KKLT an overview

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Abstract

In this thesis a detailed description of the KKLT scenario is given as well as as a comparison with later papers critiquing this model. An attempt is made to provide a some clarity in 17 years worth of debate. It concludes with a summary of the findings and possible directions for further research.

Contents

Chapter 1

Introduction

During the turn of the last century it has become apparent that our universe is expanding and that this expansion is accelerating. When attempting to describe these observations using the Friedmann equations, derived from general relativity, we find we need a energy source which behaves as a gas with negative pressure. This appears as when we add up all the known energy contributions and match it with the geometry of the universe, which we can observe from for example the cosmic microwave background, we find a miss match if we do not add such a contribution. Such an energy contribution is known as a cosmological constant.
This was somewhat of a surprise, as up until then it was generally assumed that the cosmological constant would be zero. Besides explaining the current accelerated expansion the cosmological constant also provides a way to explain the rapid expansion in the early universe known as inflation and has become a fundamental part of the Λcdm model, the current leading model in cosmology.
As mentioned above the cosmological constant behaves as a gas with negative pressure which means it has an energy contribution which does not change under expansion of the universe, contrary to contributions form matter and radiation which all fall of in density as the universe expands. This property of not diminishing in energy density even during expansion implies that the cosmological constant is a property of empty space, since the new space that is created during expansion would also carry the energy with it thereby the total per volume does not change.
This energy of empty space is very well known in quantum field theory where the groundstate, i.e. empty space, generally carries with it energy. In the Feynmann diagram picture this comes from the summation over loop diagrams, also known as quantum fluctuations of the vacuum. And

emerges in cosmology also as an explanation for the fluctuations in the cosmic micro wave background. This vacuum energy we can calculate using the standard model, our current best model of particle physics, by preforming the summation over bubble diagrams. When preformed this leads to an unexpected high energy for the vacuum, much higher than the cosmological constant observed. This is the core of the cosmological constant problem.
As the observational evidence for the cosmological constant is fundamentally gravitational, one could argue that we should not search for our explanation in the standard model as this does not include gravity. But we can not simply ignore the result from the standard model, as independent experiments, such as those measuring the Casimir force, have shown that this vacuum energy does exist. Therefor even if we do not equate the result of this calculation to the cosmological constant we still expect it to enter in to the total cosmological constant of our universe. As at this point we are attempting to explain a gravitational problem using a quantum theory, it has become clear that we'll need a theory of quantum gravity. This is maybe the largest frontier of modern physics, a theory unifying both revolutions of last century physics, general relativity and quantum mechanics. One of the most prominent candidates for quantum gravity is string theory, originally developed as a theory for the strong nuclear force but abandoned in favor of quantum chromodynamics. It's formalism was later realised to be able applicable to describe quantum gravity. The original theory describing Bosonic particles was later expanded to be applicable to Fermions as well. This required the introduction of supersymmetry, a symmetry transforming Bosons in to Fermions and vice versa. This was at the time seen as very promising as when in a supersymmetric theory one counts the bubble diagrams from both the Fermions and Bosons they cancel, resulting in a vanishing vacuum energy.* But later collision experiments failed to observe the particles predicted by supersymmetry, the so called superpartners which are the supersymmetric counterparts to known particles. This lack of superpartners and the non-vanishing cosmological constant observed mean that supersymmetry needs to be broken or abandoned in order to describe our universe. As supersymmetry provides an enormous simplification in calculations and a reduction in the number of dimensions required for string theory to be consistent, modern string theories usually are formulated such that the supersymmetry breaking and with it the mass of the superpartners put them beyond the

*In general supersymmetric string theory also allows a negative vacuum energy but not a positive one.

reach of current observations. as hinted at earlier the resulting theory still runs in to difficulties which in no small part is due to the fact that we do not have a fully consistent description of the entire string theory. The best descriptions we can write are low energy effective theories not dissimilar to the effective field theories one frequently encounters in quantum field theory. Such low energy effective string theories incorporating supersymmetry are known as supergravities. These theories differ by being different limits of the full string theory, as such there exist ways of translating between these supergravity theories. This framework of different theories being the low energy limits of a full theory was originally proposed by Edward Witten in 1995 [1] and has since become known as M-theory.

These theories however still have problems to overcome in order to solve the cosmological constant problem. Besides the issue, hinted at earlier, that supersymmetry needs to be broken in the groundstate, in order to have a non-negative vacuum energy, they also need to provide a solution to what is known as the Dine-Seiberg problem, [2]. Often stated as "When corrections are important, they are not computable, and when they are computable, they are not important." [3] (p. 125-126). Although the details of the problem are quite technical in nature the idea is quite simple. Since we expect our theory to be asymptotically free as function of a certain coupling, our potential should vanish asymptotically as function of this coupling. This means that when coming from strong coupling the potential should either approach zero from above or bellow, assuming it is not free at strong coupling. In the first case this means that to first order ours is a runaway potential to weak coupling, in the second case our potential would to first order be a runaway to strong coupling. So in order for an at least meta stable regime we need to take in to account higher order corrections, but to do this consistently we need to check all orders for relevancy which is beyond our capabilities in the regime where we would want to stabilise. So in other words our corrections are either irrelevant and we have an unstable theory or they relevant an we can't compute them.

At this point our problem seems hopeless, but string theory allows for quite a few exotic object, such as D-branes which we'll introduce in chapter 2, which possibly allow for constructions that overcome these problems. The focus of this thesis will be such a construction which attempts to overcome some of these problems. This construction, known as KKLT after the authors who originally proposed it, Kachru, Kallosh, Linde and Trivedi, [4], has been one of the most promising of such construction but is definitely not without its problems. It is this construction and the problems facing it that will be the focus of this thesis.

The structure of this thesis will be as follows, in chapter 2 we'll provide some perspective on the physical problem which motivated the research as well as giving a very brief overview and surprising results of string theory. In chapter 3 we'll go trough the model at the basis of this research while trying to highlight points which in chapter 4 we relate to critique point of this model. Then in chapter 5 we'll conclude with a summary of the research and offer some directions in which progress can be made.

Chapter 2

The Basics

2.1 Cosmological constant problem and the Expanding Universe

Often quoted as the most famous blunder of Einstein the cosmological constant, denoted as Λ, is a scalar parameter which can be added to the action of General relativity.

$$S = \int \sqrt{-g}(\frac{1}{2\kappa}(R - 2\Lambda))d^4x + S_M \tag{2.1}$$

Here g is the determinant of the metric tensor, κ is a coupling constant, R is the Ricci curvature, Λ will be the cosmological constant and S_m is the action for any matter content, which is generally depend on the metric and matter fields. The original reason for introduction, Einsteins attempt at a static universe, is dismissed in modern physics as it is unstable. However the cosmological constant remains an object of interest, particularly in cosmology. This is because the cosmological constant is closely related with the global structure of the universe. When varying (2.1) with respect to the metric we find Einsteins equation, which using the Einstein convention has the following form.

$$R_{\mu\nu} - \frac{1}{2}Rg_{\mu\nu} + \Lambda g_{\mu\nu} = \kappa T_{\mu\nu} \tag{2.2}$$

Where $T_{\mu\nu}$ is the stress energy tensor defined by:

$$T_{\mu\nu} = -\frac{2}{\sqrt{-g}}\frac{\delta S_m}{\delta g^{\mu\nu}} \tag{2.3}$$

As our universe appears to be homogeneous and isotropic on large scales we can use the FLRW metric for homogeneous and isotropic universe,

$$ds^2 = -c^2dt^2 + a(t)^2 ds_3^2 \tag{2.4}$$

where a is a scale factor and ds_3 a 3d spacial metric. Combining this metric with (2.2) we find the Friedmann equation.

$$(\frac{\dot{a}}{a})^2 - \frac{\kappa\rho}{3} - \frac{\Lambda}{3} = -\frac{k}{a^2} \tag{2.5}$$

Where we have introduced the convention $c = 1$ and ρ represents the matter density and $k = -1, 0, 1$ corresponds with a negative, flat or positive spacial curvature respectively. From observations, such as measurement of the angles cosmological sized triangles using the cosmic microwave background, we can determine the spacial curvature of our universe, which to very good approximation appears to be flat meaning $k = 0$. Thus the right hand side of (2.5) vanishes. The current expansion of our universe we can observe using type 1a supernovae and we find it to be positive, i.e. our universe is an expanding one. From observation we can also determine ρ by adding up all matter contributions, including radiation Baryonic and Dark matter. Adding all this observational evidence together we find that for our universe Λ has to be positive but extremely small. Other observational evidence for a positive cosmological constant comes from the fact that the current expansion of our universe appears to be exponential which means our current epoch is one dominated by vacuum energy. * The simplest universe with a positive cosmological constant is a de Sitter (dS) universe, this is a maximally symmetric space without matter and with a positive cosmological constant. Since our universe clearly does contain matter this is not our universe however it is a good description of what our universe looked like shortly after the big bang during the epoch of inflation and also what our universe is expected to look like in the far future as vacuum energy domination continues. So we say our universe is asymptotically de Sitter. Thus a good start for a description of our universe with it's positive cosmological constant would be to find a description of a de Sitter universe. This as outlined in the introduction, necessarily incorporates both gravity and quantum effect and therefor will require a quantum gravity theory. With our problem and goal outlined we'll spend

*These different sources for determining the cosmological constant don't agree on the precise value of the cosmological constant, this discrepancy is known as the tension in cosmology. We'll not delve in to this separate problem any further in this thesis.

the rest of this chapter on a broad stroke introduction of String theory and it's components which will be relevant to the KKLT construction.

