Singular Fluxes in Ten and Eleven Dimensions

Sources, Singularities, Fluxes and Spam

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Abstract

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The energy content of our present universe is dominated by the dark-energy, or vacuum energy, which provides accelerated cosmic expansion. Dark energy has a possible effective explanation through a positive cosmological constant. The problem present in any fundamental theory is to explain the underlying dynamics of what gives rise to the cosmological constant.

In string theory there are several scenarios that could give insight into what is behind the positive cosmological constant. One such construction uses anti-branes to achieve a net positive energy density of the vacuum. Anti-branes refers in this case to branes placed in a background with oppositely charged flux. As backreaction and localisation procedures are considered for anti-brane constructions a certain kind of singularity arise. This new type of singularity is present in the surrounding flux, which is not directly sourced by the brane.

This thesis, and the works contained, considers several aspects of this type of singularity. The first such flux singularity were discovered for the anti-D3-branes, in which the approximations and assumptions of partial smearing and perturbative expansions are used. Included in this thesis are new anti-D6-brane solutions which are placed in oppositely charged flux. It is shown that after the anti-D6-branes are localised, they display the same type of singularity. The strength of this result lies in that it is possible to show the presence of the singularity beyond partial smearing and perturbative expansions. Similar to the anti-D6-brane solutions, new anti-M2-brane solutions are presented. These solutions are also argued to display the same type of singularity.

The investigation into the presence of the singularity is just the first step. The second step is to deduce whether this singularity is acceptable and can somehow be resolved. Included in this thesis are two works that considers exactly this. One way of interpreting the singularity is through the absence of a no-force condition between the brane and the surrounding flux. This interpretation leads to the conclusion that the singularity is present due to the use of static Ansätze in a system that is inherently time dependent. Through an adiabatic approach it is here argued that this interpretation leads to a new type of instability.

Another way of arguing for a possible resolution of this singularity is whether or not the singularity can be cloaked by an event horizon. This condition have been successful in other systems with singularities. It is argued in this thesis that it is not possible to hide the flux singularity behind a horizon. This leads to one out of two conclusions, either the condition is not a necessary one and the singularity can be resolved in a static manner, or the singularity does not have a resolution.

To put these works in context the current singularities from anti-branes program is briefly reviewed to give a full overview of the current situation of these investigations.

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To my Mother

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 H. Danielsson, Daniel Junghans, Thomas Van Riet, Timm Wrase, and Marco Zagermann. *Smeared versus localised sources in flux compactifications. JHEP*, 1012:043, 2010, [arXiv:1009.1877]

II   Johan Blåbäck, Ulf H. Danielsson, Daniel Junghans, Thomas Van Riet, Timm Wrase, and Marco Zagermann. *The problematic backreaction of SUSY-breaking branes. JHEP*, 1108:105, 2011, [arXiv:1105.4879]

III   Johan Blåbäck, Ulf H. Danielsson, Daniel Junghans, Thomas Van Riet, Timm Wrase, and Marco Zagermann. *(Anti-)Brane backreaction beyond perturbation theory. JHEP*, 1202:025, 2012, [arXiv:1111.2605]

IV   Johan Blåbäck, Ulf H. Danielsson, and Thomas Van Riet. *Resolving anti-brane singularities through time-dependence. JHEP*, 1302:061, 2013, [arXiv:1202.1132]

V   Iosif Bena, Johan Blåbäck, Ulf H. Danielsson, and Thomas Van Riet. *Antibranes cannot become black. Phys.Rev.*, D87(10):104023, 2013, [arXiv:1301.7071]

VI   Johan Blåbäck. *Note on M2-branes in opposite charge. Phys.Rev.*, D89(6):065004, 2014, [arXiv:1309.2640]

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The following papers were produced during the Ph.D. studies and will not be covered in the text and should hence not be considered a part of the thesis.

• Johan Blåbäck, Ulf H. Danielsson, and Thomas Van Riet. *Lifshitz Backgrounds from 10d Supergravity. JHEP*, 1002:095, 2010, [arXiv:1001.4945]

- Johan Blåbäck, B. Janssen, T. Van Riet, and B. Vercnocke. *Fractional branes, warped compactifications and backreacted orientifold planes. JHEP*, 1210:139, 2012, [arXiv:1207.0814]
- Johan Blåbäck, Ulf Danielsson, and Giuseppe Dibitetto. *Fully stable dS vacua from generalised fluxes. JHEP*, 1308:054, 2013, [arXiv:1301.7073]
- J. Blåbäck, A. Borghese, and S.S. Haque. *Power-law cosmologies in minimal and maximal gauged supergravity. JHEP*, 1306:107, 2013, [arXiv:1303.3258]
- Johan Blåbäck, Ulf Danielsson, and Giuseppe Dibitetto. *Accelerated Universes from type IIA Compactifications. JCAP*, 1403:003, 2014, [arXiv:1310.8300]
- Johan Blåbäck, Diederik Roest, and Ivonne Zavala. *De Sitter Vacua from Non-perturbative Flux Compactifications. Submitted to PRL*, 2013, [arXiv:1312.5328]
- Johan Blåbäck, Bert Janssen, Thomas Van Riet, and Bert Vercnocke. *BPS domain walls from backreacted orientifolds. Submitted to JHEP*, 2013, [arXiv:1312.6125]

Contents

1. Introduction

String theory is a candidate *theory of everything*. A theory of *everything* might seem a bit too overwhelmingly ambitious, and also is. The goal is not as grand, or mad even, as trying to render all other fields of physics obsolete, but instead to join the four fundamental forces of nature under one description. Even saying *four* might not be a fair assessment of the problem as three of them are already unified in a common framework. So to better describe it, the goal is to unify the three forces of our microscopic universe with the one force of our macroscopic universe.

The microscopic forces are unified in a framework called *Quantum field theory* (QFT for short), this together with local (gauge) symmetries, i.e. a *gauge theory*, describe our three microscopic forces to high precision. This gauge theory is called *the Standard Model* and unifies the *Strong*, *Electromagnetic* and *Weak* forces. With the recent discovery of the so-called *Higgs particle*, awarded with the 2013 Nobel prize in physics [NFb], the standard model is more or less complete and obvious hints for physics beyond the standard model seems to be currently absent.

