String Model Building

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Abstract. In this talk I review some recent progress in heterotic and F theory model building. I then consider work in progress attempting to find the F theory dual to a class of heterotic orbifold models which come quite close to the MSSM.

Keywords: string phenomenology
PACS: 11.25.-w,11.25.Wx

RANDOM SEARCHES FOR THE MSSM

The literature includes a record of several searches for the MSSM. Among these have been random searches in the string landscape looking for features common to the MSSM. In particular, vacua with \( N = 1 \) supersymmetry [SUSY], the Standard Model gauge group and three families of quarks and leptons. These random searches have for the most part shown that the MSSM is an extremely rare point in the string landscape. For example, searches in Type II intersecting D-brane models [1] have found nothing looking like the MSSM in \( 10^9 \) tries. Searches in Gepner orientifolds have been a bit more successful finding one MSSM-like model for every 10,000 tries [2]. Even searches in the heterotic string, using the free fermionic construction, have shown that the MSSM is a very rare point in the string landscape [3]. The bottom line: if you want to find the MSSM, then a random search is not the way to go. In fact, MSSM-like models have been found in Type II D-brane vacua, BUT by directed searches [4] AND not random ones.

Virtues of SUSY GUTs

We will propose a very particular directed search. We will require that SUSY GUTs be incorporated at the first step. This constraint is motivated by the many virtues of SUSY GUTs.

1. They can “naturally” explain why the weak scale \( \mathcal{M}_Z < < \mathcal{M}_{\text{GUT}} \);
2. Explains charge quantization;
3. Predicts gauge coupling unification !
4. Predicts SUSY particles at LHC !
5. Predicts proton decay;
6. Predicts Yukawa coupling unification,
7. and with Family symmetry can explain the Fermion mass hierarchy;
8. Explains neutrino masses and mixing via See-Saw;
9. The LSP is a Dark Matter candidate, and
10. Can give a cosmological asymmetry in the number of baryons minus anti-baryons,
    i.e. baryogenesis via leptogenesis.

For all of these reasons we might suspect that SUSY GUTs are a fundamental component
of any realistic string vacuum. And thus if one is searching for the MSSM in the mostly
barren string landscape, one should incorporate SUSY GUTs at the first step.

**HETEROTIC MODEL BUILDING**

I will not attempt to discuss the many attempts to find the MSSM in the string landscape.
Instead let me just discuss some recent progress in heterotic model building, either on a
smooth Calabi-Yau 3-fold or in the context of orbifold contructions.

**Smooth Manifolds**

Bouchard et al. [5] have obtained an $SU(5)$ GUT model on a CY$_3$ with the following
properties. They have three families of quarks and leptons, and one or two pairs of Higgs
doublets. They accomplish GUT symmetry breaking and Higgs doublet-triplet splitting
via a Wilson line in the weak hypercharge direction. The CY$_3$ is defined by a double
elliptic fibration, i.e. two tori whose radii change as the tori move over the surface of a
sphere (see Fig. 1).

In addition, they obtain a non-trivial up Yukawa matrix given by [6]

$$\lambda_u = \begin{pmatrix} a & b & c \\ b & d & e \\ c & e & 0 \end{pmatrix}. \quad (1)$$

The parameters $a, \cdots, e$ are functions of the moduli. The down and charged lepton
Yukawa matrices are however zero and would require non-perturbative effects to change
this.
Orbifolds

Early work on orbifold constructions of the heterotic string was started over 20 years ago [7, 8, 9, 10, 11]. However, progress has been made recently [12, 13, 14, 15, 16, 17, 18, 19, 20, 21]. In a “mini-landscape” search of the $E(8) \times E(8)$ heterotic landscape [18] 223 models with 3 families, Higgs doublets and ONLY vector-like exotics were found out of a total of order 30,000 models or approximately 1 in 100 models searched looked like the MSSM! We called this a “fertile patch” in the heterotic landscape. Let me describe this focussed search in more detail.