2.2 String theory

Aside from the cosmological constant problem mentioned in the previous section modern physics has additional problems, such as the mass of neutrino's, chiral gauge couplings and unification. String theory has become one of the main candidate theories to address these problems. And as such the search for constructions of cosmological models in string theory is far from surprising. Should string theory really be the new physics we are looking for then we could hope that it provides an answer to the cosmological constant problem. This search has already resulted in many surprising discoveries one of which is the enormous amount of possible universes describable using string theory. An often given figure for this 10^{500} more as a way to illustrate the scale rather than to be used as a precise figure. We'll refer to this number in section 2.2.5. This section is broken down into parts each discussing different aspect of string theory relevant to our discussion of the Kachru Kallosh Linde and Trivedi model.

2.2.1 Classical relativistic strings

We'll take a look at a very simple string, namely the classical† relativistic open string, as it gives an intuitive insight into some basics of string theory.‡ As in most physics we'll be interested in formulating an action as the basis for our theory. Let's recall the action for a relativistic particle familiar from relativity. Again using the Einstein summation convention this is the following.

$$S = -mc \int_{\tau_i}^{\tau_f} \sqrt{-\eta_{\mu\nu} \frac{\partial x^\mu}{\partial \tau} \frac{\partial x^\nu}{\partial \tau}} d\tau \tag{2.6}$$

Where τ is a parameter defining a position along the worldline of a particle via the function $x(\tau)$ from the parameter space to physical space. Notice that this action is nothing more than a parametrisation invariant way to write the length of the worldline in spacetime. It is intuitively clear that the equivalent of a worldline of a particle for a string would be some surface, called a worldsheet, so one might expect that the action for a relativistic

†Classical here is meant to mean not quantum.

‡This and the next section are inspired by [5] an insightful lecture series.

string would be the parametrisation invariant area of this worldsheet. The area spanned by two vectors in Euclidian space is given by

$$dA = |dv_1 \times dv_2| = \sqrt{|dv_1|^2|dv_2|^2 - |dv_1 \cdot dv_2|^2} \tag{2.7}$$

For an infinitesimal area in parameter space spanned by τ, σ the physical space equivalent is then given by

$$dA = d\sigma d\tau \sqrt{(\frac{dx^a}{d\sigma}\frac{dx_a}{d\sigma})(\frac{dx^b}{d\tau}\frac{dx_b}{d\tau}) - (\frac{dx^a}{d\tau}\frac{dx_a}{d\sigma})^2} \tag{2.8}$$

This is however assuming physical space is Euclidean. To transform this equation to one for spacetime with metric $(-,+,+,+)$ we need to make a Wick rotation resulting in the terms in the square root changing sign. This results in the Nambu-Goto action

$$S = -\frac{T_0}{c}\int_{\tau_i}^{\tau_f} d\tau \int d\sigma \sqrt{(\frac{dx^a}{d\tau}\frac{dx_a}{d\sigma})^2 - (\frac{dx^a}{d\sigma}\frac{dx_a}{d\sigma})(\frac{dx^b}{d\tau}\frac{dx_b}{d\tau})} \tag{2.9}$$

Where T_0 has units of tension, force per unit length, and can be understood on dimensional grounds. A surprising observation at this point is that the rest energy of the string is given by $T_0 l$ where l is the length of the string, thus that the mass per unit length μ_0 is determent by the strings tension. This is rather different than the classical string like one might find on a guitar. In the classical case a string not put under tension has a mass determined by the material the string is made out of. But in the case of the relativistic string the mass of the string is given completely by the energy needed to pull it to it's length as per $E = mc^2$.
The Nambu-Goto action does not lend itself well to quantization although it is possible using the light-cone gauge. The action which is more often used when one wants to quantize the string is the Polyakov action.

$$S = \frac{T}{2}\int d^2\sigma \sqrt{-h} h^{ab} g_{\mu\nu}(X)\partial_a X^\mu(\sigma)\partial_b X^\nu(\sigma) \tag{2.10}$$

Where T is the string tension, $g_{\mu\nu}$ the metric on the target manifold, (the manifold given by the embedding of the worldsheet in spacetime), and h_{ab} is the metric on the worldsheet. This action is in the classical sense equivalent to the Nambu-Goto action, as we can recover, up to reparamaterisation, the action (2.9) from (2.10) by stabilising the later with respect to the worldsheet metric, h_{ab}. Because of this any physical solution leaves both stationary.

2.2.2 Boundary conditions and D-branes

We'll now examine what kind of solutions the action (2.9) allows. For notational convenience we'll use the canonical momenta,

$$\mathcal{P}^\tau_\mu = \frac{\partial \mathcal{L}}{\partial(\frac{dx^\mu}{d\tau})} = \frac{(\frac{dx^\nu}{d\tau}\frac{dx_\nu}{d\sigma})\frac{dx_\mu}{d\sigma} - (\frac{dx^\nu}{d\sigma}\frac{dx_\nu}{d\sigma})\frac{dx_\mu}{d\tau}}{\sqrt{(\frac{dx^\mu}{d\tau}\frac{dx_\mu}{d\sigma})^2 - (\frac{dx^\mu}{d\sigma}\frac{dx_\mu}{d\sigma})(\frac{dx^\nu}{d\tau}\frac{dx_\nu}{d\tau})}} \tag{2.11a}$$

$$\mathcal{P}^\sigma_\mu = \frac{\partial \mathcal{L}}{\partial(\frac{dx^\mu}{d\sigma})} = \frac{(\frac{dx^\nu}{d\tau}\frac{dx_\nu}{d\sigma})\frac{dx_\mu}{d\tau} - (\frac{dx^\nu}{d\tau}\frac{dx_\nu}{d\tau})\frac{dx_\mu}{d\sigma}}{\sqrt{(\frac{dx^\mu}{d\tau}\frac{dx_\mu}{d\sigma})^2 - (\frac{dx^\mu}{d\sigma}\frac{dx_\mu}{d\sigma})(\frac{dx^\nu}{d\tau}\frac{dx_\nu}{d\tau})}} \tag{2.11b}$$

With this notation the variation of (2.9) becomes

$$\delta S = \int d\tau \mathcal{P}^\sigma_\mu \delta x^\mu |_0^{\sigma_1} - \int d\tau d\sigma \delta x^\mu (\frac{\partial \mathcal{P}^\tau_\mu}{\partial \tau} + \frac{\partial \mathcal{P}^\sigma_\mu}{\partial \sigma}) \tag{2.12}$$

Requireing the term $(\frac{\partial \mathcal{P}^\tau_\mu}{\partial \tau} + \frac{\partial \mathcal{P}^\sigma_\mu}{\partial \sigma})$ to vanish we recover the equations of motion for the string as usuall.

$$(\frac{\partial \mathcal{P}^\tau_\mu}{\partial \tau} + \frac{\partial \mathcal{P}^\sigma_\mu}{\partial \sigma}) = 0 \tag{2.13}$$

For the full variation of the action to vanish the first term of (2.12) which we recognise as the boundary conditions must also vanish. In our notation the boundary values for σ are 0 and σ_1. Boundary terms like these are in quantum field theories usually taken to vanish by the assumption that the fields themselves vanish at the boundaries, which correspond to space-time infinity. But since in this case the boundaries correspond to the end of the string which need not be located at space-time infinities we cannot simply discard the boundary terms. As we are considering an open string, i.e. the endpoints are not connected, we should treat the different σ values separately. Also our choice of coordinates should not matter, this means we can consider each direction separately. So in each direction we get two conditions one for each endpoint. This leaves us with

$$0 = \mathcal{P}^\sigma_0(\tau, \sigma_*)\delta x^0(\tau, \sigma_*) \tag{2.14a}$$

$$0 = \mathcal{P}^\sigma_i(\tau, \sigma_*)\delta x^i(\tau, \sigma_*) \tag{2.14b}$$

Where we do not sum over i as each one corresponds with a different spacial direction and we have introduced the notation $\sigma_* \in \{0, \sigma_1\}$. So (2.14) represent $2D$ equations, for D the number of dimensions. Since setting

$\delta x^0(\tau, \sigma_*) = 0$ would mean fixing the endpoint in time, which is not physically valid, equation (2.14a) means that $0 = \mathcal{P}_0^\sigma(\tau, \sigma_*)$. This means that the endpoints behave as free particles along the temperal direction. Equation (2.14b) allows for two types of boundary conditions namely

$$0 = \mathcal{P}_i^\sigma(\tau, \sigma_*) \tag{2.15}$$

$$0 = \delta x^i(\tau, \sigma_*) \tag{2.16}$$

which are known as Neumann and Dirichlet boundary conditions respectively. For each endpoint and each direction one of these two types of boundary conditions needs to be satisfied.
The Neumann condition, (2.15), means it's corresponding canonical momentum vanishes. This indicates translational invariance along the i direction. Meaning the endpoint behaves as a free particle along the corresponding spacial direction. The interpretation of the Dirichlet conditions, (2.16), is surprising. This is because if $\delta x^i = 0$ the endpoint is fixed at a certain position in the i direction. Since fixing a spacial position for the endpoints means that transnational invariance is broken. But as we'll discuss in section 2.2.3, transitional invariance need only be satisfied in non compact directions.So this boundary condition can not entirely be ignored. Suppose that an endpoint satisfies the Dirichlet condition in m directions and the Neumann condition it the remaining $D - 1 - m$ directions. We can then consider this endpoint fixed to a surface which spanned the m directions corresponding to the m Dirichlet conditions, such a surface is called a Dirichlet-brane or D-brane for short. These objects are crucial in modern string theory. When one quantizes the theory these branes cease to be rigid objects and become dynamical, and as such these branes carry energy and interact. It is these branes that will play an important role in the model we'll be discussing.

