A 10d view on the KKLT AdS vacuum and uplifting

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Abstract

We continue investigating the ten-dimensional consistency of the KKLT model. We rely on the recent insight that a quartic gaugino term was missing from the ten-dimensional action. The incorporation of this term gives sensible results for the supersymmetric AdS vacuum, expressing the four-dimensional curvature in terms of quantities that appear in the four-dimensional $\mathcal{N} = 1$ supergravity description. Concerning the uplift to de Sitter using anti-branes, our results point to flattening effects and within the assumptions made we find strong evidence for a nogo against de Sitter. We explain why our results differ from other recent works in the literature that come to different conclusions.
1 Introduction

Moduli stabilisation in string theory has a long history and its understanding is crucial for the study of top-down phenomenology. One of the main breakthroughs came in 2003 with the KKLT model [1] and in 2005 with the Large Volume Scenario [2]. Both methods to stabilise moduli rely on computing leading-order corrections to ten-dimensional supergravity solutions with orientifold sources and three-form fluxes in type IIB supergravity [3–5]. Despite the long history of this field there has never been any genuine top-down understanding of these mechanisms to stabilise moduli. The arguments always involved a mixture between top-down and bottom-up viewpoints. This is not a problem per se, but it can lead to a false sense of freedom to tune parameters in the bottom-up effective field theory. It is therefore desirable to find at least one concrete top-down description of moduli-stabilisation.

Recently the field of moduli stabilisation has regained attention due to a growing suspicion that string theory may fail to accommodate any de Sitter (dS) vacuum [6–8]. This suspicion is mainly based on three things. Firstly on the lack of explicit dS examples string constructions such as classical type II supergravity and tree-level heterotic string theory [9–12] (see [7] for a review). Secondly, on non-trivial generalisations of the Dine-Seiberg argument [13] (see also [14]). And thirdly, on the old debate concerning the quantum breakdown of dS space (due to backreacting particle creation) [16–22].

This revival of interest has provided a healthy spark of ideas to construct new possible dS vacua (or more general accelerating cosmologies) [27–32] or revisit existing ones [33–36]. In this paper we direct our attention to the consistency of the KKLT model; both the construction of the AdS vacuum and its potential uplift to a meta-stable dS vacuum. In particular the uplift procedure has been heavily criticised on two fronts. First the stability of the anti-branes themselves came under intense scrutiny in [37–39]. It is our opinion that the arguments pointing towards the instability of anti-branes have been to a large extent addressed and refuted in [40–43] (see however [44, 45]). The second concern, raised in [33] (see also [46]), has to do with the interaction between the open string degrees of freedom of the anti-brane and closed string degrees of freedom of the background. This interaction is only possible in a compact model where the two effects cannot be separated indefinitely.

The authors of [33] suggested that it should be possible to study the interaction using classical ten-dimensional supergravity including the energy-momentum of a gaugino condensate on D7-branes that live in the UV. The main observation of [33] was that the ten-dimensional energy-momentum tensor behaves such that a dS vacuum is not attainable with a single gaugino condensate. From a four-dimensional viewpoint this effect can be understood as a flattening of the potential caused by a backreaction of the supersymmetry-breaking ingredients on the volume modulus.

Later, it was shown in [34] that there is a loophole in the argument of [33] due to a subtlety in the computation. Unfortunately the issue could not be settled in [34] since some of the terms in the energy-momentum tensor were divergent and a regularisation is needed that could influence the fate of the flattening effects. Recently Hamada et al [48] (see also [49]) proposed how to perform this regularisation by adding a quartic fermion term to the action. In this note we redo the computations in [33,34] including this new term. We find that indeed the four-dimensional curvature is expressed entirely in terms of regular background fluxes and

1See [15] for early suggestions of a refined form of the dS swampland bounds.
2See [23–26] for some critical remarks on these ideas.
3Other somewhat orthogonal worries about the consistency of uplifting were expressed in [36,47].
the gaugino condensate, consistent with the effective four-dimensional description of the AdS vacuum [1]. We consider this a non-trivial consistency check of the approach.

The computation can rather straightforwardly be extended to include anti-branes and we find that the flattening effects suggested by [33] persist, within the approximations made. Since the approximations gave reasonable results for the supersymmetric vacuum this can be seen as non-trivial evidence for flattening effects. We do believe that it is important to go beyond the approximations made in order to settle the discussion conclusively. There should be an effective field theory description of these flattening effects as suggested in [33] and discussed further in [50–52]. But in this paper we will not address this.

Simultaneously with this work two other works appeared [53,54] about the same problem and one of these works [54] comes to a very different conclusion than us. We comment in section 5 on the relation between our work and these two other papers.

2 The framework

The KKLT proposal [1] to stabilise moduli in a supersymmetric AdS vacuum rests on reasonable assumptions about quantum corrections to the classical flux compactifications of [3, 5]. These corrections have been argued for using four-dimensional effective field theories described within $\mathcal{N} = 1$ supergravity. The quantum correction used by KKLT is the leading non-perturbative quantum effect to the superpotential $W$. Perturbative corrections to $W$ are absent because of non-renormalisation theorems whereas it is argued that (non-)perturbative corrections in the Kähler potential $K$ can sometimes be self-consistently ignored.\(^4\) The non-perturbative quantum correction arises as a result of gauginos living on a stack of D7-branes condensing in a confining vacuum of the four-dimensional gauge theory. In what follows we assume, like KKLT, that the underlying Calabi-Yau manifold only admits a single Kähler deformation. It then follows that there is only one (holomorphic) 4-cycle wrapped by the D7-branes.

At first sight the gaugino condensation obscures a dimensional oxidation of the KKLT model. However, it has been suggested in [33] that it should nonetheless be possible using ten-dimensional supergravity with D7 probe actions included. The way the gaugino condensation couples to gravity and other closed strings could then be understood simply by keeping explicit non-zero gaugino bilinears of the fermionic part of the D7 action and compute its contribution to the ten-dimensional energy-momentum tensor [55–59].

