A Realistic World from Intersecting D6-Branes

Ching-Ming Chen,1 Tianjun Li,1,2 V. E. Mayes,1 and Dimitri V. Nanopoulos1,3,4

1 George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, TX 77843, USA
2 Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100080, China
3 Astroparticle Physics Group, Houston Advanced Research Center (HARC), Mitchell Campus, Woodlands, TX 77381, USA
4 Academy of Athens, Division of Natural Sciences, 28 Panepistimiou Avenue, Athens 10679, Greece

We briefly describe a three-family intersecting D6-brane model in Type IIA theory on the $\mathbb{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold with a realistic phenomenology. In this model, the gauge symmetry can be broken down to the Standard Model (SM) gauge symmetry close to the string scale, and the gauge coupling unification can be achieved. We calculate the supersymmetry breaking soft terms, and the corresponding low energy supersymmetric particle spectrum, which may be tested at the Large Hadron Collider (LHC). The observed dark matter density may also be generated. Finally, we can explain the SM quark masses and CKM mixings, and the tau lepton mass. The neutrino masses and mixings may be generated via the seesaw mechanism as well.

PACS numbers: 11.10.Kk, 11.25.Mj, 11.25.-w, 12.60.Jv

Introduction – During the last few years, intersecting D-brane models on Type II orientifolds [1], where the chiral fermions arise from the intersections of D-branes in the internal space [2] and the T-dual description in terms of magnetized D-branes [3] have been particularly interesting [4]. On Type IIA orientifolds with intersecting D6-branes, a large number of non-supersymmetric three-family Standard-like models and Grand Unified Theories (GUTs) were constructed in early stages [5]. However, there generically existed uncancelled Neveu-Schwarz-Neveu-Schwarz tadpoles in these models as well as the gauge hierarchy problem. To solve these problems, semi-realistic supersymmetric Standard-like and GUT models have been constructed in Type IIA theory on the $\mathbb{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [6, 7] and other backgrounds [8]. To stabilize the moduli via supergravity fluxes, flux models on Type II orientifolds have also been constructed [9, 10]. There are two main constraints on supersymmetric D-brane model building: (1) Ramond-Ramond (RR) tadpole cancellation conditions and (2) the requirement for four-dimensional $N = 1$ supersymmetric D-brane configurations.

However, there are two serious problems in almost all supersymmetric D-brane models: the absence of gauge coupling unification at the string scale, and the rank one problem in the Standard Model (SM) fermion Yukawa matrices. Thus, a comprehensive phenomenological study of a concrete model from the string scale to the electroweak scale has yet to be made. Although these problems can be solved in the flux models of Ref. [10] where the RR tadpole cancellation conditions are relaxed, those models are in the AdS vacua and the resulting flux induced superpotential for moduli is too complicated. Interestingly, we find that there is one and only one intersecting D6-brane model on the Type IIA $\mathbb{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold where the above problems can be solved [7, 10]. Therefore, it is desirable to study the phenomenological consequences of this model in great detail.

Model Building – We consider Type IIA string theory compactified on a $\mathbb{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold [6]. The $\mathbb{T}^6$ is a six-torus factorized as $\mathbb{T}^2 \times \mathbb{T}^2 \times \mathbb{T}^2$ whose complex coordinates are $z_i$, $i = 1$, 2, 3 for the $i$th two torus, respectively. The $\theta$ and $\omega$ generators for the orbifold group $\mathbb{Z}_2 \times \mathbb{Z}_2$, act on the complex coordinates of $\mathbb{T}^6$ as

$$
\theta : (z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3), \\
\omega : (z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3).
$$

