane with the appropriate signs in its wrapping numbers.

Note nevertheless that the addition of a hidden sector magnetized $D9$ brane will introduce another NSNS tadpole condition (applying eq. (3.9) to the $D9$ brane). Generically, any choice of n magnetized hidden sector branes will introduce n additional NSNS tadpole constraints on the Kähler moduli space. Although it is not necessary to satisfy these constraints, it is useful to classify solutions based on the number of constraints which arise from the solution. For every NSNS tadpole constraint which is not satisfied, supersymmetry will require a scalar charged under the appropriate gauge group to get a vev. If such a scalar is also charged under the SM, then this could be problematic for phenomenology.

The NSNS tadpole constraint of the visible sector would set $\mathcal{A}_2 = \mathcal{A}_3$, which is a constraint[7] we assume the Kähler moduli satisfy in order to minimize the risk of unacceptable gauge symmetry breaking (e.g., $\mathrm{SU}(3)_{qcd}$ charged fields condensing). When we add a hidden sector magnetized $D9$ brane to the theory and demand that it be supersymmetric, we are adding an additional FI-term constraint involving ξ, where

$$\xi \sim -\pi + \sum_i \tan^{-1}\left(\frac{|m_i|\mathcal{A}_i}{|n_i|}\right) \tag{4.8}$$

subject to $\mathcal{A}_2 = \mathcal{A}_3$. Therefore, we see that branes which are related by $(n_2, m_2) \leftrightarrow (n_3, m_3)$ will impose the same constraint on the Kähler moduli and thus can appear undeformed together in the hidden sector while preserving supersymmetry. A brane which is related to another by such an interchange of winding numbers will be referred to as a "partner" brane.

We have constructed a computer program that searches for hidden sector solutions that satisfy all the rules and constraints detailed in the previous section and which are supersymmetric on 1D-surfaces in Kähler moduli space (i.e., only one additional NSNS constraint in addition to $\mathcal{A}_2 = \mathcal{A}_3$). After an exhaustive search we have identified six unique classes of hidden sector solutions, listed in table 2.

[7] Note that this constraint appears to only be a constraint on the real Kähler moduli corresonding to the tori volumes, not on the axions which complexify them. How this constraint would be complexfied is an interesting question [45] [48].

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	magnetized $D9$ brane	N_a	$N_{\tilde{a}}$	$N_{D3,D7_i}$	$N_{\rm flux}^{\rm max}$
1	$(-2,1)(-3,1)(-4,1)$	2	2	(40,0,0,0)	1
2	$(-2,1)(-3,1)(-3,1)$	4	-	(16,0,2,2)	0
3	$(-2,1)(-2,1)(-7,2)$	2	0	(0,0,0,6)	0
4	$(-2,1)(-2,1)(-7,2)$	0	2	(0,0,6,0)	0
5	$(-2,1)(-2,1)(-5,1)$	2	2	(24,0,0,0)	0
6	$(-2,1)(-2,1)(-4,1)$	2	2	(8,0,2,2)	0

Table 2: The complete set of solutions to the hidden sector for the SM embedding which are supersymmetric along a 1D-surfacee in Kähler moduli space. Only the first model admits 3-form flux, and this model is equivalent up to trivial sign reparametrizations to the one found by [18].

From these solutions we see that there is only one possible solution with $N_{\rm flux} > 0$ and that is obtained by adding one unit of flux to the first solution in table 2. We should note that our computer search for $N_{\rm flux} = 1$ models found 109 solutions when we do not take into account the K-theory constraint, of which 65 allowed non-zero flux up to $N_{\rm flux} = 10$. Once we apply the K-theory constraint only this one solution given above is left, which through a trivial sign reparametrization is equivalent to the model presented in [18]. There appear to be no solutions satisfying all constraints for $N_{\rm flux} \geq 2$ in these constructions that satisfy supersymmetry on a 1D surface in Kähler moduli space.

