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New horizons in string theory

bubble babble in search of darkness

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Abstract
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It was discovered nearly two decades ago that we live in an accelerating universe that is dominated by dark energy. Understanding the origin of such an energy has turned out to be a very difficult open question in physics, and calls on the need for a fundamental theory like string theory. However, despite decades-long effort, string theory has proven incredibly resilient to a satisfactory construction of dark energy within its framework.

In the first part of this thesis and the included papers, we examine this problem and propose two possible solutions. The first is a construction within the framework of M-theory, the eleven dimensional cousin of string theory. Using only well-understood geometric ingredients and higher-derivative corrections to eleven dimensional supergravity, we present a new class of four dimensional vacua that contain dark energy. In the process, we also construct a new class of non-supersymmetric Minkowski vacua that were previously not known. Our second idea is a novel proposal that our universe could be embedded on the surface of an enormous spherical bubble that is expanding in a five dimensional anti de Sitter spacetime. The bubble is made of branes in string theory and its expansion is driven by the difference in the cosmological constants across it. We argued that such a construction arises naturally in string theory, and showed how four dimensional gravity arises in such a universe. We further showed that four dimensional matter and radiation arise from quantities that are innately five dimensional.

Another challenging problem in physics concerns the nature of black holes – the presence of an event horizon in particular. This poses a paradox between well understood physical principles, and requires a fundamental theory for its resolution. Towards this goal, we constructed a novel class of horizonless objects that mimics black holes, and proposed these objects as an alternative end point of gravitational collapse. Subsequently, we constructed slowly rotating versions of these "black shells" and proposed an observational signature that could distinguish them from black holes in cosmological experiments. This is discussed in the second part of the thesis and in the included papers.

Keywords: String theory, black holes, dark energy, de Sitter, cosmological constant, M-theory, braneworld, anti de Sitter

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“What makes the desert beautiful,” said the little prince,
“is that somewhere it hides a well…”

— Antoine de Saint-Exupéry, The Little Prince

List of papers

This thesis is based on the following papers, which are referred to in the text by their Roman numerals.

I Johan Blåbäck, Ulf Danielsson, Giuseppe Dibitetto, and Suvendu Giri. *Constructing stable de Sitter in M-theory from higher curvature corrections*, JHEP 09 (2019) 042, `arXiv:1902.04053`.

II Souvik Banerjee, Ulf Danielsson, Giuseppe Dibitetto, Suvendu Giri, and Marjorie Schillo. *Emergent de Sitter Cosmology from Decaying Anti–de Sitter Space*, Phys. Rev. Lett. 121 (2018) 26, 261301, `arXiv:1807.01570`.

III Souvik Banerjee, Ulf Danielsson, Giuseppe Dibitetto, Suvendu Giri, and Marjorie Schillo. *de Sitter Cosmology on an expanding bubble*, JHEP 10 (2019) 164, `arXiv:1907.04268`.

IV Souvik Banerjee, Ulf Danielsson, and Suvendu Giri. *Dark bubbles: decorating the wall*, JHEP 04 (2020) 085, `arXiv:2001.07433`.

V Ulf Danielsson, Giuseppe Dibitetto, and Suvendu Giri. *Black holes as bubbles of AdS*, JHEP 10 (2017) 171, `arXiv:1705.10172`.

VI Ulf Danielsson and Suvendu Giri. *Observational signatures from horizonless black shells imitating rotating black holes*, JHEP 07 (2018) 070, `arXiv:1712.00511`.

Contents

1. Introduction

It was discovered almost twenty years ago that our universe is not only expanding, but is also accelerating. This discovery caused a major upheaval, since prior to that it was widely believed that the expansion of the universe is slowing down. The mysterious energy that seems to be pushing the universe apart, causing it to accelerate, is called *dark energy*. Explaining this acceleration from the present theoretical understanding of cosmology is still an open problem. The most popular explanation is that there is a vacuum energy in the universe, which pushes apart spacetime itself. This is called the *cosmological constant*, and observations show that it is incredibly tiny and positive. The cosmological constant is a dimensionful quantity with dimension of inverse squared length. In general relativity, which is our experimentally verified low-energy theory of gravity, there are only two fundamental dimensionful quantities namely, the Newton's constant G and the speed of light c. There is no combination of these quantities that produces an object with the dimension of an inverse squared length. The appearance of a non-zero cosmological constant within general relativity is therefore not possible, without introducing a new dimensionful scale in the problem.

A length scale arises if one introduces quantum mechanics into the picture. This introduces another dimensionful constant – the Planck constant $\hbar$. The three fundamental constants G, c, and $\hbar$ can now be combined into a quantity $\ell_{\mathrm{Pl}} := \sqrt{G\hbar/c^3}$, called the Planck length. The Planck constant $\hbar$ is in fact associated with an energy scale in quantum mechanics, corresponding to the zero point energy. This suggests the possibility that dark energy might have an explanation in terms of the zero point energy of quantum mechanics. In terms of the Planck length, the measured value of the cosmological constant is $\Lambda \sim 10^{-120}\ell_{\mathrm{Pl}}^{-2}$, which is an incredibly tiny number. To get an intuition of how tiny this number is, note that the volume of a single grain of sand is about 10^{-90} times the volume of the entire observable universe. 10^{-120} is much much smaller than that – about as small as the volume of a single atom of hydrogen compared to the volume of the entire observable universe. This raises a deep question – why is the cosmological constant so tiny? This is the familiar question of naturalness that appears all over particle physics. The usual answer to such a question is that there is a symmetry that forces this quantity to vanish. When this symmetry is very slightly broken, the quantity acquires a tiny non-zero value. However, no such symmetry seems to exist for the cosmological constant, making it very difficult to explain its extremely fine-tuned non-zero value. This is the essence of the *cosmological constant problem*.

Since it is clear that the problem involves both quantum field theory and general relativity, one can hope that a unified theory that combines both, might provide an answer. Such a theory needs to be a quantized theory of gravity. The reason why gravity needs to be quantized can be seen, for example, from the double slit experiment. In this experiment, the wavefunction of an electron is a superposition of its wavefunction passing through both slits. If the gravitational field of the electron were classical, meaning that it can only exist at one place at one time and not in a superposition, then one could detect which slit the electron passed through, by measuring its gravitational force, contradicting the quantum nature of the electron. This strongly suggests that if anything that gravity couples to is quantized, then gravity has to be quantized as well. See Feynman's Lectures on gravitation (1996) for a delightful discussion around this, and also for the first attempt to construct a quantized theory of gravity.

String theory is a leading candidate for such a quantum theory of gravity. However, it is consistent only in ten dimensions, and requires an additional symmetry called *supersymmetry*. Supersymmetry postulates that every matter particle (fermion) has a counterpart that behaves like light (boson), and vice versa. Such a symmetry is also motivated from particle physics, where it protects the mass of certain particles, like the Higgs boson, from becoming extremely heavy. Although this is expected to be a symmetry at extremely high energies where parameters $G, \hbar$, and $1/c$ become large and cannot be neglected, it is not a symmetry that we see in everyday life. Therefore, supersymmetry must cease to exist at some energy scale, higher than our present scales of experiments, but below the Planck scale.[1] In fact, it has been proven that in ten dimensions, in the presence of supersymmetry, the only consistent theory of quantum gravity is string theory. This is a special case of a broader program that goes by the name of *string universality*.

A quantum theory of gravity, when complete, would combine general relativity with a *grand unified theory* of particle physics, and is sometimes called a *theory of everything*, indicating the fact that it is valid over the entire range of parameter space of the physical parameters G, c, and $\hbar$ mentioned above. The space of theories in this parameter space can be represented

[1]For supersymmetry to be of use in particle physics, it must be broken close to the energy scale at which the Large Hadron Collider is currently running. Failing this, it might still exist in nature, but loses much of its motivation from particle physics.

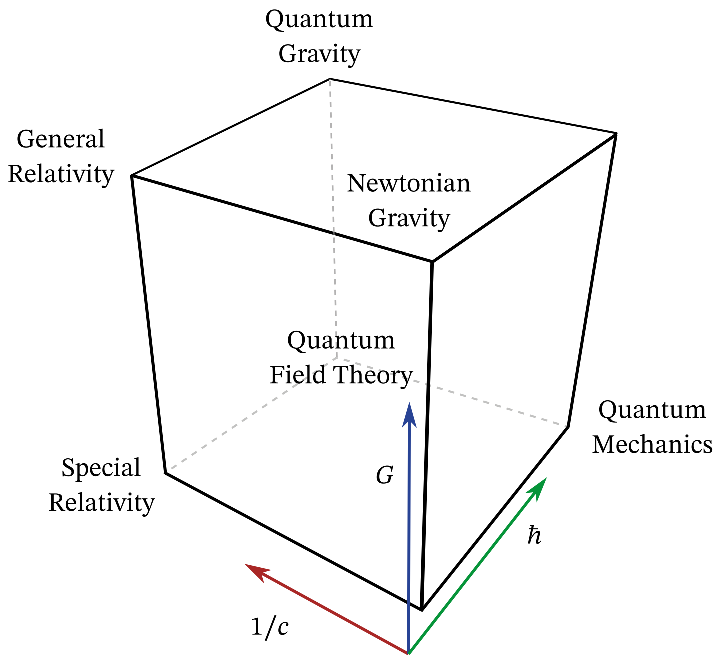

Figure 1.1. The Bronstein cube representing the space of physical theories. Starting from classical physics at the nearest corner, where $G = \hbar = 1/c \ll 1$, moving along the sides of the cube corresponds to increasing values of $G, \hbar$, and $1/c$ as indicated. Corners of the cube represent the theory that becomes relevant in that region of parameter space. The farthest corner represents the regime of quantum gravity, where all parameters are relevant. The unlabelled corner represents non-relativistic quantum gravity which is not relevant for our discussion here.

along the corners of a cube, shown in figure 1.1, which is attributed to Bronstein (1933) and Stachel (2001). Although most physical processes do not involve physics that is sensitive to all three parameters at the same time, objects with an event horizon (which is a surface that only allows one way passage through it) are an exception. A universe with a positive cosmological constant could have such a horizon very far out. Such a *cosmological horizon* gives rise to problems with unitarity in quantum mechanics, and finiteness of entropy in general relativity. Another celestial object that has an event horizon is a black hole. In this case, it gives rise to an interesting tension between the principle of unitarity in quantum mechanics and the equivalence principle in general relativity. Resolving these problems, therefore requires a quantum theory of gravity.

The goal of this thesis is to explore these two questions – the issue of dark energy, and the problem with black holes – in the context of string theory. After summarizing the current status of these problems, we will introduce solutions that we proposed to resolve them. The purpose of the thesis is to introduce the problem that we have attempted to solve, and to put our proposed solutions in the context of other work in the literature – highlighting similarities and differences. The detailed arguments and supporting computations are contained in our published papers that are a part of the thesis, and we will refrain, as much as possible, from repeating them here.

The thesis is divided into two parts. The first part discusses the problem of dark energy in string theory, while the second part deals with black holes. After introducing dark energy in cosmology, and its relation to the cosmological constant in chapter 2, we will discuss attempts to realize it in string theory in chapter 3. In chapter 4, we will do the same in M-theory, and present a new compactification to four dimensional de Sitter vacuum, which we constructed in paper I. In chapter 5, we will discuss some conjectures regarding the construction of a positive cosmological constant in string theory/M-theory. Chapter 6 revisits the braneworld construction by Randall and Sundrum and in chapter 7, we present a novel construction of a de Sitter universe in string theory, that we proposed and developed in papers II, III and IV. In the second part of the thesis, chapter 8 introduces the black hole information paradox and discusses an alternative to black holes that we developed in papers V and VI. We will summarize and conclude the thesis in Chapter 9.

Part I:
Dark energy in string theory

2. Dark energy and the cosmological constant

In this chapter, we will briefly review what dark energy is, observational evidence for it, and its relation to the cosmological constant. In this process, we will also introduce some notation that will be used throughout the thesis.

2.1 What is dark energy?

Let us start with the Einstein-Hilbert action in d dimensions,

$$S = \int \mathrm{d}^d x \sqrt{-\mathrm{g}} \left(\frac{R}{2\kappa_d^2} + \mathcal{L}_{\text{matter}} \left(\phi_i, \partial \phi_i \right) \right), \tag{2.1}$$

where $\mathcal{L}_{\text{matter}}$ is a matter Lagrangian of some matter fields (ϕ) coupled to gravity. We will work in conventions where

$$\kappa_d^2 := 8\pi G_d \equiv 8\pi m_{\text{Pl}}^{2-d} \equiv \mathrm{M}_{\text{Pl}}{}^{2-d}, \tag{2.2}$$

for G_d, m_d, and M_d being the d dimensional Newton's constant, Planck mass, and the reduced Planck mass respectively. Here, and throughout the thesis, we will work in natural units[1] ($\hbar = c = 1$), only making factors of $\hbar$ and c explicit when it adds some insight to the discussion, as in some places in part 2 of the thesis. Since we will work in both even and odd dimensions throughout the thesis, we will use the mostly plus signature for the metric, $(-, +, +, +, ...)$, so that the determinant is negative in any dimension.

Extremizing the Einstein-Hilbert action with respect to the metric gives Einstein's equations

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} = \kappa^2 \left(\mathcal{L}_{\text{matter}}\, g_{\mu\nu} - 2 \frac{\delta \mathcal{L}_{\text{matter}}}{\delta g^{\mu\nu}} \right) := \kappa^2 T_{\mu\nu}, \tag{2.3}$$

where δ is a functional derivative and $T_{\mu\nu}$ is the stress tensor. The left hand side is a function of the metric and its derivatives, thus representing pure geometry, while the right hand side involves contribution from matter fields.

[1]The corresponding four dimensional quantities, with factors of $\hbar$ and c explicit are $\kappa_4^2 := 8\pi G_4/c^4 \equiv \left(8\pi/c^4\right)\left(\hbar c/m_4^2\right) \equiv \left(1/c^4\right)\left(\hbar c/M_4^2\right)$.

It is in this sense that Wheeler famously wrote: *Spacetime tells matter how to move; matter tells spacetime how to curve.*[2] Let us solve these equations to find a solution corresponding to our universe. To do this, a good simplification is to assume that the universe is homogeneous and isotropic (which is true at large scales, anyway). In this approximation, the stress tensor takes a diagonal form, $T_\mu^\nu = \mathrm{diag}\,(-\rho, p, p, p, ...)$, where ρ is the energy density and p is the corresponding isotropic pressure. Being coupled, non-linear, partial differential equations, Einstein's equations are very difficult to solve for the metric given a matter distribution; instead, one makes an ansatz for the metric. For a homogeneous and isotropic universe, this can be written in the form of the Friedmann-Lemaître-Robertson-Walker (FLRW) metric,

$$\mathrm{d}s^2 = -d\tau^2 + a(\tau)^2 \left(\frac{\mathrm{d}r^2}{1-kr^2} + r^2 \mathrm{d}\Omega_{d-2}^2 \right), \tag{2.4}$$

where $\mathrm{d}\Omega^2$ is the metric of a unit $d-2$ dimensional sphere. This represents a time evolving universe, with the radius of a unit sphere being proportional to the *scale factor* $a(\tau) \equiv a$, for proper time τ. The spatial part of the metric represents a space of uniform curvature, which is positive, negative, or zero corresponding to $k \in \{1, 0, -1\}$ respectively. These coordinates are called *comoving coordinates.*

Since we are talking about our observable universe here, we will specialize to four dimensions in the rest of this chapter. With the FLRW ansatz for the metric and the diagonal stress tensor, Einstein's equations diagonalize and reduce to equations governing the evolution of the scale factor $a(\tau)$. These are called the *Friedmann equations*, and in four dimensions are given by

$$H^2 := \frac{\dot{a}^2}{a^2} = \frac{8\pi G_4}{3}\rho - \frac{k}{a^2}, \tag{2.5}$$

$$\frac{\ddot{a}}{a} = -\frac{4\pi G_4}{3}(\rho + 3p), \tag{2.6}$$

where the Hubble parameter, $H := \dot{a}/a$, measures the expansion rate of the universe. Here ρ and p are the total energy density and the pressure respectively, representing the sum over all types of matter in the universe, *i.e.*, $\rho = \sum_i \rho_i$ and $p = \sum_i p_i$.

While the first Friedmann equation gives the rate of expansion of the universe, the second Friedmann equation governs its acceleration in response to the matter contained in the universe. Different kinds of matter can be conveniently characterized in terms of their *equation of state* $w := p/\rho$.

[2]Throughout this chapter, we will use the word *matter* to mean not just ordinary fermionic matter, but all kinds of energy density – ordinary matter, radiation, etc.

Matter with $\rho + 3p > 0$ causes the universe to decelerate, while that with $\rho + 3p < 0$ causes it to accelerate. All observed matter in our universe is of the first kind; for example, relativistic matter (radiation) and non-relativistic pressure less matter (dust) correspond to $w = 1/3$ and $w = 0$ respectively. This suggests that our universe should be decelerating.

However, as we will see in the next section, observations show that our universe is in fact accelerating. This implies that there has to an yet unobserved energy density in the universe that contributes as $\rho + 3p < 0$ (corresponding to an equation of state $w < -1/3$) to the Friedmann equation. Since it has not yet been detected experimentally, it is referred to as a *dark energy*.

2.2 Measuring dark energy

In order to connect to observations, it is useful to rewrite the Friedmann equation in terms of the present value of the corresponding quantities. We will denote these with a subscript "0". Additionally, let us define the critical energy density $\rho_{\text{crit}} := 3H^2/(8\pi G_4)$, as the energy density for a spatially flat universe. In terms of this, we can define a dimensionless density parameter $\Omega_i := \rho_i/\rho_{\text{crit}}$, which represents the energy density of a particular kind of matter, as a fraction of the critical energy density. This allows us to rewrite equation (2.5) as

$$\Omega_{\text{total},0} \equiv \Omega_{m,0} + \Omega_{\Lambda,0} = 1 + \frac{k}{a_0^2 H_0^2}, \tag{2.7}$$

where we have written the total density parameter in terms of contributions from known matter and a dark energy, denoted by the subscripts m and Λ respectively. For a spatially flat universe $k = 0$, with $\Omega_{\text{total},0} = 1$; whereas $\Omega_{\text{total},0} \gtrless 1$ for $k \gtrless 0$ respectively. The deviation of $\Omega_{\text{total},0}$ from 1, therefore, measures the spatial curvature of our universe.

This of interest to observations because, the spatial curvature can be measured from the position of the first peak in the power spectrum of the anisotropies of the Cosmic Microwave Background (CMB) radiation. Using equation (2.7), this gives the total density parameter. The most recent data from the Planck collaboration (Aghanim et al., 2018) gives

$$|\Omega_{\text{total}} - 1| = 0.0007 \pm 0.0019 \ll 1, \tag{2.8}$$

which shows that our universe is extremely flat. The next thing that can be reliably measured is the matter density parameter Ω_m. This, combined with $\Omega_{\text{total}} \simeq 1$, gives a way of obtaining the dark energy density parameter

Ω_Λ. Below, we will briefly mention some ways of measuring Ω_m and the results obtained. Some good reviews discussing this in more detail are by Padmanabhan (2003) and by Copeland, Sami, and Tsujikawa (2006).

Type Ia supernova

At the end of its lifecycle, a main sequence star, not much heavier than the Sun, ends up expelling most of its outer material, leaving behind a hot dense core called a white dwarf. If such a white dwarf gets more massive, for example by merging with another white dwarf, it can enter a process of uncontrolled nuclear fusion and end up exploding in a supernova. These are the brightest and most uniform type of supernovae, and it is believed that they have the same peak luminosity wherever they are in the universe, *i.e.*, irrespective of their redshift. These are called type Ia supernovae and their universal peak brightness makes them useful as *standard candles*.

The apparent magnitude m (which is proportional to the logarithm of the observed flux density) and the redshift z of a supernova can be measured directly. These are related to the absolute magnitude M, by the following astrophysical relation, which can be used to find the *luminosity distance* d_{L}:

$$m - M = 5 \log\left(\frac{d_{\mathrm{L}}(z)}{M_{\mathrm{pc}}}\right) + 25. \tag{2.9}$$

Since the absolute magnitude M of all type Ia supernovae are thought to be the same, their luminosity distance can be used to fit the value of the density parameters $\Omega_{m,0}$ and $\Omega_{\Lambda,0}$. In their original data set, the Supernova Cosmology Project (Perlmutter et al., 1999) had discovered 42 supernovae of type Ia, while the High-z Supernova Search Team (Riess et al., 1998) had discovered 48 more, to find $\Omega_{m,0} \simeq 0.28$ and $\Omega_{\Lambda,0} \simeq 0.72$. Later, with the discovery of more supernovae of this type, also with higher redshifts, this was refined by Choudhury and Padmanabhan (2005) to $\Omega_{m,0} \simeq 0.31$, and $\Omega_{\Lambda,0} \simeq 0.69$.

Age of the universe

Another way of measuring the dark energy density is by integrating the Friedmann equation (2.5) to determine the age of the universe. Using the definition of the redshift, $1 + z = a_0/a$, and the Hubble parameter $(H/H_0)^2 = \Omega_{m,0}\,(a0/a)^3 + \Omega_{\Lambda,0}$, the age of the universe can be computed as

$$t_0 = \int_0^{t_0} \mathrm{d}t = \int_0^{\infty} \frac{dz}{H(1+z)} = \frac{1}{H_0} \int_0^{\infty} \frac{\mathrm{d}z}{(1+z)\sqrt{\Omega_{m,0}(1+z)^3 + \Omega_{\Lambda,0}}}$$
$$= \frac{2}{3H_0} \frac{1}{\sqrt{\Omega_{\Lambda,0}}} \operatorname{arcsinh} \sqrt{\frac{\Omega_{\Lambda,0}}{\Omega_{m,0}}}, \tag{2.10}$$

where as before, the universe has been assumed to be spatially flat. Using our current best estimates of the Hubble parameter (H_0) and the age of the universe (t_0) from the Planck 2018 data, one finds $H_0 t_0 = 0.951 \pm 0.023$. Together with the flatness constraint $\Omega_{m,0} + \Omega_{\Lambda,0} = 1$, this gives $\Omega_{m,0} \simeq 0.315$ and $\Omega_{\Lambda,0} \simeq 0.685$.

Equation of state of dark energy

Another cosmological measurement that constrains the nature of dark energy is that of baryon acoustic oscillations (BAO). In the hot plasma of the early universe consisting of photons, electrons, and baryons (protons and neutrons), photon-electron Thompson scattering created an outward pressure that was counteracted by the gravitational attraction, giving rise to *acoustic oscillations* in the plasma. As the universe cooled down, electrons and protons combined into neutral hydrogen (known as *recombination*). This caused the photons to decouple and the acoustic oscillations to cease, but their density fluctuations were frozen and are imprinted in both the CMB radiation as well as the distribution of ordinary baryonic matter. This can be measured today from the power spectrum of density fluctuations of galaxies. These measurements, together with data from type Ia supernovae and CMB measurements mentioned above, constrain the nature of dark energy to give an equation of state $w_\Lambda = -1.028 \pm 0.031$. The relevance of this will become clear in the next section.

2.3 Dark energy and the cosmological constant

To summarize the discussion so far, experiments show that our universe contains an energy density that is driving its acceleration. From the Friedmann equation, we know that such a form of energy should have an equation of state $w < -1/3$. However, ordinary matter and radiation does not have such an equation of state. What then is this dark energy and where does it come from? In this section, we will discuss the main candidate for a dark energy – the cosmological constant.

Let us add a constant $\mathcal{L}_0$ to the Einstein-Hilbert action of equation (2.1), to get

$$S = \int \mathrm{d}^d x \sqrt{-g} \left(\frac{R}{2\kappa_d^2} + \mathcal{L}_{\text{matter}} + \mathcal{L}_0 \right). \tag{2.11}$$

Such a constant does not change the equation of motion for matter (obtained by extremizing the action with respect the scalar field), but being coupled to the determinant of the metric, it does contribute to the Einstein's equations to give

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} = \kappa^2 \left(T_{\mu\nu} + \mathcal{L}_0 \, g_{\mu\nu} \right). \tag{2.12}$$

Written in this form, the constant term contributes an extra stress tensor, $\Delta T_{\mu\nu} = \mathcal{L}_0 \, g_{\mu\nu}$, to Einstein's equations. Being proportional to the metric, it has an equation of state $w = -1$. Moving this extra term to the left hand side of the equation, however, suggests a different interpretation

$$R_{\mu\nu} - \frac{1}{2} \left(R + 2\kappa^2 \mathcal{L}_0 \right) g_{\mu\nu} = \kappa^2 T_{\mu\nu}. \tag{2.13}$$

In this form, the extra term appears as a shift of the scalar curvature, without modifying the stress tensor. As indicated before, if we think of the left hand side of Einstein's equations as giving a rule for how spacetime should bend in response to stress tensor on the right hand side, this implies modifying the rule even when there is no matter on the right hand side. The modification is in fact so severe that empty flat space is no longer a solution to these equations. It is customary to write this extra term as $\mathcal{L}_0 = -\Lambda/\kappa^2$, and Λ is called the *cosmological constant*.

Going back to the Lagrangian, the addition of a constant is consistent with the symmetries of the Lagrangian, and everything that is allowed, must be present in the Lagrangian. So, generically a constant term must be included in the Einstein-Hilbert action, giving a cosmological constant in Einstein's equations

$$R_{\mu\nu} - \frac{1}{2} (R - 2\Lambda) \, g_{\mu\nu} = \kappa^2 T_{\mu\nu}. \tag{2.14}$$

This modifies the Friedmann equations to include extra contributions proportional to the cosmological constant

$$\frac{\dot{a}^2}{a^2} = \frac{8\pi G_4}{3} \rho - \frac{k}{a^2} + \frac{\Lambda}{3}, \tag{2.15}$$

$$\frac{\ddot{a}}{a} = -\frac{4\pi G}{3} (\rho + 3p) + \frac{\Lambda}{3}. \tag{2.16}$$

This shows that the cosmological constant plays an important role in determining the acceleration of the universe. When positive ($\Lambda > 0$), it serves to accelerate the expansion, thus acting like a repulsive force. Thus, the cosmological constant can drive the accelerated expansion of the universe and has an equation of state $w = -1$, which is consistent with observations. This makes it a leading candidate for dark energy.

Instead of using co-moving coordinates, one can also write the metric in *static coordinates* (so named because, in contrast to the co-moving coordinates discussed before, the metric in static coordinates does not have any time dependence). In this case, one can make a spherically symmetric ansatz for the metric to find a solution to Einstein's equation with a cosmological constant as,[3]

$$\mathrm{d}s^2 = -\left(1 - \frac{\Lambda}{3}r^2\right)\mathrm{d}t^2 + \left(1 - \frac{\Lambda}{3}r^2\right)^{-1}\mathrm{d}r^2 + r^2\mathrm{d}\Omega_2^2. \tag{2.17}$$

These spacetimes preserve all 10 Killing vectors,[4] and are therefore *maximally symmetric*.[5] These are called de Sitter (dS) and anti de Sitter (AdS) spacetimes for $\Lambda \gtrless 0$ respectively. Flat space corresponds to the absence of a cosmological constant, $\Lambda = 0$.