The theory that describes the macroscopic force of *Gravity* is *General relativity*. This is also a theory that have been put to the test and stood tall every time; the precession of elliptic planet orbits, deflection of light from massive objects, gravitational red-shift of light, and on an even greater scale, the ΛCDM-model describing the evolution of our universe. Related to this, the discovery of the accelerated cosmic expansion was rewarded with the 2011 Nobel prize in physics [NFa].

While the standard model is a quantum theory, able to describe the interactions of individual particles, no quantum description of general relativity is easily available. Some more fundamental revision of the underlying theory seems necessary. This is where string theory comes in. One of the great achievements of string theory is that it can give rise to gauge theories as well as a quantum description of gravity. Although several "kinks" still remain to be worked out. String theory comes with the concept of *supersymmetry* that mixes force mediating particles with matter particles, and specifies the number of space-time dimensions to be ten.

At an initial reflection over string theory one might conclude that, since it contains supersymmetry and extra dimensions, it might not be a realistic theory. However, one can instead argue that this could possibly add predictability to the theory, i.e. it *predicts* supersymmetry and

extra dimensions which remain to be observed experimentally. This would also mean that realistic string theory scenarios involve a description of our universe where supersymmetry and extra dimensions are effectively invisible to our current experiments. This can be achieved through breaking of supersymmetry to some energy level above current experiments as well as compactifying the extra dimensions to such a degree that they are currently undetectable.

While working with string theory itself is quite cumbersome, there exists classical low energy descriptions of string theory known as *supergravity* theories. There are several supergravity theories and only some of them are described by low energy string theory. Two of these, called type IIA and type IIB supergravity, are ten dimensional supergravity theories that will be used in this thesis.

The hope is, even though supergravity does not include all features of string theory, that it will be enough to describe realistic compactification scenarios with supersymmetry breaking. Another necessity is that a realistic supergravity solution which would describe our universe gives rise to accelerated cosmic expansion. Current observations are compatible with that the present expansion of our universe is driven by a positive *cosmological constant* Λ

$$S = \frac{1}{2\kappa_4^2} \int \star_4 \left(R^{(4)} - \Lambda \right) ; \quad \Lambda > 0 . \tag{1.1}$$

A maximally symmetric space-time, i.e. a space-time that is homogeneous and isotropic, with positive cosmological constant is called de Sitter (dS). In effective theories the cosmological constant is a parameter included by hand, while in string theory it should be possible to describe the underlying dynamics of this parameter and hence give a more fundamental understanding of the cosmological evolution of our universe.

In principle this is how a compactification would work. Consider for example the ten dimensional action of any of the type II supergravities with an Einstein-Hilbert term and other string theory related fields $\mathcal{L}_S$ and reduce this to a four dimensional action

$$\begin{aligned} S &= \frac{1}{2\kappa_{10}^2} \int \star_{10} \left(R^{(10)} + \mathcal{L}_S \right) \\ &= \frac{1}{2\kappa_4^2} \int \star_4 \left(R^{(4)} - V \right) , \end{aligned} \tag{1.2}$$

where

$$\frac{1}{2\kappa_4^2} = \frac{1}{2\kappa_{10}^2} \int \star_6 1 , \quad \text{and} \quad V = -\left(\frac{\kappa_4}{\kappa_{10}} \right)^2 \int \star_6 (R^{(6)} + \mathcal{L}_S) . \tag{1.3}$$

The potential energy V would correspond to the cosmological constant, and our universe would reside on a slightly positive, (meta-)stable, extremal point[1] of this potential. However achieving these properties of the potential turns out to be a difficult task – the generic properties of V are such that extremal points are negative and stable or positive and unstable.

There are several more or less realistic supergravity solutions present in the literature. One of the more appealing is described by [KKLT03] which is a scenario where supersymmetry is broken, an effective four-dimensional space-time arise and where the cosmological constant can be made arbitrarily small and positive, possibly compatible with observations. This description does however leave some questions unanswered needing a more detailed analysis. The aim of the articles included in this thesis is to challenge some of these questions, directly or indirectly.

More concretely, [KKLT03] describes how a supersymmetry breaking and a positive cosmological constant can be achieved by adding so-called anti-branes to a certain type of background. This is done under some seemingly reasonable approximations. One such assumption is that the anti-branes do not influence the background too much if they are few enough. This is however something that needs to be calculated in detail to see if the approximations are as reasonable as they seem.

Recently it was realised that when the anti-branes influence on the background were taken into account, i.e. their *backreaction* were calculated, they produce a previously unseen singularity. This singularity arises in the energy density of surrounding *fluxes*. These fluxes are fields that string theory introduces and are a seemingly necessary ingredients for realistic supergravity solutions.

With the discovery of this singularity, several new questions needs to be answered. Does this singularity arise in other systems? Is it possible to resolve this singularity, i.e. is it possible make the energy densities finite? Why does the singularity arise? All of these questions will, if not answered definitely, at least be considered here. The articles included in this thesis all relates to these questions. To be able to describe the included articles relation to parallel developments present in the literature a brief review of this field will be presented together with a summary of the articles.

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[1]Stability meaning that the Hessian of V with respect to the various fields that it depends on is positive definite, and extremal meaning that the first derivatives are zero. Meta-stability means that there could be extremal points of lower energies into which the elevated solution could tunnel to through quantum effects.

1.1 Outline of the thesis

The thesis is organised in the following way.

Part I includes a rough sketch of what string theory is and where the type II ten-dimensional and the 11D supergravity pictures come from. This part will also introduce the Bianchi identities and equations of motions that will be used. It will also define some terminology that will be important for the later sections and the articles included. The most relevant portion of Part I is Section 2.2.2 where several important articles will be briefly reviewed. These articles will play a very important role in Part II.

