Quantum log-corrections to swampland conjectures

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Abstract: Taking the anti-de Sitter minimum of KKLT and the large volume scenario at face value, we argue for the existence of logarithmic quantum corrections to AdS swampland conjectures. If these conjectures receive such corrections, it is natural to suspect that they also arise for other swampland conjectures, in particular the dS swampland conjecture. We point out that the proposed log-corrections are in accord with the implications of the recently proposed trans-Planckian censorship conjecture. We also comment on the emergence proposal in the context of both perturbative flux models and the KKLT construction.

Contents

1 Introduction

String theory and in particular its application to low energy physics is experiencing a period of new insights that question older arguments based on the landscape picture. The swampland program [1–3] intends to extract a set of relatively simple quantitative features that low-energy effective field theories should satisfy in order to admit a UV completion to a consistent theory of quantum gravity (see [4] for a recent review). By now several swampland conjectures have been proposed [5–17], which also induced further

new developments like the emergence [18–20] of infinite distances in field space or the appearance of towers of light strings [21–23].

In this paper we focus on two such conjectures dealing with AdS vacua. Concretely, these are the AdS/moduli scale separation conjecture (AM-SSC) [24] and the AdS distance conjecture (ADC) [15] (see also [25]). Roughly, these conjectures state that the mass of certain (towers of) modes cannot be parametrically separated from the AdS radius. The strong version of the ADC is reminiscent of observations made earlier in [26] for the class of flux vacua without negative tension objects. Moreover, we comment on the connection of our observations to the dS swampland conjecture [8, 9, 11, 12, 16]. The evidence for these conjectures derives mostly from the failure of contradicting string theory constructions. However, concerning the no-go for dS vacua in quantum gravity, there have also been alternative arguments based on the concept of quantum breaking [27–30].

The best understood string vacua examples are at tree-level, e.g. flux vacua [31–34], where fairly general results can be proven. In particular, for some tree-level type IIA flux vacua one can prove the dS swampland conjecture [35], while one can more generally exclude dS in regimes of parametric control [36]. This class of flux vacua generically gives AdS vacua, supersymmetric or not, whose cosmological constant satisfies the AM-SSC [37]. However, the supersymmetric solutions were argued to fail the strong version of the ADC [15, 33].

Stringy AdS and dS vacua can also be constructed utilizing not only tree-level ingredients, but also quantum, in particular non-perturbative effects. The most famous examples are the KKLT [38] and the large volume scenario (LVS) [39]. In both these cases, AdS minima are found in the effective 4D potential and subsequently uplifted to dS. We will focus on the AdS minima before the uplift mechanism.

Taken at face value, the KKLT model does not satisfy the two AdS conjectures. Since the tree-level vacua provide strong support for the conjectures, one might think that something is wrong or inconsistent in the quantum vacuum construction. However, it has proven to be difficult to isolate any particular shortcoming in the AdS minimum construction. Its full ten-dimensional description has been analyzed in a series of recent papers [24, 40–46] converging to the conclusion that the 4D effective KKLT description captures the main aspects of this vacuum. The uplift to a dS vacuum is more subtle and new aspects of the validity of the effective field theory in the warped throat have been investigated in [47, 48]. So while the validity of the dS vacua is still an open question, the AdS vacua seem to be true counterexamples to the AdS swampland conjectures.

In this paper we take a different approach to this issue and analyze whether the AdS swampland conjectures should rather be modified in such a way that these well-established quantum AdS vacua do satisfy them.

We will start in section 2 by recalling the AdS and dS swampland conjectures, as well as briefly describing the emergence proposal. In section 3 we briefly summarize the manifestation of the AdS swampland conjectures for the well known tree-level flux compactifications. Here one can distinguish DGKT-like models [32] with a dilute flux limit from Freund-Rubin type compactifications, where a geometric flux becomes relevant for moduli stabilization. Concrete examples of these two types are presented in appendix A.

In section 4, we take a closer look at the AdS vacua of KKLT and LVS. We show that for KKLT the relation between the lightest moduli mass and the cosmological constant receives dominant logarithmic corrections whose origin can be understood from a simple scaling argument. Taking these quantum corrections seriously, we introduce logarithmic corrections to the initial AM-SSC. Moreover, in appendix B we present two scenarios of how a small value of W_0 can be achieved and what the implied KK scales will be. These models suggest the presence of log-corrections in the ADC, as well. For the LVS we show that the two AdS swampland conjectures also receive log-corrections. As already shown in [49], for the AM-SSC these are subleading.

Since the swampland program forms an intricate set of intertwined statements, our observations should nicely align with other relevant aspects of the program. In section 5 we examine possible such connections. Once one accepts the presence of quantum corrections for the AdS conjectures, one may also expect them for the dS swampland conjecture, in high similarity to the log-corrections that have recently been proposed in the framework of the trans-Planckian censorship conjecture (TCC) [16]. We also discuss the emergence proposal both for tree-level flux models and non-perturbative AdS vacua like KKLT.

2 The swampland conjectures

Let us briefly review the conjectures that we will focus on, as well as the emergence proposal.

2.1 The AdS/moduli scale separation conjecture

The AdS/moduli scale separation conjecture (AM-SSC) [24] states that in an AdS minimum one cannot separate the size of the AdS space and the mass of its lightest mode. Quantitatively, the proposal is that the lightest modulus of non-vanishing mass has to satisfy

$$m_{\text{mod}} R_{\text{AdS}} \leq c \tag{2.1}$$

where c is an order one constant and $R_{\text{AdS}}^2 \sim -\Lambda^{-1}$ the size of AdS. A strong version of this conjecture says that this relation is saturated, i.e. $m_{\text{mod}} \sim R_{\text{AdS}}^{-1}$.

A simple, enlightening example is the 5-form flux supported $AdS_5 \times S^5$ solution of the type IIB superstring. There, the sizes R_{AdS} and R_{S^5} of AdS_5 and S^5 are equal and both related to the 5-form flux. The lightest modulus mass scales as $m_{\text{mod}} \sim R_{S^5}^{-1}$ and saturates relation (2.1).