Like [33] we compute the trace over the four-dimensional part of the (trace-reversed) Einstein equation. We will perform this computation in detail in next section but here we will discuss the form of the resulting expression,

$$R_4 = \int_6 \sqrt{g_6} \mathcal{E}[F^2, \lambda^2],$$  \hspace{1cm} (2.1)

where $R_4$ is the curvature scalar of the four large dimensions, $\int \sqrt{g_6}$ integration over the compact dimensions and $\mathcal{E}[F^2, \lambda^2]$ represents some function of the background fluxes ($F^2$) and the gaugino bilinear ($\lambda^2$). If the approach is sensible then the computation of $R_4$ should reproduce the curvature of the AdS solution of KKLT. We will find that it does so up to numerical factors and the possible reasons behind this mismatch are discussed in section 5.

\(^4\)The Large Volume Scenario [2] uses perturbative corrections to the Kähler potential to stabilise the volume modulus at exponentially large value in a non-supersymmetric vacuum.
Adding anti-branes to this computation is straightforward as explained in [33, 34]. One can show it simply adds the following term to the right hand side of equation (2.1):

\[ R_4 = \int \sqrt{g_6} \left( \mathcal{E}[F^2, \lambda^2] - 2nT\delta_6 \right), \]  

(2.2)

where \( T \) is the warped tension of a stack of anti-D3 branes and \( n \) the number of anti-D3’s. This seems a contradiction since this term contributes negatively compared to the rest, so how can it uplift? In fact [33] realised that, in the right approximation scheme, the anti-brane term is completely subdominant compared to the bulk piece \( \mathcal{E} \) and should be ignored.

This naive “paradox” arises because (2.2) computes an on-shell four-dimensional potential instead of a full off-shell expression\(^5\). The way uplifting should work from this viewpoint is rather different. Uplifting works by a shift of the volume modulus \( \rho \) induced by the supersymmetry-breaking, altering the value of \( \mathcal{E} \) and potentially flipping its sign. To make that point very explicit let us for simplicity assume that the term \( \mathcal{E} \) has two pieces

\[ \mathcal{E}[F^2, \lambda^2; \rho] = \mathcal{E}_+[F^2, \lambda^2; \rho] + \mathcal{E}_- [F^2, \lambda^2; \rho], \]  

(2.3)

where \( \mathcal{E}_- \) is negative definite and \( \mathcal{E}_+ \) positive definite and we introduced the symbolic dependence on \( \rho \). In the supersymmetric vacuum we obviously have \( \mathcal{E}_+[\rho] < -\mathcal{E}_-[\rho] \). But when \( \rho \) shifts to \( \rho + \delta\rho \) due to supersymmetry breaking we expect that the value of \( R_4 \) and therefore \( \mathcal{E}[\rho + \delta\rho] \) is less negative. A dS is reached by making sure that \( \mathcal{E}_+[\rho + \delta\rho] > -\mathcal{E}_-[\rho + \delta\rho] \).

Apart from the \( \rho \)-dependence the only quantities appearing in \( \mathcal{E} \) are the background fluxes and gaugino vev. We can assume that their values have not altered when breaking supersymmetry. This is a crucial assumption that is at the core of the KKLT model. In other words the fluxes \( F^2 \) are the ones of the GKP solution in the background. One could worry about this assumption since there are two sources for deviations from the GKP solution: 1) near the anti-D3 branes non-ISD fluxes are sourced whereas the original GKP construction only allowed for ISD fluxes [5]. But as we emphasised earlier this is a local backreaction effect, and any extra terms in the Einstein equation that are otherwise missed in the approximation will only contribute negatively. 2) Near the 7-branes also non-ISD fluxes are sourced. Even worse, these fluxes are divergent, similar to what was believed to be the case for anti-D3 branes [37,39] but later shown not too happen because of brane polarisation [41,42]. In fact the discussion of anti-D3 flux singularities parallels that of D7 gaugino-induced flux singularities. Whereas brane polarisation cures the singularities for anti-D3 branes, it was shown in [48] that quartic gaugino-terms, previously ignored, do the same for D7-branes.

The essence of this paper is the computation of the right hand side of equation (2.2) incorporating the quartic gaugino term of [48]. The computation is outlined in the next section. The outcome of it is easy to understand. We find that all local and singular backreaction effects of the D7-branes drop out of \( \mathcal{E} \) such that only background fluxes are left over, together with the gaugino vev that determine the curvature. This curvature is of the KKLT size up to a numerical factor. More importantly we find the term \( \mathcal{E}_+ \) to vanish identically if the background fluxes are taken as in the GKP solution. This would eliminate all hopes that anti-brane uplifting can generate dS vacua if the approximations are fully justified. In any case, the interpretation of this effect is a flattening of the potential caused by moduli that are not parametrically heavy with respect to the volume modulus. These moduli therefore shift during the uplift and the computation here seems to imply the effect is potentially drastic.

\(^5\)See [33,34] for a more elaborate explanation.
enough to not allow dS. The existence of meta-stable non-supersymmetric AdS is not affected.

We stress that this result depends on approximations laid out in the next sections and we keep
an open mind to the fact that going beyond these approximations could potentially change
the result. However it is rather striking that within the approximations the “no-dS” result
comes out straightforwardly.

3 The four-dimensional cosmological constant

In this section we utilise the approach of [5] to compute the four-dimensional curvature
when compactifying ten-dimensional type IIB supergravity with fluxes and D7-brane sources
with non-vanishing gaugino bilinear. In doing so, we assume a maximally symmetric four-
dimensional spacetime. The ten-dimensional metric reads

\[ ds_{10}^2 = e^{2A} ds_4^2 + e^{-2A} ds_6^2 , \]  

(3.1)

where \( ds_4^2 \) is a maximally symmetric on a four-dimensional spacetime, with a constant curva-
ture scalar \( R_4 \). The warp factor \( A \) only depends on the coordinates of the transverse manifold
\( M_6 \) equipped with the metric \( e^{-2A} ds_6^2 \). Maximal symmetry demands that the ten-dimensional
axion-dilaton \( \tau \) and 3-form field strength \( G_3 \) only have legs and dependence along \( M_6 \). The
type IIB 5-form \( F_5 \) is self-dual and given by

\[ F_5 = (1 + \ast_10) d C_4 , \quad C_4 = \alpha \text{vol}_4 , \]  