An interesting question to ask is: what happens if the matter fields appearing in the Lagrangian are not classical but are quantized instead. It was originally argued by Sakharov (1967), and later by Weinberg (1989), that taking quantum field theory into account, there would be an additional contribution coming from the constant energy density of the vacuum, given by $\langle T_{\mu\nu}\rangle = -\rho_{\text{vac}}\, g_{\mu\nu}$. Including this contribution, the semi-classical Einstein's equations become

$$R_{\mu\nu} - \frac{1}{2}(R - 2\Lambda)\, g_{\mu\nu} = \kappa^2 T_{\mu\nu} - \kappa^2 \rho_{\text{vac}}\, g_{\mu\nu}. \tag{2.18}$$

This additional contribution to the energy density, can be moved over to the left hand side, to define an effective cosmological constant

$$\Lambda_{\text{eff}} := \Lambda + \kappa^2 \rho_{\text{vac}}. \tag{2.19}$$

Since standard model matter is indeed quantized, this is the value of the cosmological constant that would be measured by cosmological observations.

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[3]An integration constant proportional to $-2G_4M/r$ also appears in the solution for Einstein's equations, which corresponds to a delta function stress tensor at the origin.

[4]The maximal number of Killing vectors for a d dimensional spacetime is $d\,(d+1)/2$.

[5]Alternatively, a maximally symmetric spacetime has $d\,(d-1)\,R_{abcd} = R\,(g_{ad}g_{bc} - g_{ac}g_{bd})$.

Quintessence

There is however, another possibility. The current accuracy of observations also allows for the equation of state of dark energy, to be slightly different from $w = -1$. As an example, imagine that the matter fields coupled to the Lagrangian discussed above, are not at the minimum of their potential, but are slowly rolling down a very flat potential instead. In that case, their vacuum expectation value would slowly change over time, and ρ_{vac} in equation (2.19) would be a time dependent function, rather than a constant. This is the main idea of *quintessence*, which was originally proposed by Wetterich (1988), and is an active line of investigation for solving the dark energy problem. A summary of this and some other proposals for explaining dark energy can be found in the review by Copeland, Sami, and Tsujikawa (2006). In this thesis, however, we will focus only on the cosmological constant.

A historical side note

The cosmological constant has a chequered history. Einstein believed the universe to be static and closed, which is why he introduced the cosmological constant in his equations. He was also of the opinion that the spatial curvature of the universe is provided by the matter contained in the universe. However, de Sitter (1917) constructed a solution to Einstein's equations representing a closed universe that did not contain any matter. This was followed by two seminal papers by Friedmann (1924) and Lemaître (1927) on non-static solutions of Einstein's equations. However, since the popular belief at the time was that the universe is static, these developments went largely unnoticed. Things changed with the discovery of an expanding universe by Hubble (1929). With this new discovery, however, Einstein no longer saw the need for a cosmological constant in his equations and vehemently argued against it. In his autobiography, Gamow (1970) wrote:

> Much later, when I was discussing cosmological problems with Einstein, he remarked that the introduction of the cosmological term was the biggest blunder he ever made in his life.

It is interesting to note that Landau and Lifshitz (1971) are among the many notable physicists at the time, who were against the cosmological constant. In fact, in their classic text, they wrote

> At the present time, however, there are no cogent and convincing reasons, observational or theoretical, for such a change in the form of the fundamental equations of the theory. We emphasize that we are talking about changes that have a profound physical significance…

Things have of course changed after the recent discovery of the dark energy domination of our universe, and the cosmological constant has regained attention.

2.4 A fine-tuning problem

Our current best estimates for the various components of the energy density in the universe from the Planck collaboration (Aghanim et al., 2018), are

$$\Omega_m = 0.315 \pm 0.007, \Omega_\Lambda = 0.685 \pm 0.007, \Omega_k = 0.0007 \pm 0.0019, \quad (2.20)$$

while the density of radiation is given by the sum of contributions from the neutrino density (Ω_ν) and CMB radiation (Ω_γ), which are $\Omega_\nu \sim 5.38\times10^{-5}$ and $\Omega_\gamma < 0.003$ respectively. This means that about 68.5% of the energy density of our universe is dark energy with the remaining 31.5% consisting of matter (both observable matter and invisible *dark* matter), while the spatial curvature is negligible. Using the current measurement of the Hubble parameter H_0, the critical density is estimated to be

$$\rho_{\text{crit},0} = \frac{3H_0^2}{8\pi G_\text{N}} \simeq 8 \cdot 10^{-46}\text{GeV}^4. \quad (2.21)$$

From this, the dark energy density can be computed using $\rho_{\Lambda,0} = \Omega_{\Lambda,0} \times \rho_{\text{crit},0}$, to give

$$\rho_{\Lambda,0}^{(\text{obs})} \approx 5 \cdot 10^{-46}\text{GeV}^4. \quad (2.22)$$

As discussed in equation (2.19), this observed value comes from the effective cosmological constant, which includes contributions from the *bare* cosmological constant, as well as the vacuum energy of all quantum fields in the universe. Of these, the latter can be computed in quantum field theory as the sum of all vacuum bubble diagrams, and is given by an integral, which turns out to be UV divergent. A naïve way to regulate this divergence is by introducing a cutoff scale for the theory, which is the mass scale up to which perturbative calculations in quantum field theory on curved space time can be trusted. Taking this to be the Planck scale, leads to the following estimate for the vacuum energy:

$$\rho_\text{vac} \sim \frac{m_\text{Pl}^4}{16\pi^2} \approx \frac{(10^{19}\text{GeV})^4}{16\pi^2} \approx 10^{74}\text{GeV}^4. \quad (2.23)$$

This is about 10^{120} orders of magnitude larger than the observed value in equation (2.22) and is the value often quoted in the literature.[6] How-

[6] One could argue against using the Planck scale as the cutoff. However, even if the cutoff scale was as low as the the scale of strong interactions (QCD), it would give $\rho_\text{vac} \sim 10^{-3}\text{GeV}^4$, which is still 49 orders of magnitude too large.

ever, such a cutoff is not Lorentz covariant, and does not lead to a stress tensor with the right equation of state ($w = -1$) for the vacuum energy. Koksma and Prokopec (2011) argued that a Lorentz covariant renormalization scheme, which gives a vacuum energy that agrees with the result from dimensional regularization, and has the right equation of state, yields

$$\rho_{\text{vac}} \sim \sum_i n_i \frac{m_i^4}{64\pi^2} \log\left[\frac{m_i^2}{\mu^2}\right] \sim 2 \cdot 10^9 \text{GeV}^4. \tag{2.24}$$

Here the sum is over all standard model particles, with n_i being the number of particles of each type, m_i being their respective masses and μ being the renormalization scale. A recent discussion of this issue can be found in an article by Danielsson (2019). The vacuum energy computed in equation (2.24) is much lower than the naïve estimate in equation (2.23), but is still about 10^{55} orders of magnitude too large compared to the observed value. This suggests that the bare cosmological constant must be equally large, but extremely fine tuned to 1 part in 10^{55}, so that it almost cancels the vacuum energy to give the tiny observed value of the cosmological constant. This extreme fine tuning problem is called the *cosmological constant problem*. See the article by Martin (2012) for a careful discussion of the various renormalization schemes.

2.5 The path forward

To summarize the discussion in this chapter, observations tell us that we live in a universe that is not only expanding but also accelerating. This acceleration is due to the presence of a *dark energy*, which constitutes over two-thirds of the total energy of the universe. Observations also tell us that the absolute value of this dark energy density is unnaturally small, and we don't have a good theoretical explanation of why this should be the case.

When faced with such a problem, a reasonable approach is to turn to a more fundamental theory for answers with the hope that this theory knows something, that our low energy theory does not. The most fundamental theory would be a quantized theory of gravity, which we do not yet have; but one of the brightest candidates is string theory. We will do exactly this in the next chapter and explore the question of a cosmological constant in the context of string theory.

3. de Sitter in string theory

String theory is a quantized theory of strings, just as quantum field theory is a quantized theory of point particles. The string, which can be open or closed, can vibrate, and the vibration modes correspond to particles. In particular, the closed string has a vibration mode corresponding to a massless spin-two particle, which was argued by Weinberg (1965) to be the particle that mediates gravitational interactions, and is called a *graviton*. While a naïve attempt to quantize gravity turns out to be perturbatively non-renormalizable in $d > 2$ (since $1/\kappa^2 \sim M^{d-2}$), string theory has the advantage of being perturbatively renormalizable, and produces finite results. Moreover, when a string propagates on a curved background, it gives rise (at sufficiently low energy) to Einstein's equations, with background fields in the theory contributing to the stress tensor. In this way, string theory (which is a quantized theory) naturally gives rise to gravity, and is hence a leading candidate for a theory of quantum gravity.

For string theory to be the right theory for our universe, it must give both the standard model of particle physics (which is the low energy theory that describes all forces in the universe other than gravity) as well as four dimensional gravity (which has two measured parameters – the Newton's constant and the cosmological constant). It turns out that (super)string theory needs $9 + 1$ dimensions to be consistent. Therefore, to be relevant for our four dimensional universe, the other six dimensions must be *compactified* in such a way that low energy four dimensional observers, such as ourselves, do not have access to them. The choice of this six dimensional manifold determines the properties of the four dimensional world. While getting the right four dimensional Newton's constant is not particularly difficult in string theory, obtaining the observed cosmological constant is an incredibly difficult open problem. Regardless of its magnitude, obtaining a positive cosmological constant, by itself, has proven to be very difficult.

In this chapter, we will discuss how string theory naturally prefers a negative cosmological constant, and what could to be done to get a positive value. We will then briefly discuss one of the best known current solutions of obtaining a positive cosmological constant in string theory, and why this is still under discussion. We will briefly comment on the fine tuning aspect of the problem. Instead of providing an introduction to string theory, we will only recapitulate some key ideas that we need here. A detailed introduction to string theory can be found in a standard textbook, such as the two volumes by Polchinski (1998a,b) or a more modern textbook, such as the one by Blumenhagen, Lüst, and Theisen (2013).

3.1 Vacuum energy in supergravity

String theory can be understood perturbatively as supergravity (which is the low energy classical limit of string theory) plus quantum corrections (g_s) and higher curvature corrections (α'). A general compactification of the ten dimensional supergravity theory on a six dimensional manifold, gives rise to a number of scalar fields in the four dimensional theory, in addition to other fields. The potential for these scalar fields determines their vacuum expectation value and gives the four dimensional vacuum energy.

Generically, some of these scalars fail to develop a potential through the compactification and remain massless. Such scalars are not allowed in our universe since they can mediate a yet unobserved fifth force. These are called *moduli* fields. Since massless scalars are not observed in the real world, a mechanism has to be found for them to develop a potential so that they can be stabilized at a minimum of the potential. Finding such a mechanism is referred to as the problem of *moduli stabilization* in string theory. We will briefly discuss vacuum energy in this section and moduli stabilization in the next.

The bosonic sector of $\mathcal{N} = 1$ supergravity in four dimensions contains complex scalars, gauge fields and the metric. The low energy interactions of the scalars are encoded in three functions: a holomorphic superpotential W, a Kähler potential K, and a gauge kinetic function f_{AB}. The Lagrangian for the scalar fields has a non-canonical kinetic term proportional to the Kähler metric $K_{i\bar{j}} := \partial_i \partial_{\bar{j}} K$ and a potential given by a sum of an F-term and a D-term potential. The D-term potential is the same as the one from global supersymmetry $V_{\mathrm{D}} \sim f_{\mathrm{R}}^{AB} D_A D_B$, where f_{R}^{AB} is the real part of the gauge kinetic function, g is the gauge coupling, and A, B are gauge indices. In particular, V_{D} is non-negative, since f_{R}^{AB} is the coefficient of the kinetic terms of the gauge fields. The F-term potential on the other hand, has two contributions. The first is the familiar positive semi-definite potential from global supersymmetry $V_{\mathrm{F}} \supset K^{i\bar{j}} D_i W \overline{D_j W}$ (recall that the kinetic term for the scalar fields is proportional to $K^{i\bar{j}}$), where $D_i W := \partial_i W + \mathrm{M_{Pl}}^{-2} W \partial_i K$, are the F-terms for scalar field indices i and j. The second contribution comes from supergravity, and is negative definite $V_{\mathrm{F}} \supset -3\,|W|^2$. This allows the F-term potential in supergravity to be negative, unlike in global supersymmetry where it is positive semi-definite. Together, the scalar potential can be written as

$$V = V_{\mathrm{F}} + V_{\mathrm{D}} = e^{K/{\mathrm{M_{Pl}}}^2}\left[K^{i\bar{j}} D_i W \overline{D_j W} - \frac{3}{{\mathrm{M_{Pl}}}^2}\left|W\right|^2\right] + \frac{g^2}{2} f_{\mathrm{R}}^{AB} D_A D_B. \quad (3.1)$$

A supersymmetric vacuum, by definition, requires vanishing of the F-terms and D-terms, *i.e.*, $D_i W \overset{!}{=} 0 \overset{!}{=} D_A$, resulting in a negative vacuum energy

$$V = -\frac{3}{{\mathrm{M_{Pl}}}^2}\left|W\right|^2 < 0. \quad (3.2)$$

The only way for this to be non-negative, is if the vacuum breaks supersymmetry and one of the F-terms or the D-terms does not vanish. However, in the absence of supersymmetry, loop corrections no longer cancel against each other, and quantum corrections become large. For the vacuum to be stable against such corrections, supersymmetry has to be broken in a controlled way, such that the quantum corrections are small and the resulting theory is calculable. We will see in section 3.3.3 how this is implemented in one of the most well discussed examples in the literature.

3.2 Moduli stabilization

Another aspect of constructing vacuum energy in string theory is that of *moduli stabilization*. In the next two sections we will use a toy example borrowed from the book by Baumann and McAllister (2015) and the review by Denef, Douglas, and Kachru (2007) to demonstrate this problem, and to see how switching on fluxes in the compact space can be used to solve it.

3.2.1 KK reduction

Let us start with the ten dimensional Einstein-Hilbert action

$$S = \frac{1}{2\kappa_{10}^2}\int \mathrm{d}^{10}x\, e^{-2\Phi}\sqrt{-g_{10}}\, R, \quad (3.3)$$

where κ_{10}^2 is defined in terms of the reduced ten dimensional Planck mass as $\kappa_{10}^2 \equiv 1/M_{10}^8$, Φ is the dilaton, and R is the Ricci scalar constructed from the spacetime metric g_{MN}. We will now perform a Kaluza-Klein reduction of this action for the following ten dimensional geometry, which is a warped product of a four dimensional space time and a *compact* six dimensional space

$$\mathrm{d}s^2 = e^{-6\varphi(x)} g_{\mu\nu}\mathrm{d}x^\mu \mathrm{d}x^\nu + e^{2\varphi(x)}\tilde{g}_{mn}\mathrm{d}y^m \mathrm{d}y^n, \quad (3.4)$$

where $\mu, \nu \in \{0,1,2,3\}$ and $m, n \in \{4, \dots, 9\}$. $\varphi(x)$ is proportional to the volume of the compact six manifold, and depends only on the coordinates of the four dimensional Minkowski space. The factor $\exp(-6\varphi)$ in front of the four dimensional metric is a convenient choice, so that the dimensionally reduced action ends up in the Einstein frame (for a constant dilaton). The dimensional reduction gives

$$S = \frac{1}{2\kappa_{10}^2} \int \mathrm{d}^6 y \sqrt{\tilde{g}_6} \int \mathrm{d}^4 x \sqrt{-g_4} e^{-2\Phi} \left(R_4 + 12 \partial_\mu \varphi \partial^\mu \varphi + e^{-8\varphi} R_6 \right), \quad (3.5)$$

where R_4 and R_6 are the Ricci scalars constructed from $g_{\mu\nu}$ and $\tilde{g}_{mn}$ respectively. $\varphi(x)$, which determines the volume of the compact manifold, appears as a scalar in the four dimensional theory and we will call it the *volume modulus*, *i.e.*, a scalar field that parametrizes the volume of the compact space. The kinetic term for φ can be canonically normalized by scaling $\varphi \mapsto \varphi/(2\sqrt{6})$ to give

$$S = \frac{1}{2\kappa_{10}^2} \int \mathrm{d}^6 y \sqrt{\tilde{g}_6} e^{-2\Phi} \int \mathrm{d}^4 x \sqrt{-g_4} \left(R_4 - \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + e^{-\frac{4}{3\sqrt{3}}\varphi} R_6 \right). \quad (3.6)$$

If the string coupling $g_s := e^\Phi$ is constant over the internal space, and R_6 is independent of y, then the six dimensional integration can be performed explicitly to give

$$S_4 = \frac{1}{2\kappa_4^2} \int \mathrm{d}^4 x \sqrt{-g_4} \left(R_4 + \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + e^{-\frac{4}{\sqrt{6}}\varphi} R_6 \right). \quad (3.7)$$

From this, the four dimensional Planck mass can be read off as the ten dimensional Planck mass scaled with the string coupling and volume of the compact space

$$M_4^2 := \frac{1}{\kappa_4^2} = \frac{\mathcal{V}_6}{g_s^2 \kappa_{10}^2}. \quad (3.8)$$

From equation (3.7), we see that the potential for the volume modulus $\varphi(x)$ goes roughly as $-R_6 \exp(-\varphi)$. For an internal manifold with negative Ricci curvature, for example H^6, the potential is monotonically decreasing towards infinity. THis causes the volume modulus to run away to infinity (*i.e.*, $\varphi \to \infty$), corresponding to a decompactification of the internal space. On the other hand, a manifold with positive scalar curvature, for example a sphere S^6, leads to a potential for the volume modulus that decreases monotonically to an unbounded negative value towards the origin. A Ricci flat manifold like $\mathbb{R}^6$, on the other hand, fails to generate a potential at all. In neither of these cases is the volume modulus stabilized. This is the essence of the problem of moduli stabilization.

3.2.2 Flux compactification

Stabilizing the volume moduli above requires competing terms in the potential. One possible source for such terms are the various electric and magnetic fluxes present in string theory. Let us consider the previous example again, but this time additionally turn on a constant three-form flux in the compact space. This modifies equation (3.3) to

$$S = \frac{1}{2\kappa_{10}^2} \int \mathrm{d}^{10}x \, e^{-2\Phi} \sqrt{-g_{10}} \left(R - e^{-6\varphi} \left| F_3 \right|^2 \right). \tag{3.9}$$

Repeating the dimensional reduction on the same warped background of equation (3.4), gives

$$S = \frac{1}{2\kappa_4^2} \int \mathrm{d}^4x \sqrt{-g_4} \left(R_4 + 12 \partial_\mu \varphi \partial^\mu \varphi + e^{-8\varphi} R_6 - e^{-12\varphi} \left| F_3 \right|^2 \right), \tag{3.10}$$

which after canonically normalizing the kinetic term for φ becomes

$$S_4 = \frac{1}{2\kappa_4^2} \int \mathrm{d}^4x \sqrt{-g_4} \left(R_4 + \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + e^{-\frac{4}{\sqrt{6}}\varphi} R_6 - e^{-\sqrt{6}\varphi} \left| F_3 \right|^2 \right). \tag{3.11}$$

The volume modulus, therefore has the following potential

$$V(\varphi) = e^{-\sqrt{6}\varphi} \left| F_3 \right|^2 - e^{-\frac{4}{\sqrt{6}}\varphi} R_6. \tag{3.12}$$

The extra contribution from the flux breaks the monotonicity of the potential and allows for stable points. The first term is strictly positive and can compete with the second term to generate a minimum, if the internal manifold has a positive scalar curvature $R_6 > 0$. This toy model shows that it is possible to stabilize a moduli by adding flux in a compact space.

3.3 The KKLT construction

Summarizing the discussion so far, the existence of a stable lower dimensional vacuum requires all moduli to be stabilized. Additionally, for this vacuum to be de Sitter, supersymmetry must be broken. However, Maldacena and Núñez (2001) showed in the form of a *no-go theorem* that under some general assumptions, there are no compactifications of ten or eleven dimensional supergravity that give a dS vacuum solution. Let us briefly summarize the theorem below.

Maldacena-Núñez no-go theorem

Maldacena and Núñez (2001) considered a compactification of D dimensional supergravity on a manifold of dimension $(D - d)$ with finite volume, for example, $D = 10$ and $d = 4$. They further assumed a static warped D dimensional metric and integrated the trace reversed Einstein equations on the internal manifold. Under the assumption that there are no negative tension sources, they proved that the warp factor is constant and no d dimensional de Sitter solutions are allowed. Furthermore, this is a statement in supergravity and assumes that there are no corrections from non-perturbative effects, string loops (g_s) or higher-derivative terms (α').

Constructing a dS vacuum, therefore requires evading one of these assumptions, for example going beyond supergravity or including negative tension sources such as O-planes. We will now briefly review a construction by Kachru, Kallosh, Linde, and Trivedi (2003, KKLT), which is one of the most discussed constructions of a dS vacuum in string theory. This evades the above no-go theorem by making use of non-perturbative corrections.

3.3.1 Flux compactification: IIB on CY_3

The starting point of the KKLT construction is type IIB string theory compactified on a *Calabi-Yau three-fold* (CY_3). A CY_3 is a three complex dimensional (equivalent to six real dimensions) Kähler manifold with a globally defined nowhere vanishing holomorphic three-form Ω. It admits a Ricci flat metric and its holonomy group is SU(3).[1] CY_3 space has several topologically non-trivial three cycles (three dimensional surfaces), the total number of which is denoted by a *Betti number* b_3. They can be written in a basis of three cycles usually denoted by A_I and B^J. A Calabi-Yau space also has a unique covariantly constant spinor, which can be used to construct a unique two-form, called the Kähler form. When the ten dimensional theory is compactified on a CY_3 manifold to get the effective four dimensional theory, these unique two- and three-forms are integrated over two cycles and three cycles to yield the *Kähler moduli* and *complex structure moduli* respectively. The Kähler moduli characterize the size of the CY_3, while the complex structure moduli correspond to its shape. For example, for a rectangular two torus $\mathbb{T}^2$ (which is an example of a CY_1) with size a and b, their product ab corresponds to the Kähler modulus, while their ratio a/b is the complex structure modulus.

Apart from the Kähler and complex structure moduli, which are universal for every Calabi-Yau compactification and correspond to the geometry, there can be additional moduli fields coming from integrating other fields over cycles of the CY manifold. For example, fluxes coming from

[1]See section 4.2 for a definition of the holonomy group of a manifold.

form fields sourced by D-branes or NS-branes can contribute to axions,[2] while position of the D-branes in the space transverse to the CY appear as D-brane moduli, to name a few.

Let us now outline the basics of flux compactification in IIB string theory following the seminal work of Giddings, Kachru, and Polchinski (2002, GKP), which KKLT takes as the starting point. The discussion in this part closely follows the presentation of Baumann and McAllister (2015). The bosonic part of the action of type IIB supergravity in Einstein frame is given by

$$\begin{aligned} S_{\mathrm{IIB}} = &\frac{1}{2\kappa_{10}^2}\int \mathrm{d}^{10}x\sqrt{-g_{10}}\left[R - \frac{|\partial\tau|^2}{2\mathrm{Im}(\tau)^2} - \frac{|G_3|^2}{2\mathrm{Im}(\tau)} - \frac{|\widetilde{F}_5|^2}{4}\right] \\ &- \frac{i}{8\kappa_{10}^2}\int \frac{C_4\wedge G_3\wedge \overline{G}_3}{\mathrm{Im}(\tau)}, \end{aligned} \tag{3.13}$$

where G_3 is defined as $G_3 := F_3 - \tau H_3$, with τ being the axio-dilaton $\tau := C_0 + ie^{-\Phi}$. $\widetilde{F}_5$ is the self dual (*i.e.*$\widetilde{F}_5 = \star_{10}\widetilde{F}_5$) five-form field strength defined as

$$\widetilde{F}_5 := F_5 - \frac{1}{2}C_2\wedge H_3 + \frac{1}{2}B_2\wedge F_3, \tag{3.14}$$

where $F_p := dC_{p-1}$. We want to look for solutions that are warped products of four dimensional Minkowski spacetime with a compact $\mathcal{M}_6$ *i.e.*,

$$\mathrm{d}s^2 = e^{2A(y)}\eta_{\mu\nu}\mathrm{d}x^\mu\mathrm{d}x^\nu + e^{-2A(y)}g_{mn}\mathrm{d}y^m\mathrm{d}y^n. \tag{3.15}$$

With this ansatz, four dimensional Poincaré invariance imposes some restrictions on the form of the fluxes G_3 and F_5. The G_3 flux should have non-zero components only along the compact space, while $\widetilde{F}_5$ takes the following special form:

$$\widetilde{F}_5 = (1+\star_{10})\,\mathrm{d}\alpha(y)\wedge\mathrm{d}x^0\wedge\mathrm{d}x^1\wedge\mathrm{d}x^2\wedge\mathrm{d}x^3, \tag{3.16}$$

where $\alpha(y)$ is an arbitrary function on $\mathcal{M}_6$. Trace of the ten dimensional Einstein equations is given by

$$\nabla^2_{(6)}e^{4A} = \frac{e^{8A}}{2\mathrm{Im}(\tau)}|G_3|^2 + e^{-4A}\left(|\partial\alpha|^2 + \left|\partial e^{4A}\right|^2\right) + 2\kappa_{10}^2 e^{2A}\mathcal{J}_{\mathrm{loc}}, \tag{3.17}$$

where $\nabla_{(6)}$ is the Laplacian on the compact space and $\mathcal{J}_{\mathrm{loc}}$ represents the energy corresponding to localized sources

$$\mathcal{J}_{\mathrm{loc}} = \frac{1}{4}\left(T^m_m - T^\mu_\mu\right). \tag{3.18}$$

[2]In the presence of topologically non-trivial cycles, fluxes can also exist without the need for brane sources, provided the fluxes satisfy the tadpole condition.

