Non-linear Supergravity and the KKLT Scenario - Report

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ABSTRACT: The KKLT scenario is the most investigated construction of de Sitter vacua from string theory that is presently known and is based on type IIB string theory. In this work we outline how this set up can be extended to type IIA theory and how there the anti-D6 brane takes the role of the uplifting contribution. We further show that this contribution can be best described by the use of constrained multiplets that realize non-linear supergravity. This starting point sparked an investigation in both type IIA and IIB set ups and led to the so called mass production process of de Sitter vacua, a predictive formalism that allows to build de Sitter vacua by starting from Minkowski solutions without the need of particular fine tuning.

Contents

1 Indroduction

This work is a report based on three scientific papers that were written during a stay at SLAC/SITP funded by a generous scholarship of the Austrian Marshall Plan Foundation. The papers are, in order of appearance:

1. **Uplifting Anti-D6-brane** [1]
2. **Mass Production of IIA and IIB dS Vacua** [2]
3. **de Sitter Minima from M theory and String theory** [3]

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Building de Sitter vacua is one of the most important applications of string theory and, by extension, its low energy limit supergravity. Due to the accelerated expansion of the universe, caused by what we call dark energy, it is known that the universe cannot be a simple flat Minkowski space. The best way to describe the large scale spacetime structure of the universe is by a de Sitter space. In terms of string theory such a space can be built by having one (or possibly multiple) scalars in the theory that are stabilized but the minimum of their potential is not at zero. The difference of zero to the the minimum of the potential then has to match the non-zero cosmological constant that we observe in our universe today.
Constructing de Sitter spaces in theories that come from string theory is not an easy task. In fact most constructions do not work out at all. The problem that usually appears are tachyons. These negative mass particles cause instabilities and prohibit a stable solution. Even in theories that do have a completely stable vacuum state there are typically remaining problems that need to be solved. Often the relation of the ten dimensional theory to the compactified effective four dimensional one is not quite clear or additional requirements from string theory are hard to satisfy.
Keeping all of this in mind the best understood and investigated working construction is the so called KKLT (Kachru-Kallosh-Linde-Trivedi) scenario [4]. This set up works by including non-perturbative corrections and extended objects in string theory in order to first stabilize all scalar particles in the theory at an anti-de Sitter minimum and then lifting this minimum to de Sitter. The KKLT scenario is built in what is called type IIB string theory. Interestingly it has not been possible thus far to achieve a similar

set up in the closely related type IIA string theory. On of the topics of this report will show how this can be achieved. Here the formalism of *non-linear supergravity* is important in order to describe general Dp-branes in terms of constrained superfields in the language of the effective low energy theory in four dimensions.

An other interesting question regards the predictive power of the KKLT scenario. A priori there is no reason that an uplift has to lead back to a stable minimum and even getting a stable minimum in anti-de Sitter is non trivial. Only in Minkowski space there is a way to show that tachyons are absent. Using this fact we were able to introduce what we call the *mass production procedure*. Starting from a stable Minkowski solution a simple 3 step process guarantees stable de Sitter solutions that do not rely on fine tuning the parameters.

In all of the above models the inclusion of non-perturbative terms is essential. Without them we were unable to find concrete solutions that give rise to stable de Sitter vacua. An investigation of models that include all tree level terms that are present in M-theory, the generalization of string theory that should include all string theories, allowed us to drastically reduce the amount of non-pertubative terms required and in one case even build a stable minimum without any non-perturbative corrections.

This report is structured as follows: In section 2 some general facts about the KKLT scenario and uplifts in general are presented. Of particular interest is non-linear supergravity as it allows the description of the uplifting branes in terms of constrained multiplets in the effective four dimensional supergravity. In section 3 the uplifting procedure in type IIA theory is introduced. We describe how an anti-D6 brane, described by constrained multiplest, take the role of the anti-D3 branes of the usual KKLT model and not only outline in detail how this procedure works but also give an explicit example. In section 4 we introduce the mass production procedure and show how, if we start from a Minkowski minimum, a stable de Sitter vacuum is guaranteed. Examples of this procedure, based on the Kallosh-Linde superpotential, are presented in section 5. In the penultimate section 6 we show how it is not necessary to include the Kallosh-Linde double exponential by including tree-level contributions and that in a certain set up it is possible to build a models without non-pertubative terms at all. Finally, section 7 contains the summary and conclusion.

2 Overview and Description of anti-Dp-brane Uplifts

In this first section we give an overview about the KKLT model of de Sitter building on which the rest of this work is based. Furthermore we show how the uplifting branes used can be described using non-linear supergravity, in particular constrained superfields. At the end of the present section we also make the connection to the observable universe in order to properly motivate the need of solid de Sitter constructions from string theory.

2.1 The KKLT Scenario

The KKLT scenario [4] is a well know mechanism to obtain dS vacua in a low energy supergravity limit of type IIB string theory compactified on some Calabi-Yau manifold. A lot of research has been conducted about this type of set up, see for some examples [5–14], and some criticism was raised, starting with [15], much of which is summarized in the review [16]. However, the fact remains that it is one of the best understood models to construct dS vacua that is known today.

On a basic level the KKLT construction works in the following way:
First one starts with a warped compactification from 10 to 4 dimensions, including non-trivial NS and RR three-from fluxes. This allows to stabilize the axio-dilaton and the complex structure moduli at a high scale.
In the second step one introduces non-perturbative corrections for the Kähler modulus. This leads to a stable AdS vacuum, see figure 1. The superpotential for this theory reads

$$W = W_0 + Ae^{\mathrm{i}aT}, \tag{2.1}$$

where W_0 comes from the flux superpotential, fixing the axio-dilaton and complex structure moduli, and the exponential term is the non-perturbative correction, stabilizing the Kähler modulus.
Lastly, in order to obtain a dS vacuum, an uplift by anti-D3-branes, located at the tip of a warped throat, is introduced, which will lift the minimum with negative vacuum energy to de Sitter by introducing a positive energy contribution into the scalar potential [17].

2.2 Non-linear Supergravity

In [18] it was shown how the contribution from the uplifting anti-D3-brane can be written consistently in 4 dimensional supergravity. This requires the use of so called

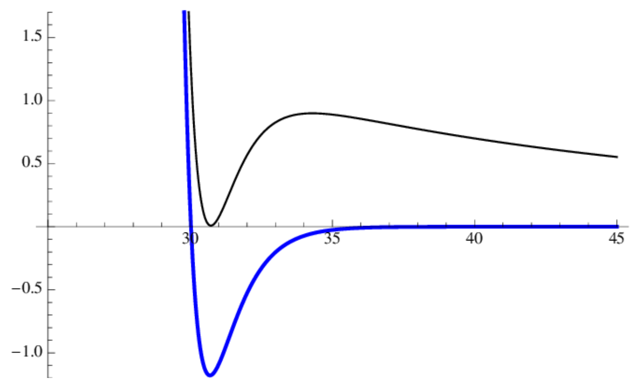

Figure 1. A basic example of a scalar potential in the KKLT scenario. The lower, blue line shows the AdS vacuum and the upper, black line the resulting de Sitter vacuum after the uplift.

nilpotent superfields and *non-linear supersymmetry*. Due to its importance we give a short description of these fields here.
A superfield X has the expansion in terms of *superspace* given by

$$X = \phi + \sqrt{2}\theta\psi + \theta^2 F\,, \tag{2.2}$$

where θ are the (anti-commuting) superspace coordinates, ϕ is a scalar field, ψ is a fermion and F is an auxiliary field, introduced such that the super-algebra closes off-shell. A good introduction into the topic of supergravity and superspace can be found in the book [19]. This expansion is exact due to the anti-commuting nature of the superspace directions θ^1 and θ^2 and the component fields ϕ, ψ and F can depend on the coordinates of *normal* space, say x. A *nilpotent* superfield of this type satisfies

$$X^2 = 0, \tag{2.3}$$

which leads to the restriction that the other component fields are given in term of the fermion:

$$X = \frac{\psi^2}{2F} + \sqrt{2}\theta\psi + \theta^2 F\,. \tag{2.4}$$

Note that the auxiliary field is not an independent degree of freedom as it is not physical at all. It also interesting that, while this breaks supersymmetry, there is still some relation between the fields remaining. This is called *non-linear* supersymmetry. We will illustrate this in the example of the Volkov-Akulov action [20] which is given as:

$$S^{VA} = -\int E^0 \wedge E^1 \wedge E^2 \wedge E^3 \qquad \text{where} \qquad E^\mu = dx^\mu + \bar{\psi}\gamma^\mu d\psi \,. \tag{2.5}$$

This action is invariant under a non-linear transformation of the fermion ψ given as:

$$\delta_\epsilon \lambda = \epsilon + \left(\bar{\psi}\gamma^\mu \epsilon\right) \partial_\mu \psi \,. \tag{2.6}$$

In [21, 22] it was noted that the Volkov-Akulov model is equivalent to

$$S^{VA} = \int d^4x \int d^2\theta \int d^2\bar{\theta} X\bar{X} + M^2 \left(\int d^2\theta X + h.c. \right) , \tag{2.7}$$

And the fermion in X exhibits the same non-linear behaviour that was outlined above. For this work this is relevant because in [18, 23] it was shown that the action of an anti-D3-brane can be written as

$$S^{VA} = -2 \int \tilde{E}^0 \wedge \tilde{E}^1 \wedge \tilde{E}^2 \wedge \tilde{E}^3 \qquad \text{where} \qquad \tilde{E}^\mu = dx^\mu + \sum_{\alpha=1}^{3} \bar{\psi}^\alpha \gamma^\mu d\psi^\alpha \,. \tag{2.8}$$

Furthermore in [24] the complete action of an anti-D3-brane at the bottom of a warped throat in a KKLT background was derived and non-linear supergravity was a key component of this description.
The in cooperation of the anti-D3-brane into the superpotential W and Kähler potential K, using non-linear supergravity, will be discussed in the the next section.

2.3 General Uplifts

Anti-D3-branes are not exclusive in their ability to uplift a scalar potential from AdS to dS. In [18] the 4d supergravity description for all viable anti-D-branes was given and it was proposed that it should be possible to use any of these branes in order to achieve an uplift. For type IIB theory it is possible to consider anti-D9-, anti-D7-, anti-D5 and anti-D3-branes. This is because these branes have to wrap around supersymmetric $(p-3)$-cycles on the compactification manifold. As it turns out in IIB theory there are 0-, 2-, 4- and 6-cycles and thus the above mentioned branes are possible. In type IIA theory on $SU(3)$ structure manifolds, the most commonly studied compactification manifolds for type IIA, however, there are only 3-cycles. Thus only anti-D6-branes are able to give a positive contribution to the scalar potential and hence lift the minimum from AdS to dS.