The orientifold projection is implemented by gauging the symmetry $\Omega R$, where $\Omega$ is world-sheet parity, and $R$ is given by

$$
R : (z_1, z_2, z_3) \rightarrow (\bar{z}_1, \bar{z}_2, \bar{z}_3).
$$

Thus, there are four kinds of orientifold 6-planes (O6-planes) for the actions $\Omega R$, $\Omega R\theta$, $\Omega R\omega$, and $\Omega R\theta\omega$, respectively. There are two kinds of complex structures consistent with orientifold projection for a two torus: rectangular and tilted [6]. If we denote the homology classes of the three cycles wrapped by the D6-brane stacks as $n'_p[a_i] + m'_p[b_i]$ and $n''_p[a''_i] + m''_p[b_i]$ with $[a'_i] = [a_i] + \frac{1}{2}[b_i]$ for the rectangular and tilted tori respectively, we can label a generic one cycle by $(n'_p, l''_p)$ in either case, where in terms of the wrapping numbers $l''_p \equiv m''_p$ for a rectangular two torus and $l''_p \equiv 2m''_p = 2m'_p + n''_p$ for a tilted two torus. Moreover, for a stack of $N$ D6-branes that does not lie on one of the O6-planes, we obtain a $U(N/2)$ gauge symmetry with three adjoint chiral superfields due to the orbifold projections, while for a stack of $N$ D6-branes which lies on an O6-plane, we obtain a $USp(N)$ gauge symmetry with three anti-symmetric
chiral superfields. Bifundamental chiral superfields arise from the intersections of two different stacks \( P \) and \( Q \) of D6-branes or from one stack \( P \) and its \( \Omega R \) image \( P' \).

We present the D6-brane configurations and intersection numbers of the model in Table I and the resulting spectrum in Table II. We put the \( a' \), \( b \), and \( c \) stacks of D6-branes on the top of each other on the third two tori, and as a result there are additional vector-like particles from \( N = 2 \) subsectors.

**TABLE I: D6-brane configurations and intersection numbers.**

| \( N \) | \( (n^1, t^1) \times (n^2, t^2) \times (n^3, t^3) \) | \( n_S \) | \( n_a \) | \( n_b \) | \( n_c \) |
|---|---|---|---|---|---|
| a | \((0, -1) \times (1, 1) \times (1, 1)\) | 0 | 0 | 0 | -3 |
| b | \((3, 1) \times (1, 0) \times (1, -1)\) | 2 | -2 | 0 | 0 |
| c | \((3, -1) \times (0, 1) \times (1, -1)\) | 2 | -2 | 1 | 0 |

Chiral superfields. Bifundamental chiral superfields arise from the intersections of two different stacks \( P \) and \( Q \) of D6-branes or from one stack \( P \) and its \( \Omega R \) image \( P' \).

We present the D6-brane configurations and intersection numbers of the model in Table I and the resulting spectrum in Table II. We put the \( a' \), \( b \), and \( c \) stacks of D6-branes on the top of each other on the third two tori, and as a result there are additional vector-like particles from \( N = 2 \) subsectors.

**TABLE II: The chiral and vector-like superfields, and their quantum numbers under the gauge symmetry**

| Quantum Number | \( Q_L \) | \( Q_{2L} \) | \( Q_{2R} \) |
|---|---|---|---|
| ab | \( 3 \times (1, 1, 1, 1, 1, 1) \) | 1 | -1 | 0 |
| ac | \( 3 \times (1, 2, 1, 1, 1, 1) \) | -1 | 0 | 1 |
| a1 | \( 1 \times (4, 1, 1, 1, 1, 1) \) | 1 | 0 | 0 |
| a2 | \( 1 \times (1, 1, 1, 1, 1, 1) \) | 1 | 0 | 0 |
| b2 | \( 1 \times (1, 2, 1, 1, 1, 1) \) | 0 | 1 | 0 |
| b1 | \( 1 \times (1, 1, 1, 2, 1, 1) \) | 0 | -1 | 0 |
| c1 | \( 1 \times (1, 1, 2, 1, 2, 1) \) | 0 | 0 | -1 |
| c3 | \( 1 \times (1, 1, 2, 1, 1, 1) \) | 0 | 0 | 1 |
| c3 | \( 1 \times (1, 2, 1, 2, 1, 1) \) | 0 | 0 | -1 |
| c3 | \( 1 \times (1, 2, 2, 1, 1, 1) \) | 0 | 0 | 2 |
| c3 | \( 1 \times (1, 1, 1, 1, 2, 1) \) | 0 | 0 | 1 |