We compactify the $E(8) \times E(8)$ heterotic string on the product of 3 two dimensional tori defined in terms of the root lattice of a $G(2)$, $SU(3)$ and $SO(4)$ root lattices (see Fig. 2). These tori are chosen since they are invariant under the symmetry $\mathbb{Z}_6 - II = \mathbb{Z}_3 \times \mathbb{Z}_2$ which defines the orbifold. We also imbed the orbifold symmetry into the $E(8) \times E(8)$ gauge lattice via a shift vector, $V_6$ and add Wilson lines, consistent with modular invariance. This has the effect of breaking the gauge group to a subgroup (without breaking the rank). Note 5 cycles on the tori are small of order the string length, while one cycle in the $SO(4)$ torus is assumed to be large.

First consider the action of the $\mathbb{Z}_3$ orbifold with a Wilson line in the $SU(3)$ torus (see Fig.3). Let us focus on the states transforming under the visible sector $E(8)$. The massless states from the untwisted sector and $(G(2), SU(3))$ twisted sectors transform as N=1 super-multiplets in 6 dimensions (or N=2 in terms of effective 4 dimensional super-multiplets) as a vector hyper-multiplet in the adjoint representation of $SU(6)$ and chiral hyper-multiplets in the $20 + 9(6 + \bar{6})$ dimensional representations. All of these states move freely in the untwisted torus. Thus they describe the degrees of freedom for a 6 dimensional $SU(6)$ orbifold GUT.

We then perform the additional $\mathbb{Z}_2$ orbifold and add a Wilson line along the long cycle in the $SO(4)$ torus (see Fig. 4). The resulting theory has the Standard Model gauge symmetry in the visible sector. The shift vector $V_2$ acts as a parity $P$ on the left, breaking $SU(6)$ to $SU(5)$. On the right, the combination of $V_2 + W_2$ gives a parity $P'$ which breaks further to the Standard Model gauge group. In addition, two light families are located at the two $SO(4)$ fixed points. In Fig. 5 this is represented in terms of an
effective 5 dimensional orbifold $SU(6)$ GUT field theory. Note, the two $SO(4)$ fixed points are very special. In terms of the full 10D string theory, these correspond to “local” $SO(10)$ fixed points. Thus families come in complete $SO(10)$ representations. This is NOT an accident, but was enforced from the very beginning. The Higgs and third family come from the 5D “bulk.” They can be seen as follows. First consider the vector hypermultiplet. In terms of 6 x 6 hermitian matrices we have the gauge multiplet on the left (Eqn. 2) and the chiral adjoint on the right. The subscripts label the charge under the parity $(P, P')$. All states include Kaluza-Klein modes beginning at the compactification scale $M_C = 1/R$, while only $(++)$ states contain massless modes. On the left, these include the $SU(3) \times SU(2) \times U(1)_Y$ SM gauge sector. On the right we have one pair of Higgs doublets. Note, the orbifolding has succeeded in splitting the Higgs doublets and triplets. This model is a string theory realization of gauge-Higgs unification.

$$
\begin{pmatrix}
V^3_{++} & V^2_{+-} & V^3_{-+} \\
V^2_{++} & V^2_{-+} & V^1_{++} \\
V^1_{-+} & V^1_{-+} & V^1_{++}
\end{pmatrix}, \begin{pmatrix}
\Phi^3_{++} & \Phi^2_{+-} & \Phi^3_{-+} \\
\Phi^2_{++} & \Phi^2_{-+} & \Phi^2_{++} \\
\Phi^1_{-+} & \Phi^1_{++} & \Phi^1_{++}
\end{pmatrix}
$$

(2)

In addition the third family is contained in [in terms of effective 4D N=1 chiral superfields]

$$
(20 + 20^c) \supset Q_3 + \bar{t} + \bar{\bar{t}} \\
2(6 + 6^c) \supset L_3 + \bar{b}
$$