For the KKLT scenario these anti-branes can be described by inclusion of a nilpotent field X into the Kähler- and superpotential:

$$\begin{aligned} K &= -3\log\left(-\mathrm{i}(T-\bar{T})\right) + X\bar{X} \\ W &= W_0 + Ae^{\mathrm{i}aT} + \mu^2 X \end{aligned} \tag{2.9}$$

where μ^2 gives the hight of the uplift and depends on the type of anti-brane that is used. An important thing to remember is that the nilpotent field X has to be set to 0 at the end of all calculations.

In general, the scalar potential in a supergravity that is given by the Kähler potential K and superpotential W is

$$V = e^K\left(D_I W K^{IJ}\overline{D_J W} + 3|W|^2\right) , \tag{2.10}$$

where I and J run over all moduli in the theory and $D_I W$ is the Kähler-covariant derivative $D_I W = \partial_I W + W \partial_I K$.

The effective contribution of anti-Dp-branes to the scalar potential of the KKLT scenario is:

$$V^{up} = +\frac{\mu^4}{(T+\bar{T})^3} \, . \tag{2.11}$$

The fact that the uplift will always have a form similar to this one simplifies calculations greatly. Once the relation to the description using nilpotent multiplets is shown one does not need to consider them during each step of the investigation but rather can add the effective uplift term after the (supersymmetric) AdS minimum was found.

How this translates in detail to other models will be discussed in the corresponding section 4.4.

2.4 The expanding Universe and dark Energy

Since the observations in [25, 26] it is an established fact that our universe is expanding at an accelerated rate. This expansion is described by the inclusion of the cosmological constant $\Lambda \cong 10^{-121}$ in the Einstein equations. The best way to accommodate this result is by considering the spacetime of our universe to be of the de Sitter type. This geometry intrinsically has the property that its volume grows with time. In the language of particle physics the source of this expansion is attributed to dark energy. It can be described by a scalar particle that sits at the minimum of its potential, however, the value of the potential at the minimum is not zero but rather some value

$\langle V \rangle$, corresponding to the behaviour of the expansion. In order to match this language with the gravity side of things we have to match the value of the minimum with the cosmological constant: $\langle V \rangle = \Lambda$. This fine tuning problem usually is explained via the anthropic principle. For us this means that, in building our de Sitter models, we should build a model such that this matching is achieved. In practice we choose not to do so for reasons of presentation. The true value of the cosmological constant is too small to be depicted by anything but zero. It is, however, clear from our constructions that the value can easily be achieved by changing some parameters which does not alter the general behaviour of the set up.

3 Anti-D6-Uplift

While the description an uplifting anti-D6-brane for type IIA models was proposed in [18], no concrete examples were know. Indeed, the mere existence of an anti-Dp-brane that can lift an AdS minimum to a (meta) stable dS vacuum does not mean it is possible to construct such a situation. The brane only gives a positive contribution to the scalar potential (see (2.10), (2.11)) but does not guarantee that all moduli are stabilized. In order to establish a KKLT-like construction in type IIA one needs not only a working description of the ant-D6-brane but also concrete models as examples. We gave first, working, examples in [1]. In this work we not only discuss the aforementioned method of construction dS vacua but also discuss an issue related to M-theory U-duality.

3.1 The STU Model

The so called STU model is a simple set up where there are only 3 independent moduli. We call the axio-dilaton S, the complex structure modulus T and the (volume) Kähler modulus U. The ten dimensional supergravity is compactified on a calibrated manifold, for example Calabi-Yau manifolds or more general SU(3) structure maniflods. In this way standard four dimensional supergravity with linearly realized $\mathcal{N} = 1$ supersymmetry will be our effective low energy theory. In addition we add a pseudo calibrate anti-D6-brane which will give a positive contribution to the scalar potential and facilitate a KKLT like uplift.

As a specific example we take a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orbifold compactification of type IIA string theory [27]. The ten dimensional metric is given by

$$ds_{10}^2 = \tau^{-2} ds_4^2 + \rho \left(\sigma^{-3} G_{ab} dy^a dy^b + \sigma^3 G_{ij} dy^i dy^j \right) . \tag{3.1}$$

In this metric τ, σ and ρ are moduli and they are identified in the following way [28]:

$$\begin{aligned}\rho &= Im(U) = (vol_6)^{1/3} \\ \tau &= Im(S)^{1/4} Im(T)^{3/4} = \mathrm{e}^{-\phi}\sqrt{vol_6} \\ \sigma &= Im(S)^{-1/6} Im(T)^{1/6}\,.\end{aligned} \tag{3.2}$$

On the other hand, G_{ab} and G_{ij} correspond to two independent three cycles around which the anti-D6-branes can be wrapped. We distinguish $N^{\parallel}_{\overline{D6}}$, corresponding to an amount of branes that are wrapped around one cycle only, and $N^{\perp}_{\overline{D6}}$, which are branes wrapped around directions along both cycles, in different combinations.

After the compactification the scalar potential of our model will be given as a sum of both the standard $\mathcal{N}=1$ supergravity potential (2.10)

$$V^{\mathcal{N}=1} = \mathrm{e}^{K}\left(K^{i\bar{j}} D_i W D_{\bar{j}}\overline{W} - 3W\overline{W}\right)\,, \tag{3.3}$$

with K the Kähler potential and W the superpotential, and the positive energy contribution from the anti-D6-brane, given effectively as:

$$V^{\overline{D6}} = \frac{\mu_1^4}{Im(T)^3} + \frac{\mu_2^4}{Im(T)^2 Im(S)}\,. \tag{3.4}$$

Here $\mu_1^4 = 2\mathrm{e}^{\mathcal{A}_1} N^{\parallel}_{\overline{D6}}$ and $\mu_2^4 = 2\mathrm{e}^{\mathcal{A}_2} N^{\perp}_{\overline{D6}}$. $\mathrm{e}^{\mathcal{A}_{1(2)}}$ corresponds to the warp factor at the position of the brane.

As Kähler and superpotential we use:

$$\begin{aligned}K &= -\log\left(-\mathrm{i}(S-\bar{S})\right) - 3\log\left(-\mathrm{i}(T-\bar{T})\right) - 3\log\left(-\mathrm{i}(U-\bar{U})\right) \\ W &= f_6 + f_4 U + f_2 U^2 + f_0 U^3 + (h_T + r_T U)T + (h_S + r_S U)S + W_{np}\,.\end{aligned} \tag{3.5}$$

Here we will set all terms in W to zero except f_6 and W_{np} which comes from non-perturbative contributions to the superpotential that take the general form

$$W^{np} = \sum_i A_i \mathrm{e}^{\mathrm{i}a_i\Phi_i} \qquad \text{where} \qquad \Phi_i = \{S, T, U\}\,. \tag{3.6}$$

For the sake of simplicity we will assume that all parameters in the potentials are constant. In principle they can depend on the moduli but under certain conditions it is reasonable to assume them to be constant.

These non-perturbative contributions can have different origins, for the S and T direction the may arise, for example, from gaugino condensation [27] or, alternatively, from Euclidean D2-branes wrapping 3-cycles. For the U direction the origin of these terms is not quite as clear, however, a good motivation for the existence of these terms is the concept of M theory U-duality. String theory exhibits both S and T-duality but not explicit U-duality. M-theory is the parent theory of all string theories and in fact does include U-duality. It is thus reasonable to expect that there are effects of U-duality present in string theory, allowing for the missing terms since they should exist in M-theory [29–31]. Concretely worldsheet instantons of $\mathcal{N}=1$ orientifold compactifications in type IIA give generically rise to the required exponential terms, see [32, 33].

3.2 Four-dimensional Action

While the effective contribution of the anti-D6-branes to the scalar potential has already been given above, we have not shown how to incorporate them into the Kähler - and superpotential. In [18] it was shown that this is indeed possible by the use of a nilpotent chiral goldstino superfield X where $X^2=0$. The Kähler potential for the STU model, including the nilpotent field and thus the contribution from the anti-D6-branes, is given as

$$\begin{aligned} K =& -\log\left(-\mathrm{i}(S-\bar{S})\right) - 3\log\left(-\mathrm{i}(T-\bar{T})\right) \\ & -\log\left(\left[-\mathrm{i}(U-\bar{U})\right]^3 - \frac{X\bar{X}}{\mathrm{e}^{\mathcal{A}_1}N_{\overline{D6}_1}\left(-\mathrm{i}(S-\bar{S})\right) + \mathrm{e}^{\mathcal{A}_2}N_{\overline{D6}_2}\left(-\mathrm{i}(T-\bar{T})\right)}\right) . \end{aligned} \tag{3.7}$$

The contribution to the superpotential has the form

$$W = \mu^2 X \tag{3.8}$$

where we identify $\mu^4_{1(2)} = \frac{1}{8}\mu^4 \mathrm{e}^{\mathcal{A}_{1(2)}} N_{\overline{D6}_{1(2)}}$. Using this formula for the Kähler potential K and adding $\mu^2 X$ to the superpotential we arrive at the same result once we derive (2.10) and use that $X^2=0$. For the remainder of this work we will use the effective description, adding the contribution of the anti-D6-brane to the scalar potential without the nilpotent field X since it is equivalent.

3.3 Requirements from String Theory

There are a number of requirements that low energy effective theories originating from string theory should satisfy for consistency. Here we give an overview about the most important ones and what kind of conditions they imply on our model.

First we need to satisfy Gauss' law in the compact space which is equivalent to satisfying the Bianchi identities for the RR fields. For our explicit example this reduces to the tadpole condition including the D6-brane charges:

$$\int dF_2 - F_0 H = -2N_{O6} + N_{D6} - N_{\overline{D6}}\,, \tag{3.9}$$

for each three-cycle independently. Since our model does not include any fluxes the contributions from $O6$-planes, $D6$-branes and anti-$D6$-branes need to cancel identically. We are free to add $D6$-branes as needed but must take care to find a geometry that is stable. In principle $D6$ and $\overline{D6}$ can annihilate. Hence it is important to find a set up that allows for a stable configuration. This might be non-trivial [34].

An additional condition on the fluxes is quantization. In principle all fluxes need to be properly quantized for consistency. Luckily there is a remaining, overall scaling symmetry of the superpotential that can be used in order to achieve any desired value of the f_6 flux. This changes the parameters A_i but not the existence nor the location of the vacua.

Another important requirement concerns higher order non-perturbative and α' corrections. All these contributions need to be small in order to not change the system. Generally the sum of corrections we consider is of form

$$\sum_{n=1}^{\infty} A_n \mathrm{e}^{\mathrm{i}na_i\Phi_i}\,, \tag{3.10}$$

for each modulus. In order to only have the first contribution of this infinite sum matter we require that $a_i Im(\Phi_i) > 1$.

The α' corrections are suppressed if the volume of the internal manifold is large, or in other words, $vol(6) = (Im(U))^3$ has to be large und thus $Im(U) \gg 1$.