These solutions have no net chiral exotics charged under $\mathrm{SU}(4)_{PS}$. This feature is related to the automatic cancellation of the $\mathrm{SU}(4)_{PS}$ cubic anomaly, which in turn arises from the satisfaction of the RR tadpole conditions [34, 27, 28]. The only contributions to the $\mathrm{SU}(4)_{PS}$ cubic anomaly arise from fermions transforming in the fundamental, anti-fundamental, symmetric or anti-symmetric representations. One can easily verify that no fermions transform in the symmetric or anti-symmetric representation. As a result, the number of fundamentals of $\mathrm{SU}(4)_{PS}$ must equal the number of anti-fundamentals, implying that there are no net chiral exotics of this type.

4.2 Solutions with 2 NSNS tadpole constraints

We may also consider solutions in which the hidden sector consists of multiple branes which impose independent NSNS tadpole constraints. If one adds no more than two NSNS tadpole constraints from the hidden sector (in addition to the contribution from the visible sector), then we might expect to be able to solve all constraints. In fact, we will find that although we have three constraints for three unknowns, the constraints nevertheless cannot be solved simultaneously in most cases.

If we choose to add a hidden sector which generates 2 NSNS tadpole conditions, then we must add two distinct hidden sector branes. The first must have negative D3-brane charge in order to cancel that RR tadpole. In order for it to be supersymmetric anywhere on moduli space, it must therefore have positive values for all three D7-brane charges.

For the second brane, we have three choices. First, we might choose another brane with negative D3-brane charge and positive D7-brane charges. Secondly, we might choose a brane with negative value for one D7-brane charge, and positive values for all other

brane a	brane b	N_a	N_b	$N_{D3,D7_i}$	$N_{\rm flux}^{\rm max}$
$(-2,1)(-5,1)(-4,1)$	$(-2,1)(-2,1)(-1,1)$	2	2	(0,0,0,4)	1
$(-4,1)(-4,1)(-4,1)$	$(0,1)(2,-1)(1,1)$	2	6	(8,0,20,0)	2
$(-3,1)(-5,1)(-5,1)$	$(-1,1)(-1,1)(-1,1)$	2	2	(0,0,2,2)	3
$(-4,1)(-6,1)(-4,1)$	$(0,1)(6,-1)(1,1)$	2	6	(8,0,40,0)	4
$(-5,1)(-6,1)(-4,1)$	$(2,1)(1,-1)(2,-1)$	2	2	(16,2,0,2)	5
$(-7,1)(-4,1)(-5,1)$	$(7,1)(2,-1)(1,-1)$	2	2	(4,8,2,2)	6
$(-6,1)(-6,1)(-4,1)$	$(2,1)(1,-1)(2,-1)$	2	2	(0,0,0,2)	7
$(-7,1)(-5,1)(-5,1)$	$(7,1)(1,-1)(1,-1)$	2	2	(24,8,2,2)	8
$(-6,1)(-6,1)(-5,1)$	$(2,1)(1,-1)(2,-1)$	2	2	(8,0,0,0)	9
$(-7,1)(-5,1)(-5,1)$	$(3,1)(1,-1)(1,-1)$	2	2	(0,0,2,2)	9
$(-6,1)(-5,1)(-6,1)$	$(4,1)(2,-1)(1,-1)$	2	2	(0,4,0,0)	9

Table 3: Illustrative set of hidden sector solutions with $N_{\rm flux} > 0$ possible, such that when combined with the visible sector of eq. (4.4) all RR tadpole constraints and K-theory constraints are satisfied. We have included in this list all three $N_{\rm flux}^{\rm max} = 9$ solutions we have found.

charges. Thirdly, we can choose a brane with positive values for two of the four charges, and with the other two charges being zero. These are the only possibilities which can be supersymmetric somewhere on Kähler moduli space, and which can solve the RR tadpole conditions.

Again, we used a computer program to search for as many such solutions as we could find. We performed a nearly exhaustive[8] search through all hidden sector brane configurations that would satisfy all constraints and introduce 2 NSNS tadpole constraints. Among the nearly 1000 classes of solutions we found, we identified several solutions in which all NSNS tadpole conditions can be solved simultaneously at a point (0D surface) in Kähler moduli space. In each of these cases, however, no flux can be turned on, which means that we have not fixed the complex structure moduli. There is one case where all NSNS tadpoles can be solved with $N_{\rm flux}$=1, but in that case one of the Kähler moduli is infinite. There were many more solutions we obtained where all such constraints cannot be satisfied.