Since the left hand side of equation (3.17) vanishes when integrated over $\mathcal{M}_6$, and the only term of indeterminate sign on the right hand side is the source term, it needs to be negative to give a non-trivial solution. Therefore, a non-trivial warped compactification requires the presence of localized sources, for at least some of which $\mathcal{J}_{\text{loc}} < 0$. For $p < 7$, these are negative tension objects, namely Op-planes. This is the essence of the Maldacena-Nùñez theorem, where in the absence of O-planes, the right hand side has to vanish and the warp factor is a constant. The presence of O-planes here allows for a non-trivial warped compactification, hence evading the no-go theorem. The Bianchi identity for $\widetilde{F}_5$ is

$$\mathrm{d}\widetilde{F}_5 = H_3 \wedge F_3 + 2\kappa_{10}^2 T_3 \rho_3^{\text{loc}}, \tag{3.19}$$

where ρ_3^{loc} is the charge density associated with D3 branes and T_3 is the D3 brane tension. Imposing the restriction on the form of $\widetilde{F}_5$ from our metric ansatz in equation (3.16), and subtracting the trace of Einstein's equation in equation (3.17) yields

$$\begin{aligned}\nabla_{(6)}^2\left(e^{4A}-\alpha\right) = \frac{e^{8A}}{24\mathrm{Im}(\tau)}\left|iG_3 - \star_6 G_3\right|^2 + e^{-4A}\left|\partial\left(e^{4A}-\alpha\right)\right|^2 \\ + 2\kappa_{10}^2 e^{2A}\left(\mathcal{J}_{\text{loc}} - \mathcal{Q}_{\text{loc}}\right),\end{aligned} \tag{3.20}$$

where $\mathcal{Q}_{\text{loc}} := T_3\rho_3^{\text{loc}}$ is the charge associated with the localized sources. Similar to equation (3.17), the left hand side of this equation integrates to zero on $\mathcal{M}_6$, implying that the right side must vanish as well. For sources satisfying a BPS-like condition $\mathcal{J}_{\text{loc}} \geq \mathcal{Q}_{\text{loc}}$, when this condition is saturated, each term on the right hand side must vanish. This implies that G_3 is *imaginary self dual* (ISD), $\star_6 G_3 = iG_3$ and the warp factor is related to the potential through $\exp(4A) = \alpha$. Solutions of this form are called *ISD solutions* and are among the best understood solutions in IIB flux compactifications. D3 branes, O3 planes and D7 branes wrapping four-cycles (so as to preserve the $\mathcal{N} = 1$ of the D3 branes) saturate the bound $\mathcal{J}_{\text{loc}} = \mathcal{Q}_{\text{loc}}$ while for $\overline{\text{D3}}$ and D5 wrapping a zero volume two cycle, it is a strict inequality $\mathcal{J}_{\text{loc}} > \mathcal{Q}_{\text{loc}}$. $\overline{\text{O3}}$ and O5 planes violate the bound, *i.e.*, $\mathcal{J}_{\text{loc}} < \mathcal{Q}_{\text{loc}}$.

In summary, from the trace of Einstein's equations, warped compactifications of the form of equation (3.15) are only possible in the presence of negative tension sources like O-planes, while positive tension sources like D-branes give a factorized metric. Combining this with the Bianchi identity for the self dual five-form tells us that for sources satisfying $\mathcal{J}_{\text{loc}} \geq \mathcal{Q}_{\text{loc}}$, only ISD solutions are possible. An important property of these GKP flux compactifications is that, on performing a KK reduction to four dimensional Minkowski space, the complex structure and the axio-dilaton get a potential as is evident from the third term in the IIB action of equation (3.13)

$$S \subset -\frac{1}{2\kappa_{10}^2}\int_{\mathbb{R}^{1,3}} \mathrm{d}^4x \int_{\mathcal{M}_6} \mathrm{d}^6y\sqrt{-\mathrm{g}_6}e^{-2A(y)}\frac{g^{pq}g^{nr}g^{ps}G_{mnp}G_{qrs}}{2\mathrm{Im}(\tau)}. \tag{3.21}$$

This involves the axio-dilaton both through the denominator and the definition of G_3, as well as the complex structure modulus through the factors of the six dimensional metric in the numerator. The effective four dimensional theory obtained from the ISD compactification above can be written in terms of a Kähler potential and a superpotential of $\mathcal{N} = 1$ supergravity. The superpotential is given by the Gukov, Vafa, and Witten (2000, GVW) superpotential

$$W = \int_{\mathcal{M}_6} G_3 \wedge \Omega. \tag{3.22}$$

The Kähler potential is

$$K = -3\ln\left(-i\left(\rho - \bar{\rho}\right)\right) - \ln\left(-i\left(\tau - \bar{\tau}\right)\right) - \ln\left(-i\int_{\mathcal{M}_6} \Omega \wedge \overline{\Omega}\right), \tag{3.23}$$

where ρ is the (universal) volume modulus of the compact manifold, the holomorphic three-form Ω depends on the complex structure moduli and G_3 depends on the dilaton. The superpotential is independent of the volume modulus, and the Kähler potential is of the form $K^{\rho\bar{\rho}}\partial_\rho K \partial_{\bar{\rho}} K = 3$ (this is called the *no-scale* structure). This implies that the scalar potential depends on the volume modulus only through the exponent of the Kähler potential to give

$$\begin{aligned} V &= e^{K/\mathrm{M_{Pl}}^2}\left[K^{i\bar{j}}D_iW\overline{D_jW} - \frac{3}{\mathrm{M_{Pl}}^2}\left|W\right|^2\right] \\ &= e^{K/\mathrm{M_{Pl}}^2}\left[K^{a\bar{b}}D_aW\overline{D_bW}\right] \geq 0, \end{aligned} \tag{3.24}$$

where i, j run over all moduli and a, b run over all except the volume modulus ρ. Therefore, all moduli except the volume modulus are stabilized in the presence of the G_3 flux. The real part of the volume modulus does not get a potential at all and remains flat, while the imaginary part gets contribution from the exponent of the Kähler potential, leading to a runaway.

3.3.2 Moduli stabilization: non-perturbative corrections

Non-perturbative corrections

In order to generate a potential for the volume modulus, it either needs to appear in the superpotential or the no-scale structure of the Kähler potential needs to be corrected. Although the superpotential and the Kähler potential considered above are at lowest order in both the perturbation parameters α' and g_s, non-renormalization theorems protect the superpotential from getting any perturbative corrections.[3] Thus, one either needs to consider quantum corrections to the Kähler potential, or look for non-perturbative corrections to the superpotential. The KKLT construction skips all perturbative g_s and α' corrections to the Kähler potential, and instead uses non-perturbative corrections to the superpotential. These corrections have two sources in IIB:

1. *Gaugino condensation*: If there is a stack of N D7 branes extending along $\mathbb{R}^{1,3}$ and wrapping a four-cycle in $\mathcal{M}_6$, the world volume theory on them is a super Yang-Mills theory with gauge group SU(N). There can be spontaneous symmetry breaking in such a theory when the gaugino bilinear $\langle\lambda\bar{\lambda}\rangle$ gets a non-zero vacuum expectation value, *i.e.*, gaugino condensation. This generates a superpotential non-perturbatively
$$W_{\text{gaugino}} = \mathcal{A}e^{-\frac{2\pi i}{N}\rho}, \tag{3.25}$$
where $\mathcal{A}$ is independent of ρ but can depend on other moduli.
2. *D3-brane instantons*: A similar non-perturbative contribution to the superpotential arises when a Euclidean D3 brane wraps a four-cycle in $\mathcal{M}_6$ and is given by
$$W_{\text{instanton}} = \mathcal{A}e^{-2\pi i\rho}, \tag{3.26}$$
where, as before, the prefactor $\mathcal{A}$ is independent of ρ but can depend on the other moduli, in general.

Together, these non-perturbative contributions can be written as
$$W_{\text{non-pert}} = \mathcal{A}e^{-ia\rho}, \tag{3.27}$$
with the inclusion of which, the corrected superpotential becomes
$$W = W_{\text{GVW}} + W_{\text{non-pert}} = W_0 + \mathcal{A}e^{-ia\rho}. \tag{3.28}$$
W_0 denotes the GVW superpotential from equation (3.22) which, as we have seen, is independent of ρ.

Supersymmetric AdS

Since all other moduli except for the volume modulus are stabilized, the relevant Kähler potential and scalar potential are given by
$$K = -3\ln\left[-i\left(\rho - \bar{\rho}\right)\right], \quad W = W_0 + \mathcal{A}e^{-ia\rho}. \tag{3.29}$$

[3]Holomorphy and shift symmetry protect against α' corrections, while the absence of perturbative g_s corrections in IIB is more delicate and was shown by Burgess, Escoda, and Quevedo (2006).

In terms of a nonzero axion $\rho = i\sigma$, the scalar potential constructed from the above becomes

$$V = \frac{a\mathcal{A}e^{-a\sigma}}{2\sigma^2}\left[\left(1+\frac{a}{3}\sigma\right)\mathcal{A}e^{-a\sigma}+W_0\right], \tag{3.30}$$

which has a supersymmetric minimum at $\sigma = \sigma_\star$ defined by $D_\rho W|_{\rho=i\sigma_\star} = 0$, with W_0

$$W_0 = -\mathcal{A}e^{-a\sigma_\star}\left(1+\frac{2a\sigma_\star}{3}\right). \tag{3.31}$$

Consequently, the potential at the supersymmetric minimum is negative.

$$V(\sigma_\star) = -\frac{a^2\mathcal{A}^2}{6}\frac{e^{-a2\sigma_\star}}{\sigma_\star} < 0. \tag{3.32}$$

3.3.3 dS uplift: anti-D3 branes

So far, we have seen that starting from a GKP compactification and adding non-perturbative corrections to the superpotential, fixes all moduli including the volume modulus and gives a four dimensional theory with a supersymmetric AdS minimum. The final step to obtaining a de Sitter vacuum is to break supersymmetry and *uplift* the minimum to dS. This is done in the KKLT construction by adding $\overline{\text{D3}}$ branes.

Considering the G_3 flux above, we realize that if it is sourced by positive tension sources like D3 branes, the internal space cannot be compact. This follows from the tadpole cancellation condition. A non-compact six manifold called the Klebanov and Strassler (2000) geometry, can be used instead. This geometry is obtained by starting from a non-compact CY_3 called the *conifold* (which is a metric cone over the base of a manifold called $\text{T}^{1,1} \simeq S^2 \times S^3$, *i.e.*, $\mathrm{d}s^2_{\text{CY}_3} = \mathrm{d}r^2 + r^2\mathrm{d}s^2_{\text{T}^{1,1}}$), and deforming it so that the S^3 does not vanish at the tip – hence called the *deformed conifold*. The coordinate going from the tip to the base of the cone is the coordinate y of equation (3.15). The volume of this manifold goes from a finite value to infinity at $y \to \infty$. This can be smoothly glued to a compact CY_3 (with the addition of O3 planes to satisfy the tadpole condition) at some $y = y_{\text{max}}$. The ISD fluxes can then be added as before, to support this geometry. The volume of the compact space therefore has a minimum value given by

$$\exp\left(A_{\text{min}}\right) \sim \exp\left(-\frac{2\pi K}{3g_s M}\right), \tag{3.33}$$

where K and M are integers giving the quantized values of the three-form fluxes

$$\frac{1}{4\pi^2\alpha'}\int_A F_3 = M, \quad \frac{1}{4\pi^2\alpha'}\int_B H_3 = K. \tag{3.34}$$

These can be chosen so that $\exp A_{\mathrm{min}} \ll 1$, which means that owing to the strong warping, physics at the tip can be considered independently of that away from it.

To break supersymmetry, it suffices to add $\overline{\mathrm{D3}}$ branes at some $y = y_0$. Since the background contains D3 branes, this breaks all supersymmetry and contributes an extra energy density

$$V_{\overline{\mathrm{D3}}} = 2g_s^{-4} e^{4A(y_0)} T_3 \frac{1}{\mathrm{Im}(\rho)^3} \equiv \frac{D}{\sigma^3}. \tag{3.35}$$

The factor $\exp[4A(y_0)]$ causes the $\overline{\mathrm{D3}}$ to be attracted to the tip of the conifold, $y = y_{\mathrm{min}}$, giving a small positive contribution to the potential proportional to D as defined above. With this, the potential of equation (3.30) becomes

$$V = \frac{a\mathcal{A}e^{-a\sigma}}{2\sigma^2}\left[\left(1 + \frac{a}{3}\sigma\right)\mathcal{A}e^{-a\sigma} + W_0\right] + \frac{D}{\sigma^3}, \tag{3.36}$$

where D is a small number, and uplifts the potential to produce a metastable dS minimum with a small positive cosmological constant.

3.4 Discussion

The KKLT construction outlined above, being one of the most concrete models of dS in string theory today, has been the subject of intense scrutiny. Several of its aspects are still a matter of active discussion and a consensus on its validity has not yet been reached. Some of the major points of discussion are the following.

A zeroth order objection: flux compactification

The first point of debate is already at the level of the GKP flux compactification, before performing moduli stabilization or the anti-brane uplift. The non-zero W_0 needed for the above construction breaks supersymmetry already in the GKP construction and one should, in principle, include perturbative string loop g_s and higher curvature corrections α', even before adding the non-perturbative corrections considered above. Since these perturbative corrections are not yet completely known in type IIB, there is a debate on whether or not these are small enough to be counteracted by the non-perturbative corrections. Sethi (2018), who pointed out this issue, argued that these corrections could lead to a runaway, while Kachru and Trivedi (2019) argued against it. A resolution of this problem requires a better understanding of instantons in time-dependent backgrounds, of which not much is currently known.

A first order objection: moduli stablization
The second point of discussion concerns the non-perturbative terms used for stabilizing the volume modulus. Moritz, Retolaza, and Westphal (2018) found a tension between the results obtained from gaugino condensation in ten dimensions and the four dimensional result by KKLT. This result was improved and elaborated on by Gautason, Van Hemelryck, and Van Riet (2019). It has been argued that the original objection was based on an incomplete treatment in ten dimensions and proposals were made to cure the problem (Hamada, Hebecker, Shiu, and Soler, 2019a,b; Kallosh, 2019). However this has recently been challenged (Carta, Moritz, and Westphal, 2019; Gautason, Van Hemelryck, Van Riet, and Venken, 2020), and it has been suggested that it might be difficult to achieve this in a controlled way.

A second order objection: dS uplift
A third problem concerns the backreaction of the $\overline{\mathrm{D3}}$ branes used for the uplift to a dS vacuum. The three-form flux associated to $\overline{\mathrm{D3}}$ branes at the tip of the KS throat are singular, and it is not yet clear if these singularities are physical. A resolution of this would be to construct the corresponding non-singular solution or to prove that no such solution exists. There has been extensive discussion around this in the literature (Bena, Graña, and Halmagyi, 2010; Blåbäck, Danielsson, and Van Riet, 2013; Danielsson and Van Riet, 2015; Bena, Blåbäck, and Turton, 2016; Danielsson, Gautason, and Van Riet, 2017, among others). It has also been argued that $\overline{\mathrm{D3}}$ brane polarization could resolve the singularity (Cohen-Maldonado, Diaz, van Riet, and Vercnocke, 2016; Armas, Nguyen, Niarchos, Obers et al., 2019).

3.5 Other dS constructions

The KKLT construction just outlined, is a compactification of type IIB string theory. Several other attempts to construct dS vacua have also been made in type IIA, IIB, and heterotic string theory. Let us briefly outline them below.

Type IIA
A feature of type IIA string theory is that the R-R and NS-NS fluxes are even and odd forms respectively, giving the possibility to stabilize all moduli classically, leading to a supersymmetric AdS vacuum. This was shown by DeWolfe, Giryavets, Kachru, and Taylor (2005). In order to promote this to a dS vacuum, one needs to evade the Maldacena-Núñez no-go theorem. To do this, one can either include quantum corrections or add classical negative tension objects like O-planes. As for the choice of internal

manifold, Hertzberg, Kachru, Taylor, and Tegmark (2007) showed that a Ricci flat manifold CY_3 ($R_6 = 0$) gives $|\nabla V|^2 / V^2 \geq 54/13$, which excludes dS extrema. dS vacua have, however, been constructed in *massive* type IIA (*i.e.*, including a zero-form field strength F_0, which appears as a parameter in the theory, and is called the *Romans' mass*) compactified on a negatively curved manifold ($R_6 < 0$), and in the presence of O6-planes. The first such constructions were by Caviezel, Koerber, Kors, Lust et al. (2009), and Flauger, Paban, Robbins, and Wrase (2009). This was followed by further constructions within this framework by Danielsson, Haque, Shiu, and Van Riet (2009) and Danielsson, Koerber, and Van Riet (2010). Subsequently, a systematic scan of possible solutions of this type was performed by Danielsson, Haque, Koerber, Shiu et al. (2011), and it was found that all dS vacua constructed using these ingredients contain tachyons. Recently, such models were also examined by Roupec and Wrase (2019). They showed that because of the intersecting O-planes necessary for these models, the ten dimensional equations of motion are not solved pointwise but only as an integral over the internal space (usually referred to as *smeared sources*).[4] Moreover, all such solutions were found numerically and have not been shown to exist in the limit of weak string coupling and small curvature, when the fluxes are properly quantized. Another class of constructions is by Córdova, De Luca, and Tomasiello (2019a,b), who numerically solved full ten dimensional supergravity equations of massive type IIA, in the presence of O8-planes; however, a string theory uplift of these solutions remains to be constructed.

Adding *non-geometric fluxes* (these are fluxes that are obtained via dualities from other perturbative string theories, but that don't have a geometric source in the theory being considered) introduces more parameters to the problem and helps in the search for a dS vacuum. The first fully stable (*i.e.*, free from tachyons) stable classical dS solutions were constructed by de Carlos, Guarino, and Moreno (2010a,b). Other tachyon free metastable dS vacua were subsequently constructed in even more general settings (Danielsson and Dibitetto, 2013; Blåbäck, Danielsson, and Dibitetto, 2013; Damian, Diaz-Barron, Loaiza-Brito, and Sabido, 2013; Blåbäck, Danielsson, Dibitetto, and Vargas, 2015).

Type IIB

The situation is quite different in type IIB, where fluxes can be used to stabilize the dilaton and the complex structure moduli, but the Kähler moduli have a no-scale structure, which keeps them massless at the classical

[4]Also see the article by Blåbäck, Danielsson, Junghans, Van Riet et al. (2010) for a discussion on smeared vs localized sources for constructing vacuum solutions in type II supergravity and its consequences for a string theory uplift.

level. There are two well known mechanisms of stabilizing them – the KKLT scenario and the *Large Volume Scenario* (LVS) by Balasubramanian, Berglund, Conlon, and Quevedo (2005). As we have just seen, in the KKLT scenario, the superpotential receives non-perturbative corrections and W_0, which is non-zero, is fine-tuned against it to give a supersymmetric AdS minimum. In contrast, W_0 is not fine-tuned in LVS. Rather, the Kähler potential receives perturbative corrections, which are fine-tuned against the non-perturbative corrections to the superpotential, to give a non-supersymmetric AdS. The volume $\mathcal{V} \sim \exp(1/g_s) \gg 1$ is exponentially large (in string units) in this minimum, hence the name *large volume*.

Several dS constructions have been proposed in type IIB that include extra ingredients. The KKLT construction presented above is one such case where the extra ingredient in an $\overline{\text{D3}}$ brane. While in this construction the uplift to dS is carried out as an extra step after stabilizing the Kähler moduli, in some other constructions, this can also be achieved at the same time as moduli stabilization. Some of the key ingredients that have been used to obtain dS are higher curvature terms by Westphal (2007), non-perturbative effects at singularities by Cicoli, Maharana, Quevedo, and Burgess (2012), a particular brane configuration called T-branes by Cicoli, Quevedo, and Valandro (2016), and a flux induced supersymmetry breaking by Gallego, Marsh, Vercnocke, and Wrase (2017), to name a few. An extensive discussion and review of several other aspects of dS construction in string theory can be found in the reviews by Danielsson and Van Riet (2018) and Cicoli, De Alwis, Maharana, Muia et al. (2019). A recent class of solutions was found numerically by Andriot, Marconnet, and Wrase (2020) but they all contain tachyons.

Heterotic string theory

Kutasov, Maxfield, Melnikov, and Sethi (2015) showed that heterotic string theory does not admit dS vacua – to all orders in α' (but not in g_s). There are several attempts to construct dS vacua using non-perturbative effects in heterotic string theory, but the explicit dS vacua found were all tachyonic (Gukov, Kachru, Liu, and McAllister, 2004; Parameswaran, Ramos-Sanchez, and Zavala, 2011; Anderson, Gray, Lukas, and Ovrut, 2011; Cicoli, de Alwis, and Westphal, 2013, among others).

3.6 What about the fine tuning problem?

The KKLT construction outlined above, tries to address the question of a positive cosmological constant in string theory but the fine tuning problem still remains. In the initial days of string theory, it was hoped that the theory would uniquely determine a vacuum that would be the correct vacuum of our universe, and would uniquely specify the values of all fundamental constants. However, it was soon realized that this is not the case and that string theory allows for an enormous number of consistent vacuum solutions. This is dubbed the *landscape* of string theory vacua and was estimated by Susskind (2003) to be of the order of 10^{100} or more.[5] These correspond to minima in a multidimensional moduli space. The universe could then start in any of these vacua and parts of it could tunnel into other vacua. This could repeat indefinitely, to give a vast universe with all possibilities realized at one place or another – an idea called the *multiverse*. In fact, Linde (2017) has argued that our universe might indeed be a part of such a multiverse. Why then do we live in precisely the vacuum that we live in? One way of arguing for this is the *anthropic principle*. A decade before the discovery of the accelerating universe in 1998, Weinberg (1989) had argued that for structure formation to be complete before the cosmological constant comes to dominate the energy density of the universe, the vacuum energy density should be roughly

$$\rho_\Lambda \sim (10 - 100)\, \rho_m^{(0)}. \tag{3.37}$$

He argued that a value slightly less than this would be the most natural, but it should not be zero, since that would require fine-tuning. This indeed turned out to be only an order of magnitude or two away from the measured value. This was an argument based on the anthropic principle – that if the cosmological constant was too different from this value, galaxies would not have formed and we would not be here to ask this question.

Several non-anthropic mechanisms have also been proposed to explain the smallness of the cosmological constant. One of these is a mechanism of *saltatory relaxation of the cosmological constant* proposed by Feng, March-Russell, Sethi, and Wilczek (2001). This uses an older idea by Brown and Teitelboim (1987, 1988) that nucleation of charged branes could lower the cosmological constant. which we will briefly review in section 7.1.3. They proposed a mechanism by which the cosmological constant can relax from a non fine-tuned value to its present value, via nucleation of a stack of branes. This decay to its present value happens quickly enough, but it then remains at its present value for a very long time.

Another proposal by Kane, Perry, and Zytkow (2003) suggests that the mixing of a large number of connected degenerate vacua in string theory will lead to a state with much lower vacuum energy than any of the individual vacua. Another radical solution to the cosmological constant problem

[5]More recent estimates, for example, by Taylor and Wang (2015) give a much larger number of possibilities from flux compactifications.

in string theory was proposed by two groups: Arkani-Hamed, Dimopoulos, Kaloper, and Sundrum (2000), and Kachru, Schulz, and Silverstein (2000). Instead of arguing for a small cosmological constant, they described a mechanism that decouples the expansion of the four dimensional universe from the value of the cosmological constant.

Solving the cosmological constant problem is an enormously difficult task, and we still don't have a string theoretic construction of a dS vacuum that is completely under control. Given this, we will only focus on the "lesser" problem of obtaining dS vacuum in the rest of the thesis, without worrying about the fine tuning problem.

4. de Sitter in M-theory

So far, we have summarized attempts of constructing de Sitter vacua in string theory and highlighted some of the difficulties involved. In this chapter, we will turn our attention to eleven dimensional M-theory and examine de Sitter constructions there. We will then present a construction that we proposed in paper I, and discuss it in the context of related work.

4.1 What is M-theory?

M-theory is a eleven dimensional theory whose low energy limit is eleven dimensional supergravity. The origin of the name *M-theory* lies in a famous article by Hořava and Witten (1996). In this section, following the original motivation by Witten (1995), we will outline how such a theory arises from considering the strong coupling limit of type IIA string theory. A more thorough, but pedagogical, introduction to M-theory can be found in the textbook by Becker, Becker, and Schwarz (2006) or the one by Ibáñez and Uranga (2012). A list of the historical benchmarks leading to the development of M-theory and following its progress up to the early 2000's, can be found in an article by Duff (2004).

String theory is a theory of fundamental strings (F1), as well as higher dimensional objects: Dp branes and NS5 branes. In terms of the two parameters of string theory, α' and g_s, the tension of these objects goes as

$$T_{\mathrm{F1}} \sim (\alpha')^{-1}, \quad T_{\mathrm{NS5}} \sim (\alpha')^{-3} g_s^{-2}, \quad T_{\mathrm{D}p} \sim (\alpha')^{-(p+1)/2} g_s^{-1}. \tag{4.1}$$

The mass scales corresponding to the tension of these objects can be read off from dimensional analysis, $M_{\mathrm{p}} \sim T_{\mathrm{p}}^{1/(p+1)}$, to give

$$M_{\mathrm{F1}} \sim (\alpha')^{-1/2}, \quad M_{\mathrm{NS5}} \sim (\alpha')^{-1/2} g_s^{-1/3}, \quad M_{\mathrm{D}p} \sim (\alpha')^{-1/2} g_s^{-1/(p+1)}. \tag{4.2}$$

Let us examine them in the weak and strong coupling regimes. In the weakly coupled limit ($g_s \ll 1$), the NS5 brane and Dp branes become extremely massive and their excitations decouple from the theory, leaving behind a theory of fundamental strings. However, in the strong coupling limit ($g_s \to \infty$), the lightest objects are not fundamental strings, but rather the lowest dimensional Dp brane. For type IIB string theory, this corresponds to a D1 brane, which is a string. This suggests that weakly coupled and strongly coupled type IIB string theory are dual to each other, *i.e.*, the $SL(2,\mathbb{Z})$ self duality of type IIB.

For type IIA string theory, however, the lightest object is a D0 brane, which is a particle and not a string. This suggests that the strongly coupled limit of type IIA theory is not a string theory, but something else. To understand what this might be, consider a bound state of n D0 branes. The mass of such a bound state goes as

$$M_{n\mathrm{D0}} \sim n\,(\alpha')^{-1/2}\, g_s^{-1} \sim \frac{n}{g_s}, \text{ with } n \in \mathbb{N}. \tag{4.3}$$

This is reminiscent of the mass spectrum of momentum modes obtained from a KK reduction on a circle of radius R, in which case one gets $m \sim k/R$, with $k \in \mathbb{N}$. This analogy suggests that type IIA string theory in the strong coupling limit, is dual to an eleven dimensional theory with the eleventh dimension compactified on a circle – the string coupling corresponding to the radius of the circle. Strong coupling, therefore, corresponds to the decompactification limit of the circle. This theory is known as *M-theory*, and has only one dimensionful scale – the eleven dimensional length scale ℓ_{11}. This should be contrasted with string theory, which has a dimensionful parameter namely, the string length $\ell_s \sim \sqrt{\alpha'}$, and a dimensionless parameter g_s that the theory can be expanded around. In terms of ten dimensional quantities, the radius of the eleventh dimension is given by $R = \ell_s g_s = \sqrt{\alpha'} g_s$, while the eleven dimensional length can be written as $\ell_{11} = g_s^{1/3}\ell_s = (\alpha' R)^{1/3}$. The absence of a dimensionless parameter makes a perturbative formulation of the theory difficult, and there is still no microscopic formulation of M-theory.

The low energy limit of M-theory, on the other hand, is well understood and is given by eleven dimensional supergravity, whose bosonic action can be written down as

$$S_{11\mathrm{d}} = \frac{1}{2\kappa_{11}^2}\int \mathrm{d}^{11}x\sqrt{-g}\left(R - \frac{|G_4|^2}{2}\right) - \frac{1}{6}\int_{\mathcal{M}_{11}} C_3 \wedge G_4 \wedge G_4. \tag{4.4}$$

The massless bosonic field content of the theory is an eleven dimensional metric g and a three-form gauge field C_3, with field strength G_4. It can be shown that performing a dimensional reduction along a circle indeed gives the action of type IIA supergravity, making the above analogy explicit. Under this dimensional reduction, the eleven dimensional metric reduces to the ten dimensional metric, the RR one-form C_1, and the dilaton; while the three-form C_3 gives the RR three-form and the NS-NS two-form of type

IIA string theory. The three-form gauge field in M-theory is sourced by a two-brane M2, while the dual six-form gauge field is sourced a five-brane M5. These are the only BPS branes in M-theory and have been explicitly constructed as solutions in supergravity.