3.4 Finding Vacua

The first step in order to find a stable dS vacuum is to stabilize all moduli in a supersymmetric anti de Sitter state. In principle it suffices to satisfy the Breitenlohner-

Freedman bound for stability in AdS but, since we want to find stable dS, we require that all moduli have masses above zero. To find such a vacuum we have to follow the F-term equations

$$D_iW = 0\,, \tag{3.11}$$

for the Kähler - and superpotential

$$\begin{aligned} K &= -\log\big(-\mathrm{i}(S-\bar{S})\big) - 3\log\big(-\mathrm{i}(T-\bar{T})\big) - 3\log\big(-\mathrm{i}(U-\bar{U})\big) \\ W &= f_6 + A_S \mathrm{e}^{\mathrm{i}a_S S} + A_T \mathrm{e}^{\mathrm{i}a_T T} + A_U \mathrm{e}^{\mathrm{i}a_U U}\,. \end{aligned} \tag{3.12}$$

Note that while $D_iW = 0$ implies $\partial_i V = 0$ the reverse is not true.
Since the $Re(\Phi_i)$ do not appear in the Kähler potential they are axions and we can set all of them to zero. The solution we find will be consistent for as long as all masses are positive. We also choose to set the location of the AdS vacuum to be at:

$$Im(S) = S_0\,, \qquad Im(T) = T_0\,, \qquad Im(U) = U_0\,, \tag{3.13}$$

and solve for the A_i. This allows us to freely choose the values of the moduli as well as f_6. After finding a critical point in this way we have to check whether or not it is a stable minimum by evaluating either the canonical mass matrix

$$m_i^j = \frac{1}{2} K^{jk} \nabla_k \partial_i V \tag{3.14}$$

or the second derivative of the scalar potential. Either one suffices for the purpose of checking stability.
After checking stability we can introduce the uplifting contribution from the anti-D6-brane, in the from of

$$V^{\overline{D6}} = \frac{\mu_1^4}{Im(T)^3} + \frac{\mu_2^4}{Im(T)^2 Im(S)}\,. \tag{3.15}$$

We can choose the values of μ_1 and μ_2 such that the minimum of the potential is above zero or even match the cosmological constant.
The uplift will slightly change the position of the minimum, depending on the magnitude of the μ. We thus need to find it's new position, typically very close to the original one. It is also necessary to check the mass matrix again after uplifting in order to make sure that the vacuum remains stable.

3.5 An explicit Example

Investigating this model we found that it is not only possible to choose the remaining free parameters such that one is able to find a suitable solution but that in fact no particular amount of fine tuning is necessary [1]. One particular solution will be given in this section.

A convenient choice for the position of the AdS vacuum is $S_0 = T_0 = 1$ and $U_0 = 10$, where the value of $Im(U)$ was chosen in order to have a large internal volume. f_6 is set equal to 2 and the $a_S = 3.1$, $a_T = 3.3$ as well as $a_U = 0.32$ in order to have all $\mathrm{e}^{-a_i \Phi_i} \lesssim 0.1$. For the uplift we chose $\mu_1^4 = \mu_2^4 = 1.34 \cdot 10^{-5}$. In table 1 the masses for AdS and dS are given and plots for this model are given in figures 2 and 3.

	m_1^2	m_2^2	m_3^2	m_4^2	m_5^2	m_6^2
AdS	$1.19 \cdot 10^{-3}$	$1.01 \cdot 10^{-3}$	$2.43 \cdot 10^{-4}$	$2.20 \cdot 10^{-4}$	$1.64 \cdot 10^{-4}$	$1.45 \cdot 10^{-4}$
dS	$8.00 \cdot 10^{-4}$	$7.40 \cdot 10^{-4}$	$1.76 \cdot 10^{-4}$	$1.63 \cdot 10^{-4}$	$1.61 \cdot 10^{-4}$	$1.50 \cdot 10^{-4}$

Table 1. The canonically normalized masses squared for both before and after the uplift are positive.

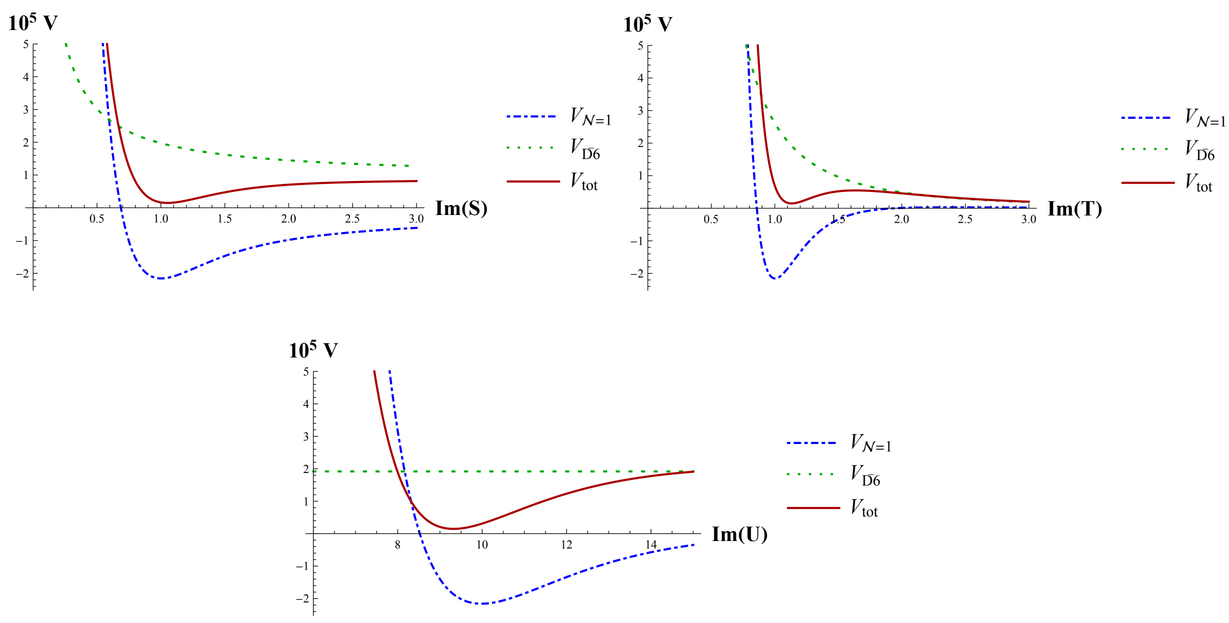

Figure 2. 2D plots of the total scalar potential V_{tot}, the anti-de Sitter potential $V_{\mathcal{N}=1}$ and the $\overline{D6}$ potential $V_{\overline{D6}}$. Starting from the top we have the Im(S) direction on the left, followed by Im(T) on the right and below that Im(U). In all plots we see clearly the anti-de Sitter and de Sitter vacua as well as the uplift term.

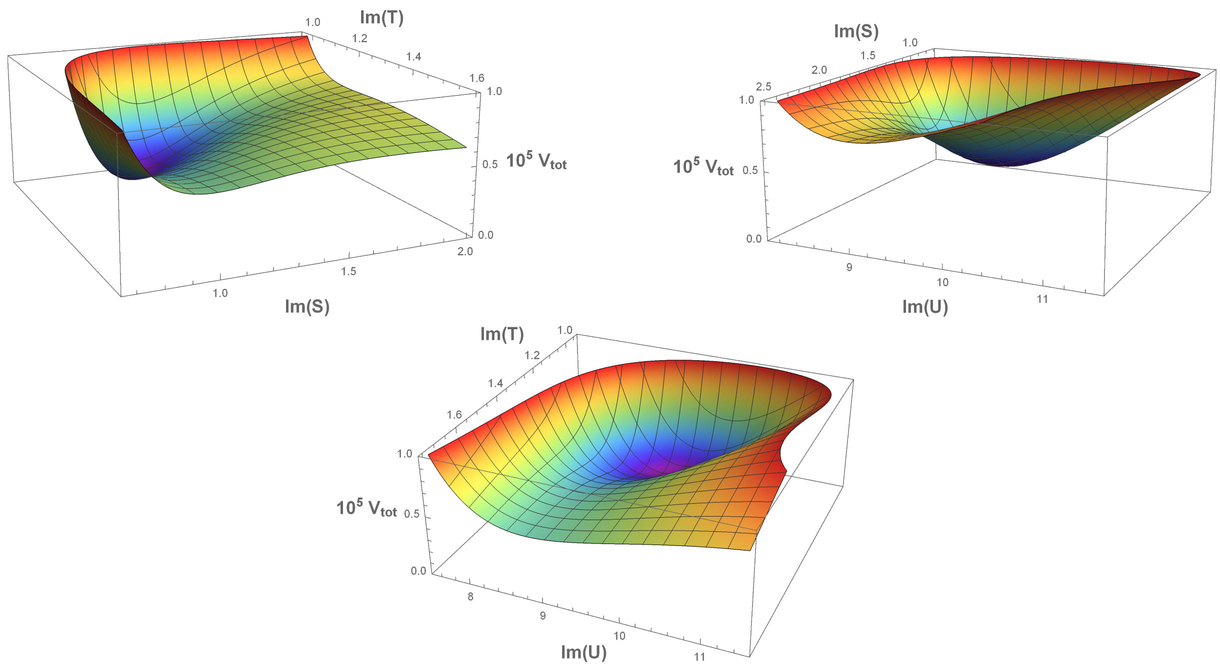

Figure 3. 3D plots of the de Sitter potential for our set of parameters. We have the following slices. Top Left: Im(S) and Im(T), Top Right: Im(S) and Im(U), Bottom: Im(U) and Im(T). In all three different plots the de Sitter minimum is clearly visible and it is metastable.

4 Mass Producing dS Vacua

While the method described above seems to work very well, at least for the class of models we investigated, there is not particular reason it has to work. In [2, 35] it was proposed that it is possible to make a prediction about the existence and success of this method by introducing an additional first step. The suggested procedure is as follows:

1. Solve the F-term equations $D_iW = 0$ as well as the Minkowski condition: $W = 0$. This will give a Minkowski vacuum without negative masses. Adjust the free parameters such that the masses are non-zero.

2. Introduce a parametrically small shift into the scalar potential: $W \to W + \Delta W$. This will change the vacuum from Minkowski to anti de Sitter. It is possible to show that the resulting AdS state is stable with all masses still positive.

3. Introduce the uplift via anti-D6-branes. For small uplifts the changes will be again small and the vacuum stable.

In the following we will outline how this procedure guarantees stable vacua. The details and more information about this procedure can be found in [35].