(3.2)

where \( \alpha \) is a function on \( M_6 \). The inclusion of three-form fluxes induces three-brane charge
that contributes to the tadpole which must be cancelled by the inclusion of local sources such
as O3/O7 planes. Taking these ingredients into account, [5] found a consistency condition
demanded by the trace-reversed Einstein equation that takes the form\(^6\)

\[ \triangle_6 \Phi^- = R_4 + \frac{e^{2A}}{\text{Im} \tau} |G_3^-|^2 + e^{-6A} |d \Phi^-|^2 , \]  

(3.3)

where

\[ \Phi^\pm = e^{4A} \pm \alpha , \quad G_3^\pm = \frac{1}{2} (\ast_6 \pm i) G_3 . \]  

(3.4)

This equation can also be derived directly from the type IIB action (together with O3/O7
source terms) upon varying with respect to the warp factor \( A \) (see appendix B). The power
of this equation becomes apparent when one integrates (3.3) over \( M_6 \). The left-hand-side
integrates to zero which implies that the four-dimensional curvature scalar \( R_4 \) is non-positive.
The GKP vacua are the ones with \( R_4 = G_3^- = \Phi^- = 0 \).

We now derive the equation corresponding to (3.3) when \( N > 1 \) D7-branes are included
with non-vanishing fermion bilinear. At this stage we will be agnostic about the value of the
fermion-bilinear, and assume that it will take some non-vanishing value at low energies where
the gauge group condenses. In the procedure outlined we combine the ten-dimensional type
IIB action with the effective action of D7-branes. The bosonic D7-brane action

\[ S_{D7} = -2 \pi N \int_{\Sigma_8} d^8 x (\text{Im} \tau)^{-1} \sqrt{|P[g]|} + 2 \pi N \int_{\Sigma_8} P[C_8] , \]  

(3.5)

\(^6\)Here \( \triangle_6 \Phi^- = - \ast_6 d \ast_6 d \Phi^- \) and all inner products of forms are ten-dimensional, i.e. they include warp
factors.
does not contribute to (3.3) since “it is BPS” with respect to the background.\footnote{BPS is used in a loose form since fluxes can already break supersymmetry of the GKP background.} One can verify this claim by using the ansatz (3.1) and noticing that the warp factor drops out. Even though the bosonic action does not play an important role in our discussion, the fermion terms do. Indeed the gauge theory living on the D7-brane world volume descends to an $\mathcal{N} = 1$ gauge theory in four dimensions with various matter couplings. In particular the eight-dimensional world-volume fermions give rise to the $\mathcal{N} = 1$ gaugini in four dimensions.\footnote{We refer to [33,34] for an explicit map between the eight-dimensional fermions and the gaugino.} The fermionic D7-brane action contains an interaction term between the gaugino bilinear $\lambda \lambda \equiv \text{tr}(\lambda^\alpha \lambda^\alpha)$ and the three-form $G_3$ [60]

$$ S_{\text{D7}^{\text{ferm}}} = \pi \int \star_{10} (G_3 \cdot I + \overline{G}_3 \cdot \bar{I}) , \quad I = \frac{e^{-4A}}{\sqrt{\text{Im} \tau}} \frac{\lambda \lambda}{16\pi^2} \Omega \delta(\Sigma_8) . \quad (3.6) $$

The contribution of the fermionic terms to (3.3) turns out to be singular due to the backreaction of the fermion bilinears on the three-form fluxes \cite{33,34}

$$ G_3 \sim (\text{Im} \tau) \bar{I} , \quad (3.7) $$

where $I$ is the source appearing in (3.6) and carries an explicit delta function. In fact, even the on-shell action is UV divergent due to this singular backreaction. In \cite{48} it was suggested that the singular backreaction is ultimately a result of omitting a four-fermion term in the D7-brane action. Currently the four-fermion terms are not known for the D-brane action and so \cite{48} suggested to fix their form such that the singularities would be cancelled. The suggestion of \cite{48} was inspired by a similar problem appearing in Hořava-Witten theory \cite{61–63} which was resolved by a four-fermion term that appears together with the M-theory 4-form in a perfect square. This was in direct analogy with a similar square structure that appears in heterotic supergravity \cite{49,64}. It is therefore reasonable to assume that the fermionic D7-brane coupling (3.6) together with the kinetic terms for $G_3$ should be combined with the four-fermion terms to form a perfect square \cite{48,49}. Indeed this goes a long way to prevent singular on-shell action as a result of backreaction the fermion bilinear.

The perfect square replaces the original action for $G_3$ and the fermionic coupling by

$$ S_3 = -\pi \int \star_{10} \left| \frac{G_3 - (\text{Im} \tau) \mathcal{P}(\bar{I})}{(\text{Im} \tau)} \right|^2 , \quad (3.8) $$

The projector $\mathcal{P}$ is defined to eliminate the coexact piece of the form it acts on. It does not affect the $G_3$ equation of motion but is the final piece of the puzzle to ensure that the on-shell action is regular \cite{48}. Upon expanding the square the action consists of three terms:

$$ S_3 = -\pi \int \star_{10} \frac{|G_3|^2}{(\text{Im} \tau)} + \pi \int \star_{10} \left( G_3 \cdot \mathcal{P}(I) + \overline{G}_3 \cdot \mathcal{P}(\bar{I}) \right) - \pi \int \star_{10}(\text{Im} \tau)|\mathcal{P}(I)|^2 . \quad (3.9) $$

The first term is the standard bulk term for $G_3$, the second term is the fermionic action of the D7-brane considered before in \cite{33,34,57}, with the exception that the three-form $I$ is now projected to the set of closed forms. The last term is the new quartic fermion term. We are now ready to re-derive the equation (3.3) with the D7-brane fermion contribution. The explicit computation is carried out in appendix B and the result is

$$ \Delta_6 \Phi^- = R_4 + e^{-6A}|d\Phi^-|^2 + e^{2A} \left( \frac{|G_3|^2}{\text{Im} \tau} - \frac{3}{4} \text{Re} \left( G_3 \cdot \mathcal{P}(I) \right) + \frac{\text{Im} \tau}{4} |\mathcal{P}(I)|^2 \right) . \quad (3.10) $$
This expression is exact and despite its appearance is completely regular.\footnote{Equation (3.10) does have delta-function sources on the right-hand-side but is regular upon integration. This should be compared to the situation without quartic fermion terms where even the integrated expression was singular due to $\delta^2$ terms appearing.} In order to demonstrate this and further analyse equation (3.10) we must now make simplifying approximations.