2.2 The AdS distance conjecture

The AdS distance conjecture (ADC) [15] states that for an AdS vacuum with negative cosmological constant Λ, the limit $\Lambda \to 0$ is at infinite distance in field space and that there will appear a tower of light states whose masses scale as

$$m_{\text{tower}} = c_{\text{AdS}} |\Lambda|^{\alpha} \tag{2.2}$$

for some constant c_{AdS} of order one and $\alpha > 0$. Moreover, for supersymmetric AdS vacua a stronger version of the AdS distance conjecture was claimed, namely that in this

case $\alpha = 1/2$. Assuming that the tower of states is just the KK tower, the strong ADC generalizes earlier Maldacena-Nuñez type obstructions [26] for scale separated type II AdS flux vacua without negative tension object and rephrases them as a swampland conjecture. In the following, we shall consider the KK tower only, leaving open the possibility of other towers appearing.

Let us again consider $AdS_5 \times S^5$. Having $\Lambda \sim -R_{AdS}^{-2}$, we are interested in R_{AdS} becoming large. Then the radius of S^5 also becomes large and the KK modes on S^5 scale as

$$m_{KK}(S^5) \sim \frac{1}{R} \sim |\Lambda|^{\frac{1}{2}}. \tag{2.3}$$

Therefore, these KK modes constitute the tower of states for the (strong) AdS distance conjecture with $\alpha = 1/2$.

2.3 The dS swampland conjecture

For completeness, let us also recall the dS swampland conjecture. In its original formulation [8] it states that

$$|\nabla V| \geq \frac{c}{M_{pl}} \cdot V, \tag{2.4}$$

where c is of order one. The refined version of the conjecture [11, 12] states that either the previous inequality or

$$\min(\nabla_i \nabla_j V) \leq -\frac{c'}{M_{pl}^2} \cdot V \tag{2.5}$$

has to hold, where $\min(\nabla_i \nabla_j V)$ is the minimal eigenvalue of the Hessian matrix and c' is also of order one.

2.4 Emergence of infinite distance

In the framework of the swampland distance conjecture it has been observed that the infinite distance emerges from integrating out the appearing tower of light states [18–20, 50]. In quantitative terms, the emergence proposal claims that the 1-loop contribution to the moduli field metric, arising from integrating out a tower of states that are lighter than the natural cut-off of the effective theory, is proportional to the tree-level metric. Here, let us briefly recall only the main relations. For more details we refer to the original literature.

Say one has an effective theory in D dimensions that has a tower of states with masses $m_n = n\Delta m(\phi)$, with a degeneracy of states at each mass level that scales like n^K. Note that the mass depends on the value of a modulus field ϕ. If N_{sp} of these states become lighter than the species scale [51]

$$\Lambda_{sp} = \frac{\Lambda_{UV}}{N_{sp}^{\frac{1}{D-2}}}, \tag{2.6}$$

they impose a one-loop correction to the field space metric of the field ϕ. Here the UV cut-off Λ_{UV} is often chosen to be the Planck scale but could in principle be lower. The number of species are given by

$$N_{sp} = \sum_{n=1}^{\Lambda_{sp}/\Delta m} n^K \approx \left(\frac{\Lambda_{sp}}{\Delta m}\right)^{K+1}. \tag{2.7}$$

The latter two relations can be inverted to give

$$\Lambda_{\text{sp}} = (\Lambda_{\text{UV}})^{\frac{D-2}{D+K-1}} (\Delta m)^{\frac{K+1}{D+K-1}}, \qquad N_{\text{sp}} = \left(\frac{\Lambda_{\text{UV}}}{\Delta m}\right)^{\frac{(K+1)(D-2)}{D+K-1}}. \tag{2.8}$$

Then the one loop-correction to the field space metric for the modulus ϕ in D dimensions can be written as

$$G_{\phi\phi}^{\text{loop}} \sim \frac{\Lambda_{\text{sp}}^{D+K-1}}{M_{\text{pl}}^{D-2}} \frac{\left(\partial_\phi \Delta m(\phi)\right)^2}{\left(\Delta m(\phi)\right)^{K+3}}. \tag{2.9}$$

In section 5, from requiring $G_{\phi\phi}^{\text{loop}} \sim G_{\phi\phi}^{\text{tree}}$ we will determine the cut-off scale Λ_{sp}. An analogous logic was followed in [48] for the effective theory in the warped throat.

3 Tree-level vacua

Before we turn to the well-known quantum AdS vacua, we review whether tree-level flux vacua comply with the aforementioned AdS conjectures.[1]

The construction of string vacua usually starts with an assumed Ricci-flat background equipped by extra fluxes and instantons. Then one looks for minima of the lowa energy effective action that stabilizes the moduli in a controlled regime. In order to determine the KK scale one has to solve an eigenvalue problem for fluctuations around the background. However, for that purpose one actually has to use the fully backreacted metric. This is often not possible and one hopes that a naive estimate using the initial background plus some control arguments give already a good estimate.[2]

3.1 Type IIA flux models

The best understood examples of AdS minima in string theory are just type IIA and type IIB flux compactifications on Calabi-Yau manifolds. In type IIA one can stabilize all closed string moduli via R-R and H_3 form fluxes. Classes of such concrete models have first been analyzed in [32] and have been called DGKT models. More flux models of this type were considered recently in [33].

A typical example for the isotropic six-torus is presented in appendix A.1, where we also take into account that in general there exists more than just a single KK scale. In these models one has a dilute flux limit that implies that the KK scales can be made parametrically larger than the masses of the moduli. In the limit $\Lambda \to 0$ some fluxes have to become infinite implying that also some of the moduli become infinite. Therefore, $\Lambda \to 0$ is reached at infinite distance in field space. As far as we can tell, all AdS flux models of this type studied in [32, 33] satisfy the relation

$$m_{\text{mod}} \sim |\Lambda|^{\frac{1}{2}} \tag{3.1}$$

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[1]We are indebted to Daniel Junghans and Timm Wrase for pointing out some misconceptions in an earlier version of this paper.

[2]That the backreaction can be essential for seeing some precise cancellations for models with geometric flux was nicely demonstrated (after this paper appeared) in [52].

between the mass of the lightest modulus and the cosmological constant. This includes also the non-supersymmetric models. Therefore, for DGKT models the AM-SSC is satisfied. However, as also claimed in [15], the relevant supersymmetric DGKT vacua do not satisfy the strong version of the ADC but only its weak form with $\alpha < 1/2$, which is also satisfied for the non-supersymmetric ones.

3.2 Geometric fluxes and Freund-Rubin models

Another well known class of AdS minima are Freund-Rubin [53] backgrounds. The standard example is the 5-form flux supported $AdS_5 \times S^5$ solution of the type IIB superstring, whose effective description we recall in the following.