Part II is divided in the following way. In Chapter 3 there will be an introduction of some of the solutions presented in Paper I and the motivation for their study. Chapter 4 is the main chapter of this thesis and summarises the result of all articles included. In addition, said chapter also includes a brief review of articles by other collaborations that also considers the issues related to the aforementioned flux singularity. The idea of this layout is to put the works included in this thesis in a broader context and emphasise their importance.

Part I:
Background material

2. Flux compactifications

The purpose of this chapter is to introduce some relevant background. Most of the work included in this thesis is performed in the so-called type II ten dimensional supergravity theories, which originate from string theory. The exception is one paper that considers the eleven dimensional supergravity theory, which has an M-theory origin. Given here is hence the most basic parts of string and M-theory including a sketch of what they are, how they are related, and what their field content is.

Later in this chapter some well known results will be presented. The text in Part II will heavily rely on these results. Therefore they are included here, not with the purpose of presenting a complete review of them, but rather serving as background material so they can be discussed more briefly when needed.

2.1 String theory

The most simple summary of the idea of string theory is to extend the point particle to a string to see what happens [Mun]. In this section the aim is to introduce the string and its world-sheet action and also to sketch the approach of finding the corresponding space-time action. The equations of motion derived from the space-time action are the set of equations that will be used through out the thesis. A lot of details are left aside but the statements made and equations presented can be found in any of the standard books on string theory: [Pol98a, Pol98b, GSW87, BBS07]. The information concerning M-theory is mostly taken from [Tow96].

The action governing the motion of a point particle in a curved background would be determined by

$$S = -m \int \sqrt{-\mathrm{d}s^2} \,, \tag{2.1}$$

where the line-element squared is given by

$$\mathrm{d}s^2 = g_{ab}\mathrm{d}X^a\mathrm{d}X^b \,. \tag{2.2}$$

This particle is embedded in the so-called *target-space*, i.e. the space-time in which it moves. The target-space has a curvature given by the

metric $g_{\mu\nu}$ with signature $(-, +, \ldots, +)$, and X^a describes the embedding. There is an equivalent action that does not contain a square root

$$S = -\frac{1}{2} \int d\tau \left(\sqrt{-h}\, h^{-1} \partial_\tau X^a \partial_\tau X^b g_{ab} - \sqrt{h} m^2 \right), \qquad (2.3)$$

where h is the one dimensional metric on the *world-line*, i.e. the line that represents the point particles path through space-time, parametrised by the proper time τ. That is, the point particles path is traced out with the parameter τ and it is embedded in a space with D dimensions, i.e. $a, b = 0, \ldots, D - 1$.

This action for the point particle can be generalised to a string

$$S_0 = -\frac{1}{4\pi\alpha'} \int d\tau d\sigma \sqrt{-h}\, h^{\alpha\beta} \partial_\alpha X^a \partial_\beta X^b g_{ab}, \qquad (2.4)$$

with a tension $T = 1/(2\pi\alpha')$. Here $h_{\alpha\beta}$ is the metric on the two-dimensional world-*sheet*. This is called the Polyakov action, although usually attributed to [BDVH76] and [DZ76]. The indices $\alpha, \ldots = \tau, \sigma$ are the world-sheet indices.

In Maxwell theory the point particle can be charged, which extends the action with the following term

$$-Q \int A_\mu dX^\mu. \qquad (2.5)$$

In a similar way one can add charges to the string. The string is charged under the antisymmetric Kalb-Ramond [KR74] field B_{ab}

$$S_B = -\frac{1}{4\pi\alpha'} \int d\tau d\sigma \sqrt{-h}\, i\epsilon^{\alpha\beta} \partial_\alpha X^a \partial_\beta X^b B_{ab}. \qquad (2.6)$$

The string could also be coupled to the gravitational interactions that takes place on the world-sheet, which is done by adding an Einstein-Hilbert term. It is also possible to couple this term to a scalar field that is called the *dilaton*, here represented by Φ. This part of the action would be

$$S_d = \frac{1}{4\pi\alpha'} \int d\tau d\sigma \sqrt{h}\, \alpha' R^{(2)} \Phi. \qquad (2.7)$$

Adding these actions together gives the world-sheet action for the *Bosonic string theory*

$$S = S_0 + S_B + S_d. \qquad (2.8)$$

What was the purpose of this? Well, the above presented action terms give a string world-sheet action, with charge and with curvature, for a string with only bosonic degrees of freedom. That is, they represent the dynamics on the world-sheet of a string. What is important for the later

parts of this thesis is the space-time action. How this action is derived is what will be outlined here.

This world-sheet action might look good, but it does not obey all the symmetries that one would expect. One symmetry that must be imposed is that the world-sheet theory is conformal. This means that the world-sheet metric can be written as

$$h_{\alpha\beta} = e^{\psi}\eta_{\alpha\beta}\,, \qquad (2.9)$$

where the $\eta_{\alpha\beta}$ is the flat metric. In other words, the world-sheet metric can be rescaled to a conformally flat metric. By adding quantum corrections to the presented actions, which is needed to match orders of the coupling constant α', it can be shown that conformal symmetry is governed by the following equations, to the lowest order in α',

$$
\begin{aligned}
0 &= R_{ab} - \frac{1}{2}|H_3^2|_{ab} + 2\nabla_a\partial_b\Phi\,, \\
0 &= \nabla_c H^c{}_{ab} - 2(\nabla_c\Phi)H^c{}_{ab}\,, \\
0 &= -\frac{1}{2}\nabla^2\Phi + |\partial\Phi|^2 - \frac{1}{4}|H_3|^2\,,
\end{aligned}
\qquad (2.10)
$$

where $H_3 = dB_2$ and the square-rules will be introduced later, see equation (2.29). These equations are not only possible to derive from the above world-sheet action, but also through the Euler-Lagrange method from a D-dimensional action

$$S_D = \frac{1}{2\kappa_0^2}\int d^D x\sqrt{-g}\,e^{-2\Phi}\left(R + 4\partial_a\Phi\partial^a\Phi - \frac{1}{2}|H|^2\right)\,. \qquad (2.11)$$

This is the space-time formulation of the bosonic string theory. These equations and the above action are only valid for $D = 26$ which is called *the critical dimension*, and is demanded for consistency – other dimensions gives rise to negative norm states or unitarity problems. This action is written in the so-called *String frame* where the action have a coupling between the dilaton and the Ricci scalar.