2.2.3 Extra Dimensions

An almost infamous property of String Theory is that in order for the theory to be consistent it needs to be considered in a higher dimensional space than we commonly observe. Exactly how high dimensional depends on the type of string theory we consider. What is called Bosonic string theory needs a total of 26 dimensions while supersymmetric theories generally need only 10 and M-theory needs 11 dimensions. The need for extra dimensions might at first be considered enough basis to abandon the theory entirely as a possible physical theory, however a second look allows for a

surprising possibility. This possibility comes from the idea of compactification. This is the idea that our universe actually is higher dimensional, but that we just haven't noticed this because the "extra" dimensions are extremely small on the order of the Planck scale. Since everyday physics happens at scales much larger we would expect that these additional dimensions are unnoticeable. It is also this reason, the smallness of the extra dimensions, that put much of string theory beyond the scope of modern experiments, as in order for an experiment to probe the length scales required we would need energies well beyond modern day reach.
To get better understanding of this we'll now take a look at the quintessential example of a compactified theory, namely Kaluza-Klein (KK) unification of gauge interactions and gravity. Along the way we'll encounter some important objects and principles which we'll use throughout the rest of this thesis. In this example we'll consider a D dimensional theory with one periodic coordinate. So let's consider a D dimensional field theory where, $D = d + 1$. The crucial element will be that we'll periodically identify one direction, say the d-th, meaning $x^d = x^d + 2\pi R$. This identification does not change the local metric but it will be useful write the part of the metric corresponding to the d-direction explicitly. When writen as such the metric becomes the following:

$$ds^2 = G^D_{MN} dx^M dx^N = G_{\mu\nu} dx^\mu dx^\nu + G_{dd}(dx^d + A_\mu dx^\mu) \tag{2.17}$$

Where $M, N = 0, \ldots, d$, $\mu, \nu = 0, \ldots, d-1$ and $A_\mu = G^{dd} G_{\mu d}$. The invariant length squared is invariant under local coordinate transformations, these include transformations in the periodic d direction of the form $x'^d = x^d + \lambda(x^\mu)$. This can be seen as follows:

$$\begin{aligned} ds^2 &= G_{\mu\nu} dx^\mu dx^\nu + G_{dd}(dx^d + A_\mu dx^\mu) && (2.18)\\ &= G_{\mu\nu} dx^\mu dx^\nu + G_{dd}(dx'^d - \partial_\mu \lambda dx^\mu + A_\mu dx^\mu) && (2.19)\\ &= G_{\mu\nu} dx^\mu dx^\nu + G_{dd}(dx'^d + A'_\mu dx^\mu) && (2.20) \end{aligned}$$

This gives us the transformation law $A'_\mu = A_\mu - \partial_\mu \lambda$. So we see that gauge transformations arise from higher-dimensional coordinate transformations.
Now it will be interesting to look at what happens to a massless scalar field $\varphi(x^M)$; for simplicity we'll assume that $G_{dd} = 1$ for now. The periodicity of x^d means that the momentum in that dimension needs to be discrete, $p_d = \frac{n}{R}, n \in \mathbb{Z}$. Now we can expand φ in terms of a complete set for it's x^d

dependence, which gives the following

$$\varphi(x^d) = \sum_{n=-\infty}^{\infty} \varphi_n(x^\mu) \exp{(in\frac{x^d}{R})} \tag{2.21}$$

We can substitute this in to the D dimensional wave equation and find that:

$$\partial_\mu \partial^\mu \varphi_n(x^\mu) = \frac{n^2}{R^2}\varphi_n(x^\mu) \tag{2.22}$$

We can interpret this with by interpreting each mode with momentum in the periodic direction as a massive mode, with mass square given by $\frac{n^2}{R^2}$ and one massless mode for corresponding to the zero momentum mode. The energy to excite these modes is related to their mass, as per $E = mc^2$, so at low energies the more massive modes cannot be excited and will therefor not be dynamical. If for example we'll only consider energies lower than $E < \frac{1}{R}$ non of the massive modes will be dynamical; we can disregard them. This is common in the kind of theories we'll consider as they'll be almost exclusively the low energy limit theories. For string theories with supersymmetry these kinds of limits are referred to as supergravity theories. Most of this thesis will deal with supergravity theories as that is the regime where calculations are somewhat tractable. We'll see in chapter 4 that precisely how tractable these calculations are, is a source of much debate.

2.2.4 Moduli

In the example above of the Kaluza-Klein compactification theory the parameter R we left unspecified. This parameter determines the size of the periodic, (or compact) dimension and we can take it as having any positive real value. But this will have consequences for our theory, as we saw when we wanted to impose an energy limit on our description the precise value of R determined what modes we would need to consider at any given limit. In principle we can promote R to a dynamical field dependent on the non-compact directions. This would mean that the masses of the modes corresponding to momenta in the compact dimension would change dependent on this field and as such the effective potential of the system could be dependent on this R field. This is an example of a modulus field, which is extremely common in supergravity theories. These moduli fields arise naturally by allowing the parameters defining a system to become dynamical fields. For example if in the Kaluza-Klein model described above we compactified an additional dimension the size of this dimension could also be left dynamical giving an additional modulus field. Then we can also consider rotations in these now two compact directions, which would transform the compact momenta in one direction in to compact momenta in the other direction. This results in the different types of masses (or charges) corresponding to the different directions transforming into one another. So by the addition of just a single modulus field the model quickly grows more complex. These moduli we would like to be fixed. This seems somewhat counter productive firstly promoting these parameters to dynamic fields and then desiring them to be fixed. The motivation behind this however is quite physical, since by allowing our moduli to be dynamical fields we impose minimal assumptions on to our system. In our example above we assume only that R is finite[§]. This way of approaching the problem means that when the moduli are stabilised, what this means we'll see shortly, this is due to the physics and not because we imposed their value beforehand.

As mentioned before moduli can in general appear in the energy of our system. In our example this occurs via the dependency of the mass of the modes being dependent on R. So as nature minimizes energies we expect that, by allowing our moduli to be dynamic, they will feel effectively a potential. So our moduli will take values where this potential is at least locally minimised. If such a minimum is at least meta-stable, in contrast to a rolling potential, we say that our modulus is stabilised at the value corresponding to that minimum. If our moduli are stabilized then we can treat the moduli as having a fixed value, namely that which it has at the minimum, meaning the physics dependent on those moduli does not change. Allowing us to analyse our system.

2.2.5 Anthropic reasoning

A possible reason for the observed value for the cosmological constant is that if the cosmological constant had a different value the universe would not be able to support life like us. So it should not be surprising that we observe it to have approximately this value since we are here to observe it in the first place.

This argumentation rests on the anthropic principle. The basis of the an-

[§]Although we do in principle allow it to run to infinity when considering it as a modulus. But since this significantly changes our physics as the dimension will no longer remain compact, we consider this a different system.

thropic principle is an argument of observation bias. This is best illustrated with an example.
Suppose we want to answer the question why life evolved on earth. A line of argumentation might go as follows. In order for life to be there needs to be something, lets call it matter, to be alive, so life can expect not to find itself in an empty universe. If the matter in the universe was completely isotropically and homogeneously distributed almost nothing would happen so there would be no life. Thus living things can expect to be in a non-homogeneous universe this leads to structure creation. Again the requirement for matter leads to the conclusion that if life exist it would be in these structures not in the voids between. Then the extremes in most of these structures, i.e. in stars or in interplanetary space, are not conducive to formation of life. Thus we expect life to evolve on a planet, with the right conditions. Therefor it is not to surprising that life evolved on earth. This kind of argumentation is almost more philosophical than it is physical. And this is the reason that it causes much debate. One might argue that it does not answer the question, our example might just be interpreted as not explaining why life evolved on earth, but just could be seen as simply asking "Well where else would it have evolved?".
The reason to mention this here is because, in the search for solutions to the cosmological constant problem, this type of reasoning is occasionally used. Particular in context with what is called the string landscape. This is the large number of possible vacua expected to be self consistently describable by supergravity theories. Suppose we find a family of different vacua each with a different cosmological constant. We could then by virtue of anthropic reasoning conclude that if this family contains vacua with our cosmological constant it describes our universe; without needing to specify a mechanism by which a selection out of this family is made. This somewhat relaxes the rigor needed for an acceptable model of our universe. The core of the debate on whether this is a good physical argument rests on this relaxation of rigor. We'll try not to take sides in this debate during this thesis. When anthropic reasoning is used we'll try to point it out leaving for the reader to decide whether the argumentation made is satisfactory.

2.3 Why KKLT?

As discussed at the start of this chapter, a good begin in providing a solution to the cosmological constant problem is to find a quantum gravity description of a de Sitter space. As mentioned before supersymmetric

groundstates, which are the only groundstates which we can with certainty compute, are either flat or anti-de-Sitter. This makes such a description somewhat of a challenge. However constructing a de Sitter is precisely what the model proposed by Kachru Kallosh Linde Trivedi, [4], attempts to provide. It is the first concerted attempt at providing a construction in string theory with a computationally controlled groundstate, which is neither flat nor Anti-de-Sitter (AdS).
Precisely how controlled this construction is has been a matter of debate since it's proposal in 2003.

