ON LARGE VOLUME MODULI STABILIZATION IN IIB ORIENTIFOLDS

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I present a brief introduction to the construction of explicit type IIB orientifold compactifications and summarize the 'Large Volume Scenario' on compact four-modulus Calabi–Yau manifolds. I discuss the relevance of this kind of setups for the physical MSSM-like model building and gravitational cosmology. These notes are based on my talk at the 'Bogolyubov Kyiv Conference 2009' on 'Modern Problems of Theoretical and Mathematical Physics'.

1. Introduction

In addition to the perturbative closed string sector, type II theories admit non-perturbative objects, the so-called *Dirichlet* branes (D-branes, in brief). These are objects, on which open strings can end. Strings can a) have both ends on the same D-brane; b) they can stretch between two different D-branes, or c) propagate (as closed strings) from one D-brane to another one. Due to their intrinsic tension, stretched strings give rise to massive excitation modes. When the distance between two or more branes is reduced, the strings allow for massless modes.

A careful analysis of the string spectrum shows that the massless modes give rise to the gauge theory on a D-brane. Consider a bunch of N D-branes on the top of each other, this configuration being often referred to as a *stack* (of N D-branes). Then the massless modes induce a non-Abelian group, more precisely a dimensional reduced Super Yang–Mills gauge theory.

D-branes are classified according to their spatial dimension. For example, a D7-brane fills the four-dimensional spacetime entirely – this accounts for the three spatial dimensions x_1, x_2, and x_3 (see Table 1) – and it wraps a four-dimensional subspace of the CY manifold (a holomorphic four-cycle extended, for instance, along the directions x_4, x_5, x_6, and x_7). The transverse directions x_8 and x_9 account for the degrees of movement of the brane.

Table 1. Spacetime extension of D-branes. Spacetime directions filled by the branes are denoted by crosses, whereas the transverse ones are left blank

	x_0	x_1	x_2	x_3	x_4	x_5	x_6	x_7	x_8	x_9
D7-brane	×	×	×	×	×	×	×	×		
D5-brane	×	×	×	×	×	×				
D3-brane	×	×	×	×						
E3-brane					×	×	×	×		

Stable configurations of D-branes underlie certain supersymmetric calibration conditions (BPS conditions). Whereby, in type IIB, D-branes have to wrap complex subspaces of the CY manifold. Configurations of other kinds turn out to be unstable. Therefore, type IIB compactifications naturally come with D3-, D5-, and D7-branes that wrap, respectively, no cycles, two-cycles, and four-cycles of the internal space.

The rest of this section presents a rough sketch of how the matter content of low-energy supergravity arises from intersecting stacks of branes. Two four-cycles generically intersect each other in a one-dimensional complex subspace, i.e. a *Riemann surface*. Let D_1 be a four-cycle, on which a stack of D7-branes is wrapped. Consider the second stack on a four-cycle D_2 such that D_1 and D_2 intersect each other. It can be shown that the chiral matter is induced at the intersection by bifundamental open strings 'stretched' between D_1 and D_2 with fluxes. Furthermore, Yukawa couplings arise in the case where three stacks of D7-branes intersect one another.

In conclusion, the type IIB compactification (with D-branes and orientifold planes) provides an efficient way to encode the entire matter content of MSSM-like theories in terms of the geometry of complex subspaces. Beside the massless modes arising from the perturbative sector in ten dimensions, we have: a) a Yang–Mills gauge theory on an eight-dimensional manifold (e.g., on a stack D_1), b) chiral matter on a six-dimensional manifold (e.g., $D_1 \cap D_2$), and c) Yukawa coupling in our spacetime (e.g., $D_1 \cap D_2 \cap D_3$).

2. The 'Large Volume Scenario'

Calabi–Yau manifolds come in families smoothly related to one another by deformation parameters called *mod-*

uli. These parameters control the shape and the size of CY manifolds. An important property of CY three-folds is that their moduli space is the product of two disjoint parts $\mathcal{M} = \mathcal{M}^{2,1}_{\rm CS} \times \mathcal{M}^{1,1}_{\rm K}$. The *complex-structure moduli* account for deformations of the shape. The *Kähler moduli*, instead, control the sizes of the three-fold and of its subspaces.