The bosonic string theory suffers from some problems, especially for phenomenological reasons. One of these problems is that there exists no fermions in this theory. To introduce fermions one reconsiders S_0 with an additional fermionic part ($T = 1$)

$$S_S = -\frac{1}{2}\int d\tau d\sigma\eta^{\alpha\beta}\left(\partial_\alpha X^a\partial_\beta X^b - i\bar{\psi}^a\rho^\alpha\partial_\alpha\psi^b\right)\eta_{ab}\,. \qquad (2.12)$$

Here the metrics are switched to a flat world-sheet ($\eta_{\alpha\beta}$) and a flat space-time (η_{ab}). There is a symmetry of this action that relates these two terms

$$
\begin{aligned}
\delta X^a &= \bar{\epsilon}\psi^a\,, \\
\delta\psi^a &= -\rho^\alpha\partial_\alpha X^a\epsilon\,.
\end{aligned}
\qquad (2.13)
$$

This symmetry is known as *supersymmetry*, which in simple terms is the symmetry that relates bosonic and fermionic fields.

What was done above for the bosonic string is that the available massless fields were entered into the string action. Then, to obey conformal symmetry for this action, the space-time actions were derived. For the *superstring*, i.e. the string with world-sheet supersymmetry, other approaches are used. What will be sketched now is how the massless modes for closed strings are identified. The work included in this thesis is mostly concerned with closed string dynamics, and therefore the emphasis will be put on the derivation of the closed string sector field content. When the massless fields are identified, space-time supersymmetry uniquely specifies the space-time actions. Since this is the only goal with this section, details concerning the quantisation of the string will be ignored. The important lesson to draw from quantisation of the string, other than that the string is possible to quantise, is that this forces the space-time dimensions to be ten dimensional. In other words, the critical dimension of the superstring is $D = 10$.

The closed strings have to obey certain boundary/periodicity conditions. For the bosonic fields these would be

$$X^a(\tau, \sigma) = X^a(\tau, \sigma + \pi), \tag{2.14}$$

while for the fermionic fields there are two possible conditions

$$\psi_\pm^a(\tau, 0) = \psi_\pm^a(\tau, \pi), \\ \psi_\pm^a(\tau, 0) = -\psi_\pm^a(\tau, \pi). \tag{2.15}$$

The first boundary condition is known as *Ramond* (R) and the second is called *Neveu-Schwarz* (NS). Here the sign labels $\genfrac{}{}{0pt}{}{+}{(-)}$ refers to left (right) movers, i.e.

$$\partial_{\genfrac{}{}{0pt}{}{+}{(-)}} \psi_{\genfrac{}{}{0pt}{}{-}{(+)}} = 0, \text{ where } \psi^a = \begin{pmatrix} \psi_-^a \\ \psi_+^a \end{pmatrix}, \text{ and } \genfrac{}{}{0pt}{}{+}{(-)} = \tau \genfrac{}{}{0pt}{}{+}{(-)} \sigma. \tag{2.16}$$

For the left and right movers on the world sheet there is hence a pairing of these boundary conditions. There are four different possible pairings of two types – same and mixed

$$(R, R), \ (NS, NS), \\ (R, NS), \ (NS, R). \tag{2.17}$$

The first type gives rise to space-time bosons, while the second type will give rise to space-time fermions. It is however not enough to only consider these parings of boundary conditions since they generally give rise to tachyons. This can be avoided using the so-called [GSO77] projection. This associates a parity with the R and NS sectors, R± and

NS±, which truncates the string spectrum such that these tachyons are no longer present. This parity are related to the chirality of the spinors. There are four ways of doing this, in order to preserve maximum supersymmetry, where only two are independent. These are

$$\text{IIA: (R+, R-), (R+, NS+), (NS+, R-), (NS+, NS+)}, \\ \text{IIB: (R+, R+), (R+, NS+), (NS+, R+), (NS+, NS+)}. \tag{2.18}$$

In type IIA where both chiralities are present, the theory is non-chiral, while type IIB is chiral. These two have a common bosonic section of space-time fields (NS+, NS+), which consists of a dilaton Φ, an antisymmetric tensor B_{ab} and the metric g_{ab}. The fermionic sectors contain a dilatino λ, which is a spin-1/2 field, and a gravitino Ψ^a, which is a spin-3/2 field. The (R,R) sector contains different ranked anti-symmetric tensors (differential forms)

$$\text{IIA: } C_1, \ C_3, \\ \text{IIB: } C_0, \ C_2, \ C_4. \tag{2.19}$$

These will act as potentials for field-strengths $\hat{F}_q = \mathrm{d}C_{q-1}$.

There is an extension of type IIA where a field-strength without potential can be added, F_0. This is called the Romans mass parameter [Rom86]. This theory have some peculiarities, while type IIA can be extended to M-theory, for *massive* type IIA no such uplift exists.

The action, or equivalently the equations of motion (which will be presented later), derived in this way are part of a perturbative expansion

$$(\text{E.o.m.})\alpha' + \mathcal{O}(\alpha'^2). \tag{2.20}$$

This means that any solution of these equations must obey $\alpha' \ll 1$.[1] There is also an expansion in terms of string loops. A string tree-level diagram would consist of a surface that is topologically a two-sphere. By adding what is called *vertex operators*, loops can be added to the tree-level. The expansion of the string action with a vertex operator is an expansion in $g_s = \mathrm{e}^{\langle\Phi\rangle}$, so to be sure to keep string loop corrections under control also $g_s \ll 1$ is necessary. The tree-level and leading order in α' expansion that will be used here is known as 10D *supergravity*. Supergravity is the theory which will be used throughout the rest of this thesis.