Chapter 3

KKLT construction

3.1 The construction in a nutshell

In this chapter we'll examine the Kachru, Kallosh, Linde, Trivedi model for constructing a de Sitter space in string theory. This model they proposed in [4] in 2003. The construction method they laid out consists of three main parts. Starting from what is known as the tadpole condition, which comes from F-theory. This we'll explain in section 3.2. F-theory is a string theory formalism which we'll not discuss in this thesis. We'll just assume the result and proceed from there. The construction is done in what is called type IIB string theory. Type IIB theory is a limit of F-theory and therefor the tadpole condition can be translated into IIB theory using a standard method which we'll not derive nor explicitly preform. However the fact that the tadpole condition is translatable to IIB theory will allow us to make some statements in the IIB theory even without going trough the explicit details.
In the IIB framework we'll firstly stabilize all but one moduli. This we discuss in section 3.3. Here we'll use a superpotential and Kähler potential that we take as given to certain order. Then adding some general higher order corrections, which we'll discuss in subsection 3.5, to the resulting potential we'll find that we can stabilise also the remaining modulus. Which we do in section 3.6. We'll see that this results in a anti-de-Sitter space. This first part of the construction is known as moduli stabilisation. During the moduli stabilisation we'll not break supersymmetry.
During the second step we'll add more flux than to our system. Then to balance the tadpole condition we'll need to add an anti-D3-brane to our set up. Such an object breaks supersymmetry.
Lastly we'll consider the effect of having added an anti-brane. And we'll

see that this has the effect of lifting the vacuum to positive value, meaning it results in a de Sitter vacuum. This last step is known as anti-brane uplifting. The resulting vacuum is meta-stable, which as explained earlier we expect all de Sitter vacua to be. So we'll conclude the chapter by a stability analysis of the resulting de Sitter vacuum.

3.2 Tadpole Condition

As mentioned above we'll start from the tadpole condition. In this section we'll give an explanation of what this condition means. But before we get to the condition we'll fist introduce some of the components which appear in the condition, namely the IIB fluxes.
Type IIB theory has what are called two sectors. These two sectors come from the fact that the fermionic part of the worldsheet action, (3.1), can satisfy two different boundary conditions.

$$S_F = \frac{1}{4\pi}\int d^2z(\varphi^\mu\bar{\partial}\varphi_\mu + \tilde{\varphi}^\mu\partial\tilde{\varphi}_\mu) \tag{3.1}$$

Where φ and $\tilde{\varphi}$ are holomorphic and anti-holomorphic anticommuting fields respectively. The two boundary conditions are:

$$\varphi^\mu(w+2\pi) = +\varphi^\mu(w) \tag{3.2a}$$
$$\varphi^\mu(w+2\pi) = -\varphi^\mu(w) \tag{3.2b}$$

With an equivalent set for $\tilde{\varphi}$. These are the Ramond (R), (3.2a), and Neveu-Schwarz (NS), (3.2b), conditions respectively. For an open string there are thus 4 possible boundary conditions. These are R-R, R-NS, NS-R, NS-NS. Due to symmetries the only ones relevant for IIB are the R-R and NS-NS cases. The resulting objects from each of these two can be separately studied. Those coming from the R-R boundary conditions are said to be from the Ramond-Ramond sector and those from the NS-NS boundary conditions are said to be from the Neveu-Schwarz-Neveu-Schwarz sector. When transform from the worldsheet to space-time objects we can still separate them in to their sections. The potentials and field strengths coming from the RR sector we'll indicate with C_p and F_{p+1} respectively. And the potentials and field strengths coming from the NSNS sector we'll indicate with B_p and H_{p+1} respectively.

The space-time IIB action then is the following:

$$S = \frac{1}{2\kappa_{10}^2}\int d^{10}x\sqrt{-g_s}\big(e^{-2\phi}(R+4(\nabla\phi)^2) - \frac{F_{(1)}^2}{2} - \frac{G_{(3)}\cdot\overline{G}_{(3)}}{12} - \frac{\tilde{F}_{(5)}^2}{4\cdot 5!}\big)$$
$$+\frac{1}{8i\kappa_{10}^2}\int e^{\phi}C_{(4)}\wedge G_{(3)}\wedge\overline{G}_{(3)} + S_{\text{loc}} \quad (3.3)$$

Where g_s is the string metric, R the Ricci curvature, ϕ the dilation field, $F_{(i)}$ are i-form fluxes, $C_{(4)}$ is a 4-form potential, $G_{(3)} = F_{(3)} - \tau H_{(3)}$ is the combined three-flux where $\tau = C_{(0)} + ie^{-\phi}$ and S_{loc} is the action of localized objects, such as branes. For more details on type IIB we'll refer the reader to [6].

As mentioned in 3.1 our starting point is an F-theory result called the tadpole condition. This is the following, [6]:

$$\frac{\chi(X)}{24} = N_{D3} + \frac{1}{2\kappa_{10}^2 T_3}\int H_{(3)}\wedge F_{(3)} \quad (3.4)$$

Here N_{D3} is the number of D3-branes minus number of anti-D3-branes present in the system. The reason that they appear here are because these branes couple to the RR sector, or in other words they carry RR charge. The $H_{(3)}, F_{(3)}$ fluxes are as explained above. On the left hand side the $\chi(X)$ is the Euler characteristic of the manifold X. This X is compact space in F-theory.

The way to think about this condition is in the sense of Gauss law. Consider for example a point charge on a sphere in familiar electromagnetism. On a compact space such as a sphere we can not have just a single point charge as the flux lines emanating form this charge would have nowhere to end. Since the branes here can be considered as charges this condition simply balances the charges in the geometry of X. Note that in the absence of branes this condition tells us that there also can't be any fluxes.

As stated in 3.1 IIB theory is a limit of F-theory. A priori the way to transform from F-theory to IIB via this limit only applies to systems without fluxes. This was however extended to also be applicable in the present of fluxes in [7] and [8].

3.3 Superpotential and complex moduli stabilisation

With our starting point clear we'll continue with N=1 supergravity in which we'll construct our vacuum. The standard potential for such a theory is the following:

$$V = e^K\Big(\sum_{a,b} g^{a\bar{b}} D_a W \overline{D_b W} - 3|W|^2\Big) \tag{3.5}$$

D_a is the Kähler covariant derivative and is defined by $D_a = \partial_a + (\partial_a K)$. In (3.5) K is the Kähler potential and W is the superpotential and the summation over a, b runs over all moduli in our system. A superpotential is a generalization of a familiar potential. It arises when considering supersymmetric theories. A Kähler potential is a mathematical object from which we can define a particular kind of manifold, known as a Kähler manifold. The compact manifolds in IIB theory need to be Calabi-Yau manifolds, which are special kinds of Kähler manifolds. So we can understand the Kähler potential as defining our compact space. The superpotential is known as a GVW superpotential, [7], and is of the form:

$$W = \int_M G_3 \wedge \Omega \tag{3.6}$$

M is the IIB compact space, G_3 is the same as $G_{(3)}$ encountered earlier and encodes the degrees of freedom of the fluxes, Ω is a unique holomorphic (3,0) form on the compact space. Ω defines the complex structure on the compact space and is therefore referred to as the complex structure modulus. The Kähler potential follows from the space-time action, (3.3), as derived in [9]. It is of the following form:

$$K = -3\ln\left(-i(\rho - \bar{\rho})\right) - \ln\left(-i(\tau - \bar{\tau})\right) - \ln\left(-i\int_M \Omega \wedge \bar{\Omega}\right) \tag{3.7}$$

Here ρ is a volume modulus and encodes the overall volume of the compact manifold. In general there could be additional Kähler moduli, but through out this thesis we'll only consider scenario's with a single Kähler modulus ρ. τ is called the axio-dilaton and is the same as encountered in section 3.2 and Ω is the unique holomorphic (3,0) form as above. This Kähler potential is a generalisation of the result given in [10] as derived in [9].

How these potentials arise is quite technical and not extremely relevant for us. There are potential corrections to be added to these potentials which

we'll ignore for now. It is however relevant to realise that these potentials are determined by the compactification and are of the given form as a result of the IIB limit we took for the original F-theory we started with and which gave us the tadpole condition. The information of this compactification is encoded in the parameters ρ, τ, Ω. The way these parameters couple to the fluxes is what is described by W and it is this coupling that means we can not simply treat the flux and background manifold separately.
When we substitute these potentials in to (3.5) we find the following

$$V = e^K \Big(\sum_{i,j} g^{i\bar{j}} D_i W \overline{D_j W} \Big). \tag{3.8}$$

Where the summation now runs over all moduli except ρ. This due to the cancellation of the ρ term in the summation with the $-3|W|^2$ term. Which is due to the fact that (3.6) does not depend on ρ meaning that $\partial_\rho W = 0$. This means that

$$g^{\rho\bar{\rho}} D_\rho W \overline{D_\rho W} = 3|W|^2 \tag{3.9}$$

Where we used that $g^{\rho\bar{\rho}} = (\partial_\rho \partial_{\bar{\rho}} K)^{-1}$ by definition of the Kähler metric $g^{i\bar{j}}$.