(3)

As a result we obtain gauge-Yukawa coupling unification with

$$
W \supset \frac{g_5}{\sqrt{\pi R}} \int_0^{\pi R} dy 20^c \Phi 20 = g_{GUT} Q_3 H_u \bar{t}
$$

(4)

The model also has a discrete non-Abelian family symmetry, $D_4$. The symmetry acts on the two light families as doublets and the third family and Higgs doublets are singlets under $D_4$. The order 8 group, $D_4$, is generated by the two operations given by $\sigma_1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}$ and $\sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$. The first is a symmetry which interchanges the two light families located at opposite sides of one cycle of the orbifolded $SO(4)$ torus. This symmetry corresponds to the geometric translation half way around the cycle which is not broken by the Wilson line lying along the orthogonal cycle. The second
is a result of so-called space-group selection rules which require an even number of states at each of these two fixed points. As a result, the theory is invariant under the action of multiplying each state located at, say, the lower fixed point by minus one. As a consequence of $D_4$ (and additional U(1) symmetries) only the third family can obtain a tree level Yukawa coupling. All others Yukawa couplings can only be obtained once the family symmetries are broken. Thus the string theory includes a natural Froggatt-Nielsen mechanism [22] for generating a hierarchy of fermion masses. [Aside, for a general analysis of discrete non-Abelian flavor symmetries obtainable in orbifold models, see [23] and for the phenomenological analysis of a $D_4$ invariant description of quark masses and flavor violation, see [24].]

**TABLE 1.** Spectrum. The quantum numbers under $SU(3) \times SU(2) \times [SU(4) \times SU(2)]^c$ are shown in boldface; hypercharge and $B - L$ charge appear as subscripts. Note that the states $s_i^\pm$, $f_i$, $\tilde{f}_i$ and $m_i$ have different $B - L$ charges for different $i$, which we do not explicitly list.

| # | irrep | label | # | irrep | label |
|---|-------|-------|---|-------|-------|
| 3 | $(3,2;1,1)_{(1/6,1/3)}$ | $q_i$ | 3 | $(3,1;1,1)_{(-2/3,-1/3)}$ | $\bar{u}_i$ |
| 3 | $(1,1;1,1)_{(1,1)}$ | $\bar{e}_i$ | 8 | $(1,2;1,1)_{(0,0)}$ | $m_i$ |
| 4 | $(3,1;1,1)_{(1/3,-1/3)}$ | $d_i$ | 1 | $(3,1;1,1)_{(-1/3,1/3)}$ | $\bar{d}_i$ |
| 4 | $(1,2;1,1)_{(-1/2,-1)}$ | $\ell_i$ | 1 | $(1,2;1,1)_{(1/2,1)}$ | $\bar{\ell}_i$ |
| 1 | $(1,2;1,1)_{(-1/2,0)}$ | $\phi_i$ | 1 | $(1,2;1,1)_{(1/2,0)}$ | $\phi_i$ |
| 6 | $(3,1;1,1)_{(1/3,2/3)}$ | $\check{\delta}_i$ | 6 | $(3,1;1,1)_{(-1/3,-2/3)}$ | $\bar{\delta}_i$ |
| 14 | $(1,1;1,1)_{(1/2,\ast)}$ | $s_i^+$ | 14 | $(1,1;1,1)_{(-1/2,\ast)}$ | $\bar{s}_i$ |
| 16 | $(1,1;1,1)_{(0,1)}$ | $\bar{n}_i$ | 13 | $(1,1;1,1)_{(0,-1)}$ | $n_i$ |
| 5 | $(1,1;1,2)_{(0,1)}$ | $\eta_i$ | 5 | $(1,1;1,2)_{(0,-1)}$ | $\eta_i$ |
| 10 | $(1,1;1,2)_{(0,0)}$ | $h_i$ | 2 | $(1,2;1,2)_{(0,0)}$ | $y_i$ |
| 6 | $(1,1;4,1)_{(0,\ast)}$ | $f_i$ | 6 | $(1,1;\bar{3},1)_{(0,\ast)}$ | $\bar{f}_i$ |
| 2 | $(1,1;4,1)_{(-1/2,-1)}$ | $f_i^-$ | 2 | $(1,1;\bar{3},1)_{(1/2,1)}$ | $\tilde{f}_i^+$ |
| 4 | $(1,1;1,1)_{(0,\pm 2)}$ | $\chi_i$ | 32 | $(1,1;1,1)_{(0,0)}$ | $s_i^0$ |
| 2 | $(3,1;1,1)_{(-1/6,2/3)}$ | $\bar{v}_i$ | 2 | $(3,1;1,1)_{(1/6,-2/3)}$ | $v_i$ |
When the SM singlet fields (Table 1) obtain VEVs, we have checked that all vector-like exotic states and unwanted U(1) gauge bosons obtain mass; leaving only the MSSM states at low energy. In addition the $\chi$ fields spontaneously break B-L leaving over a discrete $\mathbb{Z}_2$ matter parity under which all quarks and leptons are odd and Higgs doublets are even. This symmetry enforces an exact R-parity forbidding the baryon or lepton number violating operators, $\bar{U} \bar{D} \bar{D}$, $\bar{Q} \bar{L} \bar{D}$, $\bar{L} \bar{H} u$.