These moduli give rise to massless fields in the four-dimensional supergravity via a process similar to the Kaluza–Klein (KK) reduction. The dimension of the moduli space is determined by the Hodge stucture of the CY three-fold: $\dim \mathcal{M}^{2,1}_{\rm CS} = h^{2,1}$ and $\dim \mathcal{M}^{1,1}_{\rm K} = h^{1,1}$. A generic CY three-fold comes with many moduli that, after the KK reduction, lead to unwanted massless scalar fields in the four-dimensional theory. A possible way to overcome this problem is the introduction of scalar potentials that stabilize the fields at energies beyond the characteristic compactification scale. The realization of this strategy is the main issue of type IIB compactifications.

The full superpotential for the type IIB superstring theory compactified on the CY manifold X is

$$W = \int\limits_X G_3 \wedge \Omega + \sum_i A_i\,(S,U)\,e^{-a_i T_i}\ . \qquad (1)$$

The first term is the so-called Gukov–Vafa–Witten (GVW) flux superpotential [2]. It is a purely tree-level potential generated by the background fluxes $G_3 = F_3 + iSH_3$ supported on three-cycles of the CY manifold X. The quantity Ω is the holomorphic three-form of X. We stabilize the complex structure moduli U and the axion-dilaton field $S = e^{-\phi} + iC_0$, by requiring supersymmetry to hold: $D_S W = 0 = D_U W$. Note that no Kähler moduli appear in the GVW potential.

Unlike the superpotential, the Kähler potential gets corrections order-by-order in both α' and the string-loop expansion. On the top of that, there are non-perturbative effects from either the world sheet or the brane instantons. But these corrections are subdominant, and we neglect them: $K = K_{\rm tree} + K_{\rm p}\ \ (+K_{\rm np})$. The Kähler potential takes the form

$$K = -2\ln\left(\mathcal{V}\,e^{-3\phi/2} + \frac{\xi}{2g_{\rm s}^{3/2}}\right) -$$

$$-\ln\left(S+\bar{S}\right) - \ln\left(-i\int\limits_X \Omega\wedge\bar{\Omega}\right). \qquad (2)$$

Here, $g_{\rm s}$ is the string coupling, and $\xi = -\frac{\zeta(3)\chi(X)}{16\pi^3}$ encodes the leading term of perturbative corrections in terms of the Euler characteristic of the CY three-fold, whose volume is given by

$$\mathcal{V} = \frac{1}{3!}\int\limits_X J\wedge J\wedge J = \frac{1}{6}\,\kappa_{ijk}\,t^i t^j t^k\,.$$

The coefficients κ_{ijk} account for the intersection of three four-cycles D_i, D_j, and D_k. They are numbers that depend on the basis of integral two-forms $\{\eta_i\} \in H^{1,1}(X,\mathbb{Z})$, in which we decide to expand the Kähler form:

$$J = \sum_i t_i \eta_i\,, \qquad t_i \in \mathbb{R}\,. \qquad (3)$$

The 'Large Volume Scenario' (LVS), developed in [3], is a new strategy meant to stabilize Kähler moduli via non-perturbative effects due to Euclidean D3-branes (see Table 1).[1] These effects are accounted for by the second term of Eq. (1).

Consider an E3-brane wrapped on a cycle D_i of volume τ_i. In this case, the coefficient A_i depends only on the complex structure moduli and the axion-dilaton field. We assume these moduli to be already stabilized via the flux superpotential, whereby we can think of A_i as a constant. The Kähler moduli, instead, enter explicitly in the exponentials of expression (1). These moduli are related to the volume of the cycle, on which the E3-brane is wrapped in the following way:

$$T_i = \tau_i\,e^{-\phi} + i\rho_i\ . \qquad (4)$$

Here, ρ_i is the axion field originating from the Ramond–Ramond four-form C_4 supported on the four-cycle D_i of volume τ_i:

$$\tau_i = \frac{1}{2}\int\limits_{D_i} J\wedge J = \frac{1}{2}\kappa_{ijk}\,t^j t^k\,, \qquad \text{and} \qquad \rho_i = \int\limits_{D_i} C_4\ .$$

From Eq. (4), we see how the volume of the E3-brane cycle enters superpotential (1). The instantonic contribution grows exponentially as the cycle decreases in size.