Defining $\rho = R/M_{\rm pl}$ as the radius of the S^5 in Planck units, the string scale and the 5D Planck scale are related as $M_{\rm pl}^8 = M_s^8 \rho^5$. Going to Einstein-frame and performing dimensional reduction of the 10D type IIB Einstein-Hilbert term and the kinetic term for the 5-form flux on the fluxed S^5, one obtains the 5D effective potential

$$V \sim M_{\rm pl}^5 \left(-\frac{1}{\rho^2} + \frac{f^2}{\rho^5} \right) . \tag{3.2}$$

Here $f \in \mathbb{Z}$ is the quantized 5-form flux and the first term is the contribution of the internal curvature. The AdS minimum is at $\rho_0^3 = 5f^2/2$, where the cosmological constant is given by $\Lambda \sim -\rho_0^{-2} M_{\rm pl}^2$. The mass of the modulus ρ can be determined as

$$m_\rho^2 = G^{\rho\rho} \partial_\rho^2 V \big|_0 \sim \frac{M_{\rm pl}^2}{\rho_0^2} \tag{3.3}$$

with the metric on the moduli space $G_{\rho\rho} \sim \rho^{-2}$. Therefore, the mass of the ρ modulus scales in the same way as the geometric KK scale.

This seems to be a generic feature for models where curvature terms are relevant for moduli stabilization. In the framework of 4D flux compactifications this is described by turning on so-called geometric fluxes. A typical example of this kind is presented in appendix A.2. As in the example before, the $\Lambda \to 0$ limit is reached at infinite distance in field space. In these scenarios, there is no dilute flux limit and the KK scale is of the same order as the moduli mass scale. The same feature appears for the non-geometric type IIB flux models presented in [31, 34].

Therefore, irrespective of supersymmetry, these models satisfy both the AM-SSC and the strong ADC.

3.3 Generic scaling of moduli masses

We will now provide a simple argument why for classical flux compactifications the moduli masses are generally expected to scale like $|\Lambda|^{\frac{1}{2}}$. A generic contribution to the flux induced scalar potential scales like $V = A \exp(-a\phi)$, where ϕ is a canonically normalized modulus. Such terms balance against each other so that the cosmological constant is expected to also behave as $\Lambda \sim - \exp(-a\phi)$. Similarly, the masses around the minimum will be given by

$$m_{\rm mod}^2 \sim \partial_\phi^2 V \sim e^{-a\phi} \tag{3.4}$$

so that $m_{\rm mod} \sim |\Lambda|^{1/2}$ is to be expected for a generic tree-level flux compactification.

4 Non-perturbative AdS vacua

In this section, we investigate the two AdS swampland conjectures for the KKLT and the LVS. These vacua are genuinely non-perturbative, in the sense that tree-level contributions are balanced against non-perturbative effects.

4.1 The KKLT AdS vacuum

Let us first consider the KKLT AdS minimum [38] for the single Kähler modulus $T = \tau + i\theta$. Here τ measures the size of a 4-cycle and θ is an axion. After stabilizing the complex structure and axio-dilaton moduli via three-form fluxes, the effective Kähler- and superpotential of KKLT is defined by

$$K = -3\log(T + \overline{T}), \qquad W = W_0 + Ae^{-aT}. \tag{4.1}$$

Here $W_0 < 0$ is the value of the flux induced superpotential in its (non-supersymmetric) minimum and the second term in W arises from a non-perturbative effect like a D3-brane instanton or gaugino condensation on D7-branes. The resulting scalar potential after freezing the axion reads

$$V_{\text{KKLT}} = \frac{aA^2}{6\tau^2}e^{-2a\tau}(3 + a\tau) + \frac{aAW_0}{2\tau^2}e^{-a\tau}. \tag{4.2}$$

The supersymmetric AdS minimum of this potential is given by the solution of the transcendental equation

$$A(2a\tau + 3) = -3W_0e^{a\tau} \tag{4.3}$$

and leads to a negative cosmological constant

$$\Lambda = -\frac{a^2A^2}{6\tau}e^{-2a\tau}. \tag{4.4}$$

In view of the ADC, we first observe that $\Lambda \to 0$ means $\tau \to \infty$, which is at infinite distance in field space. Note that on an isotropic manifold the naive geometric KK scale can be expressed as

$$m_{\text{KK}} \sim \frac{1}{\tau} \tag{4.5}$$

and hence is exponentially larger than any scale $|\Lambda|^\alpha$ expected from the ADC (2.2).

However, one has to keep in mind that one needs an exponentially small W_0. In appendix B, we argue that this requires that the background becomes highly non-isotropic so that the naive estimate of the KK scale (4.5) is not satisfied for the lightest KK or winding modes. In fact for the toroidal example in appendix B.1 we find at leading order

$$m_{\text{w}}^2 \sim \tau e^{-2a\tau} \sim \tau^2|\Lambda| \sim \log^2(-\Lambda)\,|\Lambda| \tag{4.6}$$

and in appendix B.2 for the better controlled strongly warped throat

$$m_{\text{KK}}^2 \sim \frac{1}{\tau^2}\frac{e^{-\frac{2}{3}a\tau}}{\tau^{\frac{1}{3}}} \sim \frac{1}{\log^2(-\Lambda)}|\Lambda|^{\frac{1}{3}}. \tag{4.7}$$

In both cases the masses scale exponentially with τ and feature log-corrections. Up to these corrections, in the toroidal case the strong ADC is satisfied while the better controlled warped throat scenario only satisfies the ADC with $\alpha = 1/6$.