Thus far we have not made a choice for our fluxes, so let's do that. We'll take $F_3, H_3 \in H^3(M, \mathbb{Z})$, meaning that they are 3-form fluxes of the form such that when integrated over our manifold, M, they evaluate to integers as demanded by Dirac quantization. This choice forces G_3 to be Imaginary Self-Dual (ISD). This follows from the fact that for all objects we consider (3.10) holds, as explained in [9].

$$\frac{1}{4}(T^m_m - T^\mu_\mu)^{\text{loc}} \geq T_3 \rho_3^{\text{loc}} \tag{3.10}$$

Where T_{ab} is the stress energy tensor (considered for localised objects); the index μ runs over de compact directions and m over the non-compact directions. T_3 is the tension of D3-branes and ρ_3^{loc} is the D3 charge density from localised sources.
As discussed in 2.2.4 any modulus which enters in to the potential feels an effective potential in the presents of fluxes. So when fixing our fluxes we essentially create an effective potential for the moduli. As familiar from quantum field theories an effective potential can be viewed as generating a mass for fields. The same applies to our moduli. We recall that the

volume modulus term in the summation of (3.5) cancelled as discussed earlier. This results in all other moduli receiving a mass of the scale:

$$m \sim \frac{\alpha'}{R^3} \tag{3.11}$$

Where R is the characteristic radius of manifold, which relates to ρ. We remark that $\mathrm{Im}(\rho) \propto R^4$. So from this we conclude that the masses of all moduli are large when compared to the mass of ρ. So we can consider the low energy limit and in doing so treat all other moduli as fixed and only ρ as dynamical.
As a side note, the dependence of the masses of the fixed moduli on our choice of fluxes means that they are discretely tunable, as our flux choice is discrete.

3.4 Warping

One of the more subtle and crucial aspects of the KKLT model is that the separation of scales, the fact that we can consider the moduli other than the volume modulus fixed, depends on the warping in the system. One can think about this warping as a similar phenomenon as the warping in spacetime caused by massive objects familiar from GR. D-branes and fluxes are sources of warping and so we expect our system to exhibit warping. This allows us to write the metric of our system in Einstein frame as follows:

$$ds_{10}^2 = e^{2A(y)}\eta_{\mu\nu}dx^\mu dx^\nu + e^{-2A(y)}\tilde{g}_{mn}(y)dy^m dy^n \tag{3.12}$$

Where x coordinates along the D3 branes and y coordinates in the orthogonal direction. Since the D3 branes span the non-compact directions, as demanded by Poincare symmetry, this means that the x coordinates correspond with the non-compact directions and the y coordinates correspond with the compact coordinates.$\eta_{\mu\nu}(x)$ is the non-compact unwarped metric and $\tilde{g}_{mn}(y)$ is the compact unwarped metric on M. This warping parameterizes energy scales as it varies over the compact dimensions. Generally this is thought about as a throat expanding from the compact manifold where, as we descend in to the throat the warping increases, meaning our energy gets scaled down. This scaling should be read as scaled in relation to the value outside the throat. Although it is not exactly the same phenomena, one can again think about this in a similar manner as the red and blue shifting photons experience as they climb out resp. fall in to a potential well. The effect of objects high in the throat (more towards the unwarped compact space) are observed to have a larger energy than the same object when it sits lower in the throat, as seen from a specific location. The precise size of this warping is in general dependent on the objects in our system. One of the major results of [9] is that there it shown that using integral fluxes we find exponential warping as given by

$$e^{A_{\min}} \sim e^{-\frac{2\pi K}{3 g_s M}} \tag{3.13}$$

where K, M are called the flux quanta. The value of K, M is given by integrating the H_3, F_3 over specific subspaces of the compact manifold, see [9] for more details.
This warping is crucial for the KKLT model. As we'll see later the addition of a anti-brane without warping would have a significant effect on the shape of the potential. This effect, known as back-reaction, would make calculations in tractable. So in order to have some hope of finding a de Sitter vacuum we need exponential warping.

3.5 Corrections

In this section we'll discuss some known corrections to the no-scale potential we found in section 3.3. In order to construct a vacuum all moduli need to be stabilized. As we have seen the no-scale potential results in stabilization of the complex structure moduli as they enter in to the superpotential. But this leaves the Kähler moduli, in our case just one, ρ, unfixed. In order to have stabilization of all moduli a term needs to be added to the potential which couples to the Káhler moduli. Such a term can not be perturbative as supersymmetry forbids such terms. However there are possible non-perturbative corrections we could consider.
The precise origin of these corrections is not extremely relevant for our conclusion. As such we'll not derive these corrections in detail. An example of such a correction can be found in [11] where as a part of a broader discussion on corrections they derive a correction arising from a magnetically charged instanton.
An other example is given in [4], coming from gaugino condensation on D7-branes*.
Correction term like these generally depend on both the complex structure as well as the Kähler moduli. But since we fixed our complex structure moduli we can consider them just dependent on the Kähler modulus. This

*For details see [4] and references therein however crucially depends on the fact that the scale at which the complex structure moduli are fixed is sufficiently high. Some of the debate around KKLT revolves around this hierarchy of scales. Again the fact that we need these corrections to stabilize the Kähler modulus is crucial to the model. As both example above give a term to the super potential of the form:

$$W_{\text{correction}} = Ae^{ia\rho} \tag{3.14}$$

Where A, a parameters dependent on the details of the source of the correction. We'll assume our corrections are of this form although in general other terms could exist. An additional source of debate is the fact that in the examples above a is not a continuous parameter as it depends on the choice of fluxes, for which can only make discrete choices. However we'll treat it as continuous in the next section.

3.6 Stabilizing the volume modulus

In this section we'll consider the stabilization's we have made thus far and find that they lead to an Anti-de-Sitter vacuum. Recall that as we discussed in section 3.3, prior to adding the corrections discussed in section 3.5, we stabilized all moduli except the Kähler moduli, of which we consider there to be only one, the volume modulus. These moduli where stabilised when we fixed our flux configuration as they entered in to the superpotential. Since we have stabilized all but the volume modulus the only relevant term from the Kähler potential is

$$K = -3\ln\left(-i(\rho - \bar{\rho})\right). \tag{3.15}$$

The superpotential with the added correction term is the following:

$$W = W_0 + Ae^{ia\rho} \tag{3.16}$$

where W_0 is a tree level contribution due to the flux as coming from (3.6). The second term we assumed as the general form of corrections as we discussed in the previous section.
We'll consider the tadpole condition (3.4) to be solved by turning on only flux, which for now means that no additional D3-branes are present. For a supersymmetric vacuum it holds that $D_\rho W = 0$ where again $D_\rho = \partial_\rho + (\partial_\rho K)$. We'll write $\rho = \tau + i\sigma$. Then we find that with our Kähler and

Superpotential we get:

$$K = -3\ln\left(-i(\rho - \bar{\rho})\right)$$
$$= -3\ln(2\sigma) \tag{3.17}$$

So the real dependence drops out of the Kähler potential. And the Supersymmetry restriction gives us:

$$\begin{aligned}
0 &= D_\rho(W_0 + Ae^{ia\rho}) \\
&= D_\rho(W_0 + Ae^{ia(\tau+i\sigma)}) \\
&= \partial_\rho(W_0 + Ae^{ia(\tau+i\sigma)}) + (\partial_\rho K)(W_0 + Ae^{ia(\tau+i\sigma)}) \\
&= aAie^{ia(\tau+i\sigma)} + \frac{3i}{2}\frac{1}{\sigma}(W_0 + Ae^{ia(\tau+i\sigma)}) \\
W_0 &= -Ae^{-a\sigma}e^{ia\tau}(\frac{2}{3}a\sigma + 1)
\end{aligned} \tag{3.18}$$

Recall that the potential is given by:

$$V = e^K(G^{\rho\bar{\rho}}D_\rho W\overline{D_\rho W} - 3|W|^2) \tag{3.19}$$

Which means that with our superpotential and Kähler potential, as in (3.16) and (3.17), we find that the potential has a minimum at:

$$V_{Ads} = -3e^K W^2 \tag{3.20}$$
$$= -\frac{a^2A^2e^{2ia(\tau+i\sigma)}}{(6\sigma)} \tag{3.21}$$

The effect resulting from the imaginary part of ρ, namely σ, is the addition of a non-perturbative term to the superpotential as discussed in section 3.5.
In the original KKLT paper, [4], the axion part of the volume modulus was set to zero, $\tau = 0$, this results in the same minimum except without the factor $e^{2ia\tau}$. This has the effect of simply considering only 1 of the degenerate vacua. This degeneracy comes from the the fact that $\text{Re}(e^{2ia\tau})$ behaves as a cosine where each minimum corresponds with a different discrete vacuum. Since τ enters into the superpotential via G_3 as we saw earlier it is fixed by the choice of fluxes. This means that our choice of fluxes already break this degeneracy. This means that here we can treat the $e^{2ia\tau}$ term simply as a prefactor.
As mentioned prior the Kähler potential we use in principle receives higher

order corrections which we neglect. In order for these corrections to indeed be negligible we assumed that $\sigma >> 1$, otherwise we'd expect that the higher order terms need to be taken into account. We also require that $a\sigma > 1$ in order for contributions to the superpotential originating from the instanton to be accounted for properly in the way we have. We'll assume that this is achievable by tuning the fluxes such that $W_0 << 1$ which we can see from (3.18) has the desired result.
At this point we have stabilized all our moduli. This we did using the assumption that the masses the complex structure moduli receive is large when compared with the mass of the Kähler moduli, of which we assumed there was only 1. Thus far we have not broken supersymmetry and there for any vacuum will be supersymmetric. In figure 3.1 we see our potential plotted against σ. From both figure 3.1 as well as from (3.21) we find that this minimum is an Anti-de-Sitter one as expected.