Again, though, the solutions in which the NSNS tadpoles do not vanish are not necessarily non-supersymmetric. They instead are solutions in which the simple brane configuration must be deformed in order for the solution to become supersymmetric. We have found many of these models [46], of which we list in table 3 a representative for each value of $N_{\rm flux}^{\rm max}$ from 1 to 9, which is the maximum value of $N_{\rm flux}$ we obtained. We list all three $N_{\rm flux}^{\rm max} = 9$ solutions we found.

The existence of high flux $N_{\rm flux} = 9$ solutions is an interesting result, because we know quite generally that the number of flux vacua will grow rapidly with $N_{\rm flux}$. If the charge arising from fluxes is much larger than the number of complex structure moduli, the number

[8]We say "nearly exhaustive" because in this case, unlike the 1D Kähler surfaces case of the previous subsection, we performed a randomized Monte Carlo search for all solutions, and waited until it appeared all solutions were obtained of the general structure we were searching. However, we cannot guarantee all were found.

of flux vacua will grow as $\frac{N_{\text{flux}}^{2n+2}}{(2n+2)!}$. If it is smaller, the number of vacua is expected to grow as $e^{\sqrt{2\pi(2n+2)N_{\text{flux}}}}$ [10]. For all solutions we find, the exponential scaling is appropriate. In particular, for our highest flux model ($N_{\text{flux}} = 9$) we find that the number of vacua is approximately 10^{33}. Thus, a hidden sector with the largest value of N_{flux} will have the largest number of flux vacua, and thus is most likely to have "accidental" cancellations of phenomenological interest, such as the cancellation of the bare cosmological constant [7] against its quantum correction.

It is important to note that we are not attempting to make a probabilistic prediction — that a certain hidden sector is more likely to be chosen by nature because it is realized in more flux vacua. Instead, it is simply the statement that a hidden sector with more flux vacua is more likely to exhibit one flux vacua which, purely accidentally, happens to exhibit another property of phenomenological interest, irrespective of whether or not nature chooses this vacuum (for example, see [47]).

4.3 High flux model

We would like to analyze some details of one of our $N_{\text{flux}} = 9$ solutions to give the reader a feel for the particle content and the symmetry breaking patterns of this theory. The brane content of the hidden sector that goes along with the visible sector of eq. (4.4) is

$$\begin{aligned} N_e &= 2(-6,1)(-5,1)(-6,1) \\ N_f &= 2(4,1)(2,-1)(1,-1) \\ N_g &= 4(1,0)(0,1)(0,-1)\,. \end{aligned} \tag{4.9}$$

The exotic states that have SM quantum numbers arise from the intersections of hidden sector branes with visible sector branes. The resulting states will be in the bifundamentals of the two gauge groups from each respective brane.

In this case the chiral exotic matter is given by table 4. For simplicity we are overlapping the U(3) and $\mathrm{U}(1)_{B-L}$ brane stacks to form a $\mathrm{U}(4) = \mathrm{SU}(4) \times \mathrm{U}(1)_a$ stack. We note that the total number of chiral exotics under SU(3) is zero, which is as it should be. Thus, all SU(3) exotics have the chance of obtaining mass, which is required for phenomenological viability. The number of exotics under $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$ is even, which allows for the possibility of forming $\mathrm{SU}(2)_{L/R}$ gauge invariants to give mass to all of these exotics.

We also have 56 SU(3) exotics, but with no net chirality. As such, they can all be made massive as well. In particular, we might give vevs to the scalar of one of the 150 chiral multiplets with U(1) charges $Q_e = -1$, $Q_f = 1$, and to the scalar of one of the 54 chiral multiplets with charge $Q_f = -2$. This amounts to the brane recombination

$$[e] + [f] + [e'] + [f'] \to [j] \tag{4.10}$$

and removes all chiral SU(3) and $\mathrm{SU}(2)_{L,R}$ exotics.