It is also known that the effective mass of the Kähler modulus τ turns out to be much smaller than the naive KK-scale (4.5), in fact it is the lowest mass scale in the problem. This is then the relevant scale for the AM-SSC. In the minimum of the potential one can determine

$$m_\tau^2 = K^{T\overline{T}} \partial_\tau^2 V\big|_0 = \frac{a^2 A^2}{6\tau}\left(2 + 5a\tau + 2a^2\tau^2\right)e^{-2a\tau} \tag{4.8}$$

which indeed contains the desired factor $\exp(-2a\tau)$. Thus, one obtains the relation

$$m_\tau^2 = -\left(2 + 5a\tau + 2a^2\tau^2\right)\Lambda\,. \tag{4.9}$$

Now, neglecting log log-corrections, for large $\tau \gg 1$ one can invert (4.4)

$$a\tau = -b_1 \log(-\Lambda) + b_0 \tag{4.10}$$

with b_1 and b_0 positive constants of order one. Thus, one can express m_τ^2 as

$$m_\tau^2 = -\left(c_2^2 \log^2(-\Lambda) + c_1 \log(-\Lambda) + c_0\right)\Lambda \tag{4.11}$$

with $c_2 > 0$. After reintroducing powers of the Planck scale and working in the limit $\Lambda \to 0$, we can express the mass in the intriguing way

$$m_\tau \sim -c_2 \log\left(-\frac{\Lambda}{M_{\rm pl}^2}\right)|\Lambda|^{\frac{1}{2}}\,. \tag{4.12}$$

Note that $|\Lambda| < M_{\rm pl}$ is required for the effective theory to be controllable. Moreover, in the limit $\Lambda \to 0$ the mass scale still approaches zero.

Therefore, in comparison to the (classical) AM-SSC there appears a logarithmic correction. We propose

$$m\,R_{\rm AdS} \leq c \log(R_{\rm AdS} M_{\rm pl}) \tag{4.13}$$

to be the quantum generalization of the AM-SSC. This is a weaker bound than the classical version (2.1) so that a slight (log type) scale separation between the internal space and the light mode is allowed.

Similarly, as shown in (4.6) and (4.7) we also found log-corrections to the ADC. Therefore we summarize that for quantum vacua like KKLT, where a non-perturbative contribution is balanced against a tree-level one, it seems that there appears a logarithmic correction to the result for simple perturbative vacua.

4.2 The large volume AdS vacuum

Let us analyze another prominent example, namely the large volume scenario (LVS). Recall that here one has a swiss-cheese Calabi-Yau threefold with a large and a small Kähler

modulus, τ_b and τ_s. The precise definition can be found in [39, 54]. What is important here is the final form of the scalar potential

$$V_{\text{LVS}} = \lambda \frac{\sqrt{\tau_s} e^{-2a\tau_s}}{\mathcal{V}} - \mu \frac{\tau_s e^{-a\tau_s}}{\mathcal{V}^2} + \frac{\nu}{\mathcal{V}^3} \tag{4.14}$$

where the total volume is $\mathcal{V} \approx \tau_b^{3/2}$. Let us recall a few relevant steps from the original paper [39]. Solving the minimum condition $\partial_{\mathcal{V}} V_{\text{LVS}} = 0$ one finds

$$\mathcal{V} = \frac{\mu}{\lambda} \sqrt{\tau_s} e^{a\tau_s} \left(1 \pm \sqrt{1 - \frac{3\nu\lambda}{\mu^2 \tau_s^{3/2}}} \right) \tag{4.15}$$

and from $\partial_{\tau_s} V_{\text{LVS}} = 0$ one obtains

$$\left(1 \pm \sqrt{1 - \frac{3\nu\lambda}{\mu^2 \tau_s^{3/2}}} \right) \left(\frac{1}{2} - 2a\tau_s \right) = (1 - a\tau_s). \tag{4.16}$$

Now, one proceeds by working in the perturbative regime $a\tau_s \gg 1$, in which case the two relations can be solved analytically, yielding the values of the moduli in the LVS minimum

$$\tau_s^0 = \left(\frac{4\nu\lambda}{\mu^2} \right)^{\frac{2}{3}}, \qquad \mathcal{V}^0 = \frac{\mu}{2\lambda} \sqrt{\tau_s^0} \, e^{a\tau_s^0}. \tag{4.17}$$

However, plugging this back into the potential (4.14) one gets zero, indicating a sort of extended no-scale structure. Therefore, to find the actual non-vanishing value of the potential in the LVS minimum, one has to compute the next order in $1/\tau_s$ [55]. The only approximation we did is in the solution to (4.16). Thus there will be a correction to τ_s^0, which at leading order is a just a shift by a constant $\tau_s^0 \to \tau_s^0 + c/a$, which one can show to be positive. The value of the cosmological constant will then be

$$\Lambda \sim -\frac{3c\lambda^2 e^{-3c}}{\mu a \, \tau_s^0} e^{-3a\tau_s^0} \left(1 + O\left(\frac{1}{\tau_s} \right) \right) \tag{4.18}$$

which is indeed negative.

The lightest modulus in the game is $\mathcal{V}$, whose mass can be determined by first integrating out τ_s and taking the second derivative of the effective potential with respect to $\mathcal{V}$ (see also [49]). After solving $\partial_{\tau_s} V = 0$, we can write the effective potential as

$$V_{\text{eff}}(\mathcal{V}) = \frac{1}{\mathcal{V}^3} \left(\nu + \frac{\mu^2}{\lambda} \tau_s(\mathcal{V})^{\frac{3}{2}} \left(g(\mathcal{V})^2 - g(\mathcal{V}) \right) \right) \tag{4.19}$$

with

$$g(\mathcal{V}) = 2 \left(\frac{1 - a\tau_s(\mathcal{V})}{1 - 4a\tau_s(\mathcal{V})} \right) = \frac{1}{2} \left(1 - \frac{3}{4a\tau_s(\mathcal{V})} + \dots \right). \tag{4.20}$$

Here τ_s depends implicitly on $\mathcal{V}$. Now, using that at leading order $\partial \tau_s / \partial \mathcal{V} \approx (a\mathcal{V})^{-1}$ we realize that the leading order term (in $1/\tau_s$) again cancels[3] so that

$$m_{\mathcal{V}}^2 = K^{\mathcal{V}\mathcal{V}} \partial_{\mathcal{V}}^2 V_{\text{eff}} \big|_0 \sim \frac{\lambda^2}{\mu a \, \tau_s^0} e^{-3a\tau_s^0} \left(1 + O\left(\frac{1}{\tau_s} \right) \right). \tag{4.21}$$

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[3]We thank Joe Conlon for pointing this out to us.