Finally the mu term vanishes in the supersymmetric limit. This is a consequence of the fact that the coefficient of the $H_u H_d$ term in the superpotential has vacuum quantum numbers. Thus any product of SM singlets which can appear in the pure singlet superpotential can appear as an effective mu term. In fact both the mu term and the singlet superpotential vanish to order 6 in the product of fields. Hence in the supersymmetric vacuum the VEV of the superpotential and the mu term both vanish. As a consequence, when supergravity is considered, the supersymmetric vacuum is consistent with flat Minkowski space. We also obtain non-trivial effective Yukawa matrices. The charged fermion Yukawa matrices are

$$Y_u = \begin{pmatrix} \bar{s}_5^5 & \bar{s}_5^6 & \bar{s}_5^6 \\ \bar{s}_6^5 & \bar{s}_6^5 & \bar{s}_6^6 \\ \bar{s}_6^6 & \bar{s}_6^6 & 1 \end{pmatrix}, \quad Y_d = \begin{pmatrix} \bar{s}_5^5 & \bar{s}_5^5 & 0 \\ \bar{s}_6^5 & \bar{s}_6^5 & 0 \\ \bar{s}_6^6 & \bar{s}_6^6 & 0 \end{pmatrix}, \quad Y_e = \begin{pmatrix} \bar{s}_5^5 & \bar{s}_5^5 & \bar{s}_5^6 \\ \bar{s}_6^5 & \bar{s}_6^5 & \bar{s}_6^6 \\ \bar{s}_6^6 & \bar{s}_6^6 & 0 \end{pmatrix}.$$ (5)

where $\bar{s}_n^i$ represents a polynomial in SM singlets beginning at order $n$ in the product of fields. And we have shown that the three left-handed neutrinos get small mass due to a non-trivial See-Saw mechanism involving the 16 right-handed neutrinos and their 13 conjugates. All in all, this “benchmark” model looks very much like the MSSM!

Note, Yukawa couplings, gauge couplings and Vector-like exotic masses are functions of moduli (along SUSY flat directions). Some of these moduli are blow up modes for some, BUT NOT ALL, of the orbifold fixed points. In fact, two fixed points are NOT blown up!