4.1 Fermion and Scalar Masses

For the general case of some number of chiral matter superfields z^a and some Kähler potential $K(z^a, \bar{z}^{\bar{a}})$ as well as an arbitrary holomorphic superpotential $W(z^a)$ it is possible to show that the above works for as long as we can solve the equations outlined in 1) and find strictly positive masses.
For this section we will use the notation of [36, 37] where we use a covariantly holomorphic superpotential

$$m(z^a, \bar{z}^{\bar{a}}) = \mathrm{e}^{K/2} W \,, \tag{4.1}$$

which is related to the gravitino mass via

$$M_{3/2}^2 = |m\bar{m}| \,. \tag{4.2}$$

The Kähler covariant derivative on m is defined to be $D_a m = \partial_a m + \frac{1}{2}(\partial_a K) m \equiv m_a = \mathrm{e}^{K/2} D_a W$ and likewise for the barred quantity. m, like W, is holomorphic, meaning that:

$$\bar{D}_{\bar{a}} m = \partial_{\bar{a}} m - \frac{1}{2}(\partial_{\bar{a}} K) m = 0 \,. \tag{4.3}$$

The fermion mass matrix is given as

$$m_{ab} \equiv D_a D_b \qquad \bar{m}_{\bar{a}\bar{b}} \tag{4.4}$$

or, alternatively, as [38]

$$m_{ab} = \mathrm{e}^{K/2}(\partial_a + K_a) D_b W - \mathrm{e}^{K/2} \Gamma^c{}_{ab} D_c W \tag{4.5}$$

which, at the supersymmetric Minkowski minimum, simplifies to

$$m_{ab}^{\text{Mink}} = \mathrm{e}^{K/2} \partial_a \partial_b W \,. \tag{4.6}$$

The standard supergravity scalar potential (2.10) can be re-written in terms of this notation as

$$V = m_a K^{a\bar{b}} \bar{m}_{\bar{b}} - 3|m|^2 = |m_a|^2 - 3|m|^2 \,. \tag{4.7}$$

An extremum of the scalar potential is given at

$$\partial_a V = D_a V = -2m_a \bar{m} + m_{ab} K^{b\bar{b}} \bar{m}_{\bar{b}} = 0\,. \tag{4.8}$$

If the extremum is supersymmetric we have $m_a = \bar{m}_{\bar{b}} = 0$.
The masses at an extremum are given by the second derivative of the scalar potential and read

$$\mathcal{M}^2 = \begin{pmatrix} V_{a\bar{b}} & V_{ab} \\ V_{\bar{a}\bar{b}} & V_{\bar{a}b} \end{pmatrix}\,. \tag{4.9}$$

4.2 Minkowski Vacua

For our first step we want to find a Minkowski vacuum [37], which, in the conventions of the last section, is given by

$$m = m_a = a\,. \tag{4.10}$$

The masses in the supersymmetric Minkowski vacuum are

$$\left(\mathcal{M}^2\right)^{\mathrm{Mink}} = \begin{pmatrix} V_{a\bar{b}}^{\mathrm{Mink}} & 0 \\ 0 & V_{\bar{a}b}^{\mathrm{Mink}} \end{pmatrix}\,, \tag{4.11}$$

where:

$$V_{a\bar{b}}^{\mathrm{Mink}} = m_{ac} K^{c\bar{c}} \bar{m}_{\bar{c}\bar{b}}\,. \tag{4.12}$$

We note that the mass matrix is block diagonal and both blocks are positive definite and thus all eigenvalues will be non-negative. This means that, in Minkowski, we know for certain that there will be no tachyons to worry about.
For the remainder of this work we will require that all masses are larger than zero. In practice this is usually the case for a chosen set of parameters unless some accidental cancellations happen.

4.3 Downshift to AdS

Solving for a Minkowski vacuum means satisfying the equations $D_a W = 0$ as well as $W = 0$. When we take these solutions but add a small contribution ΔW to the superpotential ($W \to W + \Delta W$), we find that the previously found Minkowski vacuum becomes an AdS Minimum. This will change the position of the Minimum slightly and we also have to keep in mind that the supersymmetry conditions are not solved explicitly. This means that we have to check that this state remains superymmetric.

Here we will show that a parametrically small downshift ΔW will lead only to a small modification of the vacuum and the state will remain stable.
Using the covariant notation introduced above the change in the superpotential means that

$$m = \mathrm{e}^{K/2} W \neq 0 \qquad \text{and} \qquad m_a = 0\,. \tag{4.13}$$

In the following we will show the shift of the position of the minimum. We explicitly require that the shift appears only in the superpotential W and the functional dependence of both potentials on the moduli remains the same:

$$K^{AdS}(z,\bar{z}) = K^{Mink}(z,\bar{z}) \qquad \text{and} \qquad W^{AdS}(z) = W^{Mink}(z) + \Delta W\,, \tag{4.14}$$

where ΔW is a constant that does not depend on the moduli. For supersymmetry we require that $m_a = \mathrm{e}^{K/2} D_a W = 0$. This requirement tells us that the shift in z needs to be given as

$$\delta z^a = -(m_{ab})^{-1} K_b \Delta m + \cdots\,, \tag{4.15}$$

where the dots are sub leading terms. Δm is much smaller than the smallest eigenvalue m_χ of the mass matrix m_{ab}.

$$\Delta m = \mathrm{e}^{K/2} \Delta W \ll m_\chi\,. \tag{4.16}$$

In Minkowski space we found that the mass matrix is block diagonal (4.11). Now this does not hold any more and we have generically:

$$\left(\mathcal{M}^2\right)^{\mathrm{AdS}} = \begin{pmatrix} V^{\mathrm{AdS}}_{a\bar{b}} & V^{\mathrm{AdS}}_{ab} \\ V^{\mathrm{AdS}}_{\bar{a}\bar{b}} & V^{\mathrm{AdS}}_{\bar{a}b} \end{pmatrix}\,. \tag{4.17}$$

The change in the mass matrix due to the downshift turns out to be

$$V^{AdS}_{a\bar{b}} = m_{ac} K^{c\bar{c}} \bar{m}_{\bar{c}\bar{b}} - 2K_{a\bar{b}} m\bar{m}\,. \tag{4.18}$$

We observe that the first part in this formula does not change from the Minkowski case and is still positive definite. The second term, however, has an explicit minus sign. This does not pose a problem since it is parametrically small compared to the first term. For reasonably large positive masses in Minkowski space the AdS masses will be positive as well.

4.4 Lifting the Minimum to de Sitter

After finding a Minkowski minimum where all masses are positive and subsequently shifting it to anti de Sitter by introducing a *small* constant shift in the superpotential we arr no free to apply the uplifting procedure as outlined in in sections 2.3 and 3.
In the $4d$ effective theory the uplifting contribution from an anti-Dp brane is given by the inclusion of a nilpotent field X [18] that satisfies $X^2 = 0$. The complete Kähler - and superpotential are given as

$$K^{dS} = K^{AdS} + K_{X\bar{X}} X\bar{X} \qquad \text{and} \qquad W^{dS} = W^{AdS} + \mu^2 X\,. \tag{4.19}$$

The inclusion of the nilpotent field X introduces additional contributions into m that we nee to take into account. For the derivatives of m we now also need include derivatives on X and for this we define $I = \{a, X\}$. The new derivatives are:

$$\begin{aligned} m_X &= D_X m\,, \\ m_{aX} &= D_a D_X m\,, \\ m_{abX} &= D_a D_b D_X m \end{aligned} \tag{4.20}$$

and so on. For the de Sitter scalar potential one finds

$$V_{dS} = \mathrm{e}^K \left(D_I W K^{I\bar{J}} \overline{D_{\bar{J}} W} - 3|W|^2 \right) = |m_I|^2 - 3|m|^2 = |F|^2 - 3m_{3/2}^2 > 0\,. \tag{4.21}$$

Here $|F|^2 = |m_I|^2$ denotes the supersymmetry breaking terms that have to be above the scale of the gravitino.
In the end the final goal would be to build a dS vacuum where the value of the potential at the minimum matches the cosmological constant, so:

$$V^{dS}|_{min} = |m_I|^2 - 3|m|^2 \sim 10^{-120} \qquad \text{and} \qquad V'_{dS}|_{min} = 0\,. \tag{4.22}$$

We are now interested whether or not we can show that under the assumptions of section 4 we can *predict* that the vacuum will be stable after the uplift. In other words we need to check that all eigenvalues of the mass matrix are positive. The the holomorphic-holomorphic and holomorphic-anit-holomorphic part of the mass matrix are given, respectively, by

$$
\begin{aligned}
V^{dS}_{ab} &= -m_{ab}\bar{m} - m_{abI}\bar{m}^{I} \qquad \text{and} \\
V^{dS}_{a\bar{b}} &= m_{ac}K^{c\bar{c}}\bar{m}_{\bar{c}\bar{b}} - 2K_{a\bar{b}}m\bar{m} + K_{a\bar{b}}m_{I}\bar{m}^{I} - R_{a\bar{b}I\bar{I}}\bar{m}^{I}m^{\bar{I}} - m_{a}\bar{m}_{\bar{b}}\,,
\end{aligned}
\tag{4.23}
$$

where $R_{a\bar{b}I\bar{I}}$ is the moduli space curvature. Note that we do not consider the mass of X since it is not a fundamental scalar but rather given in terms of fermions as outlined in section 2.2.

The change both in the masses and also in the position of the minimum after the uplift can be related back to the anti de Sitter expressions by considering how the uplift depends on the moduli. Since the uplift is given by

$$
m_X = \mathrm{e}^{K/2}D_XW \Rightarrow |m_X|^2 = \mathrm{e}^{K}\mu^4\,, \tag{4.24}
$$

it is evident that for parametrically small μ^2 the shift in the position will be small. Indeed we will see that in all examples given in sections 5 and 6 the uplift parameter is indeed small. To make this statement more precise we remember that the dS scalar potential is the sum $V_{dS} = V_{AdS} + V_{up}$ and that the AdS part does not get modified. Now we split the complex fields in real and imaginary part $z^a = z^a_r + \mathrm{i}z^a_i$ and consider the condition for the minimum:

$$
\partial_{z^a_\alpha}\left[V^{AdS} + V^{up}\right] = 0\,. \tag{4.25}
$$

We can write the first and second part respectively as:

$$
\begin{aligned}
\partial_{z^a_\alpha}V^{AdS} &= \left(\partial_{z^a_\alpha}\partial_{z^b_\beta}V^{AdS}\right)\delta z^b_\beta\,, \\
\partial_{z^a_\alpha}V^{up} &= \mu^4\partial_{z^a_\alpha}\left(\mathrm{e}^{K}K_{X\bar{X}}\right)\,.
\end{aligned}
\tag{4.26}
$$

This allows us to write the shift in the minimum as

$$
\delta z^b_\beta = -\mu^4\partial_{z^a_\alpha}\left(\mathrm{e}^{K}K_{X\bar{X}}\right)\left(\partial_{z^a_\alpha}\partial_{z^b_\beta}V^{AdS}\right)^{-1}\,. \tag{4.27}
$$

The masses in anti de Sitter, given by $\partial_{z^a_\alpha}\partial_{z^b_\beta}V^{AdS}$ are larger than the scale of the uplift and thus the shift of the position of the minimum will be parametrically small.