Therefore, for the LVS AdS minimum we have found the relation

$$m_{\mathcal{V}}^2 \sim |\Lambda| \left(c_0 + \frac{c_{-1}}{\log(-\Lambda)} + \dots \right) , \tag{4.22}$$

which means that in the limit $\Lambda \to 0$ the LVS satisfies the strong (classical) AM-SSC. However, also for LVS there will be subleading log-corrections. In contrast to KKLT, here the first two coefficients are vanishing i.e. $c_2 = c_1 = 0$ which presumably is due to the extended no-scale structure and the perturbative stabilization of $\mathcal{V}$. For LVS one can have $W_0 = O(1)$ so that the naive estimate for the KK scale is justified

$$m_{\mathrm{KK}}^2 \sim \frac{1}{\mathcal{V}^{\frac{4}{3}}} \sim \frac{1}{\tau_s^{\frac{2}{3}}} e^{-\frac{4}{3}a\tau_s} \sim \frac{1}{\log^{\frac{2}{9}}|\Lambda|} |\Lambda|^{\frac{4}{9}} . \tag{4.23}$$

This result is similar to the KK modes (4.7) for the KKLT model in the warped throat. Thus, in the ADC one expects again extra quantum log corrections and $\alpha = 2/9$.

5 Other swampland conjectures

In this section we discuss the implications and relations of the log-correction to other swampland conjectures. First, following a similar reasoning as in section 3.1, we provide a general argument for the appearance of such corrections.

5.1 Origin of log-corrections

To see the origin of the log-corrections consider a typical non-perturbative contribution to the scalar potential, which in canonically normalized variables takes the following double-exponential form

$$V \sim A\, e^{-c\phi}\, e^{-(be^{a\phi})} + V_{\mathrm{others}} . \tag{5.1}$$

The other corrections can be perturbative or non-perturbative, depending on the nature of the model. If moduli stabilization occurs such that the first term balances the terms in V_{others}, the size of the first one is expected to set the scale for the potential and the masses in the minimum. Computing its second derivative with respect to ϕ one gets

$$m^2 \sim \partial_\phi^2 V \sim \left(c^2 + 2abc\, e^{a\phi} - ba^2\, e^{a\phi} + (ab)^2 e^{2a\phi} \right) V . \tag{5.2}$$

Inverting (5.1) one can write

$$e^{a\phi} \sim -\frac{1}{b} \log\left(\frac{V}{A} \right) = -b_1 \log|V| + b_0 \tag{5.3}$$

so that

$$m^2 \sim -\left(c_2^2 \log^2(|V|) + c_1 \log(|V|) + c_0 \right) V . \tag{5.4}$$

We observe that these terms take precisely the form of those that we found for KKLT in (4.11). One can well imagine that for a full model the potential will be more complicated so that like in LVS also further subleading corrections $\log^{-n}(|V|)$ ($n \geq 1$) will appear.

Thus, we conclude that the logarithmic corrections are genuinely related to the appearance and relevance of non-perturbative effects in the scalar potential. In the moment that such genuinely non-perturbative vacua exist in string theory, the AdS swampland conjectures are expected to receive log-corrections.

5.2 Trans-Planckian censorship conjecture

If the AdS swampland conjectures receive such corrections, it is natural to expect that also the dS swampland conjecture will be changed. Computing the first derivative of (5.1), a natural guess would be

$$|\nabla V| \geq V\Big(c_1 \log\left(|V|\right) + c_2\Big). \tag{5.5}$$

In contrast to the AdS swampland conjectures this relation is supposed to hold not only at a specific point in field space (namely the minimum) but at every point. It remains to be seen whether such a strong local bound really makes sense. In any case, it is remarkable that the right hand side could vanish for $V = \exp(-c_2/c_1)$, thus potentially allowing dS vacua.[4]

It has been recently suggested that a more "global" version might be the more general statement. In [16] an underlying quantum gravity reason for the dS swampland conjecture was proposed, namely the so-called trans-Planckian censorship conjecture (TCC). It proposes that sub-Planckian fluctuations must stay quantum and should never become classical in an expanding universe with Hubble constant H. More quantitatively it says

$$\int_{t_i}^{t_f} dt\, H < \log\left(\frac{M_{\rm pl}}{H_f}\right), \tag{5.6}$$

for more details consult [16]. Two points are to be emphasized here, namely that this conjecture is not local (as it involves an initial and final time), and the appearance of a logarithm on the right hand side. For a monotonically decreasing positive potential, the authors of [16] derived from the TCC a global version of the dS swampland conjecture

$$\left\langle \frac{-V'}{V} \right\rangle \Bigg|_{\phi_i}^{\phi_f} > \frac{1}{\Delta\phi} \log\left(\frac{V_i}{M}\right) + \frac{2}{\sqrt{(d-1)(d-2)}} \tag{5.7}$$

where $\left\langle \frac{-V'}{V} \right\rangle \Big|_{\phi_i}^{\phi_f}$ denotes the average of $-V'/V$ in the interval $[\phi_i, \phi_f]$. Here $V < M < M_{\rm pl}$ and M is a mass scale that can be lower than the Planck-scale.

Let us check that a potential of the generic form

$$V(\phi) = A e^{-c\phi}\, e^{-(b e^{a\phi})} \tag{5.8}$$

satisfies this averaged dS swampland conjecture. Note that for $a, b, c > 0$ this potential is indeed positive and monotonically decreasing. For the average value we can directly compute

$$
\begin{aligned}
\left\langle \frac{-V'}{V} \right\rangle \Bigg|_{\phi_i}^{\phi_f} &= \frac{1}{\Delta\phi} \int_{\phi_i}^{\phi_f} d\phi \left(c + ab\, e^{a\phi} \right) = c + \frac{b}{\Delta\phi}\left(e^{a\phi_f} - e^{a\phi_i} \right) \\
&> c - \frac{b}{\Delta\phi} e^{a\phi_i} > c + \frac{1}{\Delta\phi} \log\left(\frac{V_i}{A}\right).
\end{aligned}
\tag{5.9}
$$

This has precisely the form (5.7) so that we can state that non-perturbative contributions to the scalar potential induce the log-corrections in the TCC derived dS swampland conjecture (5.7). Moreover, we observe that for the three terms in the KKLT potential (4.2), one

---

[4]Utilizing quantum effects to generate stable dS vacua has been discussed in e.g. [56, 57].

gets the parameters $c \in \{\sqrt{8/3}, \sqrt{2/3}\}$ which both satisfy $c \geq \sqrt{2/3}$, the value appearing in (5.7).

We consider this connection to the TCC as further evidence for the appearance of log-corrections in the various swampland conjectures.

5.3 Emergence for KKLT

Finally, we comment on the emergence proposal. Before we come to the KKLT model let us first consider tree flux compactifications.