The key idea of the LVS lies in finding a CY three-fold such that the volume of the manifold is driven by the volume τ_l of a single large four-cycle. The remaining 'small' four-cycles contribute negatively to the overall volume:

$$\mathcal{V} \sim \tau_l^{3/2} - \sum_{s=1}^{h^{1,1}-1} \tau_s^{3/2}\ . \qquad (5)$$

[1] Euclidean D-branes are instantonic objects extended only in the spatial directions of the internal manifold. For example, an Euclidean D3-brane wraps a four-cycle of the CY manifold, and it is often denoted by E3.

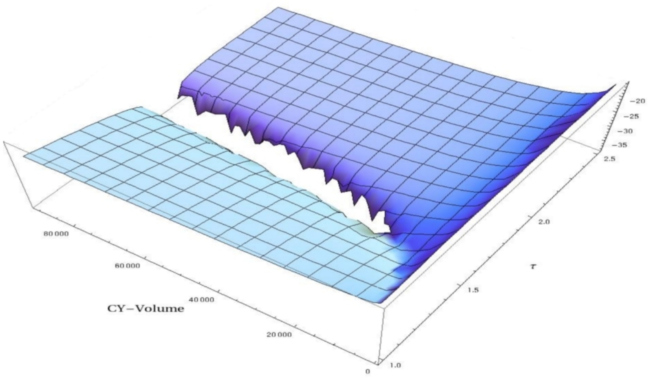

Fig. 1. $\ln(V)$ for $\mathbb{P}^4_{1,1,2,2,6}[12]/\mathbb{Z}_2 : 1\,0\,0\,0\,1$ in the large volume limit, as a function of the CY volume and the size of the E3-cycle. Here, we have set $W_0 = 5$, $A = 1$, and $g_s = 1/10$

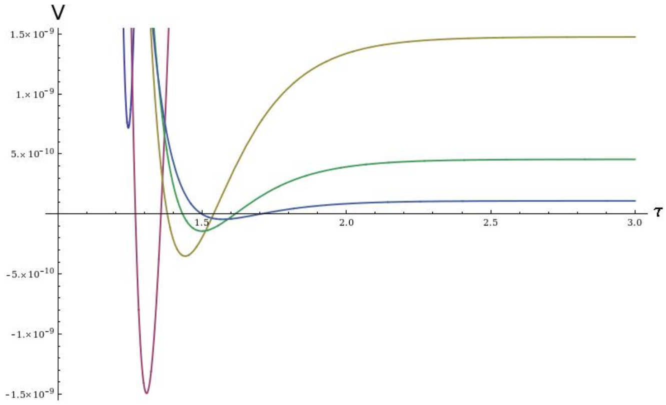

Fig. 2. V for $\mathbb{P}^4_{1,1,2,2,6}[12]/\mathbb{Z}_2 : 1\,0\,0\,0\,1$ in the large-volume limit as a function of the volume of the E3-cycle. Each curve corresponds to a different value of the CY volume. The minimum flattens and shifts to the right as the volume increases

For the obvious reason, these kinds of three-folds have been dubbed the 'Swiss-cheese' CY. This allows us to make the cycles small, while keeping the CY manifolds large. The advantage is double. First, the gauge theory, that takes place on the small cycles, decouples from the Planck scale dynamics. Second, if we choose the E3-brane to wrap one of the 'small' cycles, then we obtain instanton contributions that grow exponentially with decrease in τ_s. These give rise to an *anti de Sitter* (AdS) vacuum at a large CY volume. We devote the rest of this section to the explanation of this result.