Let us now account for the extended objects in type IIA and see how they are lifted to M-theory. Type IIA has Dp branes for even p, which are charged under the RR gauge fields, as well as F1 and NS5 branes, which are charged under the NS-NS two-form. We have seen that the D0 brane corresponds to the KK momentum modes. The D2 and $F1$ correspond to M2 branes and its compactification on the S^1 respectively. Analogously, NS5 and D4 branes correspond to the M5 brane, and its compactification on the circle respectively. This leaves the D6 brane, which has a special interpretation in M-theory.

D6 brane in IIA = KK monopole in M-theory

The D0 brane is charged under a RR one-form C_1 in type IIA, which has a corresponding two-form field strength F_2. The D6 brane is charged precisely under the magnetic dual of this two-form ($F_8 = \star_{10} F_2$, corresponding to the C_7 gauge field, under which the D6 is charged). Since the D0 brane lifts to a KK momentum mode in M-theory, we can already anticipate that the D6 brane uplifted to M-theory, could have something to do with geometry as well. This is indeed the case. More precisely, the supergravity solution of a D6 brane sources the metric, the dilaton and the RR one-form C_1. As discussed above, all of these three fields are components of the eleven dimensional metric in M-theory, and hence the M-theory lift of a D6 brane corresponds to a purely geometric background in eleven dimensions. This was worked out to be the Taub-NUT geometry, otherwise known as the KK monopole. These will play an important role in the next sections. For more details on the Taub-NUT geometry, see the book by Ortin (2015).

4.2 G_2 holonomy and G_2 structure

In order to obtain a four dimensional effective theory from M-theory, the seven extra dimensions need to be compactified. Compactifications that preserve some amount of supersymmetry in lower dimensions are stable, and do not contain tachyons. However, preserving too much supersymmetry ($\mathcal{N} \geq 2$) can place severe restrictions. For example, theories with $\mathcal{N} = 2$ do not allow for chiral fermions, making them unattractive for particle physics phenomenology. A useful approach is therefore, to consider compactifications that preserve $\mathcal{N} = 1$ supersymmetry.

The amount of supersymmetry in the four dimensional theory corresponds to the number of number of supercharges that are globally well defined on the internal seven manifold $\mathcal{M}_7$. Generically, $\mathcal{M}_7$ is curved, implying that parallel transporting a supercharge around a closed loop will bring it back to a rotated version of itself. As a consequence, the supercharge is not globally well defined. The condition for having a globally well defined supercharge is therefore the condition that the internal manifold admits a covariantly constant spinor. This is called a *Killing spinor*, and is defined as a spinor ξ that is covariantly constant, *i.e.*, $\nabla_{\mathcal{M}_7}\xi(\mathcal{M}_7) = 0$. This condition can be restated in terms of the *holonomy group* of the internal manifold – which is the group of rotations that a spinor suffers when it is parallel transported along all possible closed curves on the manifold. Generically, one would expect this to be SO(7) for the seven manifold, leaving no spinor invariant and giving a non-supersymmetric compactification. In order to preserve some supersymmetry in four dimensions, the holonomy group has to be a subgroup of SO(7), with at least one preserved spinor. The exceptional group G_2 is one such group, and compactifying on a seven dimensional manifold with G_2 holonomy leads to a $\mathcal{N} = 1$ theory in four dimensions. A CY_3 manifold, for example has SU(3) holonomy, instead of the full SO(6), which gives $\mathcal{N} = 2$ supersymmetry in four dimensions when type II string theory is compactified on it. This can be orientifolded to project out half of the supersymmetry and get $\mathcal{N} = 1$, if desired. Contrary to CY_3, a seven manifold with G_2 holonomy is much more difficult to construct. One reason for this is that, being odd dimensional, one cannot apply the tools of complex geometry to construct such manifolds. The only known explicit construction of a seven dimensional manifold with G_2 holonomy is a smooth torus orbifold $\mathbb{T}^7/\mathbb{Z}_2^3$ constructed by Joyce (1996a,b). Starting with an orbifold obtained by quotienting a square seven torus $\mathbb{T}^7$ with unit length, by three $\mathbb{Z}_2$ orbifold actions ($\mathbb{T}^7/\mathbb{Z}_2^3$), Joyce provided a constructive way to resolve the $\mathbb{Z}_2^3$ singularities to obtain a smooth G_2 holonomy manifold.

Twisted torus and G_2 structured manifold

A seven torus is defined as a seven real dimensional manifold with periodic coordinates x^m, with $m = 1, 2, \dots, 7$, where $x^m \simeq x^m + 2\pi R$. Tangent one-forms can be defined on the torus: $\eta^m = \mathrm{d}x^m$, with $\mathrm{d}\eta^m \equiv 0$. A *twisted torus* corresponds to the one-forms being linear in the coordinates, *i.e.*,

$$\mathrm{d}\eta^m = -\frac{1}{2}\omega^m_{np}\eta^n \wedge \eta^p, \tag{4.5}$$

with ω^m_{np} corresponding to the *twisting* of the coordinate x^m over the directions x^n and x^p. The twists are also referred to as *metric fluxes*. Compactifying on a twisted torus corresponds to turning on metric fluxes, which can help stabilize some of the moduli. Starting from the $\mathbb{T}^7/\mathbb{Z}_2^3$ manifold of Joyce, Dall'Agata and Prezas (2004, 2005) constructed twisted orbifolds of the seven torus and studied M-theory compactifications on it. As a result of the twisting, the holonomy group of the deformed seven manifold is no longer G_2, but the structure group of its tangent bundle is G_2 instead. Such a manifold is called a *G_2 structured manifold*. The twisted $\mathbb{T}^7/\mathbb{Z}_2^3$ is an example of such a manifold. The four dimensional theory obtained from compactifying M-theory on this manifold is $\mathcal{N} = 1$ coupled to seven chiral multiplets, the scalar sector of which consists of seven complex moduli fields: S, T_i, and U_i, for $i = 1, 2, 3$. Danielsson, Dibitetto, and Guarino (2015) showed that the four dimensional theory obtained from compactifying M-theory on a twisted $\mathbb{T}^7/\mathbb{Z}_2^3$ can resemble that obtained from compactifying type IIA string theory on an orientifolded twisted $\mathbb{T}^6/\mathbb{Z}_2^2$. In the weak coupling limit, where the type IIA description is valid, these complex moduli correspond to: the axio-dilation S, three complex structure moduli U_i, and three volume moduli T_i – fourteen real scalars in total.

The Kähler potential and the superpotential for the four dimensional $\mathcal{N} = 1$ theory obtained from compactifying M-theory on the twisted $\mathbb{T}^7/\mathbb{Z}_2^3$ is given by

$$K = -\sum_{i=1}^{7} \log\left(\mathrm{Im}\,\Phi_i\right), \qquad W = f_0 + f_i\Phi_i + f_{ij}\Phi_i\Phi_j, \tag{4.6}$$

where $\Phi_i \in \{S, T_i, U_i\}$. The constant term, f_0, in the superpotential comes from the G_7 flux, the linear terms, f_i, come from the G_4 flux and the quadratic terms correspond to the metric fluxes. Cubic order terms in the superpotential, beyond the ones included here, would correspond to non-geometric fluxes in M-theory. It was shown by Danielsson, Dibitetto, and Guarino (2015) that, among the quadratic fluxes, those corresponding to T^2 and ST are non-geometric in type IIA but geometric in M-theory. A complete map of the fluxes, and their corresponding contribution to the superpotential, can be found in table 1 of paper I.

4.3 dS in M-theory

The challenges for constructing a dS vacuum in M-theory are similar to those in string theory. Most notably, the Maldacena-Nùñez no-go theorem forbids compactifications to dS vacua in classical supergravity. Hence, constructing a dS vacuum in M-theory requires adding extra ingredients, for example, non-perturbative corrections or higher-derivative corrections. In the KKLT construction from the previous chapter, non-perturbative effects were necessary to stabilize the volume modulus, since the potential, after a GKP flux compactification, was a no-scale Minkowski with a flat volume modulus. But, as also argued in the previous chapter, the effect of adding such terms is not well understood. An alternative to non-perturbative corrections is to add non-geometric fluxes to uplift the Minkowski vacuum. However, Blåbäck, Danielsson, Dibitetto, and Vargas (2015) showed that some vacua constructed using non-geometric fluxes, can alternatively be obtained from non-perturbative terms, thus placing non-geometric fluxes and non-perturbative effects on a similar footing.

Many of the constructions mentioned in section 3.5 can be obtained from M-theory. However, most of them involve non-geometric fluxes. As an example, since the Romans' mass is a non-geometric flux in M-theory, all massive type IIA constructions are non-geometric in M-theory. Some type IIB constructions mentioned before, for example the one by Blåbäck, Danielsson, Dibitetto, and Vargas (2015) can be T-dualized to type IIA and then be lifted to M-theory. Furthermore, some fluxes (T^2 and ST in particular) that are non-geometric in type IIA, are geometric in M-theory. In addition to the dS constructions mentioned in section 3.5, many other metastable dS solutions have been constructed using non-perturbative effects (Blåbäck, Roest, and Zavala, 2014; Danielsson and Dibitetto, 2014; Guarino and Inverso, 2016). Four dimensional de Sitter solutions have also been explicitly constructed in M-theory. An example is the recent article by Cribiori, Kallosh, Linde, and Roupec (2020), who constructed a dS vacuum from a compactification of M-theory on the twisted $\mathbb{T}^7/\mathbb{Z}_2^3$, using KKLT type non-perturbative contributions to construct metastable dS solutions discussed in the previous chapter. Another construction (Acharya, Bobkov, Kane, Kumar et al., 2007; Kane and Winkler, 2019) provided an inflationary solution in M-theory compactified on an *untwisted* $\mathbb{T}^7/\mathbb{Z}_2^3$, and argued that this could be used to construct a dS vacuum via non-perturbative corrections, without the need for uplifting with antibranes. As mentioned, all of them involve either non-geometric or non-perturbative ingredients, that have been argued to not be completely understood in string theory.

4.4 dS from higher-derivative corrections

The supergravity action, should in principle, receive corrections from an infinite series of higher-derivative terms. These terms involve more derivatives of the metric, the gauge fields, or both. Such corrections have been extensively studied in type IIA string theory, where the first non-zero correction comes from terms that are quartic in the Riemann tensor (denoted by R^4). Having six more derivatives than the Ricci scalar, these terms come with a coefficient of $\ell_s^6 \sim (\alpha')^3$. Such terms do not contribute to the superpotential because of the non-renormalization theorems mentioned in the previous chapter. They do, however, contribute to the Kähler potential; for example, for type IIA string compactified on a CY_3, $\Delta K \sim (\alpha')^3 / \mathrm{vol}\,(\mathrm{CY}_3)$. Such perturbative corrections to the supergravity action provide another way to evade the Maldacena-Nùñez no-go theorem. Being dependent on the volume, they can be useful to stabilize the volume modulus. See the articles by Berg, Haack, and Kors (2005, 2006), for example, where such corrections have been extensively studied in type IIB string theory.

Higher-derivative corrections to the eleven dimensional supergravity action coming from M-theory have also been studied extensively (Antoniadis, Ferrara, Minasian, and Narain, 1997; Green, Gutperle, and Vanhove, 1997; Green and Vanhove, 1997; Russo and Tseytlin, 1997, among others). The first non-zero correction in this case comes from the eight derivative term R^4, as well as terms containing products of the Riemann tensor with the four-form field strength G_4, and its derivatives, related to each other by supersymmetry. From dimensional analysis, all such terms are proportional to ℓ_{11}^6 and contribute at the same order. The eleven dimensional Planck length can therefore be used to keep track of the order of the higher derivative terms, similar to α' in string theory. See the article by Weissenbacher (2020) for the form of currently known higher-derivative corrections in M-theory. In paper I, we used such corrections to construct a metastable dS vacuum, without using any exotic ingredients such as non-perturbative effects, non-geometric fluxes or anti-branes. We outline the construction below, highlighting crucial differences with other dS constructions in the literature. For details of the construction, please refer to paper I.

4.4.1 Non-supersymmetric Minkowski

The starting point for our dS construction is a new class of non-supersymmetric Minkowski vacua that we also constructed in paper I. These are constructed analytically using only metric fluxes in M-theory (corresponding to KK monopoles), which in the language of the STU model introduced before, implies that the superpotential only contains quadratic terms. However, this also contains T^2 and ST terms, which as mentioned before, are geometric in M-theory but non-geometric in type IIA string theory. These Minkowski vacua turn out to be free from tachyons and have only one flat direction, which is purely dilatonic. This should be compared to the no-scale Minkowski vacua obtained from compactifications of type IIB string theory à la GKP. While also free from tachyons, in terms of real scalars, they have two flat directions, corresponding to the dilatonic as well as the axionic direction. While the dilatonic direction can be made massive using higher-derivative corrections, the axionic direction is protected by virtue of the axion shift symmetry and requires non-perturbative effects to be acquire a potential. This crucial feature of our novel Minkowski vacua allows us to by pass all non-perturbative corrections, and uplift to a metastable dS minimum using only higher-derivative corrections.

4.4.2 Adding higher-derivative corrections

Next, we add higher-derivative corrections to the Minkowski vacuum obtained above. The form of such corrections is known for the G_2 holonomy manifold $\mathbb{T}^7/\mathbb{Z}_2^3$, but has not been explicitly computed in the presence of twisting. However, in the case of type IIA compactified on a CY_3, Graña, Louis, Theis, and Waldram (2015) showed that turning on SU(3) torsion does not affect the form of either α' or g_s corrections. We argued in paper I, that the situation is analogous between a twisted $\mathbb{T}^7/\mathbb{Z}_2^3$ and an untwisted $\mathbb{T}^7/\mathbb{Z}_2^3$, whose singularities have been smoothed out. From this, the form of corrections to the Kähler potential were found to be $\Delta K \sim -1/\rho^3$, for the volume modulus ρ. This contributes to the scalar potential as $\Delta V \sim -\rho^{-15/2}$ in the large volume limit where $\rho \gg 1$. To generate a dS minimum, one can turn on the seven-form field strength G_7 in M-theory, which contributes $\Delta V \sim \rho^{-21/2}$ to the scalar potential. We showed how this is achieved, both perturbatively in the large volume limit, as well as with an explicit numerical example. For the volume scaling as $\rho \sim N^{2/3}$, the potential at the minimum V, the first slow-roll parameter ϵ_V, and the second slow-roll parameter η_V were computed to be

$$
\begin{aligned}
V &\sim \frac{1}{N^3}, \\
\epsilon_V &:= \frac{K^{IJ}\partial_I V \partial_J V}{2V^2} \sim \mathcal{O}\left(N^{-4}\right), \\
\eta_V &:= \min\left(\frac{K^{IJ}\partial_I V \partial_J V}{V}\right) \sim 0.01 + \mathcal{O}\left(N^{-4}\right).
\end{aligned} \tag{4.7}
$$

Quantum corrections

A curious feature of this model is that the supersymmetry breaking scale $m_{3/2}$ is parametrically higher than the compactification scale m_{KK}. This implies that an experiment at a high enough energy would find extra dimensions before it finds supersymmetric particles. The absence of supersymmetry in the low energy theory also implies that the quantum corrections are non-zero, since boson loops are no longer cancelled against fermion loops. However, it was shown that the leading vacuum correction goes as $1/N^6$, which is parametrically much smaller than the potential at the classical dS minimum obtained above, and does not destabilize it.

4.5 Summary

To summarize, in paper I, we constructed a metastable dS vacuum in M-theory compactified on a G_2 structure manifold, namely twisted $\mathbb{T}^7/\mathbb{Z}_2^3$. To do this, we made use of metric fluxes sourced by KK monopoles, the seven-form field strength G_7, and perturbative higher-derivative corrections. The advantage of this construction over other dS constructions is that it does not use any exotic ingredients that are often not well understood and are argued to give rise to instabilities. As for the fine tuning problem, in order to get a realistic cosmological constant in this model, the coefficient of the four dimensional potential at the dS minimum goes as $\sim 1/N^3$. This has to be fine tuned, at a fixed N, so that it cancels against the vacuum energy $\sim 1/N^6$ to a high degree, to give the observed cosmological constant of our universe.

This construction, although appealing at many levels, is not yet the final answer to the dS problem. One reason for this is that the fluxes used in this construction are not quantized. Although non-quantized fluxes are perfectly valid in supergravity, all fluxes should be quantized in M-theory. We will discuss this and some other aspects in section 5.5 of the next chapter. In this context, the present construction should be viewed as leading the way for a new way to construct dS vacua using well understood ingredients. To obtain a model with quantized fluxes in this framework would require more effort, and we hope to return to this in a future work.

5. The landscape vs the swampland

We have seen in the previous chapters that constructing a vacuum with positive energy is very difficult in string theory/M-theory. We also discussed that there is a very large landscape of lower dimensional effective field theories derivable from string theory. Despite its size, Vafa (2005) argued that not all apparently consistent effective field theories can be obtained from string theory. The space of such theories forms an even larger *swampland* that surrounds the landscape of truly consistent effective field theories. In this language, the difficulty of constructing a four dimensional dS vacuum can be rephrased as asking – whether a four dimensional theory containing a dS vacuum lies in the landscape or in the swampland of string theory. In order to be certain that a positive cosmological constant can be constructed in string theory, one needs to explicitly find such a construction. However, given the practically infinite number of possibilities to exclude before finding the right solution, a parallel and arguably more productive approach could be to find features of lower dimensional theories that ensure that they can be completed in string theory. Once an elaborate set of such criteria is constructed, any effective field theory can be tested against it to determine whether or not it is possible to construct it from string theory. Apart from saving us the pain of trying to construct an incompatible theory and failing, this approach could also point out the guiding principles that prevent their string theory completion, providing a suitable way forward. This is the essence of the *swampland program* initiated by Vafa (2005). A summary and status report of the program can be found in a recent review by Palti (2019). In this chapter, we will discuss a few swampland criteria that are relevant for the problem of finding de Sitter vacua in string theory.

5.1 Weak gravity conjecture

One of the most useful swampland conjectures is that, in a theory of quantum gravity with gauge forces, gravity is the weakest force of all. In other words, there exist particles whose gauge repulsion is stronger than their gravitational attraction. This was proposed by Arkani-Hamed, Motl, Nicolis, and Vafa (2007) and is called the *weak gravity conjecture* (WGC). This can be stated more concretely for a four dimensional Einstein-Maxwell theory with a $U(1)$ coupling g,

$$S = \int \mathrm{d}x^4 \sqrt{-g} \left[\frac{M_{\mathrm{Pl}}^2}{2} R - \frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} + \cdots \right]. \tag{5.1}$$

The electric version of the conjecture states that there exists a particle with mass m and electric charge $q_{\mathrm{el}} = q$ in the theory such that

$$\frac{m}{|q|} \leq \sqrt{2} g M_{\mathrm{Pl}}. \tag{5.2}$$

This statement should also hold for magnetic monopoles in the theory, $m_{\mathrm{mag}} \lesssim \mathrm{M_{Pl}}/g$. Arkani-Hamed, Motl, Nicolis, and Vafa (2007) argued that the mass of a magnetic monopole is UV divergent and goes as $m_{\mathrm{mag}} \sim \Lambda/g^2$, for a cutoff scale Λ. Together with the bound on the mass of the monopole, this gives

$$\Lambda \lesssim g M_{\mathrm{Pl}}, \tag{5.3}$$

which is called the magnetic WGC. One way to justify the electric version of the conjecture is to argue that all black holes must be able to discharge themselves. This is because, if a charged black hole does not discharge itself completely by the time it Hawking evaporates down to the Planck scale (beyond which, our semi-classical approximation breaks down), it will end up in what is called a *charged remnant*. Such remnants would then be Planck mass objects, but have different amounts of charge depending on the charge of the black hole they were formed from. This huge degeneracy of remnants implies that they lead to pathologies like violation of entropy bounds, or they could run in loops and contribute to scattering processes. See the article by Susskind (1995) for more discussion on the trouble with remnants.

The following argument borrowed from the review by Palti (2019) illustrates the connection of black hole decay with the WGC beautifully. Imagine that a black hole of mass M and charge Q decays into a set of particles with mass m_i and charge q_i. From energy conservation, $M \geq \sum_i m_i$, and charge conservation, $Q = \sum_i q_i$, we can write down

$$\frac{M}{Q} \geq \frac{1}{Q} \sum_i \left(\frac{m_i}{q_i} \right) q_i \geq \frac{1}{Q} \left(\frac{m}{q} \right) \Big|_{\min} \sum_i q_i \geq \left(\frac{m}{q} \right) \Big|_{\min}. \tag{5.4}$$

Therefore, for a black hole to decay, there has to be a particle with a charge-to-mass ratio greater than that of the black hole. This bound is the strongest for an extremal RN black hole, which has $M = Q\sqrt{2} g M_{\mathrm{Pl}}$, giving the electric WGC of equation (5.2).

The magnetic version of the weak gravity conjecture follows from applying the electric version to magnetic monopoles, as mentioned above. Another way to motivate the magnetic WGC is by taking the limit of vanishing gauge coupling. The extremality bound, $M \geq Q\sqrt{2}gM_{\text{Pl}}$, requires that a black hole with a given mass should have charge less than

$$Q \leq \frac{M}{\sqrt{2}gM_{\text{Pl}}}. \tag{5.5}$$

In the limit of vanishing coupling $g \to 0$, the black hole can have any integral charge between 0 and $1/g \to \infty$, giving an entropy $S \sim -\log g$. However, the Bekenstein-Hawking entropy of the black hole is proportional to the area of its horizon, which is proportional to gQ. Therefore, taking a small enough coupling constant $g \to 0$, while keeping gQ fixed, makes the entropy much larger than the the Bekenstein-Hawking bound. This inconsistency arises because, in the process of taking $g \to 0$, the gauge symmetry has become a global symmetry. The above contradiction is at the core of the well known principle that there should be *no global symmetries* in a theory of quantum gravity. Taking the limit $g \to 0$ should therefore not be allowed, and equation (5.3) gives a cutoff scale for the effective theory for a given value of the gauge coupling strength. See the article by Banks, Johnson, and Shomer (2006) for more discussion along these lines.

Depending on the state of the theory that satisfies the WGC, there are two main versions of the conjecture – the *strong* version, where the lightest particle in the theory is the WGC particle, and the *mild* version, where the WGC particle is the one with the largest charge-to-mass ratio. The strong WGC when applied to axions, has implications for cosmology. Axions are usually the lightest state in a theory and the strong WGC forces their decay constant to be sub Planckian, which has implications for cosmology in general and inflationary models in particular.

5.2 Generalized WGC

Arkani-Hamed, Motl, Nicolis, and Vafa (2007) argued that the WGC should hold not only for point particle charged under a U(1) gauge field, but also for higher dimensional objects charged under general p-form gauge fields. This was made precise by Heidenreich, Reece, and Rudelius (2016) for a $(p-1)$ brane with tension T and charge q under a p-form gauge field in d dimensions, to give a generalized WGC that allows not just black holes, but also black branes to decay

$$\frac{p\,(d-p-2)}{d-2}T^2 \leq q^2 g^2 M_d^{d-2}, \tag{5.6}$$

where M_d is the Planck mass in d dimensions. This is valid for $1 \leq p \leq d-3$.

5.3 Non-supersymmetric AdS conjecture

The generalized WGC has interesting consequences when applied to an older observation by Maldacena, Michelson, and Strominger (1999). They showed that in a $(p+2)$ dimensional AdS space containing a space-time filling $(p+2)$-form flux, there are p-branes charged with respect to the flux, which can nucleate and expand out towards the boundary of the AdS space. This leaves behind an AdS space with one less unit of flux. Nucleation of multiple such branes can discharge multiple units of the flux, splitting the AdS space into two different AdS spaces separated by the stack of nucleated branes. This process is called *AdS fragmentation*. The viability of such a process depends on the ratio of the p-brane's charge-to-tension. This is because, the tension wants the brane to contract, while the charge wants it to expand. Fragmentation of AdS space therefore, requires a brane with more charge than tension. Ooguri and Spodyneiko (2017) showed that such a brane always exists in the absence of supersymmetry, and equation (5.2) becomes a strict inequality. This implies that all non-supersymmetric AdS spaces supported by flux are unstable, and will decay via brane nucleation. It is interesting to note here that the generalized WGC of equation (5.6) does not apply when $p = d-1$ as we have above. Applying it to $p = d-1$ anyway, gives a result that is trivially true. However, we showed in paper III, that energy conservation suggests an absolute value on the left hand side of the generalized WGC and this extends the result to $p = d-1$ as well.

Based on this observation, it was conjectured (Ooguri and Vafa, 2017; Freivogel and Kleban, 2016) that this is a more general result and all non-supersymmetric AdS spaces (even the ones not supported by fluxes) are unstable. This conjecture has been explored in string theory (Danielsson and Dibitetto, 2017; Danielsson, Dibitetto, and Vargas, 2017). Other aspects of this conjecture have also been explored (Banks, 2016; Aalsma and van der Schaar, 2018; Antonelli, Basile, and Bombini, 2019, among others). Recently, García Etxebarria, Montero, Sousa, and Valenzuela (2020) suggested that non-supersymmetric compactifications of string theory/M-theory could also have a bubble of nothing instability,[1] causing them to

[1]A bubble of nothing is a non-perturbative instability of a spacetime, where a compact direction shrinks to zero size leaving behind *nothing*. Such an instanton is spherically symmetric (hence, a bubble) and accelerates outward, approaching the speed of light. The existence of such instabilities was discovered by Witten (1982).

decay. This further supports the conjecture that non-supersymmetric AdS vacua should be unstable.

5.4 de Sitter conjectures

We have seen that constructing de Sitter vacua in string theory is challenging. Danielsson, Haque, Koerber, Shiu et al. (2011) summarized such constructions in type IIA string theory and performed a large scan, finding potentially stable dS vacua, and highlighting some of the difficulties involved. Similar challenges exist for constructions in type IIB, heterotic string theory, as well as M-theory as summarized in the previous chapters. One of the main difficulties is that all moduli need to be stabilized, and there are not too many well understood mechanisms for stabilizing moduli in string theory.

One could either think of these as being technical difficulties, which can be circumvented in the future given enough effort, or one could start to wonder if there are more fundamental obstructions to the construction of de Sitter vacua in string theory. The second line of thought has led to recent conjectures that string theory may not allow de Sitter vacua in the first place.