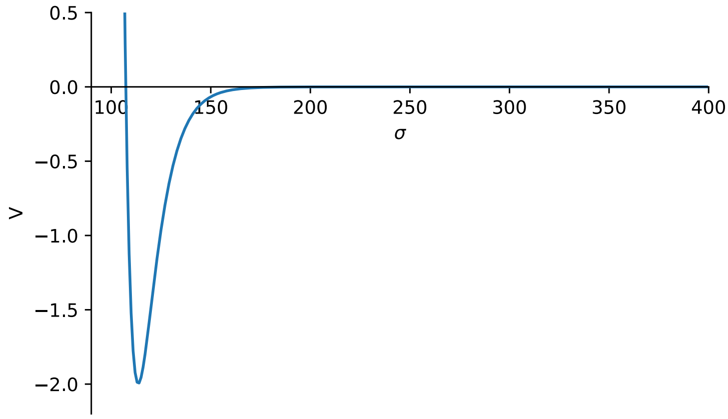

Figure 3.1: *Here we see* $V(\times 10^{15})$ *plotted for* $A = 1, a = 0.1, W_0 = -10^{-4}$

3.7 Constructing dS vacua

In the previous section we found that by only adding flux and demanding our vacuum to be supersymmetric our potential exhibited a Anti-de-Sitter minimum. Now in order to find a de Sitter vacuum we'll add additional components to our setup in order to provide what is called the uplift. This addition will break supersymmetry. This is crucial for recovering a de Sitter vacuum as these are not supersymmetric. The validity of this and the following analysis is one of the major point of debate concerning KKLT, we'll discuss some of this in the next chapter. For now we'll follow the original paper, [4], in order to compare with their results.

The uplift components we'll be adding are $\overline{D3}-$branes, which we introduce as countering additional flux in order to satisfy (3.4). These $\overline{D3}-$branes add a term to the potential proportional to $\frac{1}{\sigma^3}$† each. For the origin of this term we refer the reader to [12]. So adding multiple $\overline{D3}-$branes means we need to add $\frac{D}{\sigma^3}$ to our potential, where D depends on the number of $\overline{D3}-$branes. The factor D depends also on the warp factor at the position of the $\overline{D3}-$branes, this ensures that in a warped throat the additional term is small. In principle higher order terms would appear but these scale quadratically in D, which is exponentially suppressed due to the warp factor, so we can ignore these.

Here we assume that the $\overline{D3}-$branes do not disturb our setup to much. It is this assumption that is at the hart of much of the debates surrounding this model, which again we'll discuss in the next chapter. So assuming simply adding this term to our potential, implying that our Kähler potential and superpotential remain effectively unchanged (as side from the additional term in the superpotential) means our potential becomes:

$$\begin{aligned} V &= e^K(G^{\rho\bar{\rho}}D_\rho W\overline{D_\rho W} - 3|W|^2) + \frac{D}{\text{Im}(\rho)^3} \\ &= \frac{aAe^{-a\sigma}}{2\sigma^2}(\frac{1}{3}\sigma Aae^{-a\sigma} + W_0 + Ae^{-a\sigma}) + \frac{D}{\sigma^3} \end{aligned} \tag{3.22}$$

Where we have set $\tau = 0$ for ease of calculation as this just represents the same degeneracy we discussed in the last section. This potential has again a minimum around the same value as the not uplifted potential. For $D \geq \frac{a^2A^2\sigma^2e^{-2a\sigma}}{6}$ the minimum is no longer negative. This minimum however is now no longer global as it is positive and the potential goes to 0 as σ goes to infinity.

†Later papers use $\frac{1}{\sigma^2}$ after canceling the warp factor dependence on σ in the numerator with a σ term in the denumerator.

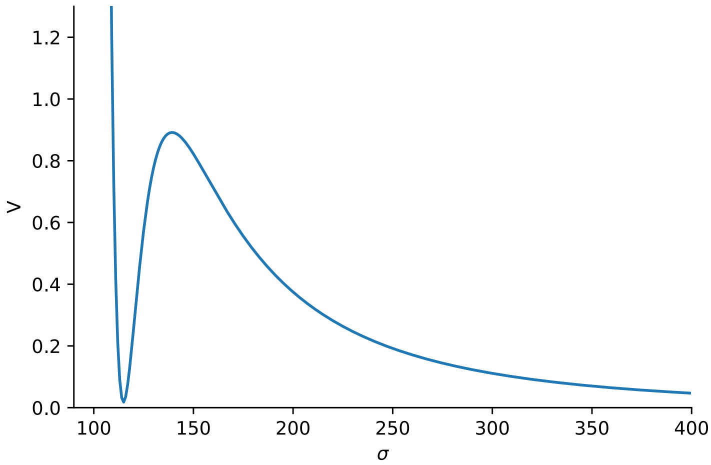

Figure 3.2: *Here we see $V(\times 10^{15})$ plotted for $A = 1, a = 0.1, W_0 = -10^{-4}, D = 3 \times 10^{-9}$.*

This new minimum, being positive, corresponds with a de Sitter vacuum. The value of the potential at this minimum is dependent on the fluxes and number of added $\overline{D3}-$branes and therefor it should be discretely tunable.

3.8 Stability of dS vacuum, Original Considerations

The uplifted vacuum we found in the last section is as mentioned not a global minimum. This means that we should expect this vacuum to be unstable, particularly this de Sitter vacuum should decay into a Minkowski vacuum in the decompactification limit of $\sigma \to \infty$. But this is not necessarily a problem for the physical relevancy of the model for if the vacuum is stable enough we may use the model as an effective description of our universe. By stable enough we mean that the time scale of the decay is larger than the expected age of the universe (so larger than 10^{10} year). On the other hand we'll also consider an upper bound as argued for in [4]. "If the decay time is longer than $t_r \sim e^{S_0}$, one may need to address the issues about the consistency of the stringy description of de Sitter space ...", [4]. Here t_r is the recurrence time and $S_0 = \frac{24\pi^2}{V_0}$ the de Sitter enthropy. This gives us a lower and upper bound on the decay time of the de Sitter vacuum. Basic tunneling decay of meta-stable states including gravitational

effects are discussed in [13] and along these lines we'll analyse the stability of our de Sitter vacuum.
We can think about this decay as a bubble of spacetime tunneling from the local de Sitter minimum in the potential to the lower value of the potential at higher σ. The probability of this occurring is related to the difference in the potential between the two configurations as well as the size of the bubble. The bubble after forming has a different vacuum than the surrounding spacetime. The vacuum of the bubble is lower in energy than the surrounding vacuum which will decay in to the lower energy vacuum. The boundary of the bubble mediates between these two vacua, which we can think about as two different phases. The mediation between two phases can be described using an instanton which is what we'll use here as well.
It is convenient for our description to switch notation to $\varphi = \sqrt{\frac{3}{2}}\ln(\sigma)$. Just as in [13] we'll look at the probability of decay per volume, $P = \frac{\Gamma}{V}$, which in the semi-classical limit they find to be approximately:

$$P = Ae^{-\frac{B}{\hbar}}(1 + \mathcal{O}(\hbar)) \approx e^{-S(\varphi)+S_0} \tag{3.23}$$

Where $S_0 = S(\varphi_0)$ is the action for the initial configuration, in our case the action at the de Sitter vacuum and $S(\varphi)$ is the Euclidean action for the tunneling trajectory. We consider spherically symmetric decay, that is to say that we assume the volume which decays is spherically symmetric. This means we can use the general $O(4)$ invariant Euclidean metric:

$$ds^2 = d\xi^2 + \rho(\xi)^2 d\Sigma^2 \tag{3.24}$$

Where ξ is the Euclidean time coordinate, $\rho(\xi)$ the scaling factor and $d\Sigma^2$ the metric of a 3-sphere. The Euclidean action is given by

$$S = \int d^4x\sqrt{g}\left(\frac{\partial_\alpha\varphi\partial^\alpha\varphi}{2} + V(\varphi) - \frac{R}{2}\right) \tag{3.25}$$

This action leads to the following EOM:

$$\varphi'' + 3\frac{\rho'}{\rho}\varphi' = \frac{dV(\varphi)}{d\varphi} \tag{3.26}$$

$$-\frac{\rho}{3}(\varphi'^2 + V(\varphi)) = \rho'' \tag{3.27}$$

Combining these with the Ricci scalar of the metric

$$R = \frac{\partial_\alpha \varphi \partial^\alpha \varphi}{2} + 4V(\varphi) \tag{3.28}$$

means that the action (3.25) simplifies to

$$S(\varphi) = -\int d^4x\sqrt{g}\big(V(\varphi)\big) = -2\pi^2 \int_0^{\xi_f} d\xi \rho(\xi)^3 V(\varphi(\xi)). \tag{3.29}$$

The initial state before the spacetime bubble forms corresponds with the limit that $\xi_f \to 0$. This equal to minus the entropy of our initial de Sitter congfiguration as formulated in, [14] and [15]. This gives us that S_0 is given by:

$$S_0 = -\boldsymbol{S_0} = \frac{24\pi^2}{V(\varphi_0)} \tag{3.30}$$

Where $\boldsymbol{S_0}$ is the de Sitter space entropy.