In order to guarantee the stability of this constructed de Sitter vacuum we require that the amount of supersymmetry breaking, introduced by the uplift, be small compared to the scale of the masses of the scalars given by the second derivative of the scalar potential. This is in addition to the condition that the gravitino has small mass when compared to the scalars:

$$m_\chi^2 \gg |F|^2 = |m_I|^2 \qquad \text{and} \qquad m_\chi^2 \gg m_{3/2}^2\,. \tag{4.28}$$

Since the potential is positive we also know that

$$|F|^2 > 3m_{3/2}^2\,, \tag{4.29}$$

Furthermore, from the formulation of the shift we conclude that breaking of supersymmetry is smaller in the chiral directions than in the direction of the nilpotent field X [39]:

$$|m_a|^2 \ll |m_X|^2\,. \tag{4.30}$$

With this we have established a mass hierarchy: $|m_a|^2 \ll |m_X|^2 \ll m_\chi^2$ and thus the amount of supersymmetry breaking if of the order of the gravitino mass, at least for a very small cosmological constant.

Finally, since all contributions that we added after finding our Minkowski vacuum lead to parametrically small changes, we conclude that, in fact, the change to the mass matrix is small after all these steps:

$$V_{a\bar{b}}^{dS} \approx V_{a\bar{b}}^{Mink}\,. \tag{4.31}$$

5 KL-Models

In [2] the process that was outlined in the previous section was expanded on and many explicit examples were built. The framework in [2] is based on so called Kallosh-Linde (KL) models [40] where the moduli are stabilized using contributions from multiple non-perturbative corrections, similar to the uplift models in [1] and as outlined here in section 3. In particular it was necessary to introduce two exponents for each modulus in order to achieve stability and facilitate the uplift. This is due to the requirement to find a *Minkowski* vacuum, something that was not done in [1] because there the starting point for the uplift procedure was an AdS state.

5.1 The Kallosh-Linde Model

The superpotential for the KL model is given as

$$W^{KL}(\Phi, X) = W_0 + \sum_i \left(A_i \mathrm{e}^{\mathrm{i}a_i\Phi^i} - B_i \mathrm{e}^{\mathrm{i}b_i\Phi^i} \right) , \tag{5.1}$$

where Φ^i runs over all moduli that we consider from compactification but does not include the nilpotent field X, which introduces the uplift. Note that here we do not sum over repeating indices. The Kähler potential and superpotential for the uplift remain unchanged in this set up. In order to find a Minkowski minimum with a KL-type superpotential one has to solve the equations

$$\partial_{\Phi^i} W^{KL} = \mathrm{i}a_i A_i \mathrm{e}^{\mathrm{i}a_i\Phi^i} - \mathrm{i}b_i B_i \mathrm{e}^{\mathrm{i}b_i\Phi^i} = 0 \qquad \forall i \tag{5.2}$$

and, the Minkowski condition

$$W^{KL} = W_0 + \sum_i \left(A_i \mathrm{e}^{\mathrm{i}a_i\Phi^i} - B_i \mathrm{e}^{\mathrm{i}b_i\Phi^i} \right) = 0 \, . \tag{5.3}$$

For as long as all masses are positive we can consistently simplify this problem by writing $\Phi^i = \phi^i + \mathrm{i}\theta^i$ and then letting all the $\theta^i = 0$. The position of the minimum is then given by

$$\Phi_0^i = \phi_0^i = \frac{\mathrm{i}}{a_i - b_i} ln \left(\frac{a_i A_i}{b_i B_i} \right) \tag{5.4}$$

and the constant and the constant part of the superpotential, W_0, is fixed to be

$$W_0 = \sum_i \left\{ -A_i \left(\frac{a_i A_i}{b_i B_i} \right)^{\frac{a_i}{b_i - a_i}} + B_i \left(\frac{a_i A_i}{b_i B_i} \right)^{\frac{b_i}{b_i - a_i}} \right\} \tag{5.5}$$

at the minimum. We see that we also have to satisfy the constraints on the parameters: $a_i > b_i$ and $a_i A_i > b_i B_i$ (ore vice versa). Solving these conditions is sufficient in order to find the stable Minkowski vacuum. The remaining free parameters have to be used in order to make sure that all masses, given, essentially, by the second derivative of the scalar potential, are strictly positive. This will the guarantee that the mass production procedure will be successful, as per the discussion of the last section.

We will now use models like this and the three step program for mass production as outlined in section 4 in order to build de Sitter vacua in both type IIA and type IIB models.

5.2 The Setup in Type IIA

First we want to study explicit models in type IIA theory. For this we consider the following Kähler and superpotential:

$$\begin{aligned} K &= -\sum_{i=1}^{m} N_i \log\left(-\mathrm{i}(\Phi^i - \bar{\Phi}^{\bar{i}}\right) \\ W &= W_0 + \sum_{i=1}^{m}\left(A_i \mathrm{e}^{\mathrm{i}a_i\Phi^i} - B_i \mathrm{e}^{\mathrm{i}b_i\Phi^i}\right) . \end{aligned} \tag{5.6}$$

The masses of the scalars are proportional to the second derivative of the scalar potential and there is a mass degeneracy between the moduli and the axions, meaning $m^2_{\phi_i\phi_i} = m^2_{\theta_i\theta_i}$. The mixed derivatives on the scalar potential are zero $\partial_{\phi_i}\partial_{\theta_j}V = 0$ and thus the mass matrix is block diagonal. For the Minkowski solution the mass matrix is actually *diagonal* meaning the only for $i = j$ the entries are non-zero. The downshift $W \to W + \Delta W$ introduces a contribution into the mass matrix that we call $\Delta_{\Phi^i\Phi^j}$. Importantly this will keep the block diagonal structure of the matrix but will introduce off-diagonal entries inside the blocks. In other words the complete mass matrix is still block diagonal but no longer diagonal after the downshift to anti de Sitter space:

$$\left(\mathcal{M}^2\right)^{\mathrm{AdS}} = \begin{pmatrix} V^{\mathrm{Mink}}_{\phi^i\phi^j} + \Delta_{\phi^i\phi^j} & 0 \\ 0 & V^{\mathrm{Mink}}_{\theta^i\theta^j} + \Delta_{\theta^i\theta^j} \end{pmatrix} . \tag{5.7}$$

In the third step an additional contribution will be added to the blocks, arising from the uplift term that is sourced by the inclusion of anti-D6 branes in type IIA:

$$\left(\mathcal{M}^2\right)^{\mathrm{dS}} = \begin{pmatrix} V^{\mathrm{Mink}}_{\phi^i\phi^j} + \Delta_{\phi^i\phi^j} + \tilde{\Delta}_{\phi^i\phi^j} & 0 \\ 0 & V^{\mathrm{Mink}}_{\theta^i\theta^j} + \Delta_{\theta^i\theta^j} + \tilde{\Delta}_{\theta^i\theta^j} \end{pmatrix} . \tag{5.8}$$

Again, the mass matrix remains block diagonal. Since we know that the contributions from downshift and uplift are both parametrically small we expect to find a stable de Sitter minimum after the mass production procedure.

5.3 Seven Moduli in Type IIA

A rather general and very interesting model in type IIA theory is a seven moduli model, originally developed in [41, 42], which is the general version of the STU model from section 3.1. In this model there is no identification for the Kähler - and complex structure moduli. Thus, we have seven different moduli appearing: one axio-dilaton S, three complex structure moduli T and three Kähler moduli U. The model is type IIA compactified on $T^6/(Z_2 \times Z_2)$ in the presence of generalized fluxes and $D6$-branes as well as $O6$-planes. Alternatively, it can be viewed as an M-theory compactification on a particular G_2 manifold with $E_{7(7)}$ symmetry [43].

This model is of particular interest because it is related to a model from M-theory and type IIA string theory in four dimensions with $\mathcal{N} = 8$ supersymmetry which can be related to B-mode detection [43, 44] and thus is relevant for phenomenology.

In this model, for $\Phi^i = \{S,\, T_1,\, T_2,\, T_3,\, U_1,\, U_2,\, U_3\}$, the superpotential that we are going to use is given by:

$$\begin{aligned} K &= -\sum_{i=1}^{7} \log\left(-\mathrm{i}(\Phi^i - \bar{\Phi}^{\bar{i}}\right), \\ W &= f_6 + \sum_{i=1}^{7}\left(A_i \mathrm{e}^{\mathrm{i}a_i\Phi^i} - B_i \mathrm{e}^{\mathrm{i}b_i\Phi^i}\right) \end{aligned} \tag{5.9}$$

which follows from equation (5.6). Only the contribution from the six-flux appears here and we set all other fluxes, and thus tree level interactions of the moduli are absent. In this model we have 29 parameters, 21 of which will remain free after solving for the Minkowski minimum at $W = \partial_i W = 0$. It is convenient to solve the seven equations for the parameters B_i as well as the flux parameter f_6. As per usual we choose the minimum to be at $\Phi^i = \mathrm{i}\Phi^i_0$, with the real parts of the moduli, which are the axions, equal to zero. We can choose the exact position of the minimum in order to satisfy stringy requirements, for example here we choose $U_0 = 5$ in order to achieve a large volume for the internal manifold. The complete list of choices for our explicit example can be found in Table 2. The choice of parameters does not require any particular

$A_S = 1$	$A_{T_1} = 3.1$	$A_{T_2} = 3.2$	$A_{T_3} = 3.3$	$A_{U_1} = 11$	$A_{U_2} = 12$	$A_{U_3} = 13$
$a_S = 2$	$a_{T_1} = 2.1$	$a_{T_2} = 2.2$	$a_{T_3} = 2.3$	$a_{U_1} = 0.41$	$a_{U_2} = 0.42$	$a_{U_3} = 0.43$
$b_S = 3$	$b_{T_1} = 3.1$	$b_{T_2} = 3.2$	$b_{T_3} = 3.3$	$b_{U_1} = 1.1$	$b_{U_2} = 1.2$	$b_{U_3} = 1.3$
$S_0 = 1$	$T_{1,0} = 1.1$	$T_{2,0} = 1.2$	$T_{3,0} = 1.3$	$U_{1,0} = 5.1$	$U_{2,0} = 5.2$	$U_{3,0} = 5.3$

Table 2. The choice for the parameters in the seven-moduli model.

amount of fine tuning. In fact many different ones will work out. This particular set was chosen for convenience and the slightly different numbers are in order ot avoid accidental cancellations. Furthermore the a_{U_i} are chosen such that the terms in the superpotential stay at approximate same order even though the value of the U_i is larger than the rest of the moduli. One thing to remember is that we have to avoid zero masses in Minkowski. While no negative values can appear, in order for guaranteed stability in anti-de Sitter and de Sitter, we have to have non-zero masses. We will give the eigenvalues of the Minkowski mass matrix together with the masses in de Sitter below, in table 3.

The next step in our procedure is the downshift to anti-de Sitter. For this we introduce a shift in the value of the six flux parameter of the form $f_6 -> f_6 + \Delta f_6$ with the value of the shift given by:

$$\Delta f_6 = -10^{-5} \,. \tag{5.10}$$

An important note here is that the sign of the downshift does not matter for as long as the downshift is actually small. For now we only considered small downshifts, however, in 5.6 we will consider larger shifts and there we will find that it is favourable to use positive shifts. One further note concerns the position of the minimum in anti-de Sitter. The shift is very small but we need to find the position of the minimum after the downshift again and check that it remains supersymemtric. For this we numerically search for the minimum and calculate the first derivative of the scalar potential at that location, which turns out to be (numerically) zero.