Stated in the context of the swampland program, all de Sitter vacua are in the swampland and do not belong to the landscape of string theory. More specifically, the de Sitter conjecture was proposed by Obied, Ooguri, Spodyneiko, and Vafa (2018), and later refined by two groups (Garg and Krishnan, 2019; Ooguri, Palti, Shiu, and Vafa, 2019). The refined version of the conjecture states that the scalar potential of any theory coupled to gravity must satisfy either of the following two bounds on derivatives with respect to scalar fields in the theory,

$$|\nabla V| \geq \frac{c}{M_{\mathrm{Pl}}} V \qquad \text{or} \qquad \min\left(G^{ij}\nabla_i\nabla_j V\right) \leq -\frac{c'}{M_{\mathrm{Pl}}^2} V. \tag{5.7}$$

c, c' are positive constants of order one in Planck units, $|\nabla V|$ is the norm of the vector of derivatives with respect to all scalar fields in the theory, and $\min\left(G^{ij}\nabla_i\nabla_j V\right)$ is the minimum eigenvalue of the mass matrix. The first condition provides a lower bound on the slope of the potential and therefore postulates that the potential, if positive, has to be steep enough not to allow any extrema. Obied, Ooguri, Spodyneiko, and Vafa (2018) motivated this criteria by considering an extension of the Maldacena-Núñez no-go theorem in eleven dimensional supergravity to show that

$$\frac{|\Delta V|}{V} \geq \frac{6}{\sqrt{(d-2)(11-d)}}. \tag{5.8}$$

However, as mentioned in the previous chapter, the no-go theorem is invalidated in the presence of O-planes, which are extensively used in type IIA de Sitter constructions. In that case however, Hertzberg, Kachru, Taylor, and Tegmark (2007) showed that

$$\frac{|\Delta V|}{V} \geq \sqrt{\frac{54}{13}}. \tag{5.9}$$

This motivates the first criteria. However, it also excludes a positive maximum for any scalar potential. This is in contradiction with the Higgs potential, which has a positive value at the "center of the hat",

$$\frac{|\nabla V|}{V} \sim \frac{10^{-55}}{\mathrm{M_{Pl}}}, \qquad \frac{\min\left(\nabla_i \nabla_j V\right)}{V} \sim -\frac{10^{35}}{\mathrm{M_{Pl}}^2}. \tag{5.10}$$

The second condition (*i.e.*, the refinement) ensures that de Sitter maxima are allowed. Combined with the first criteria, this prevent de Sitter minima from existing. Similar refinements have been proposed before (Andriot, 2018; Garg, Krishnan, and Zaid Zaz, 2019).

To further understand the dS conjecture and the non-supersymmetric AdS conjectures in string theory, a systematic study of flux compactifications at asymptotic regions of field space was done by Scalisi and Valenzuela (2019). Connections between the dS conjecture, WGC for higher dimensional objects, and the swampland distance conjecture have been explored by Lanza, Marchesano, Martucci, and Valenzuela (2020).[2]

5.5 Possible counter examples

Several counterexamples to the original dS conjecture exist in the literature, but evading the refined conjecture is much more difficult. The following two constructions, if completely under control, could however turn out to be counterexamples to the refined dS conjecture.

M-theory

The dS vacuum constructed in paper I, which we summarized in the previous chapter, clearly violates the refined dS conjecture. More specifically, equation (4.7) does not satisfy equation (5.7), making it a *potential* counter example of the above conjecture. The reason for the use of the word *potential* is that, as we discussed in the previous chapter, this model has $m_{3/2} \gg$

[2]The distance conjecture states that, in a theory coupled to gravity, as one travels infinitely far out in field space, an infinite tower of states becomes light, invalidating the effective field theory.

$m_{\rm KK}$, making it susceptible to uncontrolled quantum corrections. We argued that these quantum corrections are parametrically much smaller than the value of the potential at the dS minimum, making our potential safe against quantum corrections. However, this could be considered one potentially weak aspect of the model. Another issue is the presence of yet unknown higher-derivative corrections. We have used the state of the art higher-derivative corrections to stabilize the volume modulus to obtain a metastable dS vacuum. We have also argued how other terms at that order cannot contribute to the potential to destroy the dS minimum. However, since the explicit form of all higher-derivative corrections is not known, this leaves room for debate. Another issue is that the fluxes used in this construction are not quantized. While the solution obtained is still a perfectly valid supergravity solution, embedding it in M-theory requires the fluxes to be quantized. Although it may be possible to find a similar model with quantized fluxes, this has not yet been done.

IIA string theory
Shortly after the refined dS conjecture was proposed, Blåbäck, Danielsson, and Dibitetto (2018) constructed a possible counterexample in massive type IIA string theory compactified on a twisted $\mathbb{T}^6/\mathbb{Z}_2^2$ with orientifold planes. This makes use of only geometric fluxes, giving a tachyon free dS minimum with one flat direction that is axionic. This axionic direction can be stabilized using the non-geometric T^2 and ST fluxes. The reason for the use of the word *potential* here is again because of the unquantized fluxes used in the construction. Moreover, if non-geometric fluxes are not employed and the flat direction is kept unfixed, then perturbative quantum corrections or higher-derivative corrections could lift it. This could lead to a quintessence like runaway behavior à la Dine and Seiberg (1985), rather than a metastable dS minimum.

A comment
Known and failed examples of dS vacua are usually constructed in the regime of weak string coupling, g_s, so as to have full perturbative control. However, it is possible for dS vacua to exist away from weak coupling. Gonzalo, Ibáñez, and Uranga (2019) developed a way to examine such regimes using string dualities, but did not find a counter-example. Similar results were also found in heterotic string theory (Parameswaran, Ramos-Sanchez, and Zavala, 2011; Olguin-Trejo, Parameswaran, Tasinato, and Zavala, 2019). However, such techniques are ineffective for potentials that are invariant under dualities, for example the ones constructed by Blåbäck, Roest, and Zavala (2014). Going away from weak coupling while retaining control, is therefore still a possible way to look for dS vacua.

5.6 What does this mean for string theory?

The dS conjectures, if true, could imply that there is no compactification in string theory that gives rise to a static dS vacuum. All of the constructions discussed so far are debatable in one way or another, and there is no dS construction yet, which can claim to be completely under control. In the future, one of these constructions could turn out to be completely under control, thus invalidating the dS conjectures and providing a concrete realization of dS in string theory. However, if the swampland conjectures win, there could be a sharp tension between theory and observations that would cry out for other explanations. On the other hand, one could take a more positive attitude and try to find a clever way to evade the swampland conjecture. The review by Danielsson and Van Riet (2018) discusses these issues in detail. The article by Danielsson (2019) deals with another less discussed aspect of dS in string theory.

In the spirit of not giving up on string theory, but rather, trying to evade the dS swampland conjectures, we will review an older idea called *braneworlds* in the next chapter, and examine it from the point of view of a viable dS construction.

6. Revisiting an old idea: Braneworlds

At the turn of the millennium, Randall and Sundrum (1999a,b) made a seminal proposal that changed the way we view extra dimensions. They proposed that our four dimensional universe could be living on a three-brane (with a four dimensional world volume), embedded in a five dimensional bulk spacetime, where the fifth dimension is large (albeit warped, to have a finite volume). Although the original construction was for a flat four dimensional Minkowski universe, it was later extended to include a cosmological constant. However, given the lack of a detailed construction of such *braneworlds* in string theory, and the timely proposal by KKLT, most of the attention to constructing dS vacua in string theory shifted to the KKLT idea. Given the present debate surrounding the possibility of constructing dS vacua in string theory, it makes sense to re-examine braneworlds for lessons that it might have to offer for constructing a dS universe. We will briefly review the Randall-Sundrum construction in this chapter and discuss some aspects of its extension to a dS universe.

6.1 Braneworlds

The possibility that our universe is a four dimensional surface embedded in higher dimensional space, and its implications for particle physics and gravity have long been discussed in the literature (Akama, 1982; Rubakov and Shaposhnikov, 1983; Pavsic, 1985, 1986, 1996, 1997; Kakushadze and Tye, 1999; Gogberashvili, 2000). The novelty of the proposal by Randall and Sundrum (1999a) was that it offered a string theory inspired braneworld solution to the gauge hierarchy problem in particle physics (*i.e.*, why is there a huge gap between the electroweak scale of standard model particles and the Planck scale of gravity). This model is usually referred to as RS-I. It involves two three-branes, one with a positive tension, and the other with a negative tension, embedded in a five dimensional AdS space. Owing to the warped five dimensional AdS spacetime, particle masses on the negative tension brane are naturally much smaller than the Planck scale, providing a solution to the gauge hierarchy problem. The size of the extra dimension in this scenario is the vacuum expectation value of a modulus field that is stabilized by the Goldberger and Wise (1999) mechanism.[1]

In a subsequent article, Randall and Sundrum (1999b) introduced a second model, which is usually referred to as RS-II. This gets rid of the negative tension brane and shows how four dimensional gravity emerges on the positive tension brane, without the need for a small extra dimension. Before RS-II, it was believed that extra dimensions need to be small. This is because, in the presence of large extra dimensions, the force of gravity could permeate into the extra dimension, modifying Newton's law and leading to a contradiction with experiments. However, RS-II showed that a warped extra dimension can localize gravity, effectively confining it to four dimensions (with deviations from Newton's law being highly suppressed, and detectable only in extremely high energy experiments).

We will not talk about RS-I here, but rather focus on the RS-II model, which we will discuss in some detail in the following sections. Since this is the only RS model that we will talk about, we will use the expressions RS and RS-II interchangeably.

6.2 The model

The setup consists of a single positive tension brane and the following five dimensional action

$$S = \frac{1}{2\kappa_5^2} \int \mathrm{d}^4x \int \mathrm{d}z \sqrt{-g_5}\,(R_5 - 2\Lambda_5) + \sigma_{\mathrm{br}} \int \mathrm{d}^4x \sqrt{-g_{\mathrm{br}}}, \tag{6.1}$$

where the subscript "5" indicates quantities in the five dimensional bulk and the subscript "br" indicates quantities defined on the brane. The five dimensional cosmological constant can be written in terms of the curvature scale of AdS_5 as $\Lambda_5 = -6k^2$. Taking an ansatz for the bulk metric of the form

$$\mathrm{d}s^2 = e^{2A(z)} g_{\mu\nu} \mathrm{d}x^\mu \mathrm{d}x^\nu + \mathrm{d}z^2, \tag{6.2}$$

Einstein's equations off the brane can be solved by choosing $A(z) = \pm kz$ and $g_{\mu\nu} = \eta_{\mu\nu}$. Tension of the brane σ_{br}, is given by Einstein's equations on the brane, which can be rewritten as the *thin-shell junction conditions* as summarized below.

[1]See the article by Arkani-Hamed, Dimopoulos, and Dvali (1998) for an earlier solution to the gauge hierarchy problem with standard model particles on a three-brane and a large extra dimension, but a different bulk spacetime.

Thin-shell junction conditions

Given two spacetimes $\mathscr{M}^{\pm}$ with metrics $g^{\pm}_{\alpha\beta}$, which are both solutions to Einstein's equations, imagine constructing a composite spacetime that has $\mathscr{M}^{+}$ in one part of space and $\mathscr{M}^{-}$ in another. For the whole spacetime to be a solution to Einstein's equations, there has to be some stress energy tensor at the boundary Σ of the two spacetimes that sources the jump in the metric. The value of this stress tensor is obtained by writing down an action similar to equation (6.1), and varying it with respect to the metric to give the following thin-shell junction conditions:

$$\text{First junction condition:} \quad [h_{ab}] = 0, \tag{6.3}$$

$$\text{Second junction condition:} \quad S_{ab} = -\frac{\varepsilon}{\kappa^2}\left([K_{ab}] - K h_{ab}\right). \tag{6.4}$$

Quantities in square brackets indicate their difference across Σ, and $\varepsilon = \pm 1$ for a timelike or a spacelike hypersurface respectively. h_{ab} is the metric induced on the surface Σ and is defined as $h_{ab} := g_{\alpha\beta}\, e^{\alpha}_{a} e^{\beta}_{b}$, where the tangent vectors are given by

$$e^{\alpha}_{a} := \frac{\partial x^{\alpha}}{\partial y^{a}}. \tag{6.5}$$

K_{ab} is the extrinsic curvature of the embedded surface and is given in terms of the normal and tangent vectors as

$$K_{ab} := n_{\alpha;\beta}\, e^{\alpha}_{a} e^{\beta}_{b}, \qquad \text{where} \quad n_{\alpha} := \frac{\varepsilon\, \Phi_{,\alpha}}{\sqrt{\left|g^{\mu\nu}\Phi_{\mu}\Phi_{\nu}\right|}}, \tag{6.6}$$

and $K := K_{ab}h^{ab}$ is the trace of the extrinsic curvature K_{ab}. Derivation of the thin-shell junction conditions and more details can be found in a textbook on general relativity, like the one by Poisson (2009).

In the RS scenario, the z coordinate it taken to have a $\mathbb{Z}_2$ symmetry across $z = 0$, where the brane is placed. This amounts to choosing $A(z) = -k\,|z|$ in equation (6.2), *i.e.*, the warp factor $\exp(-2k\,|z|)$ goes to zero far away from the brane on both sides

$$\mathrm{d}s^2 = g_{ab}\mathrm{d}x^a\mathrm{d}x^b = e^{-2k|z|}\eta_{\mu\nu}\mathrm{d}x^{\mu}\mathrm{d}x^{\nu} + \mathrm{d}z^2. \tag{6.7}$$

The thin-shell junction conditions give the tension of the brane, which, for a stationary brane at $z = z_0 = 0$, is $\sigma_{\mathrm{br}} = 6k/\kappa_5^2$. This precise fine tuning between the tension of the brane and the cosmological constant in the bulk imposes four dimensional Poincaré invariance, giving a flat (Minkowski) brane.

6.3 Four dimensional Planck scale

The five dimensional action in equation (6.1) can be dimensionally reduced along the z direction for the RS metric in equation (6.7) to obtain

$$S_{\text{eff}} \subset \frac{1}{2\kappa_5^2} \int_{-\infty}^{\infty} \mathrm{d}z\, e^{-2k|z|} \int \mathrm{d}^4 x \sqrt{-\tilde{g}} \widetilde{R}, \tag{6.8}$$

where $\widetilde{R}$ is the four dimensional Ricci scalar corresponding to the four dimensional part of the metric (which we call $\tilde{g}$). From this, the four dimensional Planck scale can be read off to be

$$\frac{1}{2\kappa_4^2} = \frac{1}{2\kappa_5^2} \int_{-\infty}^{\infty} \mathrm{d}z\, e^{-2k|z|} = \frac{1}{k} \frac{1}{2\kappa_5^2} \quad \Rightarrow \quad M_4^2 = \frac{M_5^3}{k}. \tag{6.9}$$

This should be compared to a factorizable spacetime of the form $\mathrm{d}s^2 = g_{ab}(x)\,\mathrm{d}x^a \mathrm{d}x^b + \bar{g}_{\mu\nu}(y)\,\mathrm{d}y^\mu \mathrm{d}y^\nu$, where the effective Planck scale, obtained from a similar dimensional reduction, would be proportional to the volume of the extra dimensions, $M_4^2 = M_{10}^6 \mathcal{V}_6$. The internal dimensions need to be compact for the volume of the internal space $\mathcal{V}_6$ to be finite, which is then proportional to the size of these extra dimensions. In contrast, for the RS construction above, the warping keeps the internal volume finite. This gives a four dimensional Newton's constant proportional to the curvature of the extra dimension, instead of its size.

6.4 Gravity on the braneworld

Generically, a perturbation of the five dimensional metric can propagate along all five dimensions. This implies that the graviton, which is the transverse and traceless mode of the perturbation, will give rise to the five dimensional Newton's potential. In the RS construction however, the warping along the extra dimension ensures that the lowest energy state of the graviton on the brane (called the *zero mode*) actually gives rise to the four dimensional Newton's potential on the brane (plus high energy corrections). In this sense, gravity is *localized* on the braneworld. The wavefunction of this zero mode, which can be identified as the four dimensional graviton, is peaked around the brane and goes to zero away from it. In this section, we will briefly outline how this arises, using the formalism of *brane bending* developed simultaneously by two groups Garriga and Tanaka (2000), and Giddings, Katz, and Randall (2000).

To compute the wavefunction of the graviton, we start by considering perturbations of the five dimensional RS metric in equation (6.7), in response to a point mass placed on the brane. The perturbed metric is of the form $\tilde{g}_{ab} = g_{ab} + h_{ab}$, for $a, b \in \{0, \dots, 4\}$. The perturbation h_{ab}, being symmetric, has 15 degrees of freedom. Generalized coordinate transformations, corresponding to $x^\mu \mapsto \tilde{x}^\mu, z \mapsto \tilde{z}$, take away five degrees of freedom, while the differential Bianchi identities leading to covariant conservation of the Einstein tensor $\nabla^\mu G_{\mu\nu} = 0$, take away five more. This leaves five physical degrees of freedom for the five dimensional graviton, and the ten redundant degrees of freedom can be fixed by choosing a gauge. To determine what is a good gauge choice, one can write down linearized Einstein's equations for the metric perturbation h_{ab} (which is assumed to be small $h_{ab} \ll 1$), and determine which gauge choices are consistent for the given stress tensor. See the article by Giddings, Katz, and Randall (2000) for a detailed discussion. In the absence of matter away from the brane, a gauge that can be chosen everywhere in the bulk, is called the Randall-Sundrum (RS) gauge, defined by

$$h^\mu_{\nu,\mu} = 0, \quad h^\mu_\mu = 0, \quad h_{\mu z} = h_{zz} = 0. \tag{6.10}$$

For an empty bulk, the Einstein's equations are equivalent to the variation of the Ricci tensor being zero, which can be written as

$$\delta R_{ab} = -\frac{1}{2}\nabla^2 h_{ab} - R_a{}^b{}_c{}^d h_{cd} + R_{(a}{}^c h_{b)c} + \nabla_c \nabla_{(a} h^c_{b)} - \frac{1}{2}\nabla_{(a}\partial_{b)} h \overset{!}{=} 0, \tag{6.11}$$

where the curvature tensors and covariant derivatives are computed with respect to the unperturbed metric g_{ab}. In the RS gauge, this reduces to

$$\left[e^{2k|z|}\Box_{(4)} + \partial_z^2 - 4k^2\right] h_{\mu\nu} = 0, \tag{6.12}$$

where $\Box_{(4)}$ is the four dimensional scalar Laplacian operator and the other components $h_{\mu z}$ and h_{zz} vanish due to the gauge choice. However, it was shown by Giddings, Katz, and Randall (2000) that in the presence of matter on the brane (beyond its fine tuned tension), the RS gauge is inconsistent and the brane cannot be at $z = 0$; it has to be at $z = -f(x^\mu)$ as we will see in equation (6.18). Let us understand what this means. The RS brane has a fine tuned tension, is static in the RS background, and solves both bulk Einstein's equations as well as the junction conditions. If we now add an extra stress tensor on the brane (for example, a point mass), it will contribute an extra extrinsic curvature on the brane and the brane will of course not be at $z = 0$ anymore. This is called *brane bending*.

One can, however, change coordinates to $\bar{z}$, in which the brane is located at $\bar{z} = 0$. This involves relaxing the RS gauge and requiring a milder gauge choice instead, called the Gauss normal (GN) gauge

$$\bar{h}_{\mu z} = \bar{h}_{zz} = 0, \tag{6.13}$$

which fixes five degrees of freedom, leaving behind five more that can be fixed later. The stress tensor on the brane is given by the thin-shell junction conditions,

$$S_{\mu\nu} = -\frac{1}{\kappa_5^2}\left(\Delta K_{\mu\nu} - \Delta K \gamma_{\mu\nu}\right), \tag{6.14}$$

or equivalently,

$$\Delta K_{\mu\nu} = -\kappa_5^2 \left(S_{\mu\nu} - \frac{1}{3} S \gamma_{\mu\nu}\right), \tag{6.15}$$

where $\gamma_{\mu\nu}$ is the metric induced on the brane. Imposing $\mathbb{Z}_2$ symmetry across the brane and realizing that $K_{\mu\nu} = \partial_z \tilde{g}_{\mu\nu}$ gives

$$\left(\partial_z + 2k\right) \bar{h}_{\mu\nu}\Big|_{z=0^+} = -\kappa_5^2 \left(T_{\mu\nu} - \frac{1}{3} T \eta_{\mu\nu}\right), \tag{6.16}$$

where we have used that

$$S_{\mu\nu} = \sigma\eta_{\mu\nu} + T_{\mu\nu} = \frac{6k}{\kappa_5^2}\eta_{\mu\nu} + T_{\mu\nu}. \tag{6.17}$$

The junction conditions are easier to compute for a brane placed at a fixed z, which is why we have evaluated them in the GN gauge. They can however be translated back to the RS gauge via a gauge transformation that makes use of the remaining five degrees to freedom. The most general such transformation is given by

$$\hat{z} = f\left(x^\alpha\right), \tag{6.18}$$

$$\hat{x}^\mu = -\frac{1}{2k} e^{2k|z|} \eta^{\mu\nu} \partial_\nu f\left(x^\alpha\right) + Q^\mu\left(x^\alpha\right). \tag{6.19}$$

The metric perturbation in RS gauge becomes

$$h_{\mu\nu} = \bar{h}_{\mu\nu} - \frac{1}{k} f_{,\mu\nu} - 2k e^{-2k|z|} f + e^{-2k|z|} \eta_{\alpha(\mu} Q^\alpha_{\ ,\nu)}. \tag{6.20}$$

When inserted into the junction condition in the GN gauge, this gives the corresponding junction condition in the RS gauge

$$\left(\partial_z + 2k\right) h_{\mu\nu}\Big|_{z=0^+} = -\Sigma_{\mu\nu}, \tag{6.21}$$

where the stress tensor on the brane has an extra contribution proportional to the *bending* f,

$$\Sigma_{\mu\nu} = \kappa_5^2 \left(T_{\mu\nu} - \frac{1}{3} T \eta_{\mu\nu}\right) + 2 f_{,\mu\nu}. \tag{6.22}$$

Let us summarize the discussion on brane bending. In five dimensions, there are 10 gauge degrees of freedom that need to be fixed, to get the five physical degrees of freedom of the massless graviton. The RS gauge is one such gauge choice in which the brane lies at $z = 0$ in the absence of additional matter on the brane, apart from its tension. However, in the presence of extra matter, there is extra stress tensor on the brane and it cannot lie flat at $z = 0$. In order to compute this stress tensor on the brane, we need to use the thin-shell junction conditions, which are easier to compute when the brane lies at $z = 0$. So we change gauge from RS to GN and compute the stress tensor $S_{\mu\nu}$ in that gauge. However, since the GN gauge still has 5 extra degrees of freedom, we can use the gauge freedom to go back to RS gauge and find out what the stress tensor looks like in that gauge. Doing this, we realize that there is an extra contribution to the stress tensor in RS gauge proportional to the bending. The equations in RS gauge (6.12) and (6.21) can be written together as

$$\left[e^{2k|z|}\Box_{(4)} + \partial_z^2 - 4k^2 + 4k\delta(z)\right] h_{\mu\nu} = -2\Sigma_{\mu\nu}\delta(z), \tag{6.23}$$

whose solution can be formally written in terms of the retarded Green's function

$$h_{\mu\nu} = -2\int \mathrm{d}^4x' G_R\left(x, x'\right)\Sigma_{\mu\nu}\left(x'\right). \tag{6.24}$$

Since this is in the RS gauge, where the perturbation is traceless $h_\mu^\mu = 0$, it implies that so is the stress tensor $\Sigma_\mu^\mu = 0$. Combined with equation (6.22), this gives the bending f,

$$\Box_{(4)} f = \frac{\kappa_5^2}{6} T. \tag{6.25}$$

This, combined with equations (6.24) and (6.22), automatically implies $h^\nu_{\mu,\nu} = 0$. The matter part of the metric fluctuation $h^{(m)}_{\mu\nu}$ is given by equations (6.24) and (6.22) as

$$h^{(m)}_{\mu\nu} = -2\kappa_5^2 \int \mathrm{d}^4x' G_R\left(x, x'\right)\left(T_{\mu\nu} - \frac{1}{3}T\eta_{\mu\nu}\right). \tag{6.26}$$

This does not have the right tensor structure for the four dimensional graviton, which should have a $1/2$ instead of $1/3$. This is taken care of by the brane bending, and can be seen when the metric perturbation is read off in GN coordinates on the brane ($z = 0$)

$$\bar{h}_{\mu\nu} = h^{(m)}_{\mu\nu} + 2kf\eta_{\mu\nu}. \tag{6.27}$$

Using the explicit expressions for the zero mode of the Green's function and the bending f from equation (6.25),

$$G_R\left(x, x'\right) = \frac{k}{\Box_{(4)}}, \quad f = \frac{\kappa_5^2}{6}\int \mathrm{d}^4x' \frac{1}{\Box_{(4)}} T, \tag{6.28}$$

the metric perturbation on the brane $\bar{h}_{\mu\nu}$ becomes

$$\bar{h}_{\mu\nu} = -2\kappa_5^2 \int \mathrm{d}^4x' \frac{k}{\Box_{(4)}} \left(T_{\mu\nu} - \frac{1}{2} T\eta_{\mu\nu} \right). \tag{6.29}$$

Newton's potential on the brane world can now be read off from this and is given by

$$V(r) = \frac{1}{2}\bar{h}_{00} = \frac{G_5 kM}{r}\left[1 + \frac{2}{3k^2r^2} + \mathcal{O}\left(r^{-3}\right) \right] \tag{6.30}$$

With the identification $G_5k = G_4$, Newton's potential on the braneworld is correctly reproduced with subleading Yukawa type corrections for distances smaller than the AdS$_5$ scale, $r < 1/k$, which does not contradict present experimental tests of general relativity.

6.5 (A)dS braneworlds

To summarize the discussion so far, in the original RS construction, a solution to the action in equation (6.1) was found with an AdS$_5$ bulk metric, written as slices of four dimensional Minkowski space. This makes a brane at a fixed z have a Minkowski world volume, with its tension given by the thin-shell junction conditions.

However, other solutions to the action exist that are dS and AdS slicings of the AdS$_5$ metric. Placing a brane at a fixed z in this background gives a dS$_4$ or AdS$_4$ brane with a different value of the tension given again by the thin-shell junction conditions. Karch and Randall (2001) examined these solutions and showed that there is a four dimensional graviton localized on the brane, even in the presence of a non-zero cosmological constant on the brane. Solutions to Einstein's equations of the form of equation (6.2), for $g_{\mu\nu}$ being dS$_4$, Minkowski$_4$, and AdS$_4$, are given by

$$\begin{aligned}
\mathrm{dS}_4 &: A(z) = \log\left(\frac{\sqrt{\Lambda_4}}{k} \sinh(kc - k|z|) \right), \quad \sigma_{\mathrm{br}} = 6k\coth(kc), \\
\mathrm{MkW}_4 &: A(z) = kc - k|z|, \quad \sigma_{\mathrm{br}} = 6k, \\
\mathrm{AdS}_4 &: A(z) = \log\left(\frac{\sqrt{-\Lambda_4}}{k} \cosh(kc - k|z|) \right), \quad \sigma_{\mathrm{br}} = 6k\tanh(kc).
\end{aligned} \tag{6.31}$$

The Minkowski metric with $c = 0$ is the familiar result from the original RS construction that we have used so far in this chapter. For a dS braneworld, the parameter c is the distance between the brane and the horizon in these coordinates while for an AdS braneworld, it is the point where the warp factor $\exp(2A)$ turns around and increases to infinity. Just like the Minkowski braneworld, the warp factor for the dS case goes to zero away for the brane. However, for the AdS case, the warp factor decreases until $|z| = c$, before increasing again to infinity for $z \to \infty$.