First we rewrite the delta function appearing in the source term $I$ in terms of a Green’s function. Schematically we write

$$
\frac{\partial^2 G}{\partial z \partial \bar{z}} = \delta(\Sigma_8) - \frac{1}{V_2},
$$

(3.11)

where $z$ and $\bar{z}$ are complex coordinates on the transverse two-cycle and $V_2$, a volume factor of that two-cycle, is introduced to make this equation consistent. We see that the 3-form $I \sim \delta(\Sigma_8) \Omega$ is not closed and so the role of the projector $\mathcal{P}$ in this case is to add terms of the form $(\partial_z \partial_{\bar{z}} G) \Omega$ such that the resulting expression is closed. Explicitly we have \footnote{One could wonder whether in this case the orientifold projection does not eliminate the backreaction of the gaugino bilinears on the three-forms. We do not address this question here, but argue that our setup is sufficiently close to the Sen limit such that the gradients of $\tau$ do not enter without eliminating the three-form backreaction altogether.}

$$
\mathcal{P}(\bar{I}) = e^{-4A} \frac{\lambda \lambda}{\sqrt{\text{Im } \tau}} \left( \frac{1}{16\pi^2} \left( \frac{\partial G}{\partial \bar{z}} \wedge \bar{\Omega}_2 + \frac{1}{V_2} \bar{\Omega} \right) \right),
$$

(3.12)

where $\Omega_2$ is the holomorphic two-form on the holomorphic four-cycle. Notice that we have made a crucial assumption here which we will continue making in the following. Namely we have assumed that close to the seven-branes where most of our analysis takes place, we can treat the warp factor as constant. Furthermore we assume the so-called Sen limit \footnote{Under these assumptions, the Chern-Simons action does not contribute to the $G_3$ equations of motion. The details of the fate of the Chern-Simons action in our set-up are discussed in appendix C. Note that since $I^{\text{reg}}$ is harmonic with these assumptions, it can always be combined with an arbitrary harmonic contribution that can be added to a solution to the equations of motion to produce another solution to the equations of motion and we denote the entire harmonic contribution as $G_3^0$.} where the stack of 7-branes consists of an $O7$ with 4 parallel $D7$-branes on top, in this case the 7-brane stack does not source a gradient for $\tau$ and so it can be taken to be a constant.\footnote{Under these assumptions, the Chern-Simons action does not contribute to the $G_3$ equations of motion. The details of the fate of the Chern-Simons action in our set-up are discussed in appendix C. Note that since $I^{\text{reg}}$ is harmonic with these assumptions, it can always be combined with an arbitrary harmonic contribution that can be added to a solution to the equations of motion to produce another solution to the equations of motion and we denote the entire harmonic contribution as $G_3^0$.}

These assumptions allow us to determine the backreacted three-forms in the same spirit as \footnote{Under these assumptions, the Chern-Simons action does not contribute to the $G_3$ equations of motion. The details of the fate of the Chern-Simons action in our set-up are discussed in appendix C. Note that since $I^{\text{reg}}$ is harmonic with these assumptions, it can always be combined with an arbitrary harmonic contribution that can be added to a solution to the equations of motion to produce another solution to the equations of motion and we denote the entire harmonic contribution as $G_3^0$.} For notational simplicity we write

$$
\mathcal{P}(I) = I^{\text{sing}} + I^{\text{reg}},
$$

(3.13)

where $I^{\text{sing}}$ is the exact part of $\mathcal{P}(I)$ in (3.12) which is also singular, and $I^{\text{reg}}$ is the harmonic part which is regular. Finally in order to solve for $G_3$ we assume that the background close to the seven-branes is only slightly perturbed from the standard GKP background for which $\Phi^- = e^{4A} - \alpha = 0$. By our previous assumption that $e^{4A}$ is constant, this implies that the Chern-Simons terms in the action does not enter in the equations of motion for $G_3$. With these approximations, the solution to the equation of motion for $G_3$ derived from action eq. (3.8) takes a simple form\footnote{Under these assumptions, the Chern-Simons action does not contribute to the $G_3$ equations of motion. The details of the fate of the Chern-Simons action in our set-up are discussed in appendix C. Note that since $I^{\text{reg}}$ is harmonic with these assumptions, it can always be combined with an arbitrary harmonic contribution that can be added to a solution to the equations of motion to produce another solution to the equations of motion and we denote the entire harmonic contribution as $G_3^0$.}

$$
G_3 = (\text{Im } \tau) I^{\text{sing}} + G_3^0,
$$

(3.14)

where $G_3^0$ is harmonic and quantised. We can now evaluate the expression (3.10) which constrains the four-dimensional cosmological constant. Substituting the solution of the $G_3$
The equation only contains the background flux $G_3^0$ and the fermion condensate through the regular part $I_{\text{reg}}$. We notice that all singular terms have been eliminated from the expression as a result of our approximations. We expect that more generally the expression (3.10) will be regular also when the approximations are relaxed.