Finally, the uplift to de Sitter is achieved by the inclusion of the nilpotent field X in our effective description, as discussed in section 3.2. In type IIA this nilpotent field describes the presence of an anti-D6 brane which gives a positive energy contribution to the scalar potential. This nilpotent field gives, in this case, a contribution to the scalar potential of the form:

$$\begin{aligned} V^{uplift}_{\overline{D6}} = & \frac{\mu_1^4}{\mathrm{Im}(T_1)\mathrm{Im}(T_2)\mathrm{Im}(T_3)} + \frac{\mu_2^4}{\mathrm{Im}(S)\mathrm{Im}(T_2)\mathrm{Im}(T_3)} \\ + & \frac{\mu_3^4}{\mathrm{Im}(S)\mathrm{Im}(T_1)\mathrm{Im}(T_3)} + \frac{\mu_4^4}{\mathrm{Im}(S)\mathrm{Im}(T_1)\mathrm{Im}(T_2)} \,. \end{aligned}$$

The uplift parameters were chosen to be

$$\mu_1^4 = \mu_2^4 = \mu_3^4 = \mu : 4^4 = 5.49028 \cdot 10^{-15} \,. \tag{5.11}$$

These values were not chosen to match the cosmological constant but rather for illustrative purposes. Like the choice of the free parameters we do not have a fine tuning here. The only requirement is that the downshift should be *small* in some sense.
The masses in de Sitter space, given in Table 3, show how little the masses change during the whole procedure. In fact, the anti-de Sitter masses are not given because most of them are identical to the precision we give here. To conclude this section on

	Mink	dS
m_1^2	$1.80473 \cdot 10^{-3}$	$1.80465 \cdot 10^{-3}$
m_2^2	$1.80473 \cdot 10^{-3}$	$1.80465 \cdot 10^{-3}$
m_3^2	$1.37269 \cdot 10^{-3}$	$1.37262 \cdot 10^{-3}$
m_4^2	$1.37269 \cdot 10^{-3}$	$1.37262 \cdot 10^{-3}$
m_5^2	$9.96519 \cdot 10^{-4}$	$9.96472 \cdot 10^{-4}$
m_6^2	$9.96519 \cdot 10^{-4}$	$9.96471 \cdot 10^{-4}$
m_7^2	$1.30924 \cdot 10^{-4}$	$1.30911 \cdot 10^{-4}$
m_8^2	$1.30924 \cdot 10^{-4}$	$1.30911 \cdot 10^{-4}$
m_9^2	$9.41773 \cdot 10^{-5}$	$9.41667 \cdot 10^{-5}$
m_{10}^2	$9.41773 \cdot 10^{-5}$	$9.41660 \cdot 10^{-5}$
m_{11}^2	$6.37973 \cdot 10^{-5}$	$6.37888 \cdot 10^{-5}$
m_{12}^2	$6.37973 \cdot 10^{-5}$	$6.37883 \cdot 10^{-5}$
m_{13}^2	$1.89843 \cdot 10^{-5}$	$1.89809 \cdot 10^{-5}$
m_{14}^2	$1.89843 \cdot 10^{-5}$	$1.89806 \cdot 10^{-5}$

Table 3. The eigenvalues of the mass matrix for the seven-moduli type example. The mass shift is small, but noticeable, when going from Minkowski to dS. One can also notice, as predicted by the mass production procedure, that in dS the masses of scalars and pseudo scalars are not exactly equal any more, as was the case in Minkowski.

type IIA we mention that a three moduli STU model can be easily constructed from by identifying $T = T_1 = T_2 = T_3$ and $U = U_1 = U_2 = U_3$. This was explicitly done in [2] and the results are given there explicitly. The procedure works exactly as in the seven moduli case discussed here and no problems appeared.

5.4 The Setup in Type IIB

Before giving explicit examples for the procedure in type IIB models we use this section in order to show some of the basics that are different from type IIA. A much more

detailed description of the set-up in type IIB can be found in [2].
In our type IIB models we again consider a superpotential of the form

$$W = W_0 + \sum_{i=1}^{m} \left(A_i \mathrm{e}^{\mathrm{i} a_i \Phi^i} - B_i \mathrm{e}^{\mathrm{i} b_i \Phi^i} \right) , \tag{5.12}$$

where now instead of the constant contribution from a six flux we have W_0. Also of note is that the identification of the moduli in type IIB does no longer work like in type IIA. Here the axio-dilaton and the complex structure moduli are stabilized at a higher scale and do not appear in our effective theory. Instead we have different Kähler moduli appearing. For convenience we will however still call the up to three Kähler moduli S, T and U. More importantly in type IIB the Kähler potentials are different from what we considered thus far. In general the Kähler potential will be some function

$$K = K\left(-\mathrm{i}(\Phi^i - \bar{\Phi}^{\bar{i}}) \right) , \tag{5.13}$$

but where now, in some sense, more complicated forms of the potential are allowed. In particular the Kähler potential is given in terms of the internal volume of the manifold $\mathcal{V}_6$ as:

$$K = -2 \log \left(\mathcal{V}_6 \right) . \tag{5.14}$$

The volume of the compact manifold is given as a function of the Kähler moduli and is characteristic for the chosen manifold. The choices that were investigated in [2] were taken from [45] and [46].
The fact that the Kähler potential is now more complicated, when compared to the type IIA case (5.6), has effects on the Minkowski mass matrix and, by extension, anti-de Sitter and de Sitter mass matrix as well. The Minkowski mass matrix is no longer strictly diagonal, as opposed to the matrix in type IIA. However, at every stage the mass matrix will still remain block diagonal. Also, the masses in Minkowski are still positive definite, thus the mass production mechanism works as intended.
One more thing to consider is that the uplift now is performed using *anti-D3 branes*. The general mechanism remains the same, the branes still give effectively a positive contribution to the scalar potential. However, now we have two different choices for the placement of the branes [5]: We can place them either in the bulk or at the bottom

of a warped throat. Depending on the choice the Kähler potential is different

$$\begin{aligned} K_{bulk} &= -2\log\left(\mathcal{V}_6\right) + X\bar{X} \qquad \text{or} \\ K_{throat} &= -3\log\left(\mathcal{V}_6^{2/3} - \frac{1}{3}X\bar{X}\right) , \end{aligned} \tag{5.15}$$

while the contribution to the superpotential remains $\mu^2 X$. This will change the effective contribution to the scalar potential to either

$$\begin{aligned} V_{\overline{D6},\, bulk} &= \frac{\mu^4}{\mathcal{V}_6^2} \qquad \text{or} \\ V_{\overline{D6},\, throat} &= \frac{\mu^4}{\mathcal{V}_6^{4/3}} . \end{aligned} \tag{5.16}$$

Here we will only consider the placement of the brane in the bulk. Both possibilities have been considered in [2] and it was found that, other than a difference in the uplift parameter μ^4, everything works equally well in both situations.

5.5 The Fibre Inflation Model

In [2] many different models in type IIB theory were investigated: K3-fibration models, complete intersection Calabi-Yau models and so called multi-hole Swiss cheese models, based on the Fano three-fold $\mathcal{F}_{11}$. For simplicity we restrict the discussion here to one of the K3-fibrations as it is a model of particular interest: the fibre inflation model [47–49]. While building a potential suitable for inflation is not the goal of this work, it is still of particular interest for the possibility to do so.
In our notation the volume of the internal manifold for the *fibre inflation model* is given by:

$$\mathcal{V}_6 = \alpha \left[\sqrt{\left(-\mathrm{i}(S-\bar{S})\right)} \left(-\mathrm{i}(T-\bar{T})\right) - \gamma \left(-\mathrm{i}(U-\bar{U})\right)^{3/2} \right] , \tag{5.17}$$

where α and γ are some positive constant that we can choose. Besides the values for the parameters in the potentials, which are given in Table 4, we choose the value of the downshift ($W_0 \to W_0 + \Delta W$) to be:

$$\Delta W_0 = -10^{-5} \tag{5.18}$$

and the uplift parameter to be

$$\mu_{bulk}^4 = 3.10079 \cdot 10^{-10} \,. \tag{5.19}$$

These parameters lead, using our mass production procedure, to the eigenvalues of

$A_S = 1.1$	$A_T = 1.2$	$A_U = 1.3$
$a_S = 2.1$	$a_T = 2.2$	$a_U = 2.3$
$b_S = 3.1$	$b_T = 3.2$	$b_U = 3.3$
$S_0 = 1$	$T_0 = 1$	$U_0 = 1$
$\alpha = 1$	$\gamma = \frac{1}{2}$	

Table 4. One possible set of parameters for the fibre inflation model.

the mass matrix as given in Table 5. It is evident from that table that the masses in Minkowski and de Sitter are very close to each other, once more confirming that the parametrical separation during the process works out as discussed.
To illustrate the model even better we include two dimensional plots of the scalar potential for all directions in figure 4. In these plots we not only show the overall behaviour of the potential in all three directions but also a close-up of the minimum, which shows the shift of the position after the procedure. Furthermore, we include one three dimensional plot of the scalar potential, for the $Im(T)$ and $Im(U)$ direction in figure 5. These plots show that the potential is indeed stable.

	Mink	dS
m_1^2	1.01997	1.01957
m_2^2	1.01997	1.01957
m_3^2	$1.31424 \cdot 10^{-1}$	$1.31344 \cdot 10^{-1}$
m_4^2	$1.31424 \cdot 10^{-1}$	$1.31338 \cdot 10^{-1}$
m_5^2	$2.44807 \cdot 10^{-2}$	$2.44724 \cdot 10^{-2}$
m_6^2	$2.44807 \cdot 10^{-2}$	$2.44665 \cdot 10^{-2}$

Table 5. The eigenvalues of the mass matrix for the fibre inflation model, for Minkowski and dS.