Attempting to construct these braneworlds in supergravity brings up an interesting conclusion. The five dimensional metric for the Minkowski braneworld, as well as for the dS braneworld, have warp factors that go to zero at the boundary and the horizon respectively. This makes the volume of the fifth dimension finite. In particular, for dS braneworld, this is now a four dimensional static dS spacetime on a positive tension brane with the extra dimension having a finite volume. This falls right into the grips of the Maldacena-Núñez no-go theorem, which prevents it from being realized in supergravity. Other no-go results for smooth RS braneworlds in supergravity were found by Behrndt and Cvetic (2000). They found that the spacetime transverse to a domain wall that interpolates between extrema of the superpotential always approaches the boundary of AdS (instead of the center of AdS, like in the Minkowski braneworld). Kallosh and Linde (2000) also arrived at a similar conclusion. This implies that although AdS_4 braneworlds might be allowed in string theory, embedding a dS_4 braneworld in the RS setup is forbidden, unless one evades these no-go theorems by going beyond smooth domain walls in two derivative supergravity.

Recently, Karch and Randall (2020) have explored a two brane model with *mismatched* cosmologies on the branes (*i.e.*, they differ in the signs of their four dimensional cosmological constants) and found that these lead to time-dependent solutions. They argued that, combined with a mechanism to stabilize the branes, this could lead to interesting cosmological implications. Possible connections between the KKLT scenario and the mismatched RS scenario with two branes have also been explored by Randall (2019).

Given the current discussion surrounding the possibility of constructing stable dS solutions in string theory, it is pertinent to also examine other possibilities. In this context, inspired by the RS scenario, we proposed a model of a dS_4 universe, which can be potentially completed in string theory. This is discussed in the next chapter.

7. Introducing a new idea: Shellworlds

We have seen in the previous chapter that dS braneworlds describe a way to obtain a four dimensional universe with a small and positive cosmological constant. However, because of various no-go theorems, it is not clear if they can be obtained from string theory without including ingredients that go beyond two-derivative supergravity.

Inspired by the discussion so far, in paper II, we proposed a new way to realize a four dimensional dS universe, which seems to be easier to obtain from string theory. Further aspects of the model were explored in papers III and IV. In this chapter, we will point out some key features of this model and how it compares with the RS model discussed in the previous chapter. Instead of repeating the details of the model, we will point out relevant sections of the papers, and hopefully make it easier to navigate through them. The setup of this model relies on false vacuum decay, which we will briefly summarize in the next section.

7.1 False vacuum decay

7.1.1 Field theory without gravity

Consider a theory with a single scalar in a potential $V(\phi)$ with two non-degenerate minima $V_\pm$, with $V_- > V_+$, separated by an energy barrier V_b, as shown in figure 7.1,

$$S = \int \mathrm{d}^d x \left[\frac{1}{2}\partial_\mu\phi\partial^\mu\phi - V(\phi)\right]. \tag{7.1}$$

Classically, a particle at rest in the local minimum at $\phi = \phi_-$ does not have enough energy to overcome the potential barrier V_b, and escape to the global minimum at $\phi = \phi_+$. However, it can quantum mechanically tunnel through the barrier and end up in the global minimum. The local minimum is thus metastable and can *decay* through such a barrier penetration process. This vacuum is called *a false vacuum*. The physics of such decays was originally studied in two papers by Coleman (1977), and Callan and Coleman (1977). They showed that such a decay occurs via the nucleation of a spherically symmetric bubble of true vacuum in the false vacuum, and is similar to a first order phase transition in statistical mechanics. In

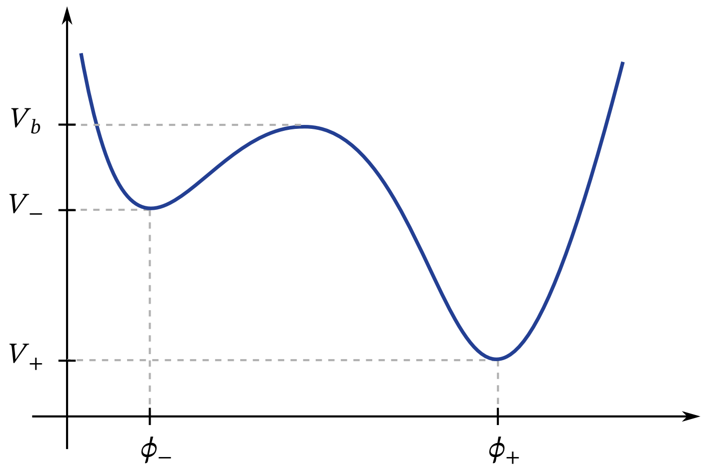

Figure 7.1. A generic potential with two non-degenerate vacua, $V_\pm$, separated by a potential barrier V_b. The false vacuum V_- can decay to the true vacuum V_+ via quantum mechanical tunneling.

the semi-classical limit ($\hbar \ll 1$), the probability of such a process (Γ) per unit volume of space (V) is given by

$$\frac{\Gamma}{V} = Ae^{-B/\hbar}\left[1 + \mathcal{O}(\hbar)\right], \quad \text{with} \quad B := S_E(\phi_I) - S_E(\phi_-), \tag{7.2}$$

where $S_E(\phi_-)$ is the Euclidean action evaluated in the false vacuum, and $S_E(\phi_I)$ is the Euclidean action for the *bounce* solution. The bounce is the instanton solution corresponding to the nucleated bubble, *i.e.*, localized solutions $\phi = \phi_I$ that extremize, and give a finite Euclidean action $S_E(\phi)$. The coefficient A depends on details of the potential $V(\phi)$. The subsequent evolution of the nucleated bubble can be studied using evolution of the particle in the potential. A particular limit in which the physics becomes simple, is when the difference in energy density between the vacua is small, *i.e.*, $V_- - V_+ \equiv \varepsilon \to 0^+$. This is called the *thin wall* limit.

7.1.2 Field theory coupled to gravity

While every false vacuum can decay in a quantum field theory, Coleman and De Luccia (1980) showed that this is not the case when the quantum field theory is coupled to gravity. Gravity can stabilize some false vacua, making them eternally stable. This can be understood intuitively by accounting for energy conservation during the nucleation event. The bubble being an instanton, is formed at a radius that minimizes the Euclidean action. In order to form, the bubble needs to spend energy (proportional to its tension times its area) to build itself. This energy comes from lowering the vacuum energy in its interior, which is proportional to the vacuum energy density times its volume. When these exactly balance each other, the bubble remains at rest after nucleation. However, if the bubble gains more energy from lowering the vacuum energy than it needs to hold itself together, then this extra energy provides the kinetic energy of expansion. Therefore, after nucleating at rest, it accelerates outward, asymptotically approaching the speed of light. This can be written symbolically as $V_{\text{before}} - V_{\text{after}} = E_{\text{wall}} + E_{\text{expansion}}$, where the last term is the kinetic energy of expansion of the bubble. The expansion energy being positive semi-definite, gives an upper bound on the tension of the wall,

$$E_{\text{wall}} \leq V_{\text{before}} - V_{\text{after}}. \tag{7.3}$$

The equality holds for a static bubble when $E_{\text{expansion}} = 0$.

As an example, let us work this out for the case of a bubble containing an AdS_d vacuum that nucleates inside a metastable AdS_d vacuum, with a different cosmological constant. Here we will generalize the four dimensional result that was worked out by Harlow (2010). Consider an AdS_d space, with cosmological constant $\Lambda = -(d-1)(d-2)/(2L^2)$, written in dS slicing

$$\mathrm{d}s^2_{\text{AdS}} = \mathrm{d}\xi^2 + f(\xi)^2 \left[-\mathrm{d}\eta^2 + \cosh^2\eta \, \mathrm{d}\Omega^2_{d-2}\right], \tag{7.4}$$

where $f(\xi) := L \sinh(\xi/L)$. Let us denote quantities outside and inside the bubble with subscripts plus and minus respectively. Further, let σ be the tension of the bubble, and let us assume that its position is given by $\xi = \text{constant} = \xi_\pm$ as seen from either side. From the thin-shell junction conditions of equations (6.3) and (6.4), we get

$$\begin{aligned} f_-(\xi_-) &= f_+(\xi_+) \equiv f, \\ f'_-(\xi_-) - f'_+(\xi_+) &= \frac{\kappa_d^2 \sigma f}{(d-2)}. \end{aligned} \tag{7.5}$$

Energy conservation implies

$$f'_-(\xi_-) - f'_+(\xi_+) = \frac{2f}{(d-1)} \frac{E(\sigma)}{\sigma}, \tag{7.6}$$

where $E(\sigma)$ is the energy of the wall. For $f_\pm(\xi_\pm) := L_\pm \sinh\left(\xi_\pm L_\pm\right)/k_\pm$ above, and defining $x_\pm := k_\pm \xi_\pm$, these become

$$\begin{aligned} L_- \sinh(x_-) &= L_+ \sinh(x_+), \\ \cosh(x_-) - \cosh(x_+) &= \frac{{}^{2}_{d}\sigma}{(d-2)} L_- \sinh(x_-), \\ \cosh(x_-) + \cosh(x_+) &= \frac{2}{(d-1)} \frac{E(\sigma)}{\sigma} L_- \sinh(x_-). \end{aligned} \tag{7.7}$$

The last two equations can be solved for functions of $x_\pm$. Defining the coefficients on the right hand side as

$$A := \frac{\kappa_d^2 \sigma}{(d-2)} L_-, \quad B := \frac{2}{(d-1)} \frac{E(\sigma)}{\sigma} L_-, \tag{7.8}$$

gives

$$\coth(x_-) = \frac{B+A}{2}, \quad \coth(x_+) = \left(\frac{L_+}{L_-}\right) \frac{B-A}{2}. \tag{7.9}$$

Multiplying the last two equations and using the first gives

$$\frac{1}{L_-^2} - \frac{1}{L_+^2} = \frac{2\kappa_d^2 E(\sigma)}{(d-1)(d-2)} \equiv AB. \tag{7.10}$$

Using this to rewrite B in terms of A, equation (7.9) has a solution only when $A \leq 1 - L_-/L_+$, which can be written as

$$\sigma \leq \frac{d-2}{\kappa_d^2} \left(\frac{1}{L_-} - \frac{1}{L_+}\right). \tag{7.11}$$

This gives an upper bound on the tension of the bubble that can nucleate.

7.1.3 Lowering the cosmological constant via bubble nucleation

Using the framework of false vacuum decays in the presence of gravity discussed above, Brown and Teitelboim (1988) demonstrated a mechanism for lowering of the cosmological constant in the presence of antisymmetric tensor fields. Consider a spacetime where the cosmological constant is given by the combination of a geometrical cosmological constant and a field strength, *e.g.*,

$$S_E \supset -\frac{1}{2\kappa_d^2} \int \mathrm{d}^d x \sqrt{-g}\,(R - 2\Lambda) - \frac{1}{d!} \int \mathrm{d}^d x \sqrt{-g}\, |F|^2 + \dots, \tag{7.12}$$

where for $F^{\mu_1 \cdots \mu_d} = \left(E/\sqrt{-g}\right) \epsilon^{\mu_1 \cdots \mu_d}$, the effective cosmological constant is a combination of Λ and E^2. They showed that similar to Schwinger pair production of charged particles in an electric field, a spherical membrane charged under the antisymmetric field (*e.g.*, a spherical Dp brane) can nucleate. By charge conservation, the space inside the bubble would contain less charge than outside, differing by precisely the amount of charge carried by the membrane. This leads to a lower value of the cosmological constant inside the bubble. In order to nucleate, such a bubble would of course have to obey energy conservation à la Coleman and De Luccia (1980), as discussed in equation (7.3).

7.2 Empty dS universe

The starting point for our model consists of an AdS_5 space, obtained from a ten or eleven dimensional compactification, whose cosmological constant is supported, at least in part, by the flux of form fields in string theory. A spherical brane instanton of co-dimension one can nucleate in this background à la Brown and Teitelboim. This brane is charged under the flux that supports the cosmological constant of the AdS_5. As a result, the space inside it is an AdS_5 with a lower cosmological constant. This is just as one would expect from the conjecture about non-supersymmetric AdS spaces discussed in section 5.3. The brane nucleates at rest and its tension dictates its evolution. For a critical value of the tension, its tendency to expand due to the difference in the cosmological constant is counteracted by its tension, and the bubble remains at rest.[1] However, for a sub-critical tension, the spherical brane nucleates at rest and expands outward towards the boundary of AdS, asymptotically approaching the speed of light. In paper II, we showed that the world volume theory on such an expanding brane with a fine tuned sub-critical tension is an empty dS universe. We call this construction the *shellworld*. This is represented in figure 7.2.

This is very different from the RS-II scenario where the brane sits at an orbifold point with $\mathbb{Z}_2$ symmetry across it. The interior of AdS is reflected along the brane and the coordinate range corresponding to $z \in (0, \infty)$ is just a copy of $z \in (-\infty, 0)$. In this sense, what should have been the *outside* of AdS from the brane to the boundary is replaced by a mirror of the *inside* from the center of the AdS to the brane. So going away from the brane on either side is like traveling to the center of AdS. We refer to this as *inside-inside*.

In contrast, there is no $\mathbb{Z}_2$ symmetry in our construction. Instead, the brane nucleates in an AdS_5 whose z coordinate goes all the way from $z \to -\infty$ at the center of the bubble to $z \to \infty$ at the boundary of AdS_5. Another way to see this is by looking at the warp factor. In RS, the warp factor is $\exp(-2k\,|z|) \equiv \exp(2kz) + \Theta(z)\exp(-2kz)$, where $\Theta(z)$ is the Heaviside theta function. This warp factor peaks at the position of the brane $z = 0$ and goes to zero away from it on both sides. In contrast, for the shellworld written in RS coordinates, the warp factor is $\exp(2k_- z) + \Theta(z)\exp(2k_+ z)$, where $k_\pm$ are the AdS_5 curvatures outside and inside the bubble respectively. This increases all the way from the center of AdS to the boundary. In this sense, the space outside the bubble is a just like true AdS before the bubble was nucleated. This is what we refer to as *inside-outside*. Let us examine the consequences that this difference has.

[1] We refer to the spherical brane including the AdS_5 in its interior as the *bubble* and the brane itself as the *shell*.

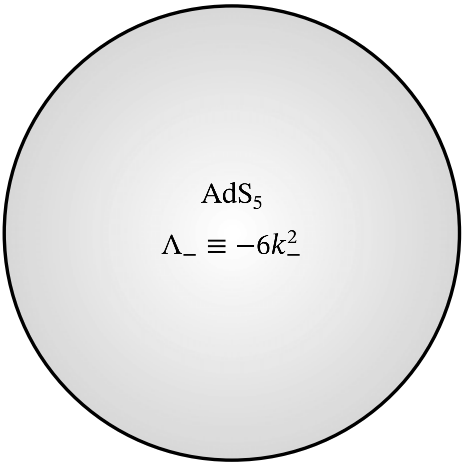

Figure 7.2. A schematic representation of a spherical bubble of true vacuum that has nucleated in a metastable AdS$_5$ with cosmological constant Λ_+. The bubble contains AdS$_5$ space in its interior with a lower cosmological constant Λ_-. The surface of the bubble is made of branes with sub-critical tension, and represents the shellworld.

1. Because the bulk metric differs between the RS and the shellworld construction, so does the critical tension needed for a stationary brane. This can be computed from the thin-shell junction conditions of equations (6.3) and (6.4). For the shellworld, this gives the following critical tension:
$$\sigma_{\text{crit}} := 3\left(k_- - k_+\right)/\kappa_5^2. \tag{7.13}$$
This should be compared to the critical tension for the RS brane, which is $\sigma_{\text{crit}} = 6k/\kappa_5^2$. Comparing the warp factors in the previous paragraph, we see that flipping the sign in front of k_+ is formally equivalent to going between the RS and the shellworld. Performing this sign flip in equation (7.13) indeed recovers the RS result. The reason that this works is because, flipping the sign between k_- and k_+ in the critical tension, corresponds to flipping the sign of the normal in the junction condition, which corresponds to changing the outside to an inside. The critical tensions for the RS braneworld and the shellworld can be seen in figure 7.3, which is taken from paper II.
2. The next point of difference between the two scenarios is in the type of brane required to get a four dimensional dS universe. Karch and Randall (2001) showed that in a dS braneworld, the cosmological constant is given by
$$\Lambda_{\text{dS}}^{\text{RS}} = k^2\left[\coth^2(kc) - 1\right]. \tag{7.14}$$

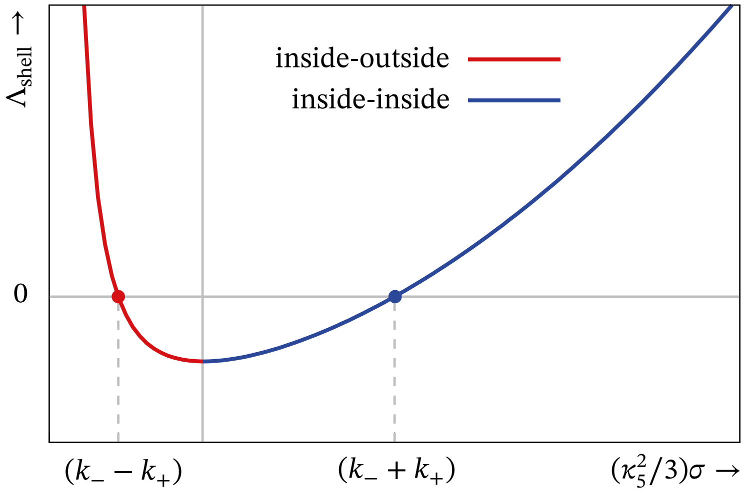

Figure 7.3. Figure (taken from paper II) showing the four dimensional cosmological constant on the shell as a function of the tension of the shell. The curve to the left of the vertical line represents the inside-outside scenario of the shellworld, while the curve to the right represents the inside-inside scenario of the RS construction. The dots represent the corresponding critical tension needed for a flat brane ($\Lambda_{\text{shell}} = 0$) in each case.

A small cosmological constant, thus requires a large kc, which implies

$$\sigma_{\text{dS}}^{\text{RS}} = \frac{6k}{\kappa_5^2}\left(1 + \mathcal{O}\left(c^{-n}\right)\right) = \sigma_{\text{crit}}^{\text{RS}}\left(1 + \epsilon\right). \tag{7.15}$$

Therefore, a super-critical brane is required to get a dS braneworld. In contrast, for a shellworld,

$$\Lambda_{\text{dS}}^{\text{shell}} = \epsilon\sigma_{\text{crit}}^{\text{shell}}, \qquad \sigma_{\text{dS}}^{\text{shell}} = \sigma_{\text{crit}}^{\text{shell}}\left(1 - \epsilon\right), \tag{7.16}$$

and a brane with sub-critical tension realizes a dS universe. This can be seen clearly in figure 7.3.

3. A third point of difference is the four dimensional Planck mass. For a flat RS braneworld, the volume of the fifth dimension is finite and the Planck mass can be read off from a dimensional reduction as discussed in section 6.3. This gives $\kappa_4^2 = k\kappa_5^2$. The same is also true for a dS braneworld. For the shellworld, however, the volume along the fifth dimension is not finite, and a dimensional reduction needs to be correctly regulated to read off the Planck mass. However, the four dimensional Planck mass can also be read off by deriving the Einstein's equations on the spherical brane on the same lines as the covariant approach by Shiromizu, Maeda, and Sasaki (2000). This involves using the Gauss-Codazzi equations to

project the five dimensional gravitational curvatures onto the four dimensional quantities on the brane. This was done in section 6 of paper III and gives

$$\kappa_4^2 \equiv 8\pi G_4 := 2\left(\frac{k_+k_-}{k_- - k_+}\right)\kappa_5^2. \tag{7.17}$$

This can also be read off both from the four dimensional graviton propagator as well as from the Friedmann equations on the expanding bubble. The graviton propagator is computed in section 3 of paper IV, while the Friedmann equations are computed in section 2 of paper II and section 3 of paper III. Of course, they all give the same result. Using the inside-inside vs inside-outside identification discussed before, this result correctly reduces to that of RS, by flipping the sign in front of k_+, giving $\kappa_4^2 = k\kappa_5^2$.

7.3 Populating the universe with matter and radiation

Let us now consider modifying the bulk AdS_5 metric by adding five dimensional matter. This backreacts on the geometry of the shell via the thin-shell junction conditions to give rise to matter and radiation on the shell, as we will summarize below.

The presence of five dimensional matter modifies the bulk geometry to five dimensional AdS Schwarzschild, which contributes as four dimensional radiation on the shell. This was shown using Friedmann equations in section 2 of paper II, and from the full four dimensional Einstein equations in section 5 of paper III. Four dimensional matter, however, arises in a much more interesting way – as end points of strings stretching in the fifth direction. This can be shown in two ways.

1. Consider a string of tension τ stretching along the fifth direction and ending on the shell. Einstein's equations require covariant conservation of the energy momentum tensor in five dimensions. It was shown in section 1 of paper IV, that this requires the end point of the string to behave like a particle. The mass of the particle is given in by the ratio of the tension τ of the string, τ, and the five dimensional AdS curvature k, *i.e.*, $m = \tau/k$.
2. An isotropic version of the above was considered in section 5 of paper III, where instead of a single string, an isotropic distribution of such strings was considered. By smearing their energy density parallel to the shell, one obtains the metric of a *cloud of strings*. This contributes like massive particles of mass τ/k to the four dimensional Friedmann equation, exactly like the single string. Not surprisingly, this can also be seen by writ-

ing down the four dimensional Einstein's equations, which was shown in section 5 of paper III.

In this way, a dS universe with matter and radiation can be realized on the shellworld, making it a model for late time cosmology.

7.4 Brane bending

To understand the effect of the stretched strings further, we can perform a linearized analysis akin to the brane bending computation in section 6.4. This was done in section 2 of paper IV. The main result is that, starting with an empty dS shellworld (*i.e.*, a brane with a finely tuned sub-critical tension in an empty AdS_5 bulk), placing a matter field directly on the brane gives the amount of bending as

$$\Box_{(4)} F = -\frac{\kappa_4^2}{6} S. \tag{7.18}$$

S is the trace of the source placed on the brane and $F := k_- f_- = k_+ f_+$ is constant across the brane by virtue of the first junction condition, although the brane appears bent by different amounts as viewed from either side ($f_\pm$ respectively).

This should be compared with the corresponding result in the RS scenario in equation (6.25). Most notably, the sign is flipped. This implies that, contrary to the RS construction, a matter field placed directly on the brane (*i.e.*, $S < 0$) causes the brane to bend inwards towards the centre of AdS. For a string stretching outward from the brane, $S \sim \alpha_+ / k_+$; while for a string stretching inward, $S \sim -\alpha_- / k_-$ with $\alpha_\pm > 0$ being the corresponding tensions. This shows that a brane stretching inwards also pulls down on the brane, while a string pulling outward has the opposite effect. Therefore, while a string stretching outward corresponds to a particle with *positive* mass on the brane, a string stretching inward, or matter placed directly on the brane contributes as a four dimensional particle with a *negative* mass. This is shown in figure 7.4, which is taken from paper IV . This is, of course, consistent with the result obtained from the Friedmann equation as well as the four dimensional Einstein equations on the shell. These equations show that a string pulling on the inside of the brane contributes as a negative mass particle, while a string pulling outward gives a positive mass in four dimensions, exactly as the bending analysis gives. Equation (7.18) is also consistent with a similar expression that was derived by Padilla (2005).

This should be compared with the brane bending in the RS construction as sketched in Garriga and Tanaka (2000). It was shown there that a positive mass in four dimensions causes the brane to bend outward just as the

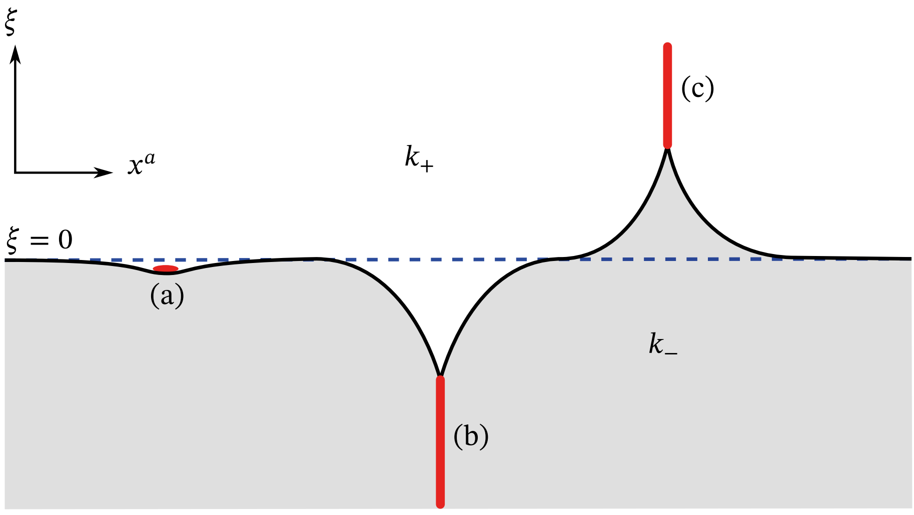

Figure 7.4. A schematic figure (taken from paper IV) showing the direction of brane bending due to sources in the bulk and on the brane. The interior of the bubble is shaded in gray, with the bold black line being the shell. The position of the brane prior to adding and mass or string on it is shown with the dotted line. Object (a) is a point source placed directly on the brane while (b) and (c) are strings ending on the brane from the inside and the outside respectively.

string stretching outward does in our model. In this regard, a particle with an effective positive mass in four dimensions, has a similar effect in our model as it has in the RS model.

7.5 Gravity on the shellworld

To study graviton modes on the shellworld, one has to perform a linearized analysis of perturbations to the five dimensional Einstein's equations. This gives the five dimensional Green's function corresponding to the graviton propagator in the bulk. Evaluating this with both points being on the shell, gives the four dimensional graviton propagator. This was computed in section 3 of paper IV. The result is that, we obtain Einstein's gravity on the shellworld for low momentum, and there are corrections when the momentum approaches the curvature scale $p \sim k_+$. This should be compared to the result obtained by Padilla (2005), who performed a similar analysis for a generalized RS like scenario, and found both UV as well as IR corrections to Einstein's equations. The crucial difference between the two setups is the presence of stretched strings for the shellworlds. As we have discussed, the presence of strings stretching outward from the brane is crucial for generating massive four dimensional particles. This allows for non-normalizable modes in the solution for the graviton Green's function. In contrast, for the case considered by Padilla (2005), the wavefunction was required to vanish at the boundary of AdS. This gives qualitatively different results.

For this mode to be the four dimensional graviton, one also has to make sure that it has the right tensor structure, as discussed in section 6.4. Similar to the RS setup, the brane bending term modifies the factor of $1/3$, to $1/2$ in the four dimensional graviton propagator. This was also computed in section 3 of paper IV. A curious feature, which might appear surprising at first is that the graviton propagator, when convoluted with the matter source has a negative sign. The minus sign appears because, as discussed before, adding a matter source directly on the brane corresponds to having a four dimensional particle of negative mass. Taking the end point of an outstretched string as the source on the brane, should remove the minus sign.

The discussion in this section highlights a truly novel feature of the shellworld construction. The presence of stretched strings is crucial for four dimensional Newton's potential to appear on the shell. The presence of these strings results in a special form of the graviton propagator – although it produces four dimensional gravity on the shellworld (upto high energy corrections), it is not localized along the fifth dimension. In fact, the five dimensional graviton mode, which produces four dimensional gravity, is non-normalizable and extends along the strings.