The four-dimensional potential for the scalar fields gets contributions from two terms: the F- and D-term potentials. At this stage of the analysis, only the F-term potential is relevant for the LVS:[2]

$$V = e^K \left(\sum_{i=T,S,U} K^{i\bar{j}} D_i W \, D_{\bar{j}} \bar{W} - 3|W|^2 \right) . \quad (6)$$

When $\mathcal{V} \gg 1$, V can be expanded in inverse powers of the three-fold volume. The potential decomposes into three parts: $V = V_{\rm np1} + V_{\rm np2} + V_{\alpha'}$. The two non-perturbative terms depend explicitly on the Kähler moduli:

$$V_{\rm np1} \sim \frac{1}{\mathcal{V}} \left(-\kappa_{ssj}\, t^j\right) e^{-2a_s\,\tau_s}\, e^K + \mathcal{O}\left(\frac{e^{-2a_s\tau_s}}{\mathcal{V}^2}\right) , \quad (7)$$

$$V_{\rm np2} \sim -\frac{a_s\tau_s e^{-a_s\tau_s}}{\mathcal{V}^2}\, e^K + \mathcal{O}\left(\frac{e^{-a_s\tau_s}}{\mathcal{V}^3}\right) . \quad (8)$$

The first term is proportional to the self-intersection of the E3-brane and is positive; $V_{\rm np2}$, instead, is negative.

[2] The D-term potential depends on the choice of the gauge bundles, which we equip the D-branes with.

There is the third contribution that depends on the α'-corrections:

$$V_{\alpha'} \sim \frac{\xi}{\mathcal{V}^3}\, e^K + \mathcal{O}\left(\frac{1}{\mathcal{V}^4}\right) . \quad (9)$$

We want this term to be positive. This is the case for negative Euler characteristics. Therefore, a necessary condition for CY three-folds to realize LVS is $h^{2,1} > h^{1,1}$ (since $\chi = 2\left(h^{1,1} - h^{1,2}\right)$).

In the general case, the right-hand side of (7) dominates the power expansion in $\mathcal{V}^{-1}$. But let us consider the special limit, in which the E3-cycle scales with the logarithm of the CY volume, i.e. $\tau_s \sim \ln \mathcal{V}$. In (7) and (8), the Kähler moduli appear in the exponentials. Therefore, in this limit, $V_{\rm np1}$ and $V_{\rm np2}$ become proportional to $\mathcal{V}^{-3}$ and compete with $V_{\alpha'}$ on the same footing.

In Fig. 1, the logarithm of the scalar potential is plotted as a function of the volume of E3-cycle τ and the CY volume. The 'void channel' corresponds to the region, where the potential becomes negative.

Figure 2 shows the potential for the same model as a function of the 'small' cycle, on which the E3-brane is wrapped. Each curve corresponds to a fixed value of the CY volume. We can choose volumes of the CY, for which the potential is minimum. But increasing $\mathcal{V}$ drives the potential to a negative minimum: this is the AdS vacuum solution. For larger volumes, the minimum flattens and shifts to the right.

The general behavior of V in the LVS can be summarized as follows. Due to the positive contributions of $V_{\rm np1}$ and $V_{\alpha'}$, the potential starts positive at small CY volumes and then, driven by $V_{\rm np2}$, reaches a negative minimum. Afterward, it approaches zero asymptotically for large values of the volume.

Remark: This behavior is of main relevance for cosmology. Such a flat potential accommodates slow-roll inflationary scenarios. The Kähler modulus plays the role of the inflaton slowly rolling down the potential. This is one of the few known setups that allow, at least in principle, a stringy inflation [4].

3. Four-Modulus 'Swiss-Cheese' Calabi–Yau Manifolds

The following section is based on my work in collaboration with Kreuzer, Collinucci, and Mayrhofer [1].

In [5], Blumenhagen *et al.* stress that it is too naive to treat the stabilization of the complex structure moduli via instanton effects and the configuration of chiral matter D7-branes, as two independent tasks. In the general case, indeed, D7-cycles intersect the E3-cycle. This gives rise to charged zero modes that could spoil the non-pertubative stabilization. This means that the moduli stabilization and the vanishing chiral intersection between D7- and E3-branes must be accounted for at the same time.

In our attempt to realize a MSSM-like scenario, we proceed in two steps. Our approach is operational: Given what we think is the minimal set of constraints for a MSSM-like configuration, we search for CY three-folds and explicit brane realizations that satisfy these conditions. The starting point is the following 'wishlist.' We look for

- a CY three-fold with 'Swiss-cheese' structure;
- two stacks of intersecting D7-branes such that they
  - give rise to a $U(N_a) \times U(N_b)$ gauge group;
  - induce bifundamental chiral fermions at the intersecting locus;
  - do not chirally intersect the E3-brane.