What's most interesting about this model is the high amount of flux. As stated above, the number of flux vacua is $\sim 10^{33}$, multiplied by other prefactors. One such prefactor arises from the integration of the vacuum density over the complex structure

Sector	$N_{\rm copies}$	$\mathrm{SU}(4)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R$	$\mathrm{U}(1)_a$	$\mathrm{U}(1)_e$	$\mathrm{U}(1)_f$
(ae)	24	$(\bar{4},1,1)$	-1	1	0
(ae')	18	$(4,1,1)$	1	1	0
(af)	10	$(4,1,1)$	1	0	-1
(af')	4	$(\bar{4},1,1)$	-1	0	-1
(be)	36	$(1,2,1)$	0	1	0
(bf)	4	$(1,2,1)$	0	0	-1
(ce)	30	$(1,1,2)$	0	1	0
(cf)	8	$(1,1,2)$	0	0	-1
(ef)	150	$(1,1,1)$	0	-1	1
(ef')	98	$(1,1,1)$	0	1	1
(eg)	30	$(1,1,1)$	0	1	0
(fg)	2	$(1,1,1)$	0	0	1
(ee')	530	$(1,1,1)$	0	2	0
(ff')	10	$(1,1,1)$	0	0	-2

Table 4: We list in this table all exotic chiral multiplets charged jointly under a SM gauge group and a hidden sector gauge group for the $N_{\rm flux}=9$ model discussed in the text.

moduli space. Another arises from the fact that only a fraction of the vacua found here will have moduli that are stabilized in a self-consistent regime (i.e., small coupling and volume larger than string scale). If we assume that these prefactors are not too small, we can still easily obtain at least one vacuum state with cosmological constant nonzero but at or below the current measured value; we need of $\mathcal{O}(10^{30})$ vacua for that. If the number of vacua is significantly larger than $\mathcal{O}(10^{30})$, it is possible that one could also have accidental fine-tuning of other observables to be consistent with current experiment or theoretical prejudice, such as gauge coupling unification, cold dark matter abundance, acceptable CP violating phases, etc. This is what makes high $N_{\rm flux}$ vacua that we have found especially interesting, since landscape statistics has a chance of enabling some good features of the model to be present simultaneously in at least one vacuum.

5. Phenomenological considerations

Our goal in the above was to construct SM embeddings within type IIB flux compactifications. We were drawn to the framework of $T^6/\mathbb{Z}_2\times\mathbb{Z}_2$ orientifold compactification with $\mathrm{SU}(3)\times\mathrm{SU}(2)_L\times\mathrm{SU}(2)_R\times\mathrm{U}(1)$ gauge group, partly due to its simplicity. The overall framework puts restrictions on the phenomenology, some of which are challenges to phenomenological viability.

Likewise, if we were to focus primarily on low-energy effective theory model building, with little consideration for how such a model could be embedded in a more complete framework like string theory, we would likely arrive at models that are challenges to string model viability. (The pure SM is of course one such model.) In this section we discussion some of the phenomenological implications of the framework we have detailed above, with the goal of gaining further understanding of what aspects are viable from both the phenomenological and string perspectives.

Exotic gauge symmetries. Additional gauge symmetries beyond the SM are common in many approaches to string model building. There is no exception here, and the general reason for this in our case is that right-handed leptons in the SM spectrum are charged only under hypercharge $U(1)_Y$ and not SU(3) color or weak $SU(2)_L$. Since matter states arise from their intersections between brane stacks, yielding bifundamental representations, we need a fourth brane stack to have intersection with a "hypercharge stack" (or equivalent) to get right-handed leptons. Therefore, there is always need for some exotic gauge symmetry arising from this fourth stack. Additional gauge symmetries can arise from more hidden sector brane stacks

There is a substantial body of literature [37] on the phenomenology of gauge bosons associated with exotic non-SM symmetries at both the TeV scale and the intermediate scale. At the level of our model building, we take no position on what scale extra symmetries are most likely to appear. We only note that their presence is required somewhere in the energy continuum.