For compactification on S^5 one needs to take heed of the degeneracy of KK modes of mass $m_n = n\Delta m = nM_{\mathrm{pl}}/\rho$. This is given by the dimensionality of the space of harmonic functions of homogeneous degree n, which for the 5-sphere goes as n^4. Applying our general result (2.9) for the one-loop correction to the field space metric and setting it equal to the tree-level metric $G_{\phi\phi}^{\mathrm{tree}} \sim \rho^{-2}$ we obtain $\Lambda_{\mathrm{sp}}^8 = M_{\mathrm{pl}}^3 (\Delta m)^5 = M_{\mathrm{pl}}^8/\rho^5$. Taking the relation between the string scale and the D dimensional Planck scale into account it follows $\Lambda_{\mathrm{sp}} \sim M_s$. We expect that this relation will appear for all tree-level flux compactifications so that the cut-off of these models is simply the string scale.

Next we will analyze the implications of the emergence proposal in the KKLT setting, both for the toroidal model and the strongly warped throat.

Emergence for toroidal KKLT. In appendix B.1 we found a (non-degenerate) tower of states with discretized masses scaling as (4.6). At leading order in τ, the 1-loop correction (2.9) to the field space metric of the modulus τ reads

$$G_{\tau\tau}^{\mathrm{loop}} \sim \frac{\Lambda_{\mathrm{sp}}^3}{M_{\mathrm{pl}}^3} \frac{a^2}{\sqrt{a\tau}} e^{a\tau} \,. \tag{5.10}$$

Imposing that this is proportional to the tree-level metric $G_{\tau\tau}^{\mathrm{tree}} \sim \tau^{-2}$, one can determine the value of the cut-off of the effective theory as

$$\Lambda_{\mathrm{sp}}^3 \sim \frac{e^{-a\tau}}{\tau^{\frac{3}{2}}} M_{\mathrm{pl}}^3 \,. \tag{5.11}$$

This is reminiscent of the dynamically generated mass scale Λ_{SQCD} of the SYM theory that undergoes gaugino condensation. This scale is usually given by $\Lambda_{\mathrm{SQCD}}^3 = e^{-a/g^2} M^3$, where M denotes a UV cut-off scale. Noting that $g^{-2} \sim \tau$ we can write the KKLT cut-off as

$$\Lambda_{\mathrm{sp}}^3 \sim e^{-\frac{a}{g^2}} (g\, M_{\mathrm{pl}})^3 \sim \Lambda_{\mathrm{SQCD}}^3 \,. \tag{5.12}$$

Thus the cut-off of the KKLT model is the scale at which the implicitly assumed gaugino condensation of the confining gauge theory occurs, while the UV cut-off of the gauge theory itself is not simply the Planck scale but rather $M \sim \Lambda_{\mathrm{UV}} \sim gM_{\mathrm{pl}}$ as suggested by the weak gravity conjecture.

Emergence for warped throat KKLT. Next we analyze our second example from appendix B.2, where the small value of W_0 is generated by a strongly warped throat. As

shown in [48], in this case there exists a tower of highly red-shifted KK modes localized at the tip of the throat with masses

$$\Delta m_{\text{KK}} \sim \frac{|Z|^{\frac{1}{3}}}{\tau^{\frac{1}{2}} y_{\text{UV}}}, \tag{5.13}$$

where Z denotes the conifold (complex structure) modulus and y_{UV} is the length of the KS throat before it reaches the bulk Calabi-Yau. It was argued in [48] that these KK modes are lighter than the cut-off of the effective theory and that their one-loop contribution corrects the second (subleading) term in the Kähler potential

$$K = -3\log(T + \overline{T}) + c\frac{|Z|^{\frac{2}{3}}}{(T + \overline{T})}. \tag{5.14}$$

Using the general relation (2.9) and setting this one-loop correction equal to the Kähler metric $G_{\tau\overline{\tau}}$ (second term) one finds for the species scale

$$\Lambda_{\text{sp}}^3 \sim \frac{|Z|}{\tau^{\frac{3}{2}} y_{\text{UV}}} M_{\text{pl}}^3. \tag{5.15}$$

Since the first term in the above Kähler potential is also present in the unwarped case, we expect it to emerge from integrating out the tower of heavier bulk KK modes (4.5) with mass scale $\Delta m_{\text{KK,h}} \sim 1/\tau$.

In the limit where the throat just fits into the warped Calabi-Yau volume, one can determine the cut-off y_{UV} as

$$y_{\text{UV}} \sim -\log\left(\frac{|Z|}{\tau^{\frac{3}{2}}}\right). \tag{5.16}$$

Now we stabilize the Kähler modulus τ via KKLT which gives the relations

$$|Z| \sim |W_0| \sim \tau e^{-a\tau}, \qquad y_{\text{UV}} \sim \tau \tag{5.17}$$

so that the species scale can be expressed as

$$\Lambda_{\text{sp}}^3 \sim \frac{e^{-a\tau}}{\tau^{\frac{3}{2}}} M_{\text{pl}}^3, \tag{5.18}$$

which is the same result (5.11) as for the toroidal setting. Therefore, also for the warped throat KKLT model the relation to Λ_{SQCD} holds.

6 Conclusions

In this paper, we have investigated the behaviour of the KKLT and LVS constructions with regard to the AdS scale separation and distance conjectures. To this end, we have identified the relevant towers of light states and provided two concrete examples of realizing an exponentially small mass scale for the KKLT model. Driven by confidence in the consistency of the aforementioned AdS vacua, we propose log-corrections to the tree-level AdS swampland conjectures. Extending our reasoning, we expect similar log-corrections

to the dS swampland conjecture. These might be in the same spirit as the log-corrections that were found for the "average" of the dS swampland conjecture in the recently proposed TCC. Whether a stronger, local version of a quantum dS swampland conjecture could make sense is left for future analysis.

Additionally, we analyzed the consequences of imposing the emergence proposal. For tree-level flux compactifications we found that the cut-off scale is simply the string scale. For the KKLT model, it is remarkable that both proposed scenarios for generating an exponentially small W_0 lead to a cut-off scale reminiscent of the dynamically generated scale for the condensing SYM theory. It is certainly encouraging that our observations seem to fit well within the broader swampland picture.