Thin-wall approximation

The decay of our de Sitter vacuum to Minkowski corresponds with a special case discussed in [13] of a positive valued false vacuum decaying to a vacuum with zero potential energy in the thin-wall approximation. There for we can use their result for the decay probability which when rewritten becomes

$$P_{\text{tunnel}} = \exp\big(-\frac{\boldsymbol{S_0}}{(1+(\frac{4V_0}{3T^2}))^2}\big) \tag{3.31}$$

Where $T = \int_{\varphi_0}^{\infty} d\varphi \sqrt{2V(\varphi)}$ is the tension of the bubble wall, equivalent to S_1 from [13]. For our model we can consider the case that $V_0 << T^2$, which allows us to reasonably expand (3.31) in terms of $\frac{V_0}{T^2}$, to zeroth order this gives us a familiar result

$$P_{\text{tunnel}} \approx \exp\left(-\boldsymbol{S_0}\right) \tag{3.32}$$

Which for realistic values of $V_0 \sim 10^{-120}$ in Planck units, means that the decay time becomes of the order of $t_{\text{decay}} \sim \exp(122)$ satisfying the lower bound. If we consider first order we get

$$P_{\text{tunnel}} \approx \exp\left(-\boldsymbol{S_0}\right) \exp\big(\frac{64\pi^2}{T^2}\big) \tag{3.33}$$

This means that the decay time becomes

$$t_{\text{decay}} = t_r \exp -(\frac{64\pi^2}{T^2}). \tag{3.34}$$

Since $\frac{64\pi^2}{T^2} > 0$ the decay time is exponentially smaller than the recurrence time. This means that our meta-stable vacuum satisfies the upper bound.

In this analysis we just considered decay via instantons as considered in [13], consideration of other instantons provide similar results leading to the conclusion that our meta-stable vacuum can be effectively considered stable.

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Chapter 4

Points of critique on KKLT

As mentioned before the Kachru, Kallosh, Linde Trivedi construction, which we examined in detail in the previous chapter, was proposed in 2003. Since it's introduction it has been one of the bases for the attempts as formulating a de Sitter space in string theory. Other methods exist, such as the Large Volume Scenario, [16], but are beyond the scope of this thesis. It then surprising that after almost 20 year there is still a large ongoing debate on whether the solutions proposed by the model are reliable or even exist at all, [17].
Ever since it's original proposal there have been critiques launched at the model ranging from minute technical details to conflicts with no-go theorems, supposedly showing that de Sitter vacua are not possible in string theory.
In this chapter we'll consider a few of these critiques. In sections 4.1 and 4.2 we go into some more detail for two discussions revolving around the probe-approximation. In the remaining sections we'll be more brief, only giving a broad outline. This chapter is not intended to be an exhaustive list of all critiques raised over the last two decades. Rather it is intended as a starting point, for the interested reader, outlining some of the more prominent critiques.

4.1 Flattening of the potential due to backreaction

The uplift in KKLT from the anti-de-Sitter vacuum to a meta-stable de Sitter vacuum needs to preserve the stability of the minimum. This is, as we mentioned in 3.7, the essence of the probe approximation. In [18] and [19]

the validity of this approximation is examined.
In principle when we add a supersymmetry breaking object to our anti-de-Sitter solution we should also account for any interactions between this object and our setup, not simply add it's energy to the potential. We however ignored these interactions in the probe-approximation.
In [19] they give sufficient criterion for the validity of ignoring these interactions, namely that the lightest scalar mass times the cosmological constant is much larger than unity (relative to the KK scale). This is simply stating that the energy added by the supersymmetry breaking effect, in our case the $\overline{D3}-$brane does not have enough energy to excite the moduli. This is necessary for us to consider the moduli as remaining fixed. Which in the notation, we used in chapter 3, this is the following statement:

$$\frac{m_\rho^2}{V_{\text{AdS}}^2} \approx 4a^2\sigma_{cr} >> 1 \tag{4.1}$$

Where we introduce m_ρ^2 as the mass of the volume modulus squared, which is approximately $m_\rho^2 \approx \frac{a^4|A|^2 e^{-2a\sigma_{cr}}\sigma_{cr}}{9}$ and σ_{cr} to indicate the value of σ at the potential minimum. This condition is not parametrically fulfilled in KKLT models as $a\sigma_{cr}$ is not arbitrarily tunable.
So a more in depth analysis of the uplift is needed. For this they use a nilpotent description. This rests on the idea of adding a nilpotent chiral superfield, S, such that $S^2 = 0$, to our set up which will play the roll of an uplift term. Then they proceed to calculate the potential in the same manner as before and at the end we put $S = 0$. It is not clear this description is accurate, see [20–22] for a discussion on this. But assuming the approach is valid we can say some thing about the uplift. Adding the nilpotent chiral superfield, S, to our description the superpotential and Kähler potential take the form

$$W = W_0 + \mathcal{A}e^{ia\rho} + e^{2A}\sqrt{24\mu_3}S \tag{4.2}$$

$$K = -3\log\left(2\text{Im}(\rho) - S\overline{S}\right) \tag{4.3}$$

Where again S is a degree 2 nilpotent superfield, i.e. $S^2 = 0$. Then calculating the potential after the uplift, then setting $S = 0$ the potential takes the naive form of

$$V_{dS} = V_{AdS} + V_{\text{uplift}} \tag{4.4}$$

This however excludes a term which arises from setting $\mathcal{A} \to \mathcal{A}(1 + cS)$, since this is not forbidden by symmetries we should include it in our EFT description. This term results in an additional term in the potential of the form:

$$V_{\text{corrections}} = \frac{e^{-a\sigma}}{12\sigma^2}(2\sqrt{24\mu_3}\text{Re}(\mathcal{A}ce^{ia\theta}) + |\mathcal{A}c|^2 e^{-a\sigma}) \tag{4.5}$$

It is clear that if c is small enough this correction term becomes negligible so lets consider where this parameter comes form. It arises from the compactification of the supersymmetry breaking object. In case of KKLT this is a $\overline{D3}-$brane in a warped throat, so naively we would expect c to be warped down meaning it indeed would be small enough. There are however a priori possibilities where this warping down does not occur, for example a gaugino condensate which lives outside the throat would have a back-reaction effect which gets blue shifted when considering it's effect at the tip of the throat. This back-reaction could there for mess with the warping down of c meaning the correction term to the potential is not negligible. In that case $V_{\text{corrections}}$ would result in a runaway potential and we would not have a meta-stable dS vacuum after the uplift, which is obviously would be a problem for the validity of the KKLT construction.

4.2 Conifold instability

Again the method of uplifing in KKLT works only if the addition of the supersymmetry breaking effect doesn't spoil the anti-de-Sitter vacuum. By this we mean that even though the minimum should get lifted it should remain a sufficiently deep minimum. A general $\overline{D3}-$brane placed anywhere in the compact manifold will not satisfy this condition, it's contribution to the potential is to large. This we discussed in 3.7 where we argued that the constant D was small due to the warp factor adding exponential suppression, without warping it would be of the order of the brane tension times the compact volume. This would violate our low energy approach, needed to consider the complex structure moduli as fixed. Therefor these anti-branes are generally placed in a region of large warping, the throat. This is not a untoward assumption as the brane is able to move and is attracted to regions with high warping, so we would expect a brane located not in the throat to fall in to the throat given enough time. Since currently no explicit construction of a compact warped manifold exist one generally assumes the throat to be glued to a flux compactification. In calculations this compact part is then integrated over and treated basically as a boundary condition for considerations made in the throat. The validity of this approach is difficult to assess as due to the lack of explicit construction there is no clear test to verify the found results. This is one of the sources

of debate on the construction. In this section we'll examine one of the arguments made in this line of thought.
Originally posed in [23], the conifold destabilisation mechanism suggests that the correct way to describe an anti-D3-brane in the throat should take in to account the conifold deformation parameter. This conifold deformation parameter can be thought of as the size of the warped throat. That this needs to be taken into account follows from a calculation done in [24] where they show that the conifold deformation parameter, which we'll denote S, is actually lighter, comparable in mass to the volume modulus, than argued for in the original KKLT paper, where it was argued to of order $\frac{1}{R^3}$ same as the complex structure moduli.
This means we can not simply integrate it out and need to treat it as a dynamical modulus. This treatment gives rise to the following potential:

$$V_{ks} = \frac{\pi^{\frac{3}{2}}}{\kappa_{10}} \frac{g_s}{(\mathrm{Im}\rho)^3} \left[c \log(\frac{\Lambda_0^3}{|S|}) + c' \frac{g_s(\alpha' M)^2)}{|S|^{\frac{4}{3}}}\right]^{-1} \left|\frac{M}{2\pi i} \log(\frac{\Lambda_0^3}{|S|}) + \frac{iK}{g_s}\right|^2 \quad (4.6)$$

Where c is a constant coming from the warp factor at UV and will be small $c << 1$, c' is a order 1 constant dependent on the warp factor. M, K are the flux quanta and Λ_0 is the UV cutoff which corresponds with the limit where the throat is glued to the flux compactification.
From this potential we can easily see that in the limit where $S \to 0$ this potential vanishes which means there is an additional minimum at 0. This minimum is not taken in to account when one assumes that S is fixed by being heavy. In order for S not to go to zero the potential barrier between the non-zero minimum and zero should be large enough, as the entire contribution when combined with the term coming from the $\overline{D3}-$branes flattens this barrier. What this means is examined in [24] in the limit where

$$\log(\frac{\Lambda_0^3}{|S|}) << \frac{g_s(\alpha' M)^2)}{|S|^{\frac{4}{3}}} \quad (4.7)$$

This limit corresponds with a large throat. Then by looking at the values where the non-zero minimum becomes an inflection point they find in terms of N the number of $\overline{D3}-$branes the flux quanta int the throat should satisfy:

$$\sqrt{g_s} M > M_{min} \quad M_{min} \propto \sqrt{N} \quad (4.8)$$

Where we explicitly keep the flux quanta as they define the number of anti-branes via the tadpole condition. This forms a lower bound on the flux quanta. The tadpole condition provides an upper bound, namely

$$MK \leq |Q_3^{\text{loc}}| \tag{4.9}$$

where Q_3^{loc} is the total D3 brane charge from local sources. When we then consider only a single $\overline{D3}-$brane and 64 $O3-$planes these bounds gives us a value for the total hierarchy between the UV cutoff and the IR to be

$$\frac{\Lambda_{IR}}{\Lambda_{UV}} = \exp\left(\frac{2\pi K}{2g_s M}\right) > 0.2 \tag{4.10}$$

Which would exclude de Sitter vacua with sufficient hierarchy. A necessary remark at this point concerning the choice for a single $\overline{D3}-$brane and 64 $O3-$planes is the following. Multiple $\overline{D3}-$branes have been argued to give additional contributions to the potential by their internal interactions which would already spoil the vacuum without accounting for the deformation parameter. And the choice for 64 $O3-$planes is some what arbitrary as there exist no known upper limit for the amount of $O3-$planes, but most examples in the literature exhibit around this number of $O3-$planes.