4 Matching to 4D EFT

We now compare our results with the four-dimensional effective theory considered by KKLT [1]. This theory is specified in terms of a Kähler potential $K$ and superpotential $W$, given by

$$K = -3 \log(2 \Im \rho) , \quad W = W_0 + A \exp(2\pi i \rho/N) ,$$

where $\rho$ is the complex volume modulus, related to the volume of the internal manifold by $\Im \rho \sim V^{2/3}$ and $W_0$ is the constant Gukov-Vafa-Witten superpotential (after integrating out the complex-structure moduli). In order to compare our ten-dimensional result (3.15) with the effective theory (4.1) we translate two results from the four-dimensional effective field theory back to ten dimensions. In particular we view the fluxes $G_3^0$ as the ISD ones, before the inclusion of the non-perturbative effects. Therefore $G_3^0$ vanishes and $G_3^0 \cdot \Omega$ can be related to $W_0$

$$\int \star_6( G_3^0 \cdot \Omega ) = i \int G_3^0 \wedge \Omega = iW_0 .$$

Secondly in the supersymmetric vacuum the four-dimensional F-term equation $D_\rho W = \partial_\rho W + W \partial_\rho K = 0$ implies that $\Re ( G_3^0 \cdot I_{\text{reg}} ) = 0$. This can be observed by relating $\lambda \lambda$ to the non-perturbative effects appearing in the four-dimensional superpotential (4.1). We use the standard gauge theory expression for the fermion bilinear in the confining vacuum

$$\langle \lambda \lambda \rangle = 16\pi i \, \partial_\rho W = -\frac{32\pi^2}{N} Ae^{2\pi \rho/N} ,$$

where, like KKLT [1] we used that the complex volume modulus $\rho$ equals the complex gauge coupling of the four-dimensional gauge theory. We note that $A$ is determined by the dynamically generated scale of the IR theory and is left unspecified. Using this map between the four-dimensional theory and the ten-dimensional quantities it is easy to obtain the desired result

$$\int \star_6 \Re ( G_3^0 \cdot I_{\text{reg}} ) = \frac{e^{-4A} \Re (iW_0 \lambda \lambda )}{16\pi^2 V_2 \sqrt{\Im \tau}} = \frac{-2e^{-4A}}{3NV_2 \sqrt{\Im \tau}} \Im \left( A e^{2\pi \rho/N} |2(3 + \frac{4\pi}{N} \Im \rho)|^2 \right) = 0 ,$$

where $W_0$ was replaced using the four-dimensional F-term equation. We expect that this will hold more generally using only the supersymmetry conditions in ten dimensions. We hope to derive this purely from ten dimensions in the near future.
Using these results our ten-dimensional equation (3.15) reduces to
\[ \int \star_6 \left( R_4 + \frac{e^{2A}}{4} (\text{Im} \, \tau) |I_{\text{reg}}|^2 \right) = 0. \] (4.5)

More explicitly, we can insert the expression for \( I_{\text{reg}} \) in (3.12) and the gauge theory expression for the fermion bilinear to find
\[ R_4 = -\frac{2}{V_2^7} \left| \frac{\lambda \lambda}{16\pi^2} \right|^2 = -\frac{8}{(\text{Im} \, \rho)N^2} |A|^2 e^{-4\pi \text{Im} \rho/N}. \] (4.6)

This final expression is everywhere negative definite. Any terms that could have given a positive contribution have disappeared. Of course, a supersymmetric vacuum should be either AdS or Minkowski, so \( R_4 \) should indeed be negative or zero. Note that we reproduce almost exactly the four-dimensional potential evaluated in the vacuum,
\[ V_{\text{KKLT}} = -3e^K |W|^2 = -\frac{2\pi^2}{3(\text{Im} \, \rho)N^2} |A|^2 e^{-4\pi \text{Im} \rho/N}. \] (4.7)

The reason for the numerical mismatch could be an artefact of the approximations made.\(^\text{12}\) The most important approximation we made is that we have put the regular harmonic piece of the 3-form flux equal to the classical GKP solution. This approximation should be very good nearly everywhere in the manifold, but most likely not close to the 7-branes. Hence one could speculate that possible extra term then arise that make the match to the AdS vacuum energy computed in KKLT exact. We leave this for future research.

Let us recall that the anti-D3 brane tension contributes only negatively to \( R_4 \). An uplifting to dS is supposed to occur through the mediation of gaugino terms that contribute positively to \( R_4 \). Such positive terms can only arise from the term \( \text{Re} \left( G_3^0 \cdot I_{\text{reg}} \right) \) being non-zero in a supersymmetry breaking vacuum.

5 Discussion

5.1 Summary and discussion of results

In this paper, we have revised the computations of [34] taking into account the flux renormalisation mechanism provided by the quartic gaugino term on D7-branes, as proposed in [14]. We must rely on a number of assumptions necessary to directly use the techniques of [14], namely constant axion-dilaton, warping, and \( C_4 \)-potential, and most importantly we used that the harmonic 3-form flux is the GKP flux. Together with four-dimensional supersymmetry constraints we have shown that there are no positive terms in \( R_4 \) consistent with the statements in [33, 34]. In the notation of equation (2.3) we found that \( \mathcal{E}_+ = 0 \). It is quite striking that a no-dS result comes out that manifestly in our approach. Of all approximations made the crucial one leading to this result is the approximation made about the harmonic three-form flux.

A worthwhile direction for future research would be to see if one can still obtain a definite result after relaxing some of these approximations. It is tempting to speculate what happens.

\(^{12}\) Another possible source for the mismatch is a subtle change of conventions between the compactified gauge theory and the four-dimensional literature.
For that we go back to equation (3.10). Note that this equation was derived without any of the simplifying approximations and is hence valid in general circumstances. We have argued that the renormalisation due to the four-fermion term is such that the only 3-form fluxes appearing in (3.10) are the background fluxes, i.e., those of the classical GKP solution. This means we can put $|G^0_{3,-}|^2 = 0$ since the flux appearing will be ISD. This is expected to be broken by the gaugino condensate but the deviations from the ISD background are mostly “self-energy” that has been removed by the renormalisation procedure. Once anti-branes are added they will also create some local non-ISD fluxes down the throat. But we will neglect these since they come anyways with a negative contribution to $R_4$ and so will not help in getting de Sitter. Secondly it is tempting to suggest that

$$\int \star_6 \text{Re}(G_3 \cdot \mathcal{P}(I)) = 0. \quad (5.1)$$

This is a non-trivial statement and, within our assumptions turned out to be a consequence of the supersymmetry of the AdS solution. The way one could check whether this is true in general, without making assumptions and without appealing to the four-dimensional supergravity is by analysing the ten-dimensional supersymmetry variations in presence of the gaugino bilinear. This is a very interesting exercise which we plan to carry out in the future. If it proves (5.1) correct then the general result for the four-dimensional curvature is simply

$$R_4 = -\frac{1}{V} \int \star_6 \frac{e^{2A}}{\text{Im} \tau} |G^0_{3,-}|^2 - \frac{1}{V} \int \star_6 \frac{e^{2A} \text{Im} \tau}{4} |\mathcal{P}(I)|^2, \quad (5.2)$$

which is manifestly negative. Anti-branes only add extra negative terms.