In addition to normal extra U(1) gauge symmetries, the generalized Green-Schwarz mechanism in these theories enables the possibility of having a low-scale global symmetry exactly preserved, with perhaps some small breaking due to a small vev, whereas the corresponding gauge bosons of the symmetry are very massive. The large mass comes from a Stueckelberg term in the potential, that can be seen after the shift is made to cancel the symmetry's anomaly. A U(1) symmetry with a Stueckelberg mass can be considered somewhat generic in this framework, and the phenomenology of this case is unique and rich [36].

Supersymmetric unification. The theories we have analyzed above are supersymmetric, in the sense that each brane stack satisfies the NSNS tadpole constraint and preserves supersymmetry somewhere in Kähler moduli space. Supersymmetry breaking can occur through a myriad of possibilities, but each model has the prospect of supersymmetry breaking [38] giving rise to soft masses anywhere between the weak scale and the Planck scale. Thus, softly broken supersymmetry is a phenomenological implication of this scenario, but it is not yet clear at what scale it should be found.

One of the attractive features of low-energy supersymmetric theories is the apparent unification of gauge couplings if the three gauge couplings of the minimal supersymmetric SM are renormalization group evolved up to the high scale. This may be a profound clue to nature or an interesting accident. In any event, being a theory with different brane stacks for different gauge groups, gauge coupling unification is by no means automatic or expected in our type IIB theories discussed above.

Pati-Salam unification is possible if we allow both the U(3) and U(1) brane stacks to be on top of each other. However, the scale at which the $SU(4)_{\rm PS}$ is recovered is a model building, or rather a model analyzing, detail that is as yet unknown. Nevertheless, it would be of interest to consider partial unification of the SM into SU(4) as a generic phenomenological outcome, and see what restrictions this would have on the spectrum[9] from a bottom up point of view.

Gauge symmetry breaking, R-parity and neutrino masses. One appealing way to break $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ down to the SM is to first condense a $(1,3,2)$ field to break $\mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L} \to \mathrm{U}(1)_Y$ and then condense a $(2,2,0)$ bidoublet field to break $\mathrm{SU}(2) \times \mathrm{U}(1)_Y$ down to $\mathrm{U}(1)_{em}$. The advantages of this approach are that the seesaw mechanism can induce small neutrino masses and R-parity can be retained as a discrete $\mathbb{Z}_2$ subgroup of $\mathrm{U}(1)_{B-L}$ [49].

The spectrum allowed in the visible sector of our D-brane configuration does not allow this symmetry breaking pattern. Instead, the first step of the symmetry breaking pattern is accomplished through an $\mathrm{SU}(2)_R$ doublet field, such as the right-handed slepton. Indeed, an equivalent description is that of brane recombination of the c and d branes, which is equivalent to veving a field charged under the two branes along a flat direction [18].

It is possible to generate viable neutrino mass spectra even restricting ourselves to the matter content available to us, by using one of the many approaches within left-right model building [50]. However, for symmetry breaking accomplished by right-slepton field condensation, R-parity is spontaneously broken since an odd lepton-number is being carried into the vacuum. Although this does not allow the lightest supersymmetric particle to be the dark matter in a straightforward way, dark matter could arise from another source such as vector-like matter with their own discrete symmetry or axions. Proton decay is also not automatically a problem, as only lepton number is spontaneously broken. Baryon number can stay preserved from the perspective of this gauge symmetry breaking pattern.

Supersymmetry breaking. In the previous sections we have searched for supersymmetric vacua which admit SM gauge group and matter content, but have also discussed some aspects of supersymmetry breaking. In any realistic string vacuum, experiment tells us that we have no moduli and no supersymmetry. In general, we would expect that supersymmetry breaking effects will generate potentials for generic scalars. But these potentials can very easily destabilize our solution, and in general it is quite difficult to maintain enough control over the calculation to be sure that we really have nonsupersymmetric solutions (e.g., see [40]). Thus, one of the central ideas of flux vacua counting has been to find supersymmetric solutions with no moduli, and then add supersymmetry breaking effects at a lower scale. If moduli can be fixed in a supersymmetric compactification, then we certainly know that our solution is stable. When we then add supersymmetry breaking effects at a lower scale, we can be confident that the solutions is not destablized because the various scalars have already been given a mass at a much higher scale. The solution may be deformed slightly by the supersymmetry-breaking effects, but it could not be destabilized.