5.6 Stability under large Shifts

Thus far we always considered the *mass production* procedure with small downshifts and subsequently small uplifts. This was shown to guarantee that, if we started from a Minkowski vacuum with all masses positive, we would arrive in a stable de Sitter

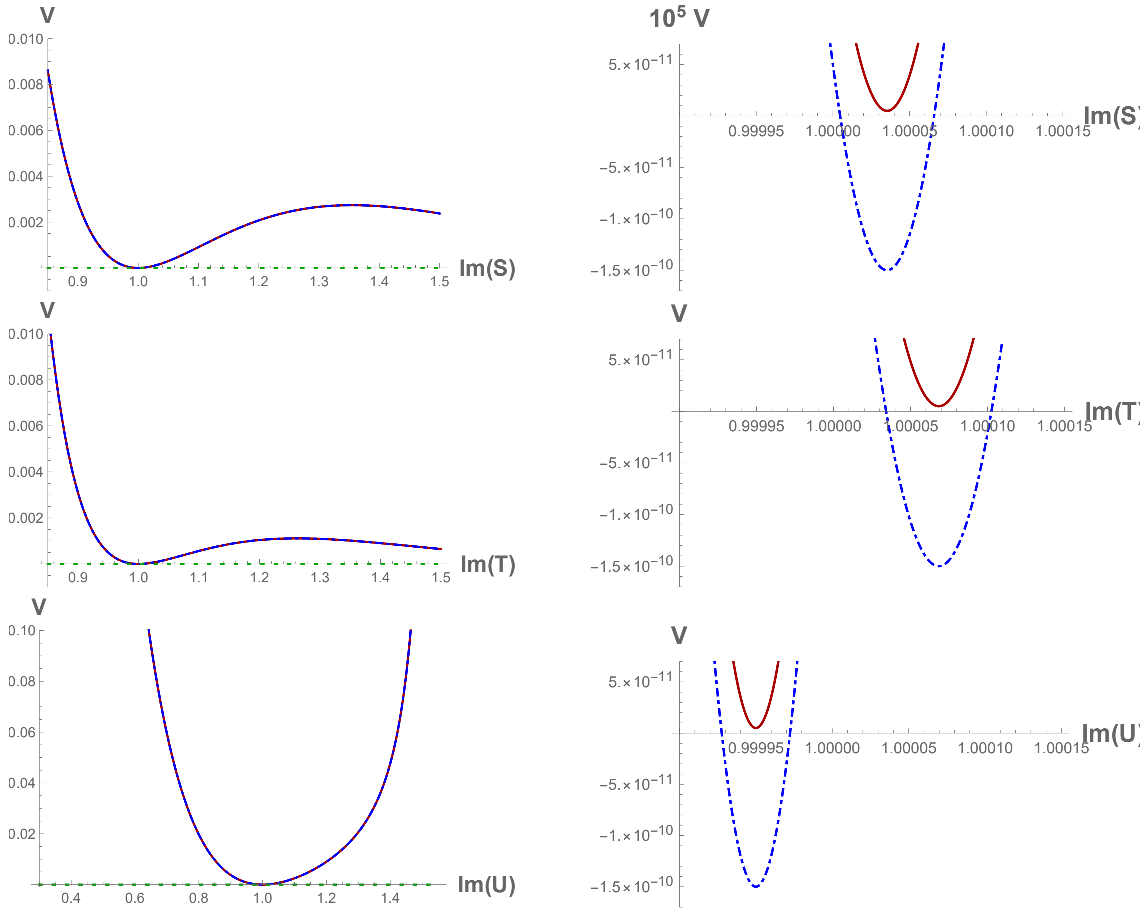

Figure 4. From top to bottom, we have the overall form of the potential on the left, for AdS and dS as well as a close-up of the minimum on the right for the directions Im(S), Im(T) and Im(U). The shift of the minimum from the initial point, with the imaginary part of the moduli set at one, is also visible.

vacuum. Admittedly, this is a vague criterion. It is therefore interesting to ask how small the shift needs to be, or if it might be possible to do large shifts. While we were only able to show that small shifts predict the procedure to be successful, there is a priori no statement about large shifts. One might expect them to not work but, in principle, there is nothing forbidding it to be possible. We devote this section to a short investigation about possible *large* shifts and show in the simple example of a one moduli set up with a Kallosh-Linde superpotential that large shifts can work out.

One important change when going to larger shifts is that the sign of ΔW becomes relevant. It turns out that for this scenario it is preferable to have a positive shift, as it will lead to stronger stabilization.

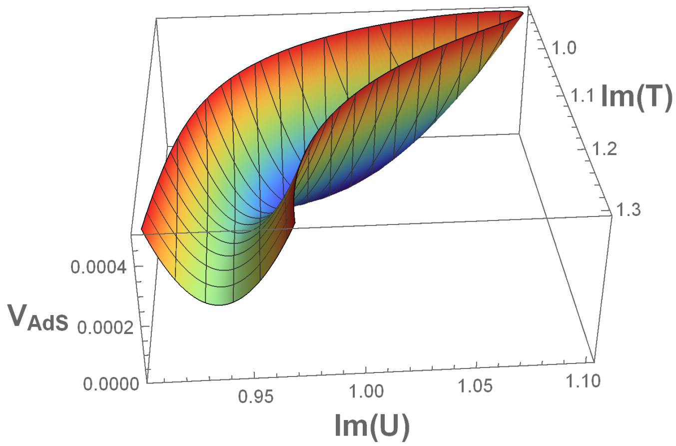

Figure 5. A 3D slice of the scalar potential in the Im(T) and Im(U) directions.

The model we want to consider is given by the Kähler - and superpotential

$$K = -3\log\left(-\mathrm{i}(T-\bar{T}\right) + X\bar{X} \qquad \text{and}$$
$$W = \mathrm{e}^{\mathrm{i}\frac{\pi T}{25}} - \mathrm{e}^{\mathrm{i}\frac{\pi T}{10}} + W_0 + \Delta W + \mu^2 X\,. \tag{5.20}$$

This model was used in [50] for inflation model building and we choose to keep the same values for the parameters as in that paper. For $\Delta W = \mu = 0$ and $W_0 = -\frac{3}{5}\left(\frac{2}{5}\right)^{2/3}$ we find a supersymmetric Minkowski minimum of the scalar potential as illustrated in figure 6. The mass of the scalar at the Minkowski minimum is $m_T = 0.013$. On the left

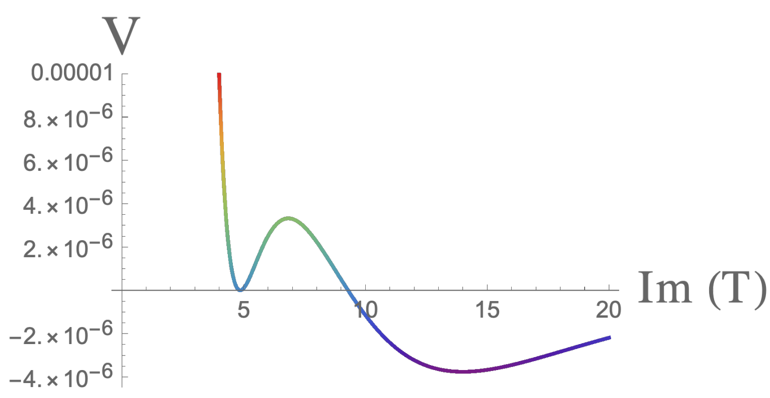

Figure 6. The potential of the KL model (5.20) in Planck units.

hand side in figure 7 we show the potential after the procedure with a shift of $\Delta W = 7$

and a large uplift of $\mu = 1.4019$, which was chosen to arrive at a de Sitter vacuum with small cosmological constant. The mass in de Sitter evaluates to $m_T = 0.27$, we see that the mass actually is higher than in Minkowski, large shifts stabilize the modulus. Another nice feature is that it is actually possible to change both ΔW and μ simultaneously in such a way that the minimum stays more or less constant. This allows to interpolate between a supersymmetric Miknowski minimum with $m_{3/2} = 0$ and a de Sitter vacuum with strongly broke supersymmetry, where the gravitino has a Planckian value of 1 for its mass. It is even possible to raise the value of the cosmological constant

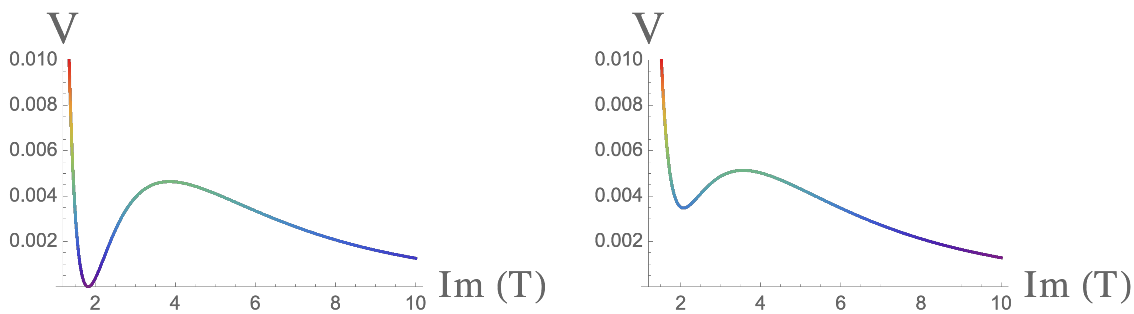

Figure 7. The potential of the KL model (5.20) after the downshift obtained by adding $\Delta W = 7$ to the superpotential. The left panel shows the potential uplifted by $\mu = 1.40199$. The right panel shows the potential uplifted by $\mu = 1.42$.

by increasing μ. On the right hand side of figure 7 this is pictured for $\Delta W = 7$ and $\mu = 1.42$. The value of the potential in the de Sitter minimum is several orders of magnitude larger than for the previous values and still the minimum is stable.
The very important conclusion here is that one can reach any desired value of supersymmetry breaking and it is possible to choose the cosmological constant with great freedom. In [2] this was also tested in the STU model in type IIA theory and it worked out in the same way as here. This gives hope that this is a general statement that should hold for any reasonable mass production model.

6 Models from M Theory

While the main focus of explicit examples thus far has been on models with a Kallosh-Linde superpotential, the mass production procedure does not explicitly require this exact type of potential. The requirement simply is that a Minkowski solution is possible. The KL-type potentials are useful because it is know that a Minkowski exists for them and also how it looks. While models with single exponents alone do not allow to construct Minkowski vacua it is possible that one can find a solution if in addition to the non-perturbative terms, generating the exponents in the superpotential, also flux contributions are present. In [3] it was investigated whether or not such tree level contributions are sufficient to use the mass production procedure. For this investigation models based on M-theroy compactified on $\mathbb{T}^6/\mathbb{Z}_2^2$ are considered. All models have seven moduli which are the coordinates of the coset space $[SL(2,\mathbb{R})/SO(2)]^7$. These models have an interpretation in terms of either type IIA or IIB string theory, however, the origin of one particular type of flux is not fully understood yet as we will discuss later.

6.1 The generalized twisted seven Torus

To start with we consider moduli stabilization on M-theory compactified on a seven torus with a G_2-structure. This is then called the *twisted seven torus*. Mathematically it is obtained by considering the toroidal orbifold $X_7 = \mathbb{T}^7/(\mathbb{Z}_2 \times \mathbb{Z}_2 \times \mathbb{Z}_2)$. Interestingly, this model leads directly to 4d, $\mathcal{N} = 1$ supergravity with seven moduli. The derivation of the twisted seven torus model from M-theory was proposed in [51] and [52]. It involves describing the set up using a pseudo action where the potentials and dual curvatures appear at the same time.