At the same time, we have to take care of consistency conditions such as

- tadpole cancellation;
- the Freed–Witten (FW) anomaly which is an open string world-sheet anomaly that arises when D7-branes are wrapped on non-spin manifolds. These are the subspaces that do not admit a spin structure.

The FW anomaly was discovered in [6] and can be compensated by means of half-integral world-volume fluxes.

We start from the complete classification of toric CY three-fold hypersurfaces in [7]. The hypersurfaces are given in terms of combinatoric data encoded in four-dimensional polytopes. Toric CY hypersurfaces are divisors, i.e. co-dimension one manifolds embedded in a toric four-fold ambient space.

The list of CY hypersurfaces with $h^{1,1} = 4$ contains 1197 objects. For simplicity, we focus on 11 models that correspond to simplicial polytopes.

Out of these 11 polytopes, there are only four inequivalent CY three-folds with the 'Swiss-cheese' structure. Each of them allows for two different D7-/E3-brane setups.

Only one of our 'Swiss-cheese' CY hypersurfaces satisfies the conditions for a MSSM-like setup *and* stabilizes (three out of four) Kähler moduli at the same time. The other CY three-folds either miss the conditions or fail to stabilize any Kähler modulus. Let us consider the (partly) successful model in more details. The following analysis is assisted by a recently enhanced version of the PALP package [8].

Consider the CY hypersurface of degree 12 embedded in the weighted projective space $\mathbb{P}^4_{1,1,2,2,6}$ orbifolded by the $\mathbb{Z}_2$-action $(1,0,0,0,1)$ identifying $x_1 \sim -x_1$ and $x_5 \sim -x_5$.

The resulting hypersurface is singular. We repair the singularities, by introducing three toric blow-ups of the ambient space. These induce three extra Kähler moduli. The resolution is shown in the following table.

We can recast Table 2 in terms of equivalence relations: every point $(x_1, \ldots, x_8)$ is equivalent to

$$\big(\lambda\mu x_1, \lambda\nu\, x_2, \lambda^2\mu^2\nu^2\rho\, x_3, \lambda^2\mu^2\nu^2\rho\, x_4,$$

$$\lambda^6\mu^5\nu^5\rho^3 x_5, \mu^2 x_6, \nu^2 x_7, \rho\, x_8\big)\ , \qquad (10)$$

where $\lambda, \mu, \nu, \rho \in \mathbb{C}^*$, and the corresponding exponents are encoded by the rows of Table 2. On the other hand, each column four-vector is associated to a global homogeneous coordinate x_i. The corresponding ambient space divisor $D_i = \{x_i = 0\}$. The divisors $D_1, \ldots, D_5$ come from the original singular ambient space, whereas D_6, D_7, and D_8 are the *exceptional* ones associated to

Table 2. Projective weights for the resolution of $\mathbb{P}^4_{1,1,2,2,6}[12]/\mathbb{Z}_2 : 1\,0\,0\,0\,1$

x_1	x_2	x_3	x_4	x_5	x_6	x_7	x_8	p
1	1	2	2	6	0	0	0	12
1	0	2	2	5	2	0	0	12
0	1	2	2	5	0	2	0	12
0	0	1	1	3	0	0	1	6

the blow-ups. The eight divisors generate four divisor classes modulo four linear equivalences, which is what we need for the four-(Kähler)-modulus model. For convenience, let us introduce the following basis: $\eta_1 = D_1$, $\eta_2 = D_2$, $\eta_3 = D_8$, and $\eta_4 = D_6$. Table 3 shows all the divisors expressed in the η-basis.

The CY divisor is $D_{\text{CY}} = D_1 + \cdots + D_8$. The smooth CY three-fold we have constructed has the Euler characteristic $\chi = -180$ and the Hodge numbers $h^{1,1} = 4$ and $h^{1,2} = 94$. PALP determines the Stanley–Reisner ideal:

$$I_{\text{SR}} = \{x_1x_2, x_1x_5, x_1x_6, x_2x_5, x_2x_7, x_5x_8,$$

$$x_3x_4x_6x_7, x_3x_4x_6x_8, x_3x_4x_7x_8\}\,.$$

This set fixes the coordinates that cannot vanish at the same time. For example, $x_1 = 0$ implies $x_2 \neq 0$.