**F THEORY / TYPE IIB**

Now let’s change directions and discuss some recent progress in F theory model building [25, 26, 27, 28, 29, 30, 31]. An $SU(5)$ GUT is obtained on a D7 “gauge” brane $S \times R^{3,1}$. D7 “matter” branes on $S' \times R^{3,1}$ also exist with chiral matter in 6D on $\Sigma \times R^{3,1}$ at the intersection of the gauge and matter branes (Fig. 6). Yukawa couplings enter at the triple intersections $\Sigma_1 \cap \Sigma_2 \cap \Sigma_3$ of matter sub-manifolds (Fig. 7).

$SU(5)$ is broken to the SM gauge group with non-vanishing hypercharge flux $\langle F_Y \rangle$. Note, this is not possible in the Heterotic string! This is because of the term in the Lagrangian $\int d^{10}x (dB + A_Y \wedge \langle F_Y \rangle)^2$ which leads to a massive hypercharge gauge boson and consequently a massive photon. In addition, $\langle F_Y \rangle$ on the Higgs brane leads to doublet-triplet splitting. Finally, spinor representations of $SO(10)$ are possible in F theory; although they are not possible in the perturbative type IIB string.

It was also demonstrated that a fermion flavor hierarchy is natural [39, 33], due to flux in the $z_2 - z_3$ surface breaking geometric flavor symmetry, with Yukawa matrices of the
The figure represents 3 complex planes labeled by $z_i, i = 1, 2, 3$. The four dimensional blue surface is the gauge brane and the matter brane is red. Open strings at the intersection give chiral matter in bi-fundamental representations.

Yukawa couplings are generated at the intersection of two quark branes with a Higgs brane. The form

$$\lambda \sim \begin{pmatrix} \varepsilon^8 & \varepsilon^6 & \varepsilon^4 \\ \varepsilon^6 & \varepsilon^4 & \varepsilon^2 \\ \varepsilon^4 & \varepsilon^2 & 1 \end{pmatrix}$$

Finally, gravity decouples (i.e. $M_{Pl} \rightarrow \infty$) with a non-compact $z_1$ direction. These are so-called “local” constructions. A bit of progress has also been made in “global” compact constructions [29, 31].

**HETEROTIC - F THEORY DUALS**

F theory defined on a CY$_4$ is dual to the heterotic string defined on a CY$_3$ (Fig. 8). We are now attempting to construct the F theory dual to our MSSM-like models [34]. The

The heterotic string (left) is compactified on the product of two tori fibered over a common 2-sphere (defining a Calabi-Yau three-fold), while F theory (right) is defined in terms of a torus (whose complex structure defines the coupling strength of a Type IIB string) fibered over the product of two 2-spheres and the additional torus fibered over a common 2-sphere (defining a Calabi-Yau four-fold).
motivation is three-fold.

1. We have also found three family MSSM-like models using an $SO(5) \times SO(5) \times SO(4)$ torus [35]. This suggests a larger class of MSSM-like models. We hope to find a more general description of MSSM-like models, i.e. all models in the same universality class.

2. It may also provide a general understanding of moduli space, since from the orbifold viewpoint we must first construct the superpotential before it is possible to identify the moduli, and

3. It may help us understand moduli stabilization and SUSY breaking.

Let me now discuss why we think the F theory duals may exist. First uplift our $E(8) \times E(8)$ orbifold models onto a smooth Calabi-Yau manifold. Recall that after the first $\mathbb{Z}_3$ orbifold plus Wilson line $W_3$ we find a 6D $SU(6)$ orbifold GUT compactified on $(T_2)^2/\mathbb{Z}_3 \times T_2$. The complete massless spectrum in this case (including the hidden sector) is given in Table 2 [36]. This spectrum satisfies the gravity anomaly constraint $N_H - N_V^6 + 29N_T = 273$, where $(N_H = 320, N_V^6 = 76, N_T = 1)$ are the number of (hyper-, vector, tensor) multiplets.