If this turns out to happen then the only angle to achieve dS is if some of the background fluxes and gaugino phases change when supersymmetry is broken. This goes against the very assumptions about the validity of the four-dimensional effective field theory, which assumes that the complex structure moduli (related to the fluxes) and the axion (related to the gaugino phase) do not shift positions.

Finally we wish to iterate the statements in [33,34]: to have certainty that uplifting AdS vacua with small supersymmetry-breaking ingredients does not lead to runaways instead of meta-stable dS vacua, one should achieve the following parametric scaling:

$$m^2 L^2 \gg 1. \quad (5.3)$$

In this equation $m^2$ is the mass squared of the lightest modulus in the AdS vacuum and $L$ the AdS length. Interestingly not a single clear top-down AdS vacuum in string theory achieves this. The only cases known to us are “stringy-inspired” like racetrack fine-tuning [66] or non-geometric fluxes [67,68]. It was conjectured in [35] that racetrack fine-tuning is in the Swampland\textsuperscript{13} whereas reference [34] conjectured this to be true for all vacua obeying (5.3). Since the conjecture of [34] forbidding (5.3) is a statement about AdS vacua and not dS vacua we believe it should be easier to verify than the actual no-dS conjecture of [6–8], and it would count as non-trivial evidence in support of the no-dS conjecture.

### 5.2 Comparison with recent papers

Simultaneous with the appearance of this work, two other papers appeared discussing the same problem [53,54]. Reference [54] set out to do similar computations as done here but

\textsuperscript{13}See [26,69] for some criticism on this.
came to a different conclusion. We differ from [54] in the way the energy-momentum tensor in ten dimensions is computed. This is also discussed in reference [53].

Our approach is as follows: we vary the D7 brane action to obtain its energy-momentum tensor. Subsequently we fill in the vev of the gaugino condensate $\langle \lambda \lambda \rangle$, whose value is set by non-perturbative physics. This is how the semi-classical limit is usually defined: one inserts the quantum vevs of operators into the classical EM tensor.

But a subtlety arises in this context which relates to the dependence of the gaugino vev on the gauge coupling (4.3). Since the coupling equals the inverse volume of the 4-cycle wrapped by the D7 branes, a dependence of $\langle \lambda \lambda \rangle$ on the ten-dimensional metric is induced. This dependence cannot be understood classically in four or ten dimensions. One could imagine reversing the order and plug the metric dependence of the gaugino condensate into the ten-dimensional action before computing the energy-momentum tensor. This introduces extra terms that match with the four-dimensional KKLT potential and reference [54] concluded there is no obstruction to uplifting to meta-stable dS vacua. We think that it is not clear how the semi-classical method really works in this case. But a straightforward method that manifestly respects ten-dimensional covariance is to have no metric dependence in $\langle \lambda \lambda \rangle$ at the level of the 10d action.

In a rather different wording reference [53] warned for this subtlety and concludes that perhaps the ten-dimensional approach is not attainable after all. We remain agnostic about that point but emphasize the following: if a ten-dimensional viewpoint is to make sense then we believe that the correct approach is the one we have followed for the following reasons. 1) Inserting the four-dimensional information about the gauge coupling dependence into the ten-dimensional action implies one is not following a ten-dimensional approach after all. 2) A semi-classical limit should be something universal. In other words when one couples classical gravity to quantum matter then the approach is that the quantum vevs are inserted into the classical EM tensor. The latter tensor is a universal object in the theory and how to compute the tensor does not depend on the background neither on the quantum effects of the matter sector. 3) It is not clear to us how ten-dimensional covariance is preserved in the approach of [54].

Finally we want to explain the reader why a ten-dimensional semi-classical approach, as we used it, is bound to give flattening effects. As explained in the appendix of [70], the trace-reversed Einstein equation with indices in four dimensions equals the definition of the on-shell effective potential $V_{\text{eff}}$. In particular the on-shell condition used to get it is $\partial_\sigma V_{\text{eff}} = 0$ with $\sigma$ the internal volume mode. The equation $\partial_\sigma V_{\text{eff}} = 0$ itself can be found from the (trace-reversed) Einstein equation with indices in the internal space. It is then obvious from the ten-dimensional approach that all dependences on the volume will always be of the power-law kind.\textsuperscript{14} Hence the ten-dimensional semi-classical limit cannot feature the sharply peaked exponential dependence of the gaugino-condensate. Potentials that are a sum of powerlaws are certainly more prone to flattening effects. Nonetheless it is not obvious why flattening effects should lead to a definite “no de Sitter” result. But that is exactly what we have found and we consider that the main result of this paper: If the ten-dimensional semi-classical limit is to make sense then flattening effects conspire against an uplift to dS space, whereas uplifts to meta-stable AdS spaces are not obviously forbidden. We emphasize once more that this does not mean that reference [54] reached a wrong conclusion, we rather think that they

\textsuperscript{14}Because dimensional reduction of the effective action gives the off-shell potential $V_{\text{eff}}$ but the same terms appear with different numerical coefficients in the internal Einstein equation, which is interpreted as $\partial_\sigma V_{\text{eff}} = 0$. 

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effectively worked in an approach that was bound to reproduce the existing four-dimensional effective theory of KKLT [1].

Finally note that, if the second part of reference [53] turns out to be correct, then dS uplifts are questionable already for a simpler reason related to having controlled uplifts and throat volumes that fit into the compact CY space.