The strings that stretch out from the shell must end somewhere. In the discussion so far, we have implicitly assumed that they stretch all the way to the boundary or close to it where they could end a second brane. However, there is no reason for this brane to be so far away, and it could very well be close to the shell. Such a scenario has its motivation in another question that we have conveniently ignored so far: *how do these strings form in the first place?* Although we have not yet addressed this question in our work, in paper III, we briefly speculated on the possibility that collisions between branes could create strings stretching between them. This could happen, for instance, when two such shellworlds collide with each other. This gives rise to an intriguing possibility sketched in figure 7.5, which is taken from paper IV. In such a scenario, gravity is truly localized between the branes via the stretched strings and whole *sandwich* could be considered as a thick brane. This also raises an associated question regarding the lifetime of such a universe. Since the bubbles are accelerating, one would expect them to collide sooner than later, spelling doom for the universe on the shell. This is indeed a reasonable outcome, and would be expected to happen instantaneously from the perspective of a five dimensional observer. However, it

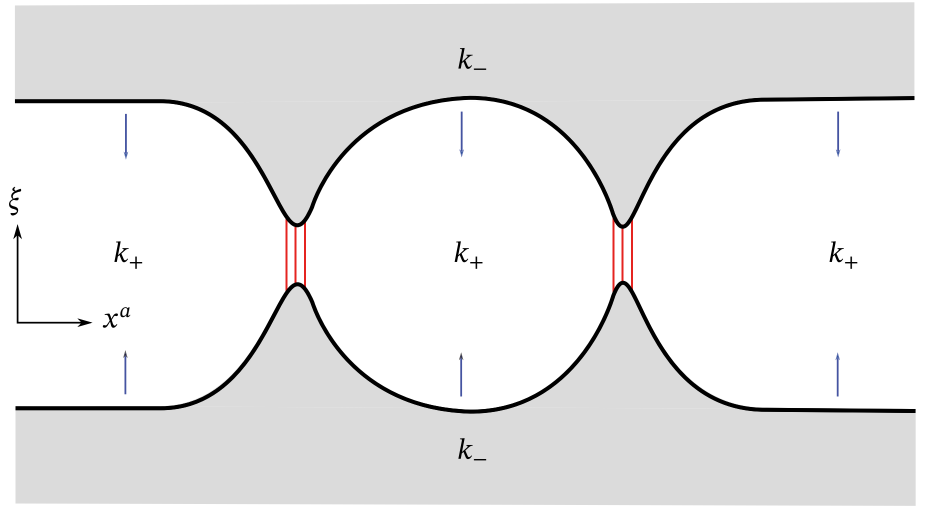

Figure 7.5. A cartoon (taken from paper IV) representing a modified shellworld construction, where the stretched strings end on an adjacent shellworld. The areas in gray represent the interior of the corresponding bubbles. In this *sandwich* model, gravity is completely localized, via the stretched strings, between the two branes.

was argued in paper III, that this could take a very long time as seen from the perspective of a shellworld observer.

7.6 Swampland conjectures

The construction presented above, is a model of late time cosmology in a dS universe. A natural question to ask is, whether there is any hope of being able to realize this in string theory. To address this question, let us briefly comment on how this model measures up against the swampland conjectures and the no-go theorems discussed in the previous chapters.

1. The dS conjecture discussed in section 5.4 is based on assuming a static compactification of a ten or eleven dimensional theory. The shellworlds, as we have discussed, are spherical instantons that nucleate at rest and accelerate outwards asymptotically approaching the speed of light. This implies that the position of the shell is time dependent and thus the full five dimensional metric is also time dependent, via the position of the shell. In this way, the shellworlds bypass the dS conjecture. Note that the four dimensional Planck scale as well as the cosmological constant, are however, time independent as explained in this chapter.
2. The Maldacena-Núñez no-go theorem also does not apply here. The reason is that it requires the volume of the extra dimensions to be finite,

which is not the case for our fifth dimension. This is also the case for an AdS braneworld in the RS construction (as pointed out by Karch and Randall, 2001), as we have discussed in the previous chapter.

3. The WGC and the non-supersymmetric AdS conjecture turn out to be helpful for us, and provide motivation for the decay of metastable AdS that our model relies on. As we have mentioned at the beginning of section 5.4, nucleation of the spherical brane instanton, that is our shellworld, is in fact motivated by these conjectures. The generalized WGC for p-form gauge fields, discussed in section 5.2, which governs such a nucleation, turns out to be trivially true for our case ($d = 5, p = 4$). Such a nucleation of a co-dimension one brane, was stated as an implicit motivation for the non-supersymmetric AdS conjecture by Vafa (2005), but has not yet been made explicit to the best of our knowledge. In section 2 of paper III, we tried to make this explicit by demanding consistency with energy conservation à la Brown and Teitelboim. Further support for the instability of non-supersymmetric AdS vacua in string theory comes from the work by Danielsson and Dibitetto (2017) and Danielsson, Dibitetto, and Vargas (2017).

7.7 Support from string theory

There does not seem to be any obvious obstruction to construct a model of shellworlds in string theory. What is required for an explicit construction is a five dimensional compactification, which has at least two non-degenerate AdS_5 vacua. The vacuum with higher energy should be non-supersymmetric, whereas the other can either be supersymmetric or non-supersymmetric. If the second vacuum is non-supersymmetric, it is likely that there will exist a third supersymmetric vacuum in the theory allowing that vacuum to decay, giving another shellworld. This cascade of decays can continue until one ends up in a supersymmetric AdS_5 vacuum. For the non-perturbative vacuum decay to occur, tension of the brane that mediates the decay between these vacua needs to be sub-extremal. In supergravity, this brane is a domain wall that sources the change in the flux-induced superpotential across it, and its tension can be calculated. To realize the shellworld construction, this tension has to be less than the critical tension, which can be computed from the vacuum energies.

In section 4 of paper II, we presented a construction in type IIB string theory that has a non-supersymmetric and a supersymmetric AdS_5 vacuum with the right hierarchy of vacuum energies. Further, we showed that the domain wall tension interpolating between these vacua was indeed subcritical. An explicit embedding of the branes needed for this construction was not addressed there, and so this is not yet an explicit construction in string theory. However, it does provide strong evidence in favor of the viability of such a construction in string theory and that the necessary ingredients are not difficult to obtain.

This should be compared with the RS construction. As we have discussed before in this chapter, four dimensional Minkowski and dS braneworld constructions in supergravity are restricted by no-go theorems and the only viable candidate for a string completion appears to be the AdS_4 braneworld. Kraus (1999) generalized this setup to a moving AdS_4 domain wall in supergravity and found that the tension of the brane needed for such a realization was much larger than those available from string theory.

7.8 What next?

To summarize, in this chapter we have presented a novel construction of a four dimensional dS vacuum in a string theory inspired setup. The dS universe sits on the world volume of an expanding bubble, and the setup is time dependent. We have explored gravitational aspects of this construction, and have shown how matter and radiation arise on the shellworld, providing a model of late time cosmology. This five dimensional construction also seems to evade the swampland conjectures, and there does not seem to be any obvious obstruction against constructing such shellworlds in string theory. This makes it a potentially viable model for the burning question of constructing a dS vacuum in string theory. In order to achieve that goal, an explicit construction of this model in string theory is required. Although a toy model was presented in paper II, showing that the requisite ingredients, with the necessary hierarchy exist in string theory, a fully explicit construction has not yet been completed. This is a very interesting direction open for future research, which, if successful, could have a significant impact.

While the model, with the stretched strings, is a model for late time cosmology, the early universe – inflation in particular – has not yet been addressed. Accommodating inflation might require adding extra ingredients in addition to the ones currently present. This was briefly pointed out in the outlook of paper III and still remains a very interesting avenue for further work.

Nature of the stretched strings, and the kind of matter represented by their endpoints, also needs further exploration. Given that the endpoint of the string represents a particle of mass τ/k, it was argued in paper III, that to get realistic standard model particles on the shellworld, these cannot be fundamental strings, but need to be low tension strings such as gauge strings or topological defects. Constructing the standard model of particle physics on the shellworld, using the stretched strings as well as gauge fields living on the brane, is a challenging and very interesting open question.

A natural question to ask is about the holographic interpretation of the expanding bubble in terms of the dual field theory on the boundary of the AdS_5 at infinity. This becomes difficult to answer because the AdS is non-supersymmetric by construction, and non-supersymmetric holography is not well understood. Heuristically, this is because the probability of bubble nucleation over the full volume of AdS is unity – implying that the field theory on the boundary decays instantaneously. However, in the presence of a brane on which the strings can end, for example in the sandwich construction proposed in figure 7.5, the second brane can serve as a holographic screen on which the dual field theory can be studied. Since radially stretched strings are the natural way to obtain massive particles in an asymptotically AdS spacetime, understanding the holographic dual of this model can lead to further insights. See the outlook of papers III and IV for more discussion on this.

The question of how four dimensional black holes are realized on the shellworld is yet another interesting open question. We will comment on this briefly in chapter 9, after the discussion of black shells in the next chapter.

Part II:
Black holes

8. Black shells

In the first part of the thesis, we have discussed how the question of dark energy brings together general relativity and particle physics, requiring us to turn to a theory of quantum gravity, like string theory, for answers. Another cosmological object which demands a treatment in quantum gravity, is a black hole. In this chapter, we will briefly sketch why this is the case. We will then highlight a model of *black shells* that we proposed, in papers V and VI, in an attempt to answer some of these questions.

8.1 Black hole thermodynamics

As an object falls past the event horizon of a black hole, its degrees of freedom become inaccessible to the outside universe. This leads to an apparent decrease in the entropy of the universe, unless the black hole itself carries an entropy that increases by an amount equal to or larger than the entropy of the infalling object. Therefore, the second law of thermodynamics leads to the notion of black hole entropy.

Furthermore, Hawking (1971) showed that in classical general relativity, under very general assumptions, the area of the event horizon is a non-decreasing function of time. This led Bekenstein (1973) to propose that black holes have an entropy proportional to their horizon area giving rise to a *second law of black hole thermodynamics*. Bardeen, Carter, and Hawking (1973) showed that to make the analogy with thermodynamics more concrete, a *first law* could be proposed if black holes carried a temperature given by the scale of their horizon radius. The discovery by Hawking (1975) that black holes have a temperature, made these ends meet, and the constant of proportionality between the entropy and the area of the horizon was found. Let us outline the derivation of the entropy below, but using a more modern approach. In the next paragraphs, we will first derive the temperature of a Schwarzschild black hole using the *Euclidean periodicity trick*, and then find the numerical constant in the entropy using the second law of thermodynamics.

Temperature

To derive the temperature of a black hole, let us start with a Schwarzschild black hole of radius r_s,

$$\mathrm{d}s^2 = -\left(1 - \frac{r_s}{r}\right)\mathrm{d}t^2 + \left(1 - \frac{r_s}{r}\right)^{-1}\mathrm{d}r^2 + r^2\mathrm{d}\Omega_2^2, \tag{8.1}$$

and write it in the near-horizon radial coordinate ε, which is defined as $r = r_s\left(1 + \varepsilon^2\right)$ for $\varepsilon \ll 1$. To simplify it further, let us ignore the angular directions. This gives the following near-horizon metric

$$\mathrm{d}s^2 = -\varepsilon^2\mathrm{d}t^2 + 4r_s^2\mathrm{d}\varepsilon^2. \tag{8.2}$$

This resembles the *Rindler metric* that describes the spacetime seen by a uniformly accelerating observer in flat space,

$$\mathrm{d}s^2 = -R^2\mathrm{d}\eta^2 + \mathrm{d}R^2. \tag{8.3}$$

This metric has a horizon at $R = 0$ and the time coordinate is periodic with an imaginary period, $\eta \sim \eta + 2\pi i$. Such a Rindler observer experiences a universe in a thermal bath, with the temperature given by the inverse of the period. The near-horizon Schwarzschild geometry can be identified with the Rindler geometry by making the identification $R = 2r_s\varepsilon$ and $\eta = t/(2r_s)$. The periodicity of $\eta \sim \eta + 2\pi i$ then implies a periodicity of time $t \sim t + 4\pi r_s i$. Analogous to the Rindler spacetime, this is just the time coordinate of a quantum field theory at a finite temperature T, *i.e.*, $t \sim t + \beta i$ for $\beta \equiv 1/T$ (in units where the Boltzmann constant is unity, $k_B = 1$). From this, the temperature of the Schwarzschild spacetime can be identified as

$$T = \frac{1}{4\pi r_s}. \tag{8.4}$$

This trick to find the temperature of a black hole spacetime can be justified by using the Euclidean path integral. See the lecture notes by Hartman (2015) for a pedagogical introduction.

Thermodynamics

Let the entropy of the black hole be given by $S = C \cdot \text{area} = C \cdot 4\pi r_s^2$, where C is the constant of proportionality that we want to determine, and $r_s \equiv 2GM$ is the Schwarzschild radius of the black hole (of mass M). Given the temperature of the black hole obtained above, its entropy can be computed from the first law of thermodynamics, $\mathrm{d}E = T\mathrm{d}S$, provided that we know the energy of the black hole. This can be computed using the ADM formalism, which gives the energy of an asymptotically flat spacetime as the boundary integral at spatial infinity. Performing the integral for the

Schwarzschild spacetime gives that the energy is equal to the mass of the black hole, *i.e.*, $E = M$. With this, the first law of thermodynamics gives

$$\frac{\mathrm{d}S}{\mathrm{d}E} = \frac{1}{T} \Rightarrow 32\pi G^2 M \cdot C = 8\pi G M \Rightarrow C = \frac{1}{4G}. \tag{8.5}$$

This gives the famous Bekenstein-Hawking entropy:

$$S = \frac{A}{4G} = \frac{A}{4\ell_{\mathrm{Pl}}^2}\frac{\hbar}{c^3}, \tag{8.6}$$

where in the last step we have expressed the Newton's constant in terms of the four dimensional Planck length, $G = \ell_{\mathrm{Pl}}^2 c^3/\hbar$. To get an idea of how big the entropy of a black hole really is, let us compute this for a solar mass black hole, which has a radius of $r_s \sim 3$ km. This gives an enormous entropy $S \sim 10^{77}$. Turning our attention back to the temperature of the black hole, we can restore all fundamental constants in the expression for temperature above, to get

$$T = \frac{\hbar c^3}{8\pi k_B G M}. \tag{8.7}$$

For a solar mass black hole, this turns out to be $T \sim 10^{-7}$K, which is smaller than the cosmic microwave background radiation by 7 orders of magnitude. A larger black hole will have an even smaller temperature and an even bigger entropy. This shows that black holes are incredibly cold objects, but carry a huge amount of entropy.

From equation (8.7), we see that the temperature vanishes both in the classical limit ($\hbar \to 0$) as well as in the Newtonian limit ($G \to 0$), indicating that the Hawking temperature and the associated Hawking radiation is a phenomenon that is fundamentally quantum mechanical as well as general relativistic. This gives an indication that black holes are objects that are interesting for understanding the physics at the intersection of both of these regimes, making them ideal test beds for understanding quantum gravity.

Given the thermodynamic nature of black holes, Bekenstein interpreted equation (8.6) as corresponding to their actual thermodynamic entropy, meaning that S gives the logarithm of the number of microstates of a black-hole of a given mass. However, general relativity (through Birkhoff's theorem) predicts that a Schwarzschild black hole of a given mass is unique, and hence should have zero entropy ($S = \ln 1 = 0$). Another way to state this tension between the two results is that, the process of black hole formation breaks time reversal at a microscopic scale, since degrees of freedom (information) of the collapsing matter is lost in the final black hole and there is no way to find out what fell into it.

This extends naturally to black holes with charge and angular momentum. The energy for such a black hole receives additional contributions from its charge Q and its angular momentum J. This gives the following first law of black hole thermodynamics

$$\mathrm{d}M = T\mathrm{d}S + \Phi\mathrm{d}Q + \Omega\mathrm{d}J, \tag{8.8}$$

for electric potential Φ, and angular velocity Ω at the horizon. The entropy also receives extra contributions from charge and spin, giving a black hole with an enormous entropy, while general relativity predicts only three degrees of freedom – charge, mass and angular momentum (*no-hair theorem*, see section 8.4).

8.2 Information paradox

Hawking (1975) showed in his celebrated paper that because of its temperature, a black hole radiates like a gray body with the rate of emission in a mode of frequency ω, harmonic l, and width $\mathrm{d}\omega$ given by

$$\frac{dE}{dt} = \frac{\omega \mathrm{d}\omega}{2\pi}\frac{T_{\omega,l}}{e^{\beta\omega} - 1}, \tag{8.9}$$

where $T_{\omega,l}$ is the gray body factor describing the departure from pure black body radiation, and $\beta = 1/T$ is the inverse of the Hawking temperature. This Hawking radiation is independent of the characteristics of the matter that fell into the black hole. This implies that any information about what fell into the black hole cannot be carried out by the Hawking radiation and is lost for ever.

This can be made precise with the following thought experiment. Let us take a pure state, say n pairs of EPR particles, and throw one of each pair into a black hole. If we wait long enough, the black hole will decay via Hawking radiation. We have seen above, that these Hawking quanta do not carry any information about the particles that fell into the black hole. When the black hole has evaporated completely, we will be left with our half of the EPR pair, while the other half that we threw into the black hole has completely disappeared into Hawking radiation. Our half of the EPR pair is, therefore, no longer entangled with anything. Something remarkable has happened – we started with a pure state of entangled pairs and are now left with a bunch of particles that are not entangled with anything anymore, *i.e.*, a mixed state. Such an evolution of a pure state to a mixed state can only occur via the action of a non-unitary operator on the state, and such operators are not allowed in quantum mechanics. Evaporation of the black hole has thus violated unitarity, which cannot happen in quantum mechanics or gravity. This, in summary, is the *black hole information paradox*. A detailed discussion of various aspects of this paradox can be found in the reviews by Harlow (2016) and Polchinski (2017).

8.3 Resolution of the information paradox

In an attempt to resolve this paradox, the various assumptions leading up to its formulation have been critically examined. In this section, we will briefly summarize some of the main lines of investigations. We will divide them into two groups: (i) black holes form normally, but evaporation requires corrections to semi-classical physics, (ii) corrections to the semi-classical theory prevent a black hole from forming in the first place.

8.3.1 Departure from semi-classicality during evaporation

Let us first examine those scenarios that offer a solution to the black hole information problem by modifying the process of black hole evaporation, without affecting their formation in general relativity.

Remnants

In the thought experiment above, the pure state went into a mixed state because the black hole was allowed to evaporate completely. However, as the black hole continues to evaporate, and gets down to a Planck sized object, the semiclassical physics used to analyze its evaporation is no longer enough. Since it is not possible to say what happens beyond this point, a reasonable stance to take is that the evaporation stops at this stage and the black hole survives for ever as a Planck mass object. Such an object is a *remnant* similar to the charged ones discussed in section 5.1. In order for remnants to provide any respite from the information paradox, they need to carry all of the entropy of the black hole that they descended from. This would suggest that all black holes end up as Planck mass remnants, but carry vastly different entropies. There would then be an infinite number of such Planck mass remnants in the universe, corresponding to every black hole that ever formed. This has implications for low energy quantum field theory. Normally, the correction to a scattering amplitude from a Planck mass object running in loops is suppressed because of its mass. However, having an infinite number of them would lead to an infinite contribution making all QFT amplitudes divergent. Apart from the breakdown of the effective field theory, remnants also lead to violation of covariant entropy bounds and are therefore believed to be forbidden. See the article by Susskind (1995) for a discussion of some issues with remnants.

Small corrections
Since Hawking's computation for the evaporation is semi-classical, one could argue that the reason we have a paradox, is that we have trusted equation (8.9) all the way to the Planck scale, where it might no longer be valid. Following this argument, it would seem that the information could be carried out in small correlations between the quanta of Hawking radiation. However, Mathur (2009) proved that such small corrections, even if they existed would not be sufficient to resolve the information paradox. Moreover, equation (8.9) is valid up to order m_p/M, and any resolution to the information paradox has to start reducing the entanglement entropy much before the black hole reaches Planck mass. See the discussion on the *Page curve*, in one of the reviews mentioned before.

Black hole complementarity
A solution to the information paradox was proposed by Susskind, Thorlacius, and Uglum (1993), who argued that causally disconnected observers could, in principle, observe mutually inconsistent versions of a event, as long as they cannot compare notes. This has interesting consequences when applied to the black hole information problem. The equivalence principle says that, an observer falling into a large black hole should not see anything special at the horizon. To prevent loss of information, a distant observer however, should see the infalling observer disappear at the horizon, and would expect her degrees of freedom to reappear as Hawking radiation. The essence of black hole complementarity is that, both of these point of views could be true at the same time. They argued that the observer who sits outside the black hole, and collects all the Hawking quanta corresponding to her in-fallen friend, cannot then jump into the black hole to see her friend again. So from her point of view, her friend has thermalized at the horizon. In this way, the equivalence principle and unitary evolution could be reconciled in an observer dependent way. This argument for black hole complementarity relies on three basic postulates: (i) Hawking radiation is in a pure state, (ii) the equivalence principle holds for an infalling observer, (iii) all information carried by the radiation is emitted near the horizon, and effective field theory holds at a microscopic distance away from the horizon.

Firewall
A decade and a half after the proposal of black hole complementarity, an inconsistency in the argument was discovered by Mathur (2009), which was

then refined by Almheiri, Marolf, Polchinski, and Sully (2013) and Marolf and Polchinski (2013), to show a sharp inconsistency among the assumptions of complementarity. The following is a caricature of their argument. Consider an old black hole formed from a pure state, *i.e.*, whose Hawking radiation forms a subsystem at least as big as the state of the black hole itself. Let us identify three subsystems: (A) the early time Hawking radiation far away from the black hole, (B) late time Hawking radiation near the black hole horizon, and (C) the interior of the black hole. The equivalence principle implies that the horizon is locally like flat space, and thus subsystems (B) and (C) should be highly entangled. However, for the external observer to not observe any loss of information, subsystem (B) should also be highly entangled with subsystem (A), implying that the subsystem (B) is highly entangled with both subsystems (A) and (C) at the same time. This violates the principle of monogamy of entanglement (*i.e.*, a system can be strongly entangled with only one other system at a time), which follows from the principle of strong sub-additivity of the von Neumann entropy, that was proved by Mathur (2009). This would still not be a problem for the complementarity argument, but they additionally showed that there is a reference frame in which a single observer can measure both of these entanglements, and thus black hole complementarity is inconsistent.

To remedy the situation, they suggested violating one of the three assumptions of black hole complementarity listed before. They argued that the most conservative approach was to drop the assumption that the infalling observer sees nothing special at the horizon. The suggestion was that the horizon of a black hole has a *firewall* instead, which is a structure where an infalling observer thermalizes instead of being able to fall into the black hole.

8.3.2 Departure from semi-classicality during formation

Another radical class of proposals suggest that a black hole never forms during a gravitational collapse, and some other horizonless structure is formed instead. This provides a way out the information paradox, since there is no paradox in the absence of a horizon. Motivating the absence of a black hole at the end of gravitational collapse requires either a modification of general relativity or another structure that can preferentially form in the place of a black hole, within the framework of general relativity. Proposals of both kinds have been made in the literature. See the review by Cardoso and Pani (2017b), especially table 1, for a classification of such proposals. In paper V, we proposed an alternative end point of gravitational collapse within the framework of general relativity, which is motivated by string theory. Here, we will briefly review different classes of such horizonless objects and point out where our proposal fits into the picture.

Such horizonless objects are classified according to their observational signatures, which are usually based on the orbits of massive and massless particles around them. Of these, the smallest possible circular orbits – both for photons as well as for massive particles – are particularly diagnostic of these objects. We will derive these below. Consider a spherically symmetric metric of the following form:

$$\mathrm{d}s^2 = -f(r)\mathrm{d}t^2 + f(r)^{-1}\mathrm{d}r^2 + r^2\left(\mathrm{d}\theta^2 + \sin^2\theta\,\mathrm{d}\phi^2\right). \tag{8.10}$$

Metric compatibility of the connection implies that

$$\epsilon = -g_{\mu\nu}\frac{\mathrm{d}x^\mu}{\mathrm{d}\lambda}\frac{\mathrm{d}x^\nu}{\mathrm{d}\lambda}, \tag{8.11}$$

is a constant of motion for an affine parameter λ. A massive particle follows a timelike trajectory and has $\epsilon = 1$, while a massless particle follows a null trajectory with $\epsilon = 0$. The Killing vectors $K = \partial_t$ and $K = \partial_\phi$ correspond to conservation of energy and angular momentum respectively:

$$E = f(r)\frac{\mathrm{d}t}{\mathrm{d}\lambda} \equiv f(r)\dot{t}^2, \quad L = r^2\frac{\mathrm{d}\phi}{\mathrm{d}\lambda} \equiv r^2\dot{\phi}^2. \tag{8.12}$$

The other two Killing vectors, corresponding to the conservation of the direction of the angular momentum, imply that motion is restricted to a plane, which we can choose to be the equatorial plane $\theta = \pi/2$. This gives

$$-f(r)\dot{t}^2 + f(r)^{-1}\dot{r}^2 + r^2\dot{\phi}^2 = -\epsilon, \tag{8.13}$$

which can be rewritten as a kinetic equation for $r(\lambda)$,

$$\dot{r}^2 = E^2 - f(r)\left(\frac{L^2}{r^2} + \epsilon\right) \equiv -V_{\mathrm{eff}}(r). \tag{8.14}$$

Written in this form, circular orbits correspond to extrema of the potential, $V'_{\mathrm{eff}}(r) = 0$,

$$V'_{\mathrm{eff}}(r) = r^2 r_s\epsilon - L^2(2r - 3r_s) \stackrel{!}{=} 0. \tag{8.15}$$

Furthermore, the orbit is stable if it corresponds to a minimum of the potential, and is unstable otherwise. This implies that a photon ($\epsilon = 0$) has circular orbits at $r_0 = 3r_s/2$ for all values of L. However, this turns out to be the maximum of the potential, implying that these orbits are unstable. This surface of unstable circular photon orbits is called the *photosphere*. A massive particle ($\epsilon = 1$), on the other hand, has two circular orbits at

$$r = \frac{L^2}{r_s}\left(1 \pm \sqrt{1 - \frac{3r_s^2}{L^2}}\right). \tag{8.16}$$

	unstable circular orbits	stable circular orbits
massless particles	$r = 3r_s/2$	—
massive particles	$3r_s/2 \leq r < 3r_s$	$r \geq 3r_s$

Table 8.1. Location of stable and unstable circular orbits for massless and massive particles in the Schwarzschild geometry. r_s is the Schwarzschild radius of the black hole.

The one at smaller r is unstable, while the other one is stable. Either of these circular orbits exist only when the angular momentum is larger than a critical value namely, $L \geq \sqrt{3}r_s$. At this critical angular momentum, the two roots coincide to give a stable orbit at $r = 3r_s$. For very large angular momentum, $L \gg 1$, the smaller root converges to an unstable orbit at $r = 3r_s/2$, analogous to the photosphere, while the larger orbit is pushed out to infinity. The smallest stable circular orbit is therefore at $r = 3r_s$, and is called the *innermost stable circular orbit* (ISCO).