A fundamental point in our search for SM flux vacua is that this entire story cannot proceed as before: if we want to get the SM, we cannot fix all moduli before breaking

[9]Tests of all unification propositions must input the known masses and couplings, which then often puts strong restrictions on the unknown masses and couplings of the model.

supersymmetry. The matter content of the SM will include chiral fermions. Their superpartners will be moduli unless supersymmetry is broken. As a result, any type of counting which focusses on vacua with no moduli must input supersymmetry breaking.

Although we cannot fix all moduli at a very high scale, we can fix many closed string moduli. This is good, because in general it has proved harder to understand whether potentials for these moduli destabilize solutions or simply deform them. On the other hand, instabilities involving open string moduli tend to resolve themselves by the annihilation of branes, without the destabilization of the space-time compactification itself.

But it will be true that the stabilization of at least some Kähler moduli will be inextricably linked to supersymmetry breaking. We require $\mathcal{A}_2 = \mathcal{A}_3$ in order to avoid undesirable deformations of the SM branes, but we generically cannot expect that this condition will arise from the F-term equations. Instead, we must demand that soft-masses arising from supersymmetry breaking give masses to the squarks and sleptons, with the equations of motion arising from the D-term potential (3.10) then fixing $\mathcal{A}_2 = \mathcal{A}_3$. This suggests that at least some of the Kähler moduli might receive a mass which, like the squark and slepton masses, is set by the supersymmetry breaking scale.

Due to these uncertainties, we have not focussed on any particular method of supersymmetry breaking. On the other hand, since it is an important part of the overall puzzle, we will list mechanisms by which supersymmetry breaking can be achieved, how the different approaches can be integrated into the string framework, and what kind of phenomenological implications they have.

The most obvious way to break supersymmetry is to add a badly misaligned brane that cannot be supersymmetric in the Kähler moduli space. The quintessential example of this is a $\overline{D3}$ brane. Supersymmetry breaking is at tree-level and is generically of order the string scale in this case. Since they contribute negatively to the $D3$-brane charge, one would be tempted to assume that it is possible to add an arbitrary number of $\overline{D3}$ branes with an arbitrary amount of flux. However, this is not the case, as the configuration is unstable to $\overline{D3}$ annihilations with flux [40]. Furthermore, given the considerations of ref. [40], which translates to $N_{\overline{D3}} \lesssim 3N_{\text{flux}}$, it is not possible to start with an $N_{\text{flux}} = 0$ model of sec. 4.1 and construct a $N_{\text{flux}} = 1$ model with added $\overline{D3}$ branes.

Nevertheless, it is possible to add a $D3/\overline{D3}$ combination to any $N_{\text{flux}} > 0$ model and obtain supersymmetry breaking. It is expected that the $\overline{D3}$ will want to locate in a warped throat, thereby possibly suppressing what would otherwise be a high-scale supersymmetry breaking mass terms to the weak scale. The supersymmetry breaking in this case would be D-term[10] and could lead to hierarchically larger scalar masses than gaugino masses.

IASD and ISD(0,3) fluxes[11] could both contribute to F-term-like supersymmetry breaking, which gives rise to soft masses of open string states connecting $D7$ branes to

[10]Note that there are many claims in the literature to the effect that $\overline{D3}$-branes break supersymmetry explicitly. However, recent work suggests that this method of supersymmetry breaking may be modelled as D-term breaking [41 – 43, 11]. A clearer resolution of this situation is desirable.

[11]For ISD(0,3) fluxes, the generation of masses is less certain, as it depends on how the no-scale structure is broken (see Camara et al. in [38]). Of course, IASD fluxes will also induce F-terms for complex structure moduli.

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other branes. Both IASD and ISD supersymmetry breaking can in principle be small, and lead to weak scale supersymmetry breaking even with the standard large hierarchy of the string scale and weak scale.