A toric divisor D_i intersects the CY hypersurface on a two-dimensional subspace: $\dim_{\mathbb{C}}(D_i \cap D_{\text{CY}}) = 2$. From the point of view of the CY hypersurface, the *ambient space* divisor D_i induces a four-cycle in $H_4(D_{\text{CY}})$ that we denoted again by D_i for simplicity.[3]

On the CY, three CY divisors intersect one another at points. In our case, the resulting triple intersection numbers in the η-basis are encoded as coefficients of a polynomial as follows:

$$I_3 = -78\eta_4^3 - 6\eta_3\eta_4^2 - 6\eta_3^2\eta_4 + 2\eta_3^3 +$$

$$+36\eta_2\eta_4^2 + 6\eta_2\eta_3\eta_4 + \eta_2\eta_3^2 - 18\eta_2^2\eta_4 -$$

$$-3\eta_2^2\eta_3 + 9\eta_2^3 + \eta_1\eta_3^2 - 3\eta_1^2\eta_3 + 9\eta_1^3\,. \qquad (11)$$

T a b l e 3. Toric divisors of $\mathbb{P}^4_{1,1,2,2,6}/\mathbb{Z}_2 : 10001$

Divisor	η-basis
D_1	η_1
D_2	η_2
D_3	$2\eta_2 + \eta_3 + \eta_4$
D_4	$2\eta_2 + \eta_3 + \eta_4$
D_5	$\eta_1 + 5\eta_2 + 3\eta_3 + 2\eta_4$
D_6	η_4
D_7	$-2\eta_1 + 2\eta_2 + \eta_4$
D_8	η_3
D_{CY}	$12\eta_2 + 6\eta_3 + 6\eta_4$

[3] This abuse of notation is justified since the four divisor classes of the CY three-fold descend from the four divisor classes of the toric ambient space X. In general, it might happen that a toric divisor intersects the D_{CY} at two or more loci. On the CY hypersuface, such a divisor reduces into a sum of disjoint non-intersecting four-cycles. More precisely, the number of disconnected parts of the cycle D is counted by $h^{0,0}(D)$.

With help of an appropriate recast, the volume of the CY three-fold exhibits the wanted Swiss-cheese structure of formula (5):

$$\mathcal{V} = \frac{\sqrt{2}}{9}\left(\frac{3}{2\sqrt{6}}\tau_a^{3/2} - \frac{1}{2\sqrt{2}}\tau_b^{3/2} - \tau_c^{3/2} - \tau_d^{3/2}\right). \qquad (12)$$

Here, $\tau_a, \ldots, \tau_d$ are the volumes of the cycles

$$D_a = \eta_1 + 5\eta_2 + 3\eta_3 + 2\eta_4\,,$$

$$D_b = \eta_1 + \eta_2 + 3\eta_3\,,$$

$$D_c = \eta_1\,,$$

$$D_d = \eta_2\,.$$

The convenience of the diagonal basis becomes obvious by rewriting the triple intersections from expression (11):

$$I_3 = 24D_a^3 + 72D_b^3 + 9D_c^3 + 9D_d^3\,. \qquad (13)$$

The divisors D_b, D_c, and D_d contribute negatively to the overall volume. Indeed, we will check if these cycles are shrinkable in such a way that the LVS accommodates.

The *Kähler cone* is a subspace of the space of parameters t_i, for which the integral of the Kähler form (3) over all *effective curves* is positive: $\int_{\text{curve}} J > 0$. We obtain the Kähler cone starting from the *Mori cone*. The Mori cone is the cone of (numerically) effective curves. It is determined, by using the Oda–Parks algorithm, as described in [9]. The result is a set of constraints for the volumes of the cycles:

$$\sqrt{\tau_a} - 3\sqrt{\tau_b} > 0\,, \quad 2\sqrt{\tau_b} - \sqrt{\tau_d} > 0$$

$$2\sqrt{\tau_b} - \sqrt{\tau_c} > 0\,, \quad \sqrt{\tau_c} > 0\,, \quad \sqrt{\tau_d} > 0\,.$$

We can arbitrarily shrink D_b, D_c, and D_d while keeping D_a large. In terms of the η-basis, η_1, η_2, and η_3 are the small directions, whereas η_4 is the large one.