The anomalies from three global \( U(1) \)s of \( U(4)_C \), \( U(2)_L \) and \( U(2)_R \) are cancelled by the Green-Schwarz mechanism, and the gauge fields of these \( U(1) \)s obtain masses via the linear \( B \times F \) couplings. Thus, the effective gauge symmetry is \( SU(4)_C \times SU(2)_L \times SU(2)_R \). In order to break the gauge symmetry, on the first torus, we split the \( a \) stack of D6-branes into \( a_1 \) and \( a_2 \) stacks with 6 and 2 D6-branes, respectively, and split the \( c \) stack of D6-branes into \( c_1 \) and \( c_2 \) stacks with two D6-branes for each one. In this way, the gauge symmetry is further broken to \( SU(3)_C \times SU(2)_L \times U(1)_T \times U(1)_B \). Moreover, the \( U(1)_T \times U(1)_B \) gauge symmetry may be broken to \( U(1)_Y \) by giving vacuum expectation values (VEVs) to the vector-like particles with the quantum numbers \((1, 1, 1, 2, -1)\) and \((1, 1, -1, 2, 1)\) under the \( SU(3)_C \times SU(2)_L \times U(1)_T \times U(1)_B \) gauge symmetry from \( a_2 \) \( Q'_L \) intersections [7, 10].

Since the gauge couplings in the Minimal Supersymmetric Standard Model (MSSM) are unified at the GUT scale \( \sim 2.4 \times 10^{16} \) GeV, the additional exotic particles present in the model must necessarily become superheavy. To accomplish this it is first assumed that the \( U(2)_L \) and \( U(2)_R \) stacks of D6-branes lie on the top of each other on the first torus, so we have two pairs of vector-like particles with \( U(2)_L \times U(2)_R \) quantum numbers \((2, 2)\). These particles can break \( U(2)_L \times U(2)_R \) down to the diagonal \( U(2)_L \times U(2)_R \) near the string scale, and then states arising from intersections \( a_1 \) and \( a_2 \) may obtain vector-like masses close to the string scale. Moreover, we assume that the \( T'_L \) and \( S'_R \) obtain VEVs near the string scale, and their VEVs satisfy the D-flatness of \( U(1)_L \). To preserve the D-flatness of \( U(1)_L \), we assume that the VEVs of \( T'_L \) is TeV scale. We also assume that there exist various suitable high-dimensional operators in the effective theory. With \( T'_L \) and \( S'_R \), we can give the GUT-scale masses to the particles from the intersections \( c_1 \), \( c_3 \), and \( c_5 \) via high-dimensional operators. The remaining states and adjoint chiral superfields may also obtain GUT-scale masses via high-dimensional operators by the Higgs mechanism and from strong dynamics since all of the \( U(2)_R \) have negative beta functions as shown in Table III. To have one pair of light Higgs doublets, it is necessary to fine-tune the mixing parameters of the Higgs doublets. In particular, the \( \mu \) term and the right-handed neutrino masses may be generated via the following high-dimensional operators:

\[
W = \frac{y_{ijkl}^{ijkl}}{M_{St}} S_{L,L}^{ij} S_{R, R}^{ij} H_{i}^{j} H_{i}^{j} + \frac{y_{Nij}^{ijkl}}{M_{St}} T_{R}^{T} T_{R}^{T} T_{R}^{T} T_{R}^{T} \Phi_{j} F_{R}^{j} F_{R}^{j} + \Phi_{i} F_{R}^{i} F_{R}^{i} .
\]

where \( y_{ij}^{ijkl} \) and \( y_{Nij}^{ijkl} \) are Yukawa couplings, and \( M_{St} \) is the string scale. Thus, the \( \mu \) term is TeV scale and the right-handed neutrino masses can be in the range \( 10^{10} \text{ to } 10^{14} \) GeV for \( y_{ij}^{ijkl} \sim 1 \) and \( y_{Nij}^{ijkl} \sim 10^{-7} \text{ to } 10^{-3} \).