### Table 2

| Multiplet Type | Representation | Number |
|----------------|----------------|--------|
| tensor         | singlet        | 1      |
| vector         | $(35,1,1) \oplus (128,1)$ | 35 + 28 |
|                | $\oplus (11,8) \oplus 5 \times (1,1,1)$ | 8 + 5 |
| hyper          | $(20,1,1) \oplus (18_{1+c+s},1) \oplus 4 \times (1,1,1)$ | 20+24+4 |
|                | $\oplus 9 \times \{(6,1,1) \oplus (\bar{6},1,1)\}$ | 108 |
|                | $\oplus 9 \times \{(1,1,3) \oplus (1,1,\bar{3})\}$ | 54 |
|                | $\oplus 3 \times (1,8_{1+c+s},1)$ | 72 |
|                | $\oplus 36 \times (1,1,1)$ | 36 |
|                | SUGRA singlets | 2 |

However, using the results of Bershadsky et al. [37] we show that the $E(8) \times E(8)$ heterotic string compactified on a smooth $K_3 \times T_2$, with instantons imbedded into $K_3$, is equivalent to the orbifolded theories. For example, with 12 instantons imbedded into an $SU(3) \times SU(2)$ subgroup of the first $E(8)$ leaves an $SU(6)$ 6D GUT with the massless hypermultiplets $(20 + cc) + 18(6 + c.c.)$. Then imbedding 12 instantons into an $E(6)$ subgroup of the second $E(8)$ plus additional higgsing leaves an unbroken $SO(8)$ gauge symmetry with the massless hypermultiplets $4(8v + 8s + 8c + c.c.)$. This is identical to the massless spectrum of the orbifold GUT, IF we neglect the additional $SU(3) \times U(1)^5$ symmetry which must be broken when going to the smooth limit. In fact, it is expected that the blow up modes necessary to smooth out the orbifold singularities will carry charges under some of the orbifold gauge symmetries; spontaneously breaking these symmetries. Therefore $K_3 \times T_2$ with instantons is the smooth limit of $T_4/\mathbb{Z}_3 \times T_2$ orbifold plus Wilson line. In addition, it was shown that F theory compactified on a Calabi-
At the 6D level we have a Calabi-Yau three-fold times a torus. The gauge branes are located at the points \( z_1 = 0 \) and \( \infty \). The matter branes are located schematically at the solid point in \( Z_2 \).

At the 4D level we fiber the last torus over the 2-sphere and retain two orbifold fixed points. These two fixed points are where the two light families are conjectured to be located. In addition, as long as the fixed points remain, the Wilson line wrapping the last torus is stable.

Yau 3-fold [defined in terms of a torus \( T_2 \) fibered over the space \( F_n \times T_2 \)] is dual to an \( E(8) \times E(8) \) heterotic string compactified on \( K_3 \times T_2 \) with instantons [37] (see Fig. 9).

Pictorially we see that the \( SU(6)[SO(8)] \) gauge branes are localized at the upper [lower] points on the \( z_1 \) 2-sphere (Fig. 9). These 7 branes wrap the four dimensional surface \( S = (z_2, z_3) \). The matter 7 branes intersect the gauge branes at points in \( z_2 \) and wrap the four dimensional surface \( S' = (z_1, z_3) \). The intersection of the matter and gauge branes is along the two dimensional surface \( \Sigma = (z_3) \).

We now need to break the 6D \( SU(6) \) GUT to \( SU(5) \) and then to the Standard Model. At the same time we must break the N=1 SUSY in 6D to N=1 in 4D. This is accomplished by acting with the \( Z_2 \) orbifold on the torus and the 2-sphere. A \( U(1) \) flux in \( SU(6) \) on the gauge and matter branes breaks \( SU(6) \) to \( SU(5) \). The breaking to the Standard Model requires a Wilson line on the torus. However, we now encounter a possible obstruction to finding the F theory dual of our Heterotic orbifold model. We need to keep two orbifold fixed points (Fig. 10) -

1. otherwise hypercharge gets mass [38],
2. and the Wilson line shrinks to a point, since \( T_2/Z_2 \) is topologically equivalent to a 2-sphere;
3. and blow up modes on the heterotic side leave two orbifold fixed points.