Before moving on some notes should be made about *M-theory*. M-theory is an eleven-dimensional theory with a low energy description, in the sense of $\alpha' \ll 1$, that is 11D supergravity. This theory is the

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[1]The characteristic curvature radius, call it $\mathcal{R}$, also influence this expansion. Effectively the expansion parameter is $\alpha'/\mathcal{R}^2$.

strongly coupled regime of string theory where $g_s \to \infty$. The radius of the eleventh dimension is hidden from string theory perturbation theory and is related to g_s as

$$R_{11} = g_s^{2/3} \,. \tag{2.21}$$

In the weak coupling regime $g_s \ll 1$ this is invisible while at $g_s \to \infty$ the eleventh dimension plays an equally important role as the other ten [Tow96].

M-theory does not contain strings and should not be referred to as a string theory. Instead, on the same basis as string theory has D-branes, that will be covered shortly, M-theory has M2- and M5-branes. To the content of the 11D supergravity belongs also a metric, a four-form G_4 and a gravitino Ψ_a. The bosonic sector of the 11D supergravity will be introduced later together with the type II supergravities.

2.1.1 String and M-theory sources

The two types of sources that will be most relevant to what is discussed in this thesis are the Dirichlet branes (D-branes) of a certain spatial dimension p (Dp-branes), and the Orientifold planes (O-planes/Op-planes). Although the use of these sources in this thesis will be only distinguishable up to a sign change, their dynamics and properties are very different. D-branes are objects which have dynamics on their world-volume in terms of open strings that end on the brane. The O-planes are instead non-dynamical objects, usually defined to be the set of fixed points of an orientifold involution that it creates on the space in which it is embedded.

The Dp-branes and Op-planes influence the ten dimensional action through their Dirac-Born-Infeld (DBI) action and the Wess-Zumino (WZ) action[2]

$$S_{\text{DBI}} = -T_{\text{D}p} \int d^{p+1}\xi \sqrt{|g^{(p+1)}|} \, e^{(p-3)\phi/4} \,,$$
$$S_{\text{WZ}} = Q_{\text{D}p} \int C_{p+1} \,, \tag{2.22}$$

$T_{\text{D}p}$ is the tension and couples to the Einstein equation and the equation for the dilaton. Whereas $Q_{\text{D}p}$ is the charge and couples its corresponding field-strength, F_{8-p}, through the Bianchi identity. The g refers to the determinant of the space-time metric, pulled back onto the world-volume, and its superscript refers to the dimension. Written above

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[2]The sign convention here is taken from [Koe11] since the same convention (up to a sign change of the NSNS potential) will be used for the equations of motion that will be introduced later.

are the actions for a Dp-brane. To get the Op-plane actions simply change the sign of the tension and charge.[3] As mentioned earlier, the D-branes have dynamic properties on its world-volume, the actions presented above are given in the lowest order expansion of this dynamics, i.e. the world-volume gauge fields have been neglected here.

To be able to write a complete ten dimensional action, the $(p+1)$ dimensional world-volume actions presented above needs to be extended in the transversal space. Both actions are extended by adding a trivial integral over the space transversal to the world-volume, that is

$$1 = \int \delta(\mathrm{D}p) \star_{9-p} 1 = \int \delta_{9-p}(\mathrm{D}p) \,. \qquad (2.23)$$

One important point that should be made here, which will be important later, is that equation (2.23) shows that the sources used are *localised*. The action written terms in (2.22) are localised to the world-volume and the delta function introduced above, $\delta(\mathrm{D}p)$, represents the position of the source (or a sum of sources) in the transversal directions.

The Neveu-Schwarz five brane (NS5-brane) will also be relevant for this thesis. While the D-branes couple to the RR-sector fields though its WZ-action, the NS5-brane couples to the NSNS-sector potential B_6. More specifically, the action of the NS5-brane is

$$S_{\mathrm{DBI}} = -\frac{\tau_{\mathrm{NS5}}}{g_s^2} \int \mathrm{d}^{p+1}\xi \sqrt{\left|\det\left(g_{\mu\nu}^{(6)} + 2\pi g_s \mathcal{F}_{\mu\nu}\right)\right|} \,,$$
$$S_{\mathrm{WZ}} = \mu_{\mathrm{NS5}} \int B_6 \,. \qquad (2.24)$$

$\mathcal{F}_{\mu\nu}$ is an object describing the world-volume fluxes. The relation between the two potentials B_6 and B_2 is given by

$$\mathrm{d}B_6 = \frac{1}{g_s^2} \star_{10} \mathrm{d}B_2 \,. \qquad (2.25)$$

The NS5-brane is hence the (electromagnetic) dual brane to the fundamental string and is therefore present in both type IIA and type IIB.

In M-theory the fundamental object is the M2-brane (also called (super)membrane). When the eleventh dimension of M-theory is wrapped by one direction of this brane it reduces to the type IIA string, and when no directions of the brane wraps the eleventh dimension it reduces to the type IIA D2-brane. The M2-brane has a dual brane called the M5-brane.

---

[3]There is a quantisation condition for the D-brane and O-plane charge that relates to the D-brane charge with a certain magnitude factor. However these details, that relates to quantisation, will not be relevant to anything presented in this thesis and will not be discussed further.

2.2 Type II and 11D supergravity compactifications

So far a short introduction of string theory has been presented and a sketch of the origin of the type II and 11D supergravities. In the present section a summary of all equations of motion will be presented so that they can be easily referred to in the text avoiding unnecessary repetition.