**Phenomenological Consequences** – In the string theory basis, we have the dilaton \( S \), three Kahler moduli \( T^i \), and three complex structure moduli \( U^i \) [12]. The \( U^i \) for the present model are

\[
U^1 = 3i , \quad U^2 = i , \quad U^3 = -1 + i .
\]

The corresponding moduli \( s \), \( t^i \) and \( u^i \) in the supergravity theory basis are related to the \( S \), \( T^i \) and \( U^i \) mod-
\[
\text{Re}(s) = \frac{e^{-\phi_4}}{2\pi} \left( \frac{\sqrt{U_1^2 U_2^2 U_3^2}}{|U^1 U^2 U^3|} \right), \quad \text{Re}(t') = \frac{i \sigma^I}{T^I},
\]
\[
\text{Re}(u^I) = \frac{e^{-\phi_4}}{2\pi} \left( \frac{\sqrt{U_2^1 U_3^1 U_1^1}}{|U^2 U^3 U^1|} \right),
\]
where \(\phi_4\) is the four-dimensional dilaton, \(U_I^j\) is the imaginary part of \(U^i\), and \(j \neq k \neq l \neq j\).

The holomorphic gauge kinetic function for a generic \(P\) stack of D6-branes which does not lie on one of O6-planes, is given by
\[
f_P = \frac{1}{8} \left( 2n_p \, n_s \, n_3 \, s - n_p \, l_p^2 \, l_p^3 \, u^1 - n_p^2 \, l_p^1 \, l_p^2 \, u^2 - 2n_p \, l_p^1 \, l_p^3 \, u^3 \right).
\]
Thus, the gauge couplings for \(SU(4)_C, SU(2)_L\) and \(SU(2)_R\) in our model are unified at the string scale. For simplicity, we neglect the little hierarchy between the string scale and the GUT scale, which may be explained via threshold corrections. Assuming the value of the unified gauge coupling in the MSSM, we obtain
\[
e^{-\phi_4} = 20.1.
\]
Thus, the string scale is \(\sim 2.1 \times 10^{17}\) GeV for \(M_{\text{St}} = \pi^{1/2} e^{\phi_4} M_{\text{Pl}}\) where \(M_{\text{Pl}}\) is the reduced Planck scale.

The Kähler metric for the chiral superfields from the intersections of the \(P\) and \(Q\) stacks of D6-branes is
\[
\tilde{K} \supset e^{\phi_4 + \gamma_E \sum_{i=1}^{3} \theta^{i}_{PQ} \sum_{j=1}^{3} \left[ \frac{\Gamma(1 - \theta^{i}_{PQ})}{\Gamma(\theta^{i}_{PQ})} (t^j + \bar{t}^j)^{-\theta^{i}_{PQ}} \right]},
\]
where \(\gamma_E\) is the Euler-Mascheroni constant, and \(\theta^{i}_{PQ}\) is the suitable positive angle between the \(P\) and \(Q\) stacks of D6-branes on the \(i^{th}\) two torus in units of \(1\), and can be written as a function of \(s, u^I\), and the wrapping numbers for the \(P\) and \(Q\) stacks of D6-branes.

The Kähler metric for the vector-like chiral superfields from the intersections of the \(P\) and \(Q\) stacks of D6-branes that are parallel on the \(j^{th}\) two torus and intersect on the \(k^{th}\) and \(l^{th}\) two torus is given by
\[
\tilde{K} \supset \left[ (s + \bar{s})(u^I + \bar{u}^I)(t^{k} + \bar{t}^{k})(t^{l} + \bar{t}^{l}) \right]^{-1/2}.
\]