In addition, on the heterotic side the two light families are located at 4D orbifold fixed points. We expect that on the F theory side they will be located on \( D_3 \) branes fixed at the two remaining 4D fixed points.
CONCLUSION

Of course, there are other model building considerations which I have no time to discuss. These include,

- Gauge coupling unification,
- Proton decay from dimension 5 and 6 operators,
- Supersymmetry breaking and sparticle masses (see [39, 40]),
- Moduli stabilization,
- Cosmological constant and the possible $10^{500}$ vacua; or
- Cosmology.

In conclusion, I believe that it is important to test string theory. In this talk I have discussed several new ideas in string model building. I have made the case that orbifold and local GUTs may be necessary ingredients for finding the MSSM in the string landscape. Finally, global F theory constructions may open a new window onto the general MSSM landscape. The first critical test comes with data at the LHC and the possible discovery of supersymmetry.

ACKNOWLEDGMENTS

I would like to acknowledge the hospitality and partial support of the Stanford Institute for Theoretical Physics where this work was begun. I also received partial support from DOE grant, DOE/ER/01545-883.

REFERENCES

1. F. Gmeiner, R. Blumenhagen, G. Honecker, D. Lust and T. Weigand, JHEP 0601, 004 (2006) [arXiv:hep-th/0510170].
2. P. Anastasopoulos, T. P. T. Dijkstra, E. Kiritsis and A. N. Schellekens, Nucl. Phys. B 759, 83 (2006) [arXiv:hep-th/0605226].
3. K. R. Dienes, M. Lennek, D. Senechal and V. Wasnik, Phys. Rev. D 75, 126005 (2007) [arXiv:0704.1320 [hep-th]].
4. R. Blumenhagen, M. Cvetic, P. Langacker and G. Shiu, Ann. Rev. Nucl. Part. Sci. 55, 71 (2005) [arXiv:hep-th/0502005].
5. V. Bouchard and R. Donagi, Phys. Lett. B 633, 783 (2006) [arXiv:hep-th/0512149].
6. V. Bouchard, M. Cvetic and R. Donagi, Nucl. Phys. B 745, 62 (2006) [arXiv:hep-th/0602096].
7. L. J. Dixon, J. A. Harvey, C. Vafa and E. Witten, Nucl. Phys. B 261, 678 (1985).
8. L. J. Dixon, J. A. Harvey, C. Vafa and E. Witten, Nucl. Phys. B 274, 285 (1986).
9. L. E. Ibanez, H. P. Nilles and F. Quevedo, Phys. Lett. B 187, 25 (1987).
10. L. E. Ibanez, J. E. Kim, H. P. Nilles and F. Quevedo, Phys. Lett. B 191, 282 (1987).
11. A. Font, L. E. Ibanez, F. Quevedo and A. Sierra, Nucl. Phys. B 331, 421 (1990).
12. T. Kobayashi, S. Raby and R. J. Zhang, Phys. Lett. B 593, 262 (2004) [arXiv:hep-ph/0403065].
13. T. Kobayashi, S. Raby and R. J. Zhang, Nucl. Phys. B 704, 3 (2005) [arXiv:hep-ph/0409098].
14. W. Buchmuller, K. Hamaguchi, O. Lebedev and M. Ratz, Phys. Rev. Lett. 96, 121602 (2006) [arXiv:hep-ph/0511035].