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A Conventions

The type IIB supergravity action (in units where $2\pi\ell_s = 1$) is

$$S = 2\pi \int \star_{10} \left( R_{10} - \frac{d\tau \cdot d\bar{\tau}}{2(\text{Im } \tau)^2} - \frac{G_3 \cdot \bar{G}_3}{2(\text{Im } \tau)} - \frac{1}{4} |F_5|^2 \right) - \frac{\pi}{2i} \int \frac{C_4 \wedge G_3 \wedge \bar{G}_3}{\text{Im } \tau}, \quad (A.1)$$

where $R_{10}$ is the ten-dimensional Ricci scalar calculated using the metric $g_{MN}$ with determinant $g_{10}$. The axion-dilaton is denoted by $\tau = C_0 + ie^{-\phi}$ and the NSNS and RR 3-form field strengths have been combined into a single complex 3-form

$$G_3 = F_3 - ie^{-\phi} H = dC_2 - \tau dB_2 \ , \quad H = dB_2 \ , \quad F_3 = dC_2 - HC_0 \ . \quad (A.2)$$

Throughout the paper we use short-hand notation to denote form contractions, let $\omega_p$ and $\psi_p$ denote two $p$-forms, then

$$\star_{10} \omega_p \wedge \psi_p = \omega_p \cdot \psi_p \star_{10} 1 \ , \quad \omega_p \cdot \psi_p = \frac{1}{p!} \omega_{M_1M_2...M_p} \psi_{M_1M_2...M_p} \ , \quad |\omega_p|^2 = \bar{\omega}_p \cdot \omega_p \ . \quad (A.3)$$

The complex-conjugation in the last expression is to allow for the possibility that $\omega_p$ is a complex $p$-form.

B Variation of the action

In this appendix we derive eq. (3.10) step-by-step. We also fill in some of the gaps of the subsequent calculations in section 3. First we show that we can obtain equation (3.3) directly by varying the ten-dimensional type IIB supergravity action with respect to the warp factor $A$, as defined in (3.1). After establishing this, we include the D7-branes and redo the same computation.
The ten-dimensional action for the Ansatz (3.1, 3.2) can be found to be:

\[ S_{\text{IIB}} = 2\pi \int d^{10}x \sqrt{-g_4} \sqrt{g_6} e^{-2A} \mathcal{I}, \]  

where

\[ \mathcal{I} = e^{-2A} R_4 + e^{2A} (R_6 + 2\Delta_6 A - 8|dA|^2) - e^{2A} \frac{|d\tau|^2}{2(\text{Im}\tau)^2} - e^{6A} \frac{|G_3|^2}{2\text{Im}\tau} + e^{-6A} \frac{|d\alpha|^2}{2}. \]  

(B.2)

The \(|\cdots|^2_6\) denote metric contractions using purely the six-dimensional metric \(ds_6^2\) without the warp factor. We now compute a variation with respect to \(A\). We find:

\[ 0 = R_4 - \Delta_6 e^{4A} + e^{-6A} |de^{4A}|^2 + e^{2A} \frac{|G_3|^2}{2\text{Im}\tau} + e^{-6A} |d\alpha|^2. \]  

This is simply the trace-reversed ten-dimensional Einstein-equation traced over the four-dimensional indices. Adding the Bianchi identity, we recover equation (3.3):

\[ \Delta_6 \Phi^- = R_4 + e^{2A} \frac{|G_3|^2}{2\text{Im}\tau} + e^{-6A} |d\Phi^-|^2. \]  

We did not include any sources, but including O7/O3 planes does not modify the final result. Now we would now like to add D7-branes to the configuration. As explained in the main text, the inclusion of the D7-brane implies that we should perform the replacement (3.8) and vary that with respect to the warp factor. We have two crucial remarks. The careful reader will notice that \(I\) as defined in (3.6) contains a warp factor dependence of \(e^{-4A}\). This factor should not be varied when deriving the equivalent to (B.3), since it does not arise from a metric in the action. On the other hand, the delta function and the holomorphic three-form do have a warp factor dependence. The holomorphic three-form goes like \(e^{-3A}\) and the delta function like \(e^{2A}\). We can therefore write the perfect square action as

\[ S_3 = -2\pi \int d^{10}x \sqrt{-g_4} \sqrt{g_6} e^{4A} \frac{e^{2A} |G_3 - e^{-A}(\text{Im}\tau)\hat{P}(I)|^2}{2\text{Im}\tau}, \]  

where we have made the warp-factor dependence explicit by momentarily writing \(\mathcal{P}(I) = e^{-A}\hat{P}(I)\). Using this we find

\[ -\frac{e^{4A}}{8\pi \sqrt{-g_4} \sqrt{g_6}} \frac{\delta S_3}{\delta A} = e^{2A} \frac{|G_3 - (\text{Im}\tau)\mathcal{P}(I)|^2}{2\text{Im}\tau} + e^{2A} \frac{3}{4} \text{Re} \left( (G_3 - (\text{Im}\tau)\mathcal{P}(I)) \cdot \mathcal{P}(I) \right) \]  

\[ = e^{2A} \left( \frac{|G_3|^2}{2\text{Im}\tau} - \frac{3}{4} \text{Re} \left( G_3 \cdot \mathcal{P}(I) \right) + \frac{\text{Im}\tau}{4} |\mathcal{P}(I)|^2 \right). \]  

(B.6)

We should replace the \(|G_3|^2\)-term in (B.3) by (B.7). This modifies (B.4) to exactly (3.10) which reads:

\[ \Delta_6 \Phi^- = R_4 + e^{-6A} |d\Phi^-|^2 + e^{2A} \left( \frac{|G_3|^2}{\text{Im}\tau} - \frac{3}{4} \text{Re} \left( G_3 \cdot \mathcal{P}(I) \right) + \frac{\text{Im}\tau}{4} |\mathcal{P}(I)|^2 \right). \]  