For simplicity, we assume that only the F terms of the complex structure moduli \(u^I\) break supersymmetry and are parametrized as follows
\[
F^{u^I} = \sqrt{3} m_{3/2}(u^I + \bar{u}^I)\Theta_i, \quad \text{for } i = 1, 2, 3\,,
\]
where \(m_{3/2}\) is the gravitino mass, and \(\Theta_i\) are real numbers and satisfy \(\sum_{i=1}^{3} |\Theta_i|^2 = 1\). Then, we can calculate the gaugino masses \((M_i)\), the universal scalar masses \(m_{FL}\) and \(m_{FR}\) respectively for the left-handed and right-handed SM fermions, the universal scalar mass \(m_H\) for Higgs fields \(H_u\) and \(H_d\), and the universal trilinear soft term \(A_Y\) at the string scale \(13\). Choosing \(m_{3/2} = 1100\) GeV, \(\Theta_1 = -0.6\), \(\Theta_2 = 0.293\), \(\Theta_3 = 0.744\), \(\text{Ret}_1 = 1/6.6\), and \(\text{Ret}_2 = \text{Ret}_3 = 0.5\), we obtain the string-scale supersymmetry breaking soft terms given in Table III. Using the code SuSpect \(14\), we calculate the low energy supersymmetric particle spectrum. An example for \(\tan \beta = 46\) and positive \(\mu\) is shown in Table IV. This spectrum is consistent with all the known experiments and can be tested at the LHC. Finally, using the code MicrOMEGAs \(15\), we obtain a dark matter density \(\Omega h^2 = 0.117\) which is very close to the observed value.

\[
\begin{array}{|c|c|c|c|c|c|c|c|}
\hline
M_1 & M_2 & M_3 & m_{FL} & m_{FR} & m_H & A_Y \\
\hline
477.4 & 271.9 & 1987.8 & 1047 & 524.7 & 451.7 & 732.6 \\
\hline
\end{array}
\]

\text{TABLE III: Supersymmetry breaking soft terms (in GeV) at the string scale.}

\[
\begin{array}{|c|c|c|c|c|c|c|}
\hline
A^0 & H^0 & A^0 & H^+ & \tilde{g} & \chi^+_1 & \chi^-_1 \\
\hline
121.3 & 1016 & 1017 & 1020 & 2192 & 192.3 & 1406 \\
\hline
\chi^0_1 & \chi^0_2 & \tilde{t}_1 & \tilde{t}_2 & \tilde{t}_3 & 1947 \\
\hline
1404 & 1405 & 1542 & 1912 & 1948 & 2144 & 1763 \\
\hline
d_1/s_1 & d_2/s_2 & \tilde{g}_1 & \tilde{g}_2 & \tilde{g}_3 & 1947 \\
\hline
2146 & 234.4 & 1010 & 1000 & 550.2 & 1059 & 1056 \\
\hline
\end{array}
\]

\text{TABLE IV: Low energy supersymmetric particles and their masses (in GeV).}

\text{The SM Fermion Masses and Mixings – Because all the SM fermions and Higgs fields arise from the intersections on the first torus, we will only consider it for simplicity. The up-type quark mass matrix \(M^U\) at the GUT scale is \(16\).}
\[
\begin{pmatrix}
A^U v^1_u + E^U v^6_u & B^U v^2_u + F^U v^6_u & D^U v^3_u + C^U v^5_u \\
C^U v^3_u + D^U v^5_u & A^U v^1_u + E^U v^2_u & B^U v^2_u + F^U v^4_u \\
F^U v^4_u + B^U v^5_u & C^U v^5_u + D^U v^4_u & A^U v^1_u + E^U v^6_u
\end{pmatrix},
\]
where \(v^i_u = \langle H^+_u \rangle\), and \(e^U\) is a constant which includes the quantum corrections and the contributions to the Yukawa couplings from the second and third two tori. The theta functions \(A^U, B^U, C^U, D^U, E^U, \) and \(F^U\) are
\[
A^U = \bar{\theta} \left[ e^{U_1} \phi^{(1)} \right] (\kappa^{(1)}), \quad B^U = \bar{\theta} \left[ e^{U_1 + \frac{1}{3}} \phi^{(1)} \right] (\kappa^{(1)}) ,
\]
\[
C^U = \bar{\theta} \left[ e^{U_1 - \frac{1}{3}} \phi^{(1)} \right] (\kappa^{(1)}), \quad D^U = \bar{\theta} \left[ e^{U_1 + \frac{1}{3}} \phi^{(1)} \right] (\kappa^{(1)}),
\]
\[
E^U = \bar{\theta} \left[ e^{U_1 - \frac{1}{3}} \phi^{(1)} \right] (\kappa^{(1)}), \quad F^U = \bar{\theta} \left[ e^{U_1 - \frac{1}{3}} \phi^{(1)} \right] (\kappa^{(1)}),
\]
where