15. W. Buchmuller, K. Hamaguchi, O. Lebedev and M. Ratz, Nucl. Phys. B 785, 149 (2007) [arXiv:hep-th/0606187].
16. O. Lebedev, H. P. Nilles, S. Raby, S. Ramos-Sanchez, M. Ratz, P. K. S. Vaudrevange and A. Wingerter, Phys. Lett. B 645, 88 (2007) [arXiv:hep-th/0611095].
17. O. Lebedev, H. P. Nilles, S. Raby, S. Ramos-Sanchez, M. Ratz, P. K. S. Vaudrevange and A. Wingerter, Phys. Rev. Lett. 98, 181602 (2007) [arXiv:hep-th/0611203].
18. O. Lebedev, H. P. Nilles, S. Raby, S. Ramos-Sanchez, M. Ratz, P. K. S. Vaudrevange and A. Wingerter, Phys. Rev. D 77, 046013 (2008) [arXiv:0708.2691 [hep-th]].
19. J. E. Kim and B. Kyae, arXiv:hep-th/0608085.
20. I. W. Kim, J. E. Kim and B. Kyae, Phys. Lett. B 647, 275 (2007) [arXiv:hep-ph/0612365].
21. J. E. Kim, J. H. Kim and B. Kyae, JHEP 0706, 034 (2007) [arXiv:hep-ph/0702278].
22. C. D. Froggatt and H. B. Nielsen, Nucl. Phys. B 147, 277 (1979).
23. T. Kobayashi, H. P. Nilles, F. Ploger, S. Raby and M. Ratz, Nucl. Phys. B 768, 135 (2007) [arXiv:hep-ph/0611020].
24. P. Ko, T. Kobayashi, J. h. Park and S. Raby, Phys. Rev. D 76, 035005 (2007) [Erratum-ibid. D 76, 059901 (2007)] [arXiv:0704.2807 [hep-ph]].
25. C. Beasley, J. J. Heckman and C. Vafa, JHEP 0901, 058 (2009) [arXiv:0802.3391 [hep-th]].
26. C. Beasley, J. J. Heckman and C. Vafa, JHEP 0901, 059 (2009) [arXiv:0806.0102 [hep-th]].
27. R. Donagi and M. Wijnholt, arXiv:0802.2969 [hep-th].
28. R. Donagi and M. Wijnholt, arXiv:0808.2223 [hep-th].
29. R. Blumenhagen, V. Braun, T. W. Grimm and T. Weigand, Nucl. Phys. B 815, 1 (2009) [arXiv:0811.2936 [hep-th]].
30. C. M. Chen and Y. C. Chung, Nucl. Phys. B 824, 273 (2010) [arXiv:0903.3009 [hep-th]].
31. J. Marsano, N. Saulina and S. Schafer-Nameki, JHEP 0908, 030 (2009) [arXiv:0904.3932 [hep-th]].
32. J. J. Heckman and C. Vafa, arXiv:0811.2417 [hep-th].
33. L. Randall and D. Simmons-Duffin, arXiv:0904.1584 [hep-ph].
34. K. Bobkov, S. Raby, A. Westphal and T. Weigand, work in progress.
35. S. Raby, A. Wingerter and P. Vaudrevange, unpublished.
36. B. Dundee, S. Raby and A. Wingerter, Phys. Rev. D 78, 066006 (2008) [arXiv:0805.4186 [hep-th]]; Addendum, Phys. Rev. D 79, 047901 (2009) [arXiv:0811.4026 [hep-th]].
37. M. Bershadsky, K. A. Intriligator, S. Kachru, D. R. Morrison, V. Sadov and C. Vafa, Nucl. Phys. B 481, 215 (1996) [arXiv:hep-th/9605200].
38. S. G. Nibbelink, J. Held, F. Ruehle, M. Trapletti and P. K. S. Vaudrevange, JHEP 0903, 005 (2009) [arXiv:0901.3059 [hep-th]].
39. J. J. Heckman and C. Vafa, arXiv:0811.2417 [hep-th].
40. L. Aparicio, D. G. Cerdeno and L. E. Ibanez, JHEP 0807, 099 (2008) [arXiv:0805.2943 [hep-ph]].