If one now writes \(\mathcal{P}(I) = I^{\text{sing}} + I^{\text{reg}}\) as in (3.13) and substitutes the solution to the equation of motion \(G_3 = (\text{Im}\tau)I^{\text{sing}} + G_0^3\) (while making the approximation that \(\Phi^- = 0\) and...
warping, $C_4$-potential and axion-dilaton are constant) into (B.6), one finds together with the other terms in the IIB action and the Bianchi identity the following equation:

\[ 0 = R_4 + \frac{e^{2A}}{2\text{Im} \tau} |G_3^0 - (\text{Im} \tau)I_{\text{reg}}|^2 + \frac{e^{2A}}{4\text{Im} \tau} \text{Re} \left[ (G_3^0 - (\text{Im} \tau)I_{\text{reg}}) \cdot (\text{Im} \tau) \left( I_{\text{sing}} + I_{\text{reg}} \right) \right] \\
+ \frac{e^{2A}}{2\text{Im} \tau} \left( |G_3^{-0}|^2 + 2\text{Re} \left( G_3^{-0} \cdot \bar{I}_{\text{sing},-} \right) - |G_3^{-0}|^2 - 2\text{Re} \left( G_3^{0,+} \cdot \bar{I}_{\text{sing},+} \right) \right) \]

(B.8)

The substitution already took into account that the Bianchi identity does not contribute $\delta^2$ singularities, which is discussed in appendix C. Recall that $I_{\text{sing}}$ is by definition exact, so its inner product with a harmonic form ($G_3^0$ or $I_{\text{reg}}$) is therefore zero when the expression is integrated over (the (A)ISD-parts vanish in the same fashion). In this way, we arrive at

\[ 0 = \int \star 6 \left( R_4 + \frac{e^{2A}}{\text{Im} \tau} |G_3^0 - (\text{Im} \tau)I_{\text{reg}}|^2 - \frac{3e^{2A}}{4} \text{Re} \left( G_3^0 \cdot I_{\text{reg}} \right) + \frac{e^{2A}}{4} (\text{Im} \tau)|I_{\text{reg}}|^2 \right) . \]  

(B.9)

### C (Non)renormalisation of the Chern-Simons action

The effective bulk action of IIB string theory contains a Chern-Simons term,

\[ S_{\text{CS}} = -\frac{\pi}{2i} \int \frac{C_4 \wedge G_3 \wedge \bar{G}_3}{\text{Im} \tau} . \]

(C.1)

This term is topological and does not affect the Einstein equation directly. Still one can be concerned and wonder whether we should renormalise the Chern-Simons action or not. If we should, the renormalisation could introduce metric dependence and affect our results.

Under the assumptions of constant $C_4$ and axion-dilaton made in the main body of the text, we should not renormalise the Chern-Simons term. Not renormalising produces no disasters and thus there is perhaps no obvious need for renormalisation. Let us check this.

At first, it may appear that filling the solution for $G_3$ into the Chern-Simons action will produce a badly divergent action. However, this is not the case. Filling the on-shell value of $G_3$ into the Chern-Simons action yields

\[ S_{\text{CS, on-shell}} = -\frac{\pi g_s}{2i} \int \left\{ C_4 \wedge G_3^0 \wedge \bar{G}_3^0 + (\text{Im} \tau)C_4 \wedge G_3^0 \wedge \bar{I}_{\text{sing}} + (\text{Im} \tau)C_4 \wedge I_{\text{sing}} \wedge \bar{G}_3^0 \right\} \\
- \frac{\pi g_s}{2} \int (|I_{\text{sing}}^-|^2 - |I_{\text{sing}}^+|^2) \left( \text{Im} \tau \right)^2 C_4 \wedge \tilde{x}_6 1, \]

(C.2)

where we introduced the symbols $I_{\text{sing}}^\pm = \frac{1}{2}(\tilde{x}_6 \pm i)I_{\text{sing}}$. The last two terms in this expression seem divergent while all other terms are clearly finite upon evaluation of the integral. Interestingly the two divergent terms exactly cancel against each other and the on-shell action is well-behaved. To see this, note that near a D7-brane we can locally write the internal manifold as $\mathcal{M} = \Sigma_2 \times \Sigma_4$, with $\Sigma_2$ transverse and $\Sigma_4$ parallel to the D7-brane. We then have

\[ I_{\text{sing}} \sim d \left( \frac{\partial G}{\partial z} \right) \wedge \bar{\Omega}_2 = \left( \frac{\partial^2 G}{\partial z \partial \bar{z}} \right) d\bar{z} \wedge \bar{\Omega}_2, \]

(C.3)
with $G$ the Green’s function, which varies transverse to the D7-brane, and $dz$ and $d\bar{z}$ are defined on $\Sigma_2$. It is clear that the first of these terms is $(1, 2)$ while the second is $(0, 3)$. First we notice that the two terms have the same magnitude:

$$ \int_6 * \left| \left( \frac{\partial^2 G}{\partial z \partial \bar{z}} \right) dz \wedge \bar{\Omega}_2 \right|^2 = \int_6 * \left| \left( \frac{\partial^2 G}{\partial z \partial \bar{z}} \right) d\bar{z} \wedge \bar{\Omega}_2 \right|^2 , \tag{C.4} $$

which follows from integration by parts and the reality of $G$. Then we use that the first term in (C.3) corresponds to the IASD part of $I^{\text{sing}}$, while the second term corresponds to the ISD part. It is immediate that the $(0, 3)$ term is ISD as the unique $(0, 3)$ form at our disposal is $\bar{\Omega}$ and this is ISD in our convention. That the $(1, 2)$ form is IASD follows from the fact that it is primitive and primitive $(1, 2)$ forms are IASD in our convention. The primitivity can be seen by decomposing $J$ as

$$ J = J_{\Sigma_2} + J_{\Sigma_4} , \tag{C.5} $$

with the two terms defined in the obvious way. Using this, evaluation yields

$$ \left( \frac{\partial^2 G}{\partial z \partial z} \right) dz \wedge \bar{\Omega}_2 \wedge J = 0 . \tag{C.6} $$

We thus see that $I^{\text{sing}+}$ and $I^{\text{sing}−}$ have the same magnitude. Therefore, the two problematic terms in (C.2) cancel against each other and the on-shell Chern-Simons action is finite without the introduction of any renormalisation counterterms.

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