Following [51] and [53] we have the potentials in this model as

$$\begin{aligned} K &= -\sum_{i=1}^{7} \log\left(-\mathrm{i}(\Phi^i - \bar{\Phi}^i)\right) \qquad \text{and} \\ W_{pert} &= g_7 + G_i \Phi^i + \frac{1}{2} M_{ij} \Phi^i \Phi^j \,, \end{aligned} \tag{6.1}$$

where the Φ^i are the seven moduli, g_7 is a seven flux contribution and we choose $G_i = 0$ in order to have only quadratic terms in the superpotential. Finally, M_{ij} is symmetric and its diagonal entries vanish. These interactions arise from geometric fluxes. Adding non-perturbative corrections, arising for example from wrapped M2-branes [54], the superpotential we consider is:

$$W = g_7 + \frac{1}{2} M_{ij} \Phi^i \Phi^j + \sum_{i=i}^{7} A_i \mathrm{e}^{\mathrm{i} a_i \Phi^i} \,. \tag{6.2}$$

Our conventions remain such that all parameters in the superpotential are real.

Once again we follow the *mass production procedure*, as outlined in section 4. We have to solve the equations $\partial_i W = 0$ and $W = 0$ in order to find a supersymmetric

Minkowski vacuum. First, considering $\partial_i W$ we find that

$$A_i = \mathrm{i} a_i^{-1} \mathrm{e}^{-\mathrm{i} a_i \Phi_i} M_{ij} \Phi^j \, . \tag{6.3}$$

As always we choose to use $\Phi^i = \theta^i + \mathrm{i}\phi^i$ and set $\theta^i = 0$ because only the ϕ^i appear in the Kähler potential. Plugging this back into the superpotential it is easy to solve for $W = 0$ and we arrive at the required Minkowski solution. In the earlier sections we always used a KL-type double exponent in order to arrive at a Minkowski minimum, here we achieve this by the inclusion of fluxes instead. Furthermore, we will find in the discussion down below that it is even possible to completely get rid of the non-perturbative corrections in some directions. This is interesting because in some directions these corrections are better motivated than in others.

In [3] a variety of different models were investigated and explicit numerical examples were given. Here we will list all these examples, however, we choose to omit the numbers as they are not very illuminating. For all models it was not only found that the procedure works and produces stable de Sitter minima but also no particular amount of fine tuning was necessary. In other words a large parameter space is viable in order to construct these models.

For the discussion of the models we will again use the notation where we split $\Phi^i = \{S, T_I, U_I\}$, $I = 1, 2, 3$, as in the earlier sections, regardless of the theory we are in. Some of the models do have interpretations in terms of type IIA/IIB theory and there this nomenclature is meaningful. For us it serves mainly as a presentation tool. In this notation we can represent the quadratic terms as [53]:

$$\frac{1}{2} M_{ij} \Phi^i \Phi^j = S b^K U_K + U_I C^{IJ} T_j + a^I \frac{U_1 U_2 U_3}{U_I} + c^I \frac{T_1 T_2 T_3}{T_I} + S d^K T_K \, . \tag{6.4}$$

The models we will consider differ in which ones of the parameters in this equation and which non-perturbative terms are non-zero.

Model 1, all exponents

First off, we can consider a model where non-perturbative corrections in all directions are present. For this we consider the superpotential (6.2) with the expansion (6.4) where we set $a^I = c^I = d^I = 0$. This reduces the superpotential to:

$$W_1 = g_7 + b^k S U_K + C^{IJ} U_I T_J + A_S \mathrm{e}^{\mathrm{i} a_S S} + \sum_I A_{T_I} \mathrm{e}^{\mathrm{i} a_{T_I} T_I} + \sum_I A_{U_I} \mathrm{e}^{\mathrm{i} a_{U_I} U_I} \, . \tag{6.5}$$

There are 19 parameters in this potential, 8 of which will be fixed by the Minkowski conditions. In particular we choose g_7 and the A_i. The rest of the parameters can be used in order to tune the masses to our liking.

Model 1, all exponents, with satisfied tadpole conditions
Here we consider the same set up as above but we also solve the tadpole conditions

$$\begin{aligned} b^I C^{IJ} + b^J C^{II} &= 0 \\ C^{IJ} C^{JK} + C^{IK} C^{JJ} &= 0 \end{aligned} \tag{6.6}$$

without introducing sources on the right hand side of the equation. This is usually problematic and thus typically one has to include sources like O-planes on the right hand side of these equations. In this specific model, however, wer were able to satisfy all the relevant tadpole conditions in terms of the C^{IJ} with $I \neq J$.

Model 2, without S exponent
Again we consider a similar set up as before but now we set $A_S = a_S = 0$ from the very beginning:

$$W_2 = g_7 + b^k S U_K + C^{IJ} U_I T_J + \sum_I A_{T_I} \mathrm{e}^{\mathrm{i} a_{T_I} T_I} + \sum_I A_{U_I} \mathrm{e}^{\mathrm{i} a_{U_I} U_I} \, . \tag{6.7}$$

We solve the Minkowski equations for g_7, A_{T_I}, A_{U_I} and now in addition b^1. This does not prohibit a solution and an explicit example was found.

Model 3, without U exponents
Even more interesting is the case where we do not include U-exponents into the superpotential at all. For this we are, however, required to include the tree-level term $a^I \frac{U_1 U_2 U_3}{U_I}$ into the potential:

$$W_3 = g_7 + b^k S U_K + C^{IJ} U_I T_J + a^I \frac{U_1 U_2 U_3}{U_I} + A_S \mathrm{e}^{\mathrm{i} a_S S} + \sum_I A_{T_I} \mathrm{e}^{\mathrm{i} a_{T_I} T_I} \, . \tag{6.8}$$

Once more an explicit solution for this model is possible. This is important because the exponents in the U-directions are somewhat less well established than the other directions.

Model 4, without T and U exponents
If we decide to include even more tree level contributions into the superpotential we are able to arrive at a solvable model that includes only one exponent in the S-direction:

$$W_4 = g_7 + b^k S U_K + C^{IJ} U_I T_J + a^I \frac{U_1 U_2 U_3}{U_I} + c^I \frac{T_1 T_2 T_3}{T_I} + A_S \mathrm{e}^{\mathrm{i} a_S S} \,. \qquad (6.9)$$

Now we have 21 parameters in the superpotential and we solve for the parameters g_7, A_S, a^I and c^I.
One important remark about the relation of these models to type IIA string theory is in order. Models 1-3 can be related to type IIA on a generalized twisted six-torus. There we can identify the ingredients with standard contributions also present in IIA. Model 4 on the other hand includes the term $c^I \frac{T_1 T_2 T_3}{T_I}$ which has no analogue in IIA and this model thus can only be considered in M-theory.

6.2 A possible model in type IIB

The standard ingredients in the superpotential of type IIB theory include contributions from F-flux, H-flux and Q-flux [55–57]. These are the standard contributions that are usually considered. However, already in [55] it was conjectured that so called P-fluxes might also be present and in [56] it was shown that such terms can arise in the superpotential naturally as components of *gauged* supergravity in four dimensions.
Keeping again only even terms in the moduli the superpotential for this case reads

$$\begin{aligned} W_5 = \;& a_0 + a^I \frac{U_1 U_2 U_3}{U_I} + S\left(b^I U_I + b_3 U_1 U_2 U_3\right) \\ & + T_K\left(C^{IK} U_I - c^K U_1 U_2 U_3\right) - S T_K\left(d^K - D^{IK}\frac{U_1 U_2 U_3}{U_I}\right) . \end{aligned} \qquad (6.10)$$

Conveniently, we find that we do not need the term proportional to D^{IK} for our purposes of moduli stabilization. Note that we do not have included *non-perturbative* contributions in this potential. It turns out that including the P-fluxes suffices in order to achieve a de Sitter vacuum through the mass production procedure. In this case we choose to solve for a_0, a^I, b_3 and the c^K and are able to find a stable de Sitter minimum after the procedure.

7 Summary and Conclusion

Building semi-realistic models of de Sitter vacua motivated by string theory has been an ongoing endeavour since the discovery of Dark Energy. Still today, the most promising example of such a construction is the so called KKLT scenario [4]. Here, based on [1–3, 35], we expand on several aspects of this mechanism.

The KKLT scenario is based in type IIB theory, using anti-D3-branes. In type IIA theory no analogous uplifting procedure was known until the description of Dp-branes in terms of nil-potent multiplets was investigated in [18]. There the authors also proposed that it should be possible to use their formalism for uplifting using branes different from anti-D3. In [1] we showed that this is indeed the case for type IIA theory and gave explicit examples. The anti-D6 brane takes the role of the anti-D3 brane of KKLT and supplies a positive energy contribution to the minimum of the potential, lifting it from a stable anti-de Sitter to de Sitter. In further analogy to the KKLT scenario non-perturbative corrections need to be included in order to make this procedure work. Importantly, it was found that the slightly unconventional U-fluxes are needed in order to achieve stability in a simple model. The availability of an uplifting procedure in type IIA opens up many new possibilities for model building in this theory and will be helpful in further endeavours of dS constructions.

Another problem in constructions of de Sitter vacua that utilize an uplift is that there is a priori no predictive power that the mechanism will work out. At no point is it clear that an uplift from a stable anti-de Sitter vacuum will lead to stable de Sitter with all masses positive. In [2, 35] it was shown that, if one starts by building a Minkowski vacuum first, it is guaranteed that the resulting de Sitter vacuum is stable if one follows a three step program. First, solve the Minkowski equations $\partial_\Phi W = 0 = W$. This yields a Minkowski vacuum where no tachyons will appear. The free parameters are adjusted such that there are also no flat directions. Secondly one introduces a *parametrically small* shift in the superpotential: $W \to W + \Delta W$, after which one finds a stable anti-de Sitter minimum. It can be shown that this is always the case. In the third and final step we use and anti-Dp brane in order to lift this minimum to de Sitter which, again, is guaranteed to be stable. This *mass production procedure* is useful not only for its predictive power but also because the Minkowski conditions are usually relatively easy to solve. In more conventional constructions fine tuning or numerical optimization is an integral part of the procedure in order to even arrive at a suitable minimum. In this regard the procedure proposed here is a huge step forward. However, whether this

method is applicable to more general and realistic set ups remains to be seen.
An interesting observation that was made in [2] regards the possibility of large downshifts and uplifts. While the idea was not fully explored in that work it was further investigated in [58]. There this concept was applied in the KKLT scenario and pushed to its extreme. The author found that it is not necessary to have a stable anti-de Sitter progenitor for the Minkowski vacuum and the uplifting procedure still works out. This could potentially be important since it allows for stronger vacuum stabilization. During the early universe the energy scales can be too large for usual models and the vacuum could de-stabilize.
Finally, based on [3] more general set ups were investigated, originating ultimately from M-theory. While most of the models do have an analogue in one of the string theory compactifications the models can be understand most easily coming directly from M-theory. We have shown that it might be worthwhile to investigate M-theory compactifications on its own due to the way the additional tree level terms help in building de Sitter vacua without the inclusion of all non-perturbative corrections. While M-theory compactifications are certainly less well understood and investigated than their string theory counterparts and there are many open questions it is also true that non-perturbative corrections, for example from instantons, lack a complete description in string theory. For this reason our M-theory models could be an alternative that are worthy for further investigation.

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