4.3 IIB backgrounds with de Sitter space and time-independent internal manifold are part of the swampland

A term common in the literature of string vacua is the swampland. This is the idea that some of the vacua that one can construct in low energy theories, such as supergravities, are not vacua of the full string theory. Similar as how in effective field theories we sometimes find problems when considering UV energy scales. There are general arguments made that certain types of vacua are in the swampland, i.e. they are not really vacua of the full theory. These are generally referred to as swampland conjectures.
Swampland conjectures are mostly based on calculations in regimes where top-down calculations can be made. There for it is possible that these conjectures miss out on phenomena which only are possible in other regimes. For example anti-brane back reaction on the world volume which leads to the the uplift in KKLT. So a closer look at possible quantum corrections appears to be necessary to expand or circumvent these conjectures. This is precisely what is investigated in [25]. This paper gives an extensive anal-

ysis of different possible corrections applicable to the construction of de Sitter vacua in IIB theory. We'll not be able to discuss in depth the entire analysis, so we'll just mention the results relevant to our discussion. The first of these is that in [25] they find that a IIB background with 4d de Sitter isometries, a time independent 6 dimensional internal manifold, with a metric of the form (4.11), and time independent fluxes are necessarily in the swampland.

$$ds^2 = \frac{1}{\Lambda(t)\sqrt{h}}(-dt^2 + dx_1^2 + dx_2^2 + dx_3^2) + \sqrt{h}g_{mn}(y)dy^m dy^n \tag{4.11}$$

These kind of solutions would include the original version of KKLT as we described it in Chapter 3. However a further result of [25] suggest that certain alterations to KKLT might provide a de Sitter space. This result is that they find that when one allows for the fluxes and internal manifold to be time dependent the resulting theory need not be part of the swampland as previous considerations no longer holds. Such a time dependent setup would have a metric of the form

$$ds^2 = \frac{1}{\Lambda(t)\sqrt{h}}(-dt^2 + dx_1^2 + dx_2^2 + dx_3^2) + \sqrt{h}(F_1(t)g_{\alpha\beta}(y)dy^\alpha dy^\beta + F_2(t)g_{mn}(y)dy^m dy^n) \tag{4.12}$$

Where $\alpha, \beta = 4, 5$ and $m, n = 6, 7, 8, 9$, which is to say the internal manifold has the structure of a product manifold consisting of a 2d and 4d part. * So this argument suggest an alteration to the KKLT model is needed in order to proceed.

4.4 Global compatibility

As pointed out by Liam Mcallister in his talk at Strings 2019, [17], the components of KKLT are quite well examined seperately. Altough at this point it should be clear that the consensus on the validity of each part is not unanimous in the literature. With this in mind we'll examine his point anyway. As Mcallister points out we can consider the different components of the KKLT senario seperately and should pose some questions regarding these components. Concerning Moduli stabilisation we should ask whether there exit global models with:

*Cross terms between these manifolds could in general also exist but we'll not discuss these.

- Quantized fluxes giving small classical superpotenial
- Incorporating D7-brane stacks supporting gaugino condensation
- Klebanov Strassler throat regions

And concerning the antibrane uplifting we can wonder whether:

- there is a supersymmetric action that describes $\overline{D3}-$branes?
- decompactification is the only important instability coming from the $\overline{D3}-$branes?
- de Sitter can vacua be described in 10d supergravity manner?

4.4.1 Quantized fluxes giving small classical superpotenial

By this we mean that the W_0 from chapter 3 is indeed $<< 1$ in string units. This we required in order to be able to neglect all moduli except the Kähler as these would be fixed at large mass. Examples of constructions satisfying this condition have been found, for example in [26].

4.4.2 Incorporating D7-brane stacks supporting gaugino condensation

The presence of gaugino condensation is an example of a correction that in chapter 3 gave rise to the exponential part in (3.16) needed for the AdS_4 vacuum prior to uplift. Although other corrections resulting in a similar term to superpotential might be substituted instead means the existence of gaugino condensation is some sense less crucial. However examples in the literature, such as [27], show that such substitutions while remaining interesting are not necessary. By which we mean that the addition of gaugino condensation is enough to lead to the AdS_4 vacuum. Gaugino condensation is the most often considered source of these correction terms but other sources exist. So gaugino condensation in particular is not required but there has to be something that gives rise to such a correction term to formulate a anti-de-Sitter vacuum.

4.4.3 Klebanov Strassler throat regions

As discussed in section 4.2 the existence of such a warped part of our manifold is crucial to the warping needed for the scale separation. Luckily the existence of these regions are well established, at least in the non-compact space where they where originally formulated in [28]. The method of gluing of this region to a compact CY is still an open point of discussion and could benefit greatly from an explicit formulation of a CY metric. Besides the method of gluing there is even still debate on if such a throat is possible in on a compact manifold.

4.4.4 Supersymmetric action describing $\overline{D3}-$branes

A supersymmetric action which describes $\overline{D3}-$branes is use full as it sheds light on the coupling of the branes to the background. This as we have seen is quite relevant. It is these couplings which are argued to ruin the uplift as in certain cases we saw that the back reaction due to the coupling is supposedly stronger than expected, which would indicate that the probe approximations is not valid. A extensive explicit Susy action for $\overline{D3}-$branes has been formulated in [29] to quadratic order in the Fermion fields. So such a description is not beyond our reach although it non trivial to preform the analysis using this action. And the obvious route of going beyond quadratic order also remains an avenue for improvement.

4.4.5 Is decompactification the leading $\overline{D3}-$branes instability

As we mentioned earlier when working under certain restrictions the $\overline{D3}-$branes can be found to cause singularities when considered in the flux background. These singularities would indicate further instabilities besides solely decompactification. However as shown by [30] these singularities do not occur when these restrictions are relaxed as then the $\overline{D3}-$branes puff up in to $NS5-$branes as per the Myers effect. Leaving decompactification as the leading instability as expected.

4.4.6 10d supergravity of dS

Most of the analysis that has been made of our model is in the framework of 4d EFT. But fundamentally we should be able to do this in 10d as well and when we do the 10d calculation we better find the same result as we did in the 4d. This comparison has been done where possible in the literature, for example in [18]. Thus the EFT analysis remains valuable. There however still various corrections one could consider in 10d which are integrated out in the 4d EFT, but these are generally found to be subleading. So it appears that the 10d calculation agrees with the 4d argumentation.

4.4.7 All at once

Thus far no example has been constructed that exhibits all these properties simultaneously. There for there is no guarantee that this actually exist but so far no conclusive proof exist that this impossible. But subtleties might arise in the construct which means these necessary features can not all be realised simultaneously. So might the D7 branes change the back-reaction of the anti-brane in the throat, dependent on the way the throat is glued to the compact manifold. And so there are many more possible point of failure that need to be checked for mutual compatibility to be sure that the construction is valid in general. If one could construct a explicit model with all these properties an analysis can be made, by checking this example for the existence of the de Sitter, which might prove that there exist de Sitter vacua in string theory. But this obviously is easier said than done.

Chapter 5

Conclusion

To conclude we have discussed in chapter 3 a construction proposed by Kachru Kallosh Linde and Trivedi to recover a de Sitter vacuum in supergravity. Which consisted of formulating a stabilized anti-de-Sitter vacuum and adding an anti-brane, which broke supersymmetry and provided an uplift to the vacuum. This resulted in a de Sitter vacuum. We discussed the main assumptions made during this construction an analysed the stability of the resulting vacuum.
In chapter 4 we saw that this construction is not without flaws. In the 17 years since it's proposal alterations have been constructed tweaking the model here and there but fundamentally these remain the same model. Extensive examination of the model showed many a priori problems with the construction. However all of these augmentations and no-go theorems have been found to miss certain subtleties or remain dependent on the inability to formulate an explicit example. The later therefor we found to be the main challenge left for the construction. There is however not a clear method to proceed in this direction. And there for it would be presumptuous to simply call it a technical detail as some authors have argued. As with all open questions it remains impossible to predict what we might find along the way. But even if it turns out that due one reason or an other the KKLT construction does not work the debate about the model has already provided numerous insights in to the deeper workings of string theory models applicable beyond attempted cosmological models.
In conclusion we have found no clear answer to the question of whether the Kachru Kallosh Linde Trivedi model provides a de Sitter space in supergravity. But we have gained some insight in to the obstacles to find a conclusive answer to this question.
It is clear at this point that currently there is no definitive method to solve

the cosmological constant problem in string theory using the KKLT construction. There are other potential methods, such as the large volume scenario, which could provide an alternative should KKLT be found to be invalid, but these we have not discussed.

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