Type II conventions

Let us quickly summarise the constituents of the type II supergravities and their relations. The type II supergravities consists of the following Ramond-Ramond (RR) form fields

$$\text{RR:} \quad F_q \,, \tag{2.26}$$

where q is even (odd) for type IIA(B). They have a common Neveu-Schwarz (NSNS) sector consisting of

$$\begin{aligned} \text{metric:} \quad & g_{ab} \,, \\ \text{dilaton:} \quad & \phi \,, \\ \text{NSNS 3-form:} \quad & H_3 \,. \end{aligned} \tag{2.27}$$

The equations of motion that govern their interactions, in the ten dimensional Einstein frame[4]. These are the trace reversed Einstein equation

$$\begin{aligned} R_{ab} = {}& \frac{1}{2}|\mathrm{d}\phi|_{ab}^2 + \frac{1}{2}\mathrm{e}^{-\phi}\left(|H_3|_{ab}^2 - \frac{1}{4}g_{ab}|H_3|^2\right) \\ & + \sum_{q \leq 5} \frac{1}{2(1+\delta_{q5})}\mathrm{e}^{\frac{5-q}{2}\phi}\left(|F_q|_{ab}^2 - \frac{q-1}{8}g_{ab}|F_q|^2\right) \\ & + \frac{1}{2}\left(T_{ab}^{\ell} - \frac{1}{8}g_{ab}T^{\ell}\right) \,. \end{aligned} \tag{2.28}$$

The Kronecker delta have been introduced in this expression to account for the self-duality of the F_5; $F_5 = (1 + \star_{10})\mathcal{F}_5$. The squares are taken according to the following rule

$$|A_q|_{ab}^2 = \frac{1}{(q-1)!}A_{a\,a_2,\ldots a_q}A_b{}^{a_2,\ldots a_q}, \quad |A_q|^2 = \frac{1}{q!}A_{a_1,\ldots a_q}A^{a_1,\ldots a_q}\,, \tag{2.29}$$

for a q-form, where contractions have been done with the inverse metric g^{ab}. The part of the stress tensor representing the localised sources is given by

$$T_{ab}^{\ell} = -\mathrm{e}^{\frac{p-3}{4}\phi}T_{\mathrm{D}p}g_{\mu\nu}\delta_{ab}^{\mu\nu}\delta(\mathrm{D}p)\,. \tag{2.30}$$

---

[4]In previous sections equations have been written in the string frame. From now on a ten dimensional Einstein frame will be used where the relation between the two frames are given by $g_{ab}^{\mathrm{E}} = \mathrm{e}^{-\phi/2}g_{ab}^{\mathrm{S}}$.

Ten dimensional indices are labelled $a, b, \ldots = 0, 1, \ldots 9$ and world-volume indices $\mu, \nu, \ldots = 0, 1, \ldots p$, while the transversal directions are labelled with $i, j, \ldots = (p+1), \ldots 9$, unless otherwise specified for each section.

The dilaton equation of motion is given by

$$\nabla^2 \phi = -\frac{1}{2} e^{-\phi} |H_3|^2 + \sum_{q \leq 5} \frac{5-q}{4} e^{\frac{5-q}{2}\phi} |F_q|^2 - \frac{p-3}{4(p+1)} T^\ell , \qquad (2.31)$$

with the Laplace operator defined as

$$\nabla^2 \cdot = \frac{1}{\sqrt{|g_{10}|}} \partial_a \left(\sqrt{|g_{10}|} g^{ab} \partial_b \cdot \right) . \qquad (2.32)$$

The form fields obeys the following Bianchi identities

$$\begin{aligned} \mathrm{d}H_3 &= 0 , \\ \mathrm{d}F_{8-p} &= H_3 \wedge F_{6-p} + Q_{\mathrm{D}p} \delta_{9-p}(\mathrm{D}p) . \end{aligned} \qquad (2.33)$$

The equations of motion for the RR sector fields are given by

$$\begin{aligned} \mathrm{d} \left(e^{\frac{p-1}{2}\phi} \star F_{6-p} \right) = &-e^{\frac{p-3}{2}\phi} H_3 \wedge \star_{10} F_{8-p} \\ &+ (-1)^{\frac{(6-p)(5-p)}{2}} Q_{\mathrm{D}(4-p)} \delta_{5+p}(\mathrm{D}(4-p)) , \end{aligned} \qquad (2.34)$$

where every RR field satisfies the following duality relation

$$e^{\frac{5-q}{2}\phi} F_q = (-1)^{\frac{(q-1)(q-2)}{2}} \star_{10} F_{10-q} . \qquad (2.35)$$

By assuring that both Bianchi identities and equations of motion are satisfied for each RR q-form it is sufficient to only consider $q \leq 5$. Keeping all the forms for all q is called *the democratic formalism*, which will not be used here. For the NSNS 3-form

$$\mathrm{d} \left(e^{-\phi} \star H_3 \right) = -\sum_q e^{\frac{5-q}{2}\phi} \star_{10} F_q \wedge F_{q-2} . \qquad (2.36)$$

The above are given in not-necessarily extremal Dp-brane charges and tensions; $T_{\mathrm{D}p} = |T_{\mathrm{D}p}|$ and $Q_{\mathrm{D}p} = |Q_{\mathrm{D}p}|$. Extremality, on the other hand, is given by $T_{\mathrm{D}p} = Q_{\mathrm{D}p}$. To convert to any other source Table 2.1 can be used.

The type of solutions that is considered in this thesis all have a common structure. Generally considered here are (anti-)Dp-branes or Op-planes that sources a F_{8-p} field-strength. These solutions are also surrounded by a flux built up by the RR-flux F_{6-p} and the NSNS-flux H_3. Only few exceptions to this is present in this thesis and will be emphasised in each corresponding section.

	Tension	Charge
Dp	$T_{\mathrm{D}p}$	$Q_{\mathrm{D}p}$
$\overline{\mathrm{D}p}$	$T_{\mathrm{D}p}$	$-Q_{\mathrm{D}p}$
Op	$-T_{\mathrm{D}p}$	$-Q_{\mathrm{D}p}$
$\overline{\mathrm{O}p}$	$-T_{\mathrm{D}p}$	$Q_{\mathrm{D}p}$

Table 2.1. *The different sources relation to each other in terms of tension and charge.*

11D supergravity conventions

The 11D supergravity is much more concise since there are nothing equivalent to an RR- and NSNS-sector, nor does there exist a dilaton. Instead there exists a 4-form G_4 that obeys the following equation of motion and Bianchi identity

$$\mathrm{d} \star_{11} G_4 = \frac{1}{2} G_4 \wedge G_4 + Q_{\mathrm{M2}} \delta_8(\mathrm{M2}) \,,$$
$$\mathrm{d} G_4 = 0 \,, \tag{2.37}$$

where $Q_{\mathrm{M2}} = |Q_{\mathrm{M2}}|$ is the charge of a M2-brane. Einstein's equation is given by

$$R_{ab} = \frac{1}{2} \left(|G_4|^2_{ab} - \frac{1}{3}|G_4|^2 g_{ab} \right) + \frac{1}{2} \left(T^l_{ab} - \frac{1}{9} T^l g_{ab} \right) \,, \tag{2.38}$$

where the stress-tensor for the localised sources is

$$T^\ell_{ab} = -T_{\mathrm{M2}} g_{\mu\nu} \delta(\mathrm{M2}) \delta^{\mu\nu}_{ab} \,. \tag{2.39}$$

Some terminology

Let us specify the meaning of some of the terminology frequently used throughout this thesis.