$$\epsilon^{U_1} = \frac{\epsilon^{U_1}_a - \epsilon^{U_1}_b - 2\epsilon^{U_1}_c}{6}, \quad \kappa^{(1)} = \frac{6J^{(1)}}{\alpha^2},$$

$$\phi^{(1)} = \theta^{(1)}_c - \theta^{(1)}_b - 2\theta^{(1)}_a,$$

(10)

where $\epsilon^{U_1}_a$, $\epsilon^{U_1}_b$, and $\epsilon^{U_1}_c$ are the shifts of $a$, $b$, and $c$ stacks of D6-branes, $J^{(1)}$ is the Kähler modulus, and $\theta^{(1)}_c, \theta^{(1)}_b$, and $\theta^{(1)}_a$ are the Wilson line phases for the $a$, $b$, and $c$ stacks on the first two torus, respectively.

At the GUT scale, the down-type quark mass matrix $M^D$ is obtained from the above up-type quark mass matrix $M^U$ by changing the upper index $U$ and lower index $a$ to $D$ and $d$, respectively. The lepton mass matrix $M^L$ is obtained from $M^D$ by changing the upper index $D$ to $L$. To generate the suitable SM fermion masses and mixings at the GUT scale, we choose $\epsilon^{U_1} = \epsilon^{U_1} = 0$, $\epsilon^{D_1} = 0.061$, and $\kappa^{(1)} = 39.6i$. And for Higgs VEVs, we choose $v^{U}_u = 0.000266, v^{U}_d = 0.236, v^{U}_\theta = 0.999, v^{D}_u = 0.981, v^{D}_d = 0.00481, v^{D}_\theta = 0.0345, v^{L}_u = 0.00224, v^{L}_d = 0, v^{L}_\theta = 1.58, v^{L}_\mu = 0, v^{L}_\tau = 0.0445$, and $v^{L}_\nu = 0.0001$. Then, with suitable $c^U_\nu, \tau^D_\mu$, and $c^L_\nu$, we obtain the SM fermion mass matrices at the GUT scale

$$M^U \simeq m_t \begin{pmatrix} 0.000266 & 0.00109 & 0.00481 & 0.0310 \\ 0.00109 & 0.00481 & 0.0310 & 0.000747 \\ 0.00047 & 0.0310 & 0.999 & 0.00747 \\ 0.999 & 0.000747 & 0.00747 & 0.000199 \end{pmatrix}$$

$$M^D \simeq m_b \begin{pmatrix} 0.00141 & 0.000025 & 4 \times 10^{-6} \\ 0.000155 & 0.028 & 0 \\ 0 & 2.2 \times 10^{-7} & 1 \\ 0 & 0 & 0 & 0 \end{pmatrix}$$

$$M^L \simeq m_\tau \begin{pmatrix} 0.00142 & 3 \times 10^{-6} & 2.8 \times 10^{-8} \\ 3 \times 10^{-6} & 0.0282 & 1.4 \times 10^{-9} \\ 2.8 \times 10^{-8} & 1.4 \times 10^{-9} & 1 \end{pmatrix}$$

The above mass matrices can produce the correct quark masses and CKM mixings, and the correct $\tau$ lepton mass at the electroweak scale $[17]$. The electron mass is about 6.5 times larger than the expected value, while the muon mass is about 40% smaller. Similar to the GUTs $[18]$, we have roughly the wrong fermion mass relation $m_c/m_\mu \simeq m_d/m_\mu$, and the correct electron and muon masses can be generated via high-dimensional operators $[11]$. Moreover, the suitable neutrino masses and mixings can be generated via the seesaw mechanism by choosing suitable Majorana mass matrix for the right-handed neutrinos.