To summarize, in the Schwarzschild geometry, a massless particle only has unstable circular orbits. These are located at the photosphere, $r = 3r_s/2$. Massive particles on the other hand, have unstable circular orbits between $3r_s/2 \leq r < 3_s$, and stable circular orbits for $r \geq 3r_s$. This is summarized in table 8.1. Therefore, the two radii $r = 3r_s/2$ and $r = 3r_s$ are special for the Schwarzschild geometry.

By Birkhoff's theorem, the geometry outside any uncharged, non-rotating object is given by the Schwarzschild metric. So any object dense enough to have a radius smaller than $r = 3r_s$, would have an ISCO at $r = 3r_s$, whereas bigger objects would not. Such objects smaller than $r = 3r_s$ are called *compact objects*, and the presence of an ISCO would observationally distinguish them from their non-compact counterparts. The densest known objects in classical general relativity are neutron stars, which by the above definition are compact objects, and any object with a higher density is expected to collapse into a black hole. Therefore, any object that is compact but not a neutron star, must be made of *exotic* matter, and is called an *exotic compact object (ECO)*.

On the other end of the range of special radii lies the photosphere. Being the closest circular light orbit to a black hole, the region within the photosphere is expected to look dark when imaging a black hole in the electromagnetic spectrum. This dark region is called the *shadow* of the black hole. An object with a radius smaller than the photosphere is therefore expected to look similar to a black hole in the electromagnetic spectrum and is called an *ultra-compact object (UCO)*. The dynamics of light close to the photosphere is however, vastly more interesting than that of the ISCO. This is because, contrary to the ISCO, the photosphere is an unstable extremum of the potential and orbits of light very slightly inside the photosphere will fall inwards with time. For a black hole, such a light ray will encounter vacuum all the way to the horizon where it is absorbed. However, for an UCO, it will find the surface of the object before hitting the horizon and can scatter back to the photosphere. This is true for both electromagnetic as well as gravitational waves, and the properties of the photosphere can help to distinguish between a black hole and a horizonless UCO. However, Cardoso and Pani (2017a,b) showed that if the radius is close enough to the Schwarzschild radius, most of the waves trapped at the photosphere will have relaxed away by the time the reflected wave makes it way back to the photosphere. For an object of radius $r = r_s(1+\epsilon)$, they showed that this corresponds to $\epsilon \lesssim \epsilon_{\text{crit}} \sim 0.0165$. An object satisfying this limit, will therefore have a photosphere very similar to that of a black hole, and is called a *Clean Photosphere Object (ClePhO)*. Another special value of the radius is the *Buchdahl radius*. For an object made of an incompressible isotropic fluid, whose pressure does not decrease as one goes towards the center of the object, compressing it beyond $r = 9r_s/8$ makes the pressure at the center divergent. This limit was first found by Schwarzschild (1916) when computing the metric inside an incompressible fluid sphere, and was later generalized by Buchdahl (1959) to any matter distribution with these properties. Therefore, any compact object smaller than the Buchdahl radius must be made of matter that does not obey these assumptions.

To summarize, we have discussed four special values of the radius: the ISCO at $r = 3r_s$, the Buchdahl radius at $r = 9r_s/8$, the photosphere at $r = 3r_s/2$, and the ClePhO radius at $r = 1.0165r_s$. The significance of these radii is that they can be used to classify ECOs, based on their observational signatures. A comprehensive summary of ECOs can be found in reviews by Cardoso and Pani (2017b, 2019). Here, we will only highlight some aspects of two particular ECOs, namely gravastars and black shells. Based on the above reviews, we present a classification of ECOs, depending on their radii, in figure 8.1.

Gravastars

Gravastars are short for *gravitational vacuum stars*, and were constructed by Mazur and Mottola (2001) as an alternate end point of gravitational collapse. They are non-singular, spherically symmetric horizonless objects, consisting of a layer of stiff matter ($p = \rho$) sandwiched between two thin shells with dS space in the interior. The metric is continuous across the thin shells, with energy densities on them given by the thin-shell junction conditions. By replacing the two shells and the enclosed matter with a single

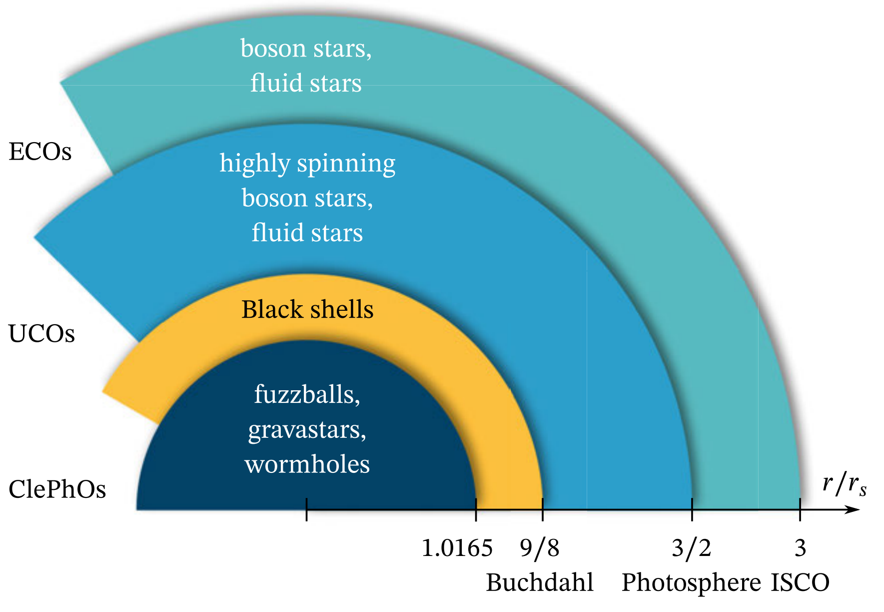

Figure 8.1. A schematic classification of compact objects following Cardoso and Pani (2017b). The black shells proposed in paper V are UCOs and can be as small as $r = r_s$ (for charged extremal blackshells) to as large as the Buchdahl radius (for uncharged black shells).

shell, Visser and Wiltshire (2004) constructed a *thin-shell* gravastar. These thin-shell gravastars share essentially the same properties as Mazur and Mottola's gravastars, but are much easier to analyse. Unlike Schwarzschild black holes, which have a negative specific heat and are thermodynamically unstable, gravastars are thermodynamically stable. Visser and Wiltshire analysed dynamical stability of thin-shell gravarstars against radial perturbations and showed that they can be stable under some physically reasonable equation of state.

Black shells

Motivated by gravastars, we proposed a novel construction in paper V, as non-singular horizonless objects, which are thin spherical shells containing AdS space inside. We call these objects *black shells*. The main motivation for replacing the dS core of the gravastar with AdS, is to construct such objects from a quantum theory of gravity, like string theory. As we have discussed in the first part of this thesis, construction of dS vacuum in string theory is a difficult challenge, and whether or not such a construction is possible is still under debate. However, there are several well understood examples of supersymmetric AdS vacua in string theory. We showed how such a black shell can be constructed out of one such vacuum, which is one of the sixteen supersymmetric AdS vacua in type IIA string theory compactified on a twisted torus orbifold $\mathbb{T}^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ that were found by Dibitetto, Guarino, and Roest (2011). Let us briefly outline the construction of these shells below.

The idea behind black shells is to construct a spherical shell enclosing an AdS_4 space in its interior, and having an effective Schwarzschild metric in the exterior. The physical properties of such a shell are given by the thin-shell junction conditions, as discussed in equations (6.3) and (6.4). Just as for the thin-shell gravastars of Visser and Wiltshire (2004), naïvely, the junction conditions do not give a sensible equation of state for the matter that makes up the shell. However, since our shells should follow from a construction in string theory, details of such a construction provide the physical properties of the matter that make up the shell.

The main idea is to consider D0 branes in four dimensions. These could also be Dp branes wrapped along p internal directions, which look like D0 branes in four dimensions. These branes can polarize into spherical D branes à la Myers (1999), with the D0 brane charge dissolved in them. Non-extremal black holes in general, and a Schwarzschild black hole in particular, would have both D0 and $\overline{\text{D0}}$ branes such that the total charge is given by the difference in their numbers. The D0-$\overline{\text{D0}}$ system could be stabilized thermally, similar to the discussion by Danielsson, Guijosa, and Kruczenski (2001) and could exist without annihilating. Physical properties of matter present on the spherical brane can be deduced from its world volume action, which consists of the DBI and Chern-Simons terms. This was discussed in detail in paper V, with the conclusion that there are three main contributions to the stress tensor on the brane, which in the limit of a microscopic AdS radius are:

(i) tension of the spherical brane, with equation of state $p = -\rho = -\tau$,
(ii) dissolved D0 branes that behave like matter with a stiff equation of state $p = \rho$, and
(iii) a gas of open strings stretching between the D0 branes, with an equation of state $p = \rho/2$.

These physical properties of the matter constituting the shell allow for a solution to the thin-shell junction conditions, and fix the radius of the shell. For an uncharged shell, this turns out to be the Buchdahl radius $r = 9r_s/8$, while for a charged shell, the radius approaches the Schwarzschild radius as the charge-to-mass ratio approaches extremality. Thus, in the classification scheme for compact objects, the uncharged shell is an UCO, and the charged shell can be anything from a ClePhO to an UCO depending on how close to extremality it lies. This is shown in figure 8.1.

Since we have used high energy objects to construct the shell, it is interesting to see how the energy balance works. The shell has an enormous tension, which is balanced by the huge negative energy of the AdS_4 in its interior. The energy of the dissolved D0 particles and the strings stretching between them contribute a small effective mass, which is visible from outside as the mass of this object.

In order for such an object to be an alternative to a black hole, there has to be a mechanism for the shell to form at the end of a gravitational collapse. The shell is a spherical instanton that can nucleate à la Brown and Teitelboim, if the Minkowski vacuum that we live in is unstable to such a tunneling process. The probability for such an instanton tunneling event is extremely small and is suppressed by the exponent of the square of its radius $\Gamma/V \sim \exp\left(-r^2\right)$ making the metastable Minkowski extremely long lived. However, in the presence of collapsing matter, entropy of the infalling matter can compensate for this suppression leading to a huge enhancement of the probability of nucleation. This is how such shells would form. When matter falls onto such a shell, its degrees of freedom would likely get absorbed into open string degrees of freedom of the gas on the shell, making such shells highly absorbing, hence the name *black* shells. It was argued in paper V that these objects are stable against small radial perturbations, with the stability arising from an exchange of energy between the various components.

8.4 Spinning black objects

Discovery of the Kerr solution was followed by the development of uniqueness theorems in the 1960s and 1970s, which led to the famous phrase by Ruffini and Wheeler (1971) – *a black hole has no hair*. An overview can be found in the review by Chrusciel, Lopes Costa, and Heusler (2012). Roughly, the no-hair theorem states that, *a stationary asymptotically flat spacetime, which is non-singular at and outside a connected event horizon, is uniquely characterized by its mass, charge and angular momentum*. This is a powerful theorem because, although the spacetime outside a spinning black hole could have an infinite number of multipole moments, it can only have three degrees of freedom namely mass, charge, and angular momentum, *i.e.*, the multipole moments are not all independent. However, all of this relies on the presence of a horizon, in the absence of which, no such uniqueness theorem exists and the geometry outside a spinning axisymmetric horizonless object can be characterized by multipole moments different from that of the Kerr solution. Apart from trying to solve the information paradox, another motivation for considering alternatives to Kerr black holes is that the Kerr geometry has closed time like curves close to $r = 0$.

Many of the ECOs in figure 8.1 have rotating counterparts, *e.g.*, rotating thin-shell gravastars were constructed by Pani (2015), rotating fuzzball solutions were constructed by Jejjala, Madden, Ross, and Titchener (2005). A rotating version of black shells was constructed in paper VI for small spin (up to second order) where, as expected, the spacetime in the exterior differs from Kerr in having an additional quadrupole moment. The results for a fast spinning black shell have not yet been computed and we can only speculate on what could happen. One reasonable expectation is that it would behave similar to a charged black shell. The charged black shell shrinks down to the size of its horizon as it approaches extremality. If this were also true for the spinning black shell in the limit of extremal spin, it would be expected to reduce down to the size of a Kerr black hole, restoring the horizon and the uniqueness theorem, making the extra quadrupole moment disappear. If this is the case, it would be very difficult to distinguish between fast spinning black shells and Kerr black holes. On the other hand, the radius could remain well outside the horizon and the deviation of the quadrupole moment could persist even at higher spins.

Ergoregion instability

Let us make a few comments about an instability that can effect horizonless spinning objects. It was discovered by Friedman (1978), and recently proved rigorously by Moschidis (2018), that spinning objects with an ergosphere but no horizon, are unstable against scalar, electromagnetic and gravitational perturbations. However, it was shown by Maggio, Pani, and Ferrari (2017) that stability can be ensured if the surface of the object has an absorption rate greater than or equal to 0.4%. Let us see what this means for our black shells. A slowly rotating black shell has a radius bigger than the ergosphere and so there is no ergoregion instability. As we speculated above, if the object approaches close to the horizon size for faster spins, an ergoregion would appear. However, the surface of the shell is highly absorptive due to the large number of degrees of freedom coming from the gas of open strings on it, and the ergoregion instability is expected to be *quenched* à la Maggio, Pani, and Ferrari (2017).

8.5 Observational tests

Electromagnetic and gravitational wave observations can distinguish between black holes and some exotic compact objects. A comprehensive review going into this in detail is by Barack et al. (2019). Here, following the review, we will highlight some of the observational signatures that will be relevant for us.

8.5.1 Electromagnetic signature

Since light trajectories around compact objects can only be traced back to the photosphere, observations in the electromagnetic spectrum cannot be used to distinguish ClePhOs (which have very similar photospheres) from black holes with sufficient accuracy. They are however, good probes for distinguishing UCOs from black holes. When a UCO is illuminated by an accretion disc, it is expected to appear as a dark shadow in the middle of a bright disc, corresponding to the photosphere. For a spinning UCO like the black shell above, the photosphere is no longer spherically symmetric and is expected to have an asymmetry proportional to its spin. For a Kerr black hole, this is of the order a^3. However, the additional quadrupole moment for the black shell contributes an extra asymmetry of order a^2, thus providing an observational test that would distinguish it from a Kerr black hole. There is, in fact, an ongoing cosmological experiment that aims to measure black hole horizons with sufficient precision to be able to detect such a deviation. This is known as the Event Horizon Telescope (EHT), and is a very large baseline interferometery (VLBI) array that images horizon scale structures around the supermassive black object at the center of the Milky Way, and in the elliptic galaxy M87. The first images of the black hole in M87 were recently published by the EHT collaboration (Akiyama et al., 2019). Although the current M87 images do not have the resolution needed to distinguish between these effects, future observations of the black hole at the center of the Milky Way are expected to have enough resolution to do so.

8.5.2 Gravitational waves signature

Unlike light, gravitational waves (GW) are not restricted to the photosphere and can be used to probe the structure of UCOs beyond their photosphere, making them valuable for distinguishing ClePhOs from black holes. For two black objects coalescing into one, if one or both of the objects is different from a black hole, the GW signal can differ in all three stages of the merger: inspiral, merger and ringdown.

(i) In the inspiral phase, the objects are far from each other and can be analysed with perturbation theory. If the merging objects are rotating black shells instead of black holes, their extra quadrupole moment could produce detectable signatures in the gravitational waves emitted during this phase.

(ii) Merger is a highly non-linear event and does not yield itself to a perturbative analysis. It requires full blown numerical simulations, and many techniques for analyzing this phase are under active development.

(iii) Next comes the ringdown phase where the two objects, post merger, are settling down into one. This phase is well described by a set of quasi-normal modes (QNM). Various kinds of ECOs have distinctly different QNMs, and this phase provides one of the best signatures to distinguish between them. The QNMs of an object depend on the reflectivity of its surface as well as its internal structure. Although we have not computed QNMs for the black shells, they are expected to be different from a black hole and the GW signal during ringdown can distinguish one from the other. Another feature of UCOs discovered by Cardoso, Hopper, Macedo, Palenzuela et al. (2016) is presence of *echoes* in the GW ringdown signal. The region between the surface of the UCO and the photosphere can support quasi-bound trapped modes, which contribute these echoes to the GW signal. The detection of such echoes would be indicative of a UCO.

Since the LIGO collaboration studies collisions of stellar mass black holes, it is unlikely that it will have the sensitivity needed to probe deviations from the Kerr quadrupole moment. However, the space based interferometer LISA will be sensitive to lower frequencies. It will therefore be able to measure gravitational waves from the merger of super massive black holes, as well as gravitational waves from stellar mass objects orbiting a supermassive black hole. This can provide a definitive test that will either confirm or rule out the existence of rotating black shells in our universe.

8.6 What next?

To summarize, in paper V, we constructed a novel horizonless object called a black shell. We argued that this is a viable end point of gravitational collapse, and that such an object is well motivated from string theory. We explored the possibility that black holes with horizons may not exist in nature, and that they could all be black shells instead. In paper VI, we constructed slowly spinning black shells, and discovered that they differ from Kerr black holes in their quadrupole moment, therefore providing an observational signature. There are two particularly interesting aspects of black shells that we have not yet explored but would like to, in the future.

First is the extension of our construction to fast spinning black shells. This is particularly relevant for observations because real world black holes are expected to be spinning fast, *e.g.*, the black hole at the center of the Milky Way is estimated to have close to maximal spin. In this chapter, we have speculated on how our results could apply to fast spinning black shells, and it would be very interesting to see which of these speculations holds.

The second aspect, also extremely relevant for observations, is the computation of quasi-normal modes for the black shells. Being dependent on the internal structure of the object, different UCOs have differing, and often very specific quasi-normal mode signatures. This could serve to distinguish them from each other if such a gravitational wave signal is detected in the future. The same is true for gravitational wave echoes, which could have variable amplitudes and frequencies depending on the structure of the ECO. Echoes have been computed for other rotating Kerr-like objects (Bueno, Cano, Goelen, Hertog et al., 2018; Wang and Afshordi, 2018) as well as for specific non-rotating ECOs. See the review by Barack et al. (2019) and the references within, for a summary of such computations.

9. Conclusion

In this thesis, we have presented our results along two main lines of research – dark energy in string theory, and black holes. To tackle the dark energy problem, we have presented two new ideas: one being a new way to obtain de Sitter vacuum from a M-theory compactification, and the other being a novel time dependent construction of de Sitter vacuum, called *shellworlds*. Towards solving the black hole information problem, we have constructed a novel horizonless object called a *black shell*, that could serve as an alternative to a black hole.

In our M-theory compactification, we have constructed a new class of non-supersymmetric Minkowski vacua using only geometric fluxes, and which possess only one flat direction. This flat direction is uplifted to a metastable de Sitter minimum, using higher-derivative corrections to eleven dimensional supergravity, without the need for including any non-perturbative effects, or other exotic ingredients like anti-branes.

The main idea of the *shellworld* model is that, our universe could be sitting on the surface of a bubble that is expanding in a five dimensional anti de Sitter space. This could explain the observed dark energy in our universe. Our model is motivated by string theory, which contains all the right ingredients to construct such a shellworld. We showed that end points of strings stretching out along the extra dimension of the anti de Sitter space represent massive objects on the shellworld, while radiation arises from the presence of mass in five dimensions. We further demonstrated how gravitational attraction between these massive particles on the shellworld arises from interaction between the stretched strings.

Black shells discussed in the second part of the thesis, are proposed as an alternate end point for gravitationally collapsing objects. These are higher dimensional spherical branes coming from string theory that don't have a horizon, but have other properties similar to black holes. We pondered on the possibility that there are no black holes in string theory and every collapsing object forms a black shell instead. We discovered that spinning black shells have a quadrupole moment that is different from a Kerr black hole in general relativity, and proposed observational tests that could distinguish them from black holes.

A natural question to ask is: if our universe really is a shellworld, how does a four dimensional black hole look like? Given that massive particles are endpoints of stretched strings, a black hole formed from a collapsing shell of matter would correspond to the stretched strings converging together. Naïvely, one would expect the strings to coalesce into a black string à la Chamblin, Hawking, and Reall (2000). Four dimensional sections of such a black string corresponds to a Schwarzschild black hole, and so this appears like a plausible way to get four dimensional black holes. However, Gregory and Laflamme (1993) showed that such black strings are unstable, making it difficult to realize a black hole with them. Additionally, they do not correspond to strings of uniform tension along the fifth dimension.

Interestingly, there is another possible fate for the converging strings. They do not need to coalesce into a black string, but can instead tunnel into a brane with cylindrical symmetry, when they get close to each other. This would be like a five dimensional analogue of the black shells. Normally, such a nucleation process would be exponentially suppressed, but in the presence of the stretched strings, the huge entropy of these strings could provide the necessary compensation, just as in the case of fuzzballs and black shells, to form a *black tube*. Four dimensional sections of this object on the shellworld would correspond to the black shells, thus providing a natural uplift of the black shells to the shellworld picture. This is a very intriguing possibility and is something that we hope to investigate further in a future work.

To conclude, the ideas presented here open up very exciting avenues for further investigation and we hope to explore them in our future research. It would also be very interesting to see if these ideas survive observational tests and the test of time.

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Svensk sammanfattning

Det har gått nästan två årtionden sedan man upptäckte att vårt universum domineras av en mörk energi som får det att expandera allt snabbare. Den överraskande upptäckten ändrade helt vår uppfattning om universum och belönades med 2011 års Nobelpris. Fysiker har försökt att förstå den mörka energins ursprung inom ramen för gängse teorier. Man har kopplat den mörka energin till den kosmologiska konstanten i Einsteins fältekvationer för den allmänna relativitetsteorin. Den kosmologiska konstanten är enligt experimentella mätningar numeriskt liten och positiv.

En förklaring till mörk energi kräver inte bara fysik på en galaktisk skala, utan också fysik på skalor mycket mindre än atomens storlek. Detta kräver en teori som inkluderar den allmänna relativitetsteorin, som beskriver objekt på stora skalor; samt också kvantmekaniken som beskriver små objekt på små skalor. En sådan teori skulle kunna förklara fenomen på alla skalor och är således en *teori för allting*. Vi har hittills inte lyckats hitta en sådan teori.

Strängteorin är dock den bästa kandidaten för en sådan fundamental teori. Strängteorin antar att universum är uppbyggt av små strängar (kanske en miljard miljard gånger mindre än protonen) och att partiklar är vibrationer hos sådana strängar. Teorin kräver tio rumtidsdimensioner (vilket kan jämföras med vårt universums fyra rumtidsdimensioner). Strängteorin har också en elvadimensionell kusin känd som M-teorin, vilket är en teori bestående av membran istället för strängar. Strängteorin innehåller dessutom tunga högredimensionella membran som är relaterade till membranen i M-teorin. Eftersom extra rumsdimensioner inte har observerats experimentellt måste dessa sex extra dimensioner i strängteorin (eller sju i M-teorin) vara extremt små och kompakta. Det finns olika sätt att kompaktifiera dimensionerna, vilket ger upphov till distinkta teorier i våra fyra rumtidsdimensioner. Till följd av de många olika sätten att kompaktifiera de extra dimensionerna, finns det ett enormt antal av möjliga fyrdimensionella teorier som strängteorin kan ge upphov till. Alla dessa fyrdimensionella teorier beskriver ett distinkt universum. Svårigheten är att hitta vårt universum, med rätt mängd mörk energi. Denna utmaning har sysselsatt fysiker de senaste två årtiondena.

Trots de otaliga möjliga universumen så har detta visat sig svårare än väntat. Att hitta en tillfredsställande modell av mörk energi är fortfarande ett öppet problem inom strängteorin. Till följd av många oväntade hinder har vissa forskare även gått så långt som att föreslå att det kan vara omöjligt att konstruera en positiv kosmologisk konstant inom strängteorin.

Strängteorins matematik kan bara hanteras om man gör vissa approximationer. Utmaningen med att konstruera en positiv kosmologisk konstant är att det ofta krävs ingredienser där dessa approximationer inte längre är giltiga. Många av de konstruktioner som föreslagits är av denna typ. I denna doktorsavhandlingen presenterar jag en konstruktion vid lågenergigränsen av M-teorin och formulerar en modell för mörk energi där approximationerna är under kontroll. Detta är ett lovande nytt resultat som banar vägen för ytterligare konstruktioner som inte bara är giltiga vid lågenergigränsen av M-teorin, utan vid alla energier.

Ett annat spännande resultat som jag presenterar i denna avhandling är en helt ny konstruktion av mörk energi inom strängteorin. Jag antar att vårt universum är inbäddat i ytan av en sfärisk bubbla som expanderar i fem dimensioner. En sådan bubbla skulle bestå av högredimensionella membran från strängteorin, där den femte rumsdimensionen är inte nödvändigtvis liten. Detta är en radikal ny modell som skulle kunna förklara mörk energi. Materiepartiklar svarar mot strängar som sträcker sig ut från bubblan i den femte dimensionen.

Ett annat öppet problem inom fysiken rör svarta hål. Dessa svarar mot kroppar vars gravitationskraft är så stor att inte ens ljuset kan ta sig därifrån. Svarta hål omsluts av en händelsehorisont som utgör en enkelriktad väg för ljus och materia. Närvaron av en händelsehorisont skapar en konflikt mellan två grundläggande fysikaliska principer: den ena ifrån den allmänna relativitetsteorin och den andra ifrån kvantmekaniken. Antag att en nationalencyklopedi kastas in i ett svart hål. Stephen Hawking – en av de mest kända fysikerna någonsin – upptäckte att svarta hål inte kan existera för evigt. Istället dunstar svarta hål så småningom bort, vilket gör att också informationen som ramlar in försvinner och aldrig kommer tillbaka. Vid första anblick förefaller detta inte vara ett så stort problem. Borde inte information försvinna om du eldar upp en bok? Även om så verkar vara fallet, är det faktiskt möjligt att samla in all aska, rök och ljus från elden och återskapa informationen i boken. Det må kräva avancerad utrustning och hårt arbete, men det är i princip möjligt. Problemet med svarta hål är att det inte är möjligt, ens i princip, att återskapa informationen. En sådan informationsförlust går emot kvantmekanikens fundamentala principer. Detta kallas informationsparadoxen. Troligen krävs en ny fundamental teori, som strängteorin, för att man skall kunna lösa problemet.

I avhandlingen presenterar jag ett nytt försök att lösa gåtan, där jag använder strängteorin för att konstruera en nytt slags objekt med egenskaper som liknar de hos svarta hål men utan en händelsehorisont. Ett sådant objekt består av strängteoretiska sfäriska membran i form av ett svart skal och har en härva av vibrerande strängar på ytan. Allt som dras till det svarta skalet träffar till slut ytan och blir till en del av stränghärvan. Om svarta hål i själva verket är svarta skal skulle informationsparadoxen lösas upp eftersom ingen händelsehorisont existerar. I avhandlingen konstruerar jag också roterande versioner av de svarta skalen som efterliknar roterande svarta hål. Fastän de svarta skalens egenskaper liknar svarta håls, så finns det viktiga skillnader. Dessa skillnader skulle kunna göra det möjligt att skilja svarta skal från svarta hål genom astronomiska observationer. Förhoppningen är att man kommer att nå den nödvändiga precisionen inom en nära framtid.

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