Let us take a closer look at the topology of the 'small' cycles. The Euler characteristic $\chi(D_i)$, the holomorphic Euler characteristic

$$\chi_h(D_i) = \sum_{p=0}^{2} (-1)^p\, h^{0,p}(D_i)\,,$$

and the Hodge numbers are listed in Table 4.

T a b l e 4. Topological quantities of the divisors $D_1, \ldots, D_8$

Divisor	D_1	D_2	D_3	D_4	D_5	D_6	D_7	D_8
χ	3	3	46	46	108	–30	–30	10
χ_h	1	1	4	4	11	–9	–9	1
$h^{0,0}$	1	1	1	1	1	1	1	1
$h^{0,1}$	0	0	0	0	0	10	10	0
$h^{0,2}$	0	0	3	3	10	0	0	0
$h^{1,1}$	1	1	38	38	86	8	8	8

The cycles η_1 and η_2 turn out to be both $\mathbb{CP}^2$. These are Del Pezzo surfaces of the dP_0 type. The anticanonical line bundle $-K_D$ of a Del Pezzo surface D is ample by definition. This implies that $-K_D$ must have positive intersections with any effective curve C on the surface: $\int_D (-K_D) \wedge C > 0$. The right-hand side of expression (11) yields $\eta_2\eta_3^2 = +1$, whereby, for the curve $C : \{x_8 = 0\} \cap \{x_2 = 0\}$, the following holds: $\int_{\eta_3} (-\eta_3) \wedge C = -1$. This means that η_3 is not Del Pezzo and, hence, is not shrinkable.

4. Conclusions and Outlook

In [1], we have analyzed two different configurations of MSSM D7-branes with chiral matter for the resolved $\mathbb{P}^4_{1,1,2,2,6}/\mathbb{Z}2[12] : 1\,0\,0\,0\,1$. The most relevent results are as follows.

Not every rigid surface D_i that contributes negatively to the CY volume can be stabilized at small values of τ_i. Indeed, we find that only Del Pezzo surfaces can be stabilized at small values. This is a general result. In our model, three out of the four Kähler moduli can be fixed via instantonic effects in the LVS. Cicoli suggests in [10] that, if g_s corrections are taken into account in (2), a full moduli stabilization might be achieved even if not all the cycles satisfy the topological constraint of being Del Pezzo.

The physical and consistency conditions we have taken into account turn out to be very restrictive and significantly reduce the list of Swiss-cheese CY hypersurfaces that can accommodate MSSM-like configurations. This is true, in particular, if the FW anomaly cancellation is required. This forces D7- and E3-branes to be magnetized in such a way that the condition of zero chiral intersections between the instantonic brane and the matter branes, on the one hand, and between the hidden D7-branes – that we eventually need to introduce for tadpole saturation – and the physical sector, on the other hand, is hard to satisfy.

While we did not succeed so far in finding a model with all the Kähler moduli stabilized, it appears likely that the systematic analysis of a larger list of toric examples will provide phenomenologically attractive models.

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ЩОДО СТАБІЛІЗАЦІЇ МОДУЛІВ ДЛЯ ВЕЛИКИХ ОБ'ЄМІВ В IIБ ОРІЄНТОВАНИХ МНОЖИНАХ

Н.-О. Уоллизер

Р е з ю м е

Подано короткий вступ до побудови явних компактифікацій IIБ орієнтованих множин та розглянуто "сценарій великих об'ємів" для компактних чотиримодульних Калабі–Яу многовидів. Обговорено доречність таких схем для побудови фізичної моделі МССМ-типу та для гравітаційної космології. Основою роботи є доповідь на "Боголюбовській київській конференції 2009" з "Сучасних проблем теоретичної та математичної фізики".