**Conclusions** – We have briefly presented a three-family intersecting D6-brane model where the gauge symmetry can be broken down to the SM gauge symmetry and the gauge coupling unification can be realized at the string scale. We have calculated the supersymmetry breaking soft terms, and obtained the low energy supersymmetric particle spectrum within the reach of the LHC. Our model may also generate the observed dark matter density. Finally, we can explain the SM quark masses and CKM mixings, and the tau lepton mass. The neutrino masses and mixings may be generated via the seesaw mechanism as well.

**Acknowledgments** – This research was supported in part by the Mitchell-Heep Chair in High Energy Physics (CMC), by the Cambridge-Mitchell Collaboration in Theoretical Cosmology (TL), and by the DOE grant DE-FG03-95-er-40917 (DVN).

[1] J. Polchinski and E. Witten, Nucl. Phys. B 460, 525 (1996).
[2] M. Berkooz, M. R. Douglas and R. G. Leigh, Nucl. Phys. B 480 (1996) 265.
[3] C. Bachas, hep-th/9503030.
[4] R. Blumenhagen, M. Cvetić, P. Langacker and G. Shiu, Ann. Rev. Nucl. Part. Sci. 55, 71 (2005), and the references therein.
[5] R. Blumenhagen, L. Görlich, B. Körs and D. Lüst, JHEP 0010, 006 (2000); C. Angelantonj, I. Antoniadis, E. Dudas and A. Sagnotti, Phys. Lett. B 489, 223 (2000).
[6] M. Cvetić, G. Shiu and A. M. Uranga, Phys. Rev. Lett. 87, 201801 (2001); Nucl. Phys. B 615, 3 (2001).
[7] M. Cvetić, T. Li and T. Liu, Nucl. Phys. B 698, 163 (2004).
[8] R. Blumenhagen, L. Görlich and T. Ott, JHEP 0301, 021 (2003); G. Honecker, Nucl. Phys. B 666, 175 (2003); G. Honecker and T. Ott, Phys. Rev. D 70, 126010 (2004) [Erratum-ibid. D 71, 069902 (2005)].
[9] J. F. G. Cascales and A. M. Uranga, JHEP 0305, 011 (2003); R. Blumenhagen, D. Lüst and T. R. Taylor, Nucl. Phys. B 663, 319 (2003); G. Villadoro and F. Zwirner, JHEP 0506, 047 (2005); P. G. Camara, A. Font and L. E. Ibáñez, JHEP 0509, 013 (2005).
[10] C.-M. Chen, T. Li and D. V. Nanopoulos, Nucl. Phys. B 740, 79 (2006).
[11] C.-M. Chen, T. Li, V. E. Mayes and D. V. Nanopoulos, in preparation.
[12] D. Lüst, P. Mayr, R. Richter and S. Stieberger, Nucl. Phys. B 696, 205 (2004); G. L. Kane, P. Kumar, J. D. Lykken and T. T. Wang, Phys. Rev. D 71, 115017 (2005); A. Font and L. E. Ibáñez, JHEP 0503, 040 (2005).
[13] A. Brignole, L. E. Ibáñez and C. Muñoz, hep-ph/9707209.
[14] A. Djouadi, J. L. Kneur and G. Moortkta, hep-ph/0211331.
[15] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, Comput. Phys. Commun. 174, 577 (2006).
[16] D. Cremades, L. E. Ibáñez and F. Marchesano, JHEP 0307, 038 (2003); M. Cvetić and I. Papadimitriou, Phys. Rev. D 68, 046001 (2003) [Erratum-ibid. D 70, 029903 (2004)]; S. A. Abel and A. W. Owen, Nucl. Phys. B 663, 197 (2003).
[17] H. Fuseoka and Y. Koide, Phys. Rev. D 57, 3986 (1998).
[18] J. R. Ellis and M. K. Gaillard, Phys. Lett. B 88, 315 (1979); D. V. Nanopoulos and M. Srednicki, Phys. Lett. B 124, 37 (1983).