Localisation: The string theory sources used here are multidimensional objects with a world-volume. Their spatial dimension will occasionally be labelled p for an unspecified number of dimensions and their corresponding world-volume is then $p + 1$ dimensional. These sources have lower world-volume dimension than the space-time dimension of the supergravity and hence there are *transversal* directions to the source. The position of the brane in the space of the transversal directions will be specified with an object $\delta_{(9-p)}(\mathrm{D}p)$ for a Dp-brane. This is essentially a $(9 - p)$-form proportional to the Dirac delta function with certain normalisations that will be specified later. When the source position in the transversal space is specified by this object, i.e. as a point, the source is *localised*.

Smearing: A *smeared* source refers to an object that have more spatial directions than specified by p. If the extra directions are less than $9 - p$, such that it does not cover the whole ten dimensional space-time, the source is considered to be *partially* smeared. To make a consistent smearing the delta function, introduced above, is exchanged for the corresponding integrated value of the source term. This integration takes place in the direction in which smearing is considered.

Probe: The sources that are considered here will influence the form-fields present in a setup. If a brane is added to an already existing background, and the effect of the branes interaction with the fields can be neglected, the brane is said to be in the *probe* approximation.

Backreaction: The *backreaction* of a source refers to when one accounts for the influence of a source on the other objects present in a solution.

The Ansätze used for the metric will differ in between sections, however some general terminology can be introduced here. A common Ansatz for the metric related to a source of p spatial dimensions will be of the form

$$\mathrm{d}s_D^2 = \mathrm{e}^{2A}\left(\mathrm{d}\tilde{s}_{p+1}^2\right) + \mathrm{e}^{2B}\left(\mathrm{d}\tilde{s}_{D-1-p}^2\right) , \qquad (2.40)$$

where the factors included here are:

A: The *warp-factor*. This is a general function on the internal coordinates, coefficient to the world-volume of the source considered.

B: The *conformal factor*. This is a general function on the internal coordinates, coefficient to the transversal directions of the source considered.

Tilde: Tilde over any object means that warping, conformal and any other added factors have been explicitly accounted for. For example

$$g_{\mu\nu} = \mathrm{e}^{2A}\tilde{g}_{\mu\nu} , \ \star_d F_q = \mathrm{e}^{dA - 2qA}\tilde{\star}_d F_q , \ \text{and} \ |F_q|^2 = \mathrm{e}^{2qA}|\tilde{F}_q|^2 , \quad (2.41)$$

in these examples F_q is a q-form with only space-time components.

The two general metrics will be referred to as

$$\begin{aligned} \textit{External: } \mathrm{d}s_{p+1}^2 &= \tilde{g}_{\mu\nu}\mathrm{d}x^\mu \mathrm{d}x^\nu , \\ \textit{Internal: } \mathrm{d}s_{D-1-p}^2 &= \tilde{g}_{ij}\mathrm{d}y^i \mathrm{d}y^j . \end{aligned} \qquad (2.42)$$

For some sections another factor will be added, the metric then has the form

$$\mathrm{d}s_D^2 = \mathrm{e}^{2A}\left(-\mathrm{e}^{2f}\mathrm{d}t^2 + \mathrm{d}s_p^2\right) + \mathrm{e}^{2B}\left(\mathrm{e}^{-2f}\mathrm{d}r^2 + \mathrm{d}s_{D-2-p}^2\right) . \qquad (2.43)$$

This new factor is:

f: The *blackening factor*, representing a horizon, dependent on r.

The field-strength related to a source of p spatial dimensions is the F_{8-p} RR-flux, for type II. Also including 11D supergravity, the common Ansatz for this form will be

$$F_{D-2-p} = \mathrm{e}^{X} \star_{D-1-p} \mathrm{d}\alpha \,. \tag{2.44}$$

The factor X has no physical relevance, but the important field here is α, i.e. the *potential associated to the field-strength*. There are some special cases for this Ansatz. For $(D, p) = (10, 3)$ this is the F_5-form which is self dual according to (2.35), the Ansatz is then extended by taking $F_{D-2-p} \to (1 + \star_D) F_{D-2-p}$. Since the democratic formalism is not used here, for $(D, p) = (10, 2)$ the field-strength is F_6, which is dual to F_4 which is the highest form kept. Similarly for 11D supergravity, $(D, p) = (11, 2)$ there does not exist a seven-form. Hence in both these cases the field-strength will instead be added to the F_4/G_4 according to

$$\star_D \mathrm{e}^{X} \star_{D-1-p} \mathrm{d}\alpha = (\text{Four-form for } p = 2) \,. \tag{2.45}$$

2.2.1 Smearing versus localisation

As was introduced earlier, the position of a localised brane is specified by the object $\delta_{(9-p)}(\mathrm{D}p)$. This object enters for example the Bianchi identity for the field-strength corresponding to such a source

$$\mathrm{d}F_{8-p} = H_3 \wedge F_{6-p} + Q_{\mathrm{D}p}\delta_{(9-p)}(\mathrm{D}p) \,. \tag{2.46}$$