We remark that in the case of ISD/IASD fluxes contributing to supersymmetry breaking, the FI D-terms arising from misaligned supersymmetric branes in Kähler moduli space will no longer be able to zero themselves out completely by appropriate choice of scalar vevs, and so they will contribute to the overall supersymmetry breaking accounting in the low-energy phenomenology. One should think of these D-terms as simply constraints that involve both the Kähler moduli and open-string moduli. Given other constraints on the Kähler and open-string moduli arising from F-terms, it may not be possible to simultaneously solve the F-term and D-term equations. However, the size of that supersymmetry breaking is controlled by the F-terms: as we tune the F-terms to zero the FI-induced D-terms must go to zero.

Lastly, gaugino condensation is another potential source of supersymmetry breaking whenever the hidden sector has a large enough gauge group and sufficiently small matter content. This supersymmetry breaking is F-term. Although gaugino condensation is not necessarily a crucial ingredient of supersymmetry breaking (fluxes can do all the work), it might nevertheless play an important role in fixing Kähler moduli through its non-perturbative dynamics.

6. Conclusions

One of the major lessons of this exercise has been the importance of the exact method by which the open string and Kähler moduli are fixed. We have seen that this is linked to the mechanism by which supersymmetry is broken. In SM flux vacua of the type we discuss, the masses given to complex structure moduli and the dilaton can be made numerically much larger than the scale of supersymmetry breaking. But the masses given to Kähler moduli may be of the same scale as those given to open-string moduli, and these are of the order of the supersymmetry breaking scale. Thus it seems that these three problems (breaking supersymmetry, fixing Kähler moduli and fixing open string moduli) may have to be dealt with simultaneously.

One can consider these various types of hidden sectors, and under the assumption that any embedding leaves all Kähler moduli fixed (without overconstraining the moduli), compare the number of flux vacua. In that case, as follows our intuition from flux vacua counting [10, 11], we see that the number of vacua scales as $e^{\sqrt{2\pi(2n+2)N_{\rm flux}}}$ where $n = 51$ is the number of complex structure moduli and $N_{\rm flux}$ is the amount of three-form flux turned on as part of the hidden sector. We thus see that we can get insight into the nature of flux vacua counting in this theory, but to have complete control over the counting, we should understand the nature of the non-perturbative corrections to the superpotential. Furthermore, it would be interesting to compare the number of flux vacua with this SM embedding to the total number of flux vacua for this orientifold compactification. Perhaps one can estimate the total number of flux vacua using techniques similar to those used in [20].

Our discussion has thus far been for a choice of discrete torsion which yields 51 complex structure moduli and 3 Kähler moduli. However, for a different choice of discrete torsion (B-field turned on the shrunken cycles) [23], we would instead find 3 complex structure moduli and 51 Kähler moduli. This would significantly reduce the scaling of the number of flux vacua with N_{flux}. In this case, we would have many Kähler moduli that can participate in constraints arising from non-perturbative superpotential corrections, but not from NSNS tadpole constraints (which only affect the 3 "toroidal" Kähler moduli). This may affect how easy it is to fix all the Kähler moduli, even if one has the right number of constraints (i.e., the toroidal set of 3 might be over-constrained, while the non-toroidal set of 48 is under-constrained).

We have identified many solutions with one or two NSNS tadpole constraints. It would be interesting to do a systematic search for hidden sectors with three or more NSNS tadpole constraints. As we discussed in the text, such solutions are likely to be much more difficult to find, as higher numbers of branes will struggle to satisfy the RR tadpole constraints. Nevertheless, one should search for their existence, and if there are solutions, determine their $N_{\text{flux}}^{\text{max}}$.

Another subtlety which we do not address is the possibility of wrapping branes on the shrunken cycles of the orientifold. Although we do not know if such branes can participate in a SM embedding, they may affect flux vacua counting within our choice of orientifold.

The approach we have used in this study can be generalized to other string compactifications beyond $T^6/\mathbb{Z}_2 \times \mathbb{Z}_2$. The important data needed for analysis of the theory are the intersection numbers and the D-brane charges, and how they interact with the tadpole constraints. In our case, the analysis was simplified by manipulating the brane wrapping numbers. In other compactifications different manipulation techniques are required, but the general procedure is the same.

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