Acknowledgments

A KK scales in type II flux compactifications

In this appendix we present two simple though typical type IIA flux compactifications on the isotropic T^6. The first one only contains R-R and H_3-form fluxes and as expected features a dilute flux limit that allows to separate the KK scale from the moduli mass scale. The second one also contains geometric fluxes, in which case there will be no dilute flux limit and the moduli masses are of the same scale as the KK modes. Freund-Rubin type compactifications are fully fledged 10D uplifts of such effective models.

A.1 A typical DGKT model

Let us consider type IIA orientifolds with fluxes on an isotropic six-torus. Here one has three chiral superfields $\{S, T, U\}$ whose real parts are defined as

$$\tau = r_1 r_2, \qquad s = e^{-\phi} r_1^3, \qquad u = e^{-\phi} r_1 r_2^2. \tag{A.1}$$

The axions do not play any role in the following and will in all examples be stabilized at vanishing value. The Kähler potential is given as

$$K = -3\log(T + \overline{T}) - 3\log(U + \overline{U}) - \log(S + \overline{S}). \tag{A.2}$$

Now we turn on just R-R fluxes and H_3-form flux so that the flux induced superpotential reads

$$W = if_0 T^3 - 3if_4 T + ih_0 S + 3ih_1 U. \tag{A.3}$$

Then there exist both supersymmetric and non-supersymmetric AdS minima. For instance in the supersymmetric vacuum, the saxionic moduli are stabilized at

$$\tau = \kappa \frac{f_4^{\frac{1}{2}}}{f_0^{\frac{1}{2}}}, \qquad s = \frac{2\kappa}{3} \frac{f_4^{\frac{3}{2}}}{f_0^{\frac{1}{2}} h_0}, \qquad u = 2\kappa \frac{f_4^{\frac{3}{2}}}{f_0^{\frac{1}{2}} h_1} \tag{A.4}$$

with $\kappa = \sqrt{5/3}$. For the non-supersymmetric minima only the numerical prefactors change. The effective masses of the moduli all scale in same way as

$$m_{\text{mod}}^2 \sim -\Lambda \sim \frac{f_0^{\frac{5}{2}} h_0 h_1^3}{f_4^{\frac{9}{2}}} M_{\text{pl}}^2 \,. \tag{A.5}$$

Therefore the model satisfies the AM-SSC. The two KK scales are

$$m_{\text{KK},1}^2 = \frac{M_s^2}{r_1^2} = \frac{M_{\text{pl}}^2}{s \, \tau \, u} = \frac{f_0^{\frac{3}{2}} h_0 h_1}{f_4^{\frac{7}{2}}} M_{\text{pl}}^2 \,,$$

$$m_{\text{KK},2}^2 = \frac{M_s^2}{r_2^2} = \frac{M_{\text{pl}}^2}{\tau \, u^2} = \frac{f_0^{\frac{3}{2}} h_1^2}{f_4^{\frac{7}{2}}} M_{\text{pl}}^2 \,. \tag{A.6}$$

To be in the perturbative regime ($g_s \ll 1$) we now choose $f_0, h_0, h_1 = O(1)$ and $f_4 \gg 1$. In this regime the KK scales are parametrically larger than the moduli masses and one has

$$m_{\text{KK},i}^2 \sim |\Lambda|^{\frac{7}{9}} \,, \tag{A.7}$$

thus satisfying the ADC with $\alpha = 7/18$, both in the supersymmetric and non-supersymmetric case. This reflects the fact that type IIA flux compactifications admit a dilute flux limit (where $f_4 \to \infty$).

A.2 A type IIA model with geometric flux

The Freund-Rubin background $AdS_5 \times S^5$ is an example of a flux compactification, where for moduli stabilization the curvature of the internal space is essential (as seen in e.g. the effective potential eq. (3.2)). Such backgrounds can be described by turning on geometric flux ω in the effective theory.

As a typical simple model we consider the superpotential

$$W = f_6 + 3f_2 T^2 - \omega_0 ST - 3\omega_1 UT \,. \tag{A.8}$$

Then the saxions are stabilized in a supersymmetric AdS minimum at

$$\tau = \frac{1}{3} \frac{f_6^{\frac{1}{2}}}{f_2^{\frac{1}{2}}} \,, \qquad s = 2 \frac{f_2^{\frac{1}{2}} f_6^{\frac{1}{2}}}{\omega_0} \,, \qquad u = 2 \frac{f_2^{\frac{1}{2}} f_6^{\frac{1}{2}}}{\omega_1} \tag{A.9}$$

and receive masses that scale as

$$m_{\text{mod}}^2 \sim -\Lambda \sim \frac{\omega_0 \, \omega_1^3}{f_2^{\frac{1}{2}} f_6^{\frac{3}{2}}} M_{\text{pl}}^2 \,. \tag{A.10}$$

In this case the two KK scales are

$$m_{\text{KK},1}^2 = \frac{\omega_0 \, \omega_1}{f_2^{\frac{1}{2}} f_6^{\frac{3}{2}}} M_{\text{pl}}^2 \,, \qquad m_{\text{KK},2}^2 = \frac{\omega_1^2}{f_2^{\frac{1}{2}} f_6^{\frac{3}{2}}} M_{\text{pl}}^2 \tag{A.11}$$

which satisfy $m_{\text{KK},1}^2 \sim m_{\text{mod}}^2/\omega_1^2$ and $m_{\text{KK},2}^2 \sim m_{\text{mod}}^2/(\omega_1\omega_2)$. It was shown in [52] that taking the backreaction of the fluxes onto the metric into account, the geometric fluxes in the denominator also cancel and parametrically one indeed finds $m_{\text{KK}}^2 \sim m_{\text{mod}}^2$. Therefore, in such models with geometric flux there is no parametric separation of the KK scale and the moduli masses, the same behavior that occurs for Freud-Rubin compactifications. Here both the AM-SSC and the strong ADC are satisfied.

B KK scales for KKLT

The ultimate question is what happens with the KK scale in the KKLT scenario. As we have seen, the naive KK scale (4.5) is not exponentially small in the Kähler modulus τ, so that there seems to be no way that the ADC can hold. However, we have to keep in mind that for the KKLT scenario to work one needs an exponentially small value of W_0 after stabilizing the complex structure moduli and the dilaton à la GKP [58].