The source also influences the dilaton equation of motion and the Einstein equation. This object is normalised such that when integrated over the transversal space, call it M^{9-p}, it obeys

$$\int_{M^{9-p}} \delta_{(9-p)}(\mathrm{D}p) = 1 \,. \tag{2.47}$$

The full expression for this object is

$$\delta_{(9-p)}(\mathrm{D}p) = \delta(\mathrm{D}p) \star_{9-p} 1 = \tilde{\delta}(\mathrm{D}p)\tilde{\star}_{9-p}1 \,, \tag{2.48}$$

where $\tilde{\delta}(\mathrm{D}p)$ is the ordinary delta function on the $9 - p$ dimensional internal space. This means that a consistent smearing procedure involves the following substitution

$$\tilde{\delta}(\mathrm{D}p) \to \tilde{v}^{-1} \,, \tag{2.49}$$

with $\tilde{v}$ being the internal volume, without any conformal factor. This smearing procedure also makes the corresponding field-strength vanish and the warping and conformal factor are zero

$$F_{8-p} = 0\,, \quad A = 0\,, \quad B = 0\,. \tag{2.50}$$

For *partial* smearing, that is, when the brane's position is only specified in some directions, the procedure is similar. The full object $\delta_{(9-p)}(\mathrm{D}p)$ is integrated over a $q < 9-p$ dimensional subspace of the internal space and for those direction the $9 - p$ dimensional delta function $\tilde{\delta}(\mathrm{D}p)$ is replaced by a delta function in $9 - p - q$ dimensions. A partial smearing retains a profile for the field-strength, warping and conformal factor, in the directions in which the source is still localised.

For the cases of complete smearing, the Bianchi identities and equations of motion reduce to algebraic equations, and hence significantly eases the effort of finding solutions. An effort which in general involves solving coupled non-linear partial differential equations.

The smearing procedure is heavily used in the literature. One of the most common uses is in the field of lower dimensional effective field theories where indeed the dynamics of the internal space have been integrated out, and substituted by the values of the corresponding integral.

2.2.2 Review of some legendary papers

In the present section some particular articles will be briefly reviewed. The results presented here will all be significant for what will be considered in Part II. The reason for introducing this work is to be able to emphasise the relevance of the work presented in this thesis, and to also give a broader perspective on the complete field.

Some of these works are covered in modern textbooks which considers this topic, e.g. [BBS07].

Klebanov-Tseytlin [KT00]
[KT00] is a type IIB solution with the following field content

$$\begin{aligned} \text{RR:} \quad & F_5,\ F_3\,, \\ \text{NSNS:} \quad & H_3,\ g_{\mu\nu},\ \phi\,. \end{aligned} \tag{2.51}$$

Present is also one or several D3-branes. The field F_5 is the field-strength associated with a D3-brane and F_3 and H_3 are surrounding fluxes. The D3-branes are so-called *space-filling* which means that it covers the entire external space, i.e. the space-time part of the metric. Topologically the ten dimensional space-time is
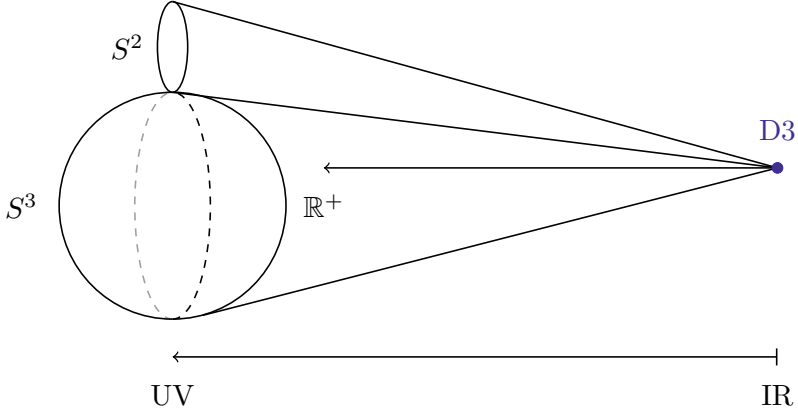
$$\mathrm{Mink}_4 \times \mathrm{CY}_6\,, \tag{2.52}$$

Figure 2.1. The [KT00] conifold, where the two spheres of the base ends in a singular point in the IR where D3-branes are positioned.

where CY_6 is a non-compact Calabi-Yau manifold of six real dimensions. This Calabi-Yau has a conical structure, meaning that it has a base, denoted $T^{1,1}$, and a radial direction $\mathbb{R}^+$. The base of the conifold has the topology $S^2 \times S^3$. The fields (2.51) are then positioned in the internal space as

$$
\begin{aligned}
F_5: &\quad T^{1,1}\,, \\
F_3: &\quad S^3 \subset T^{1,1}\,, \\
H_3: &\quad \mathbb{R}^+ \times S^2 \subset \mathrm{CY}_6\,,
\end{aligned}
\tag{2.53}
$$

and the dilaton, ϕ, is a constant. The 3-form fluxes F_3 and H_3 are organised in such a way in this solution that if one considers the following complex combination

$$
G_3 = F_3 - i\mathrm{e}^{-\phi} H_3\,,
\tag{2.54}
$$

then they obey

$$
\star_6 G_3 = iG_3\,,
\tag{2.55}
$$

where $\star_6$ is the Hodge-dual of the internal space. This condition will be referred to as the *ISD condition*, meaning *imaginary self-dual*. There is a corresponding real statement of this duality which looks like

$$
H_3 = \mathrm{e}^{\phi} \star_6 F_3\,.
\tag{2.56}
$$

However, this will still be referred to as the ISD condition.

The conifold geometry can be described by four complex coordinates z_α that satisfies the following relation

$$
\sum_{\alpha=1}^{4} z_\alpha^2 = 0\,.
\tag{2.57}
$$

This conifold gives rise to a naked singularity at the bottom of the conifold – the warp factor diverges without being shielded by a horizon. The D3-branes present are positioned at the singular point of the conifold. A depiction of the background can be found in Figure 2.1.

Herzog-Klebanov [HK01]

The [HK01] solutions are generalisation of [KT00] to various dimensions. This means that they consider space-filling Dp-branes in type IIA(B) for p being even (