B.1 A toroidal example

Let us consider again a simple toroidal type IIB model, for which we can easily compute the KK scales directly. In this case, the real parts of the chiral superfields $\{S, T, U\}$ are defined as

$$\tau = e^{-\phi}r_1^2 r_2^2, \qquad s = e^{-\phi}, \qquad u = r_2/r_1. \tag{B.1}$$

We turn on F_3 and H_3 form flux such that the superpotential is

$$W = ifU + ihSU^2 \tag{B.2}$$

with f, h positive. This freezes the axions completely and the saxions have to satisfy

$$us = \frac{f}{h}, \tag{B.3}$$

leading to the value $W_0 = 2ifu$ of the superpotential along the minimum. Therefore, the superpotential becomes very small for $u \ll 1$ while the dilaton stays in the perturbative regime, i.e. $e^{\phi} \ll 1$. The KK and winding scales can be computed in terms of u and τ as

$$m_{\text{KK},\pm}^2 = \frac{M_{\text{pl}}^2}{\tau^2}u^{\pm 1}, \qquad m_{\text{w},\pm}^2 = \frac{M_{\text{pl}}^2}{s\tau}u^{\pm 1}. \tag{B.4}$$

In the regime of interest $s \sim u^{-1}$, the large radius regime, where the KK scale is smaller than the winding scale, is given by $\tau u \gg 1$. As long as τ is not stabilized this can always be satisfied by choosing τ large enough. However, in KKLT τ is fixed as $W_0 \sim u \sim \tau \exp(-a\tau)$, which implies $\tau u \ll 1$. Therefore, the lightest tower of states in KKLT is given by the winding modes

$$m_{\text{w},+}^2 \sim \frac{M_{\text{pl}}^2}{\tau}u^2 \sim \frac{M_{\text{pl}}^2}{\tau}|W_0|^2. \tag{B.5}$$

In the KKLT minimum, using (4.3) and (4.4), the winding mass becomes

$$\begin{aligned} m_{\text{w},+}^2 &\sim \frac{e^{-2a\tau}}{\tau}(4a^2\tau^2 + 12a\tau + 9)\,M_{\text{pl}}^2 \\ &\sim (\log^2(-\Lambda) - 6\log(-\Lambda) + 9)|\Lambda|\,, \end{aligned} \tag{B.6}$$

which is of the same form as (4.11). This satisfies the strong ADC up to log corrections.

Thus we have seen that, if we want to have an exponentially small value of W_0, the torus becomes highly non-isotropic and towers of states become exponentially light $m_{\mathrm{w},+}^2 \sim M_{\mathrm{pl}}^2 \tau^{-1} \exp\left(-\sqrt{8/3}\,\phi_u\right)$ in the canonically normalized field $\phi_u = -\sqrt{3/2}\log u$ corresponding to the complex structure modulus u. This is nothing else than the tower of states that must become light in the large ϕ_u regime due to the swampland distance conjecture.

One relevant concern is that for such large excursions in the complex structure u, the effective theory that we used for computing complex structure moduli stabilization is not under control anymore. For instance, as we have seen we do not satisfy $\tau u \gg 1$ so that one radius of the torus becomes significantly smaller than the string length. As a consequence, higher derivative terms in the effective action might not be negligible. Moreover, the mass scale of these winding modes is lighter than the mass $m_{u,s}^2 \sim \tau^{-3}$ of the stabilized complex structure moduli. Therefore, this simple toroidal model can certainly not serve as a completely convincing flux GKP compactification with $W_0 \ll 1$. Nevertheless, it exhibits an important feature, namely that $W_0 \ll 1$ goes along with the occurrence of a tower of states whose mass scales as $m \sim \exp(-a\tau)$.

B.2 The warped throat

Another option was proposed in [48], namely that a superpotential involving the complex structure modulus Z governing the appearance of a conifold singularity can also generate an exponentially small value for W_0. If $|Z| \ll 1$ the three-cycle of the conifold becomes very small and locally the geometry is described by a Klebanov-Strassler(KS) throat. For our purpose we only need a couple of relations. First the superpotential in the minimum is given by

$$|W_0| \sim |Z| \sim \exp\left(-\frac{2\pi h}{g_s f}\right) \tag{B.7}$$

where f, h are F_3 and H_3 fluxes supporting the strongly warped KS throat. It was shown in [48] that there exists a tower of light KK modes localized close to the tip of the conifold with masses

$$m_{\mathrm{KK}}^2 \sim \frac{1}{y_{\mathrm{UV}}^2}\left(\frac{|Z|}{\mathcal{V}}\right)^{\frac{2}{3}} M_{\mathrm{pl}}^2 \tag{B.8}$$

where $\mathcal{V} = \tau^{3/2}$ denotes the warped volume of the threefold. (Note that one must have $\mathcal{V}|Z|^2 \ll 1$ in the strongly warped throat.) Moreover, y_{UV} denotes the length of the warped KS throat before it goes over to the bulk Calabi-Yau manifold. In the limit that the throat just fits into the Calabi-Yau volume one can relate y_{UV} to the other quantities as

$$y_{\mathrm{UV}} \sim -\log\left(\frac{|Z|}{\mathcal{V}}\right) \tag{B.9}$$

(see [48] for further details). It was also found that the mass scale of these KK modes is of the same order as the mass of the complex structure Z. Thus, one is (still) at the limit of control of the utilized effective theory. In that respect, this scenario is better controlled than the toroidal model discussed before.

Now using again the KKLT minimum condition (4.3) we get[5] $y_{\text{UV}} \sim a\tau$ and can express this exponentially small KK scale as

$$
\begin{aligned}
m_{\text{KK}}^2 &\sim \frac{|W_0|^{\frac{2}{3}}}{a^2\tau^3} M_{\text{pl}}^2 \sim \left(\frac{(2a\tau+3)^2}{a^8\tau^8} \right)^{\frac{1}{3}} \left(\frac{a^2 e^{-2a\tau}}{\tau} \right)^{\frac{1}{3}} M_{\text{pl}}^2 \\
&\sim \left(\frac{1}{\log^2(-\Lambda)} - \frac{2}{\log^3(-\Lambda)} + \dots \right) |\Lambda|^{\frac{1}{3}} M_{\text{pl}}^2 .
\end{aligned}
\tag{B.10}
$$

Up to the log-term this satisfies the ADC with $\alpha = 1/6$.

We note that $\mathcal{V}|Z|^2 \sim \tau^{7/2} \exp(-2a\tau)$ which for large τ is indeed much smaller than one. Therefore, stabilizing the Kähler modulus via KKLT is self-consistent with using the effective theory in the warped throat. In this respect, it also behaves better than the toroidal model from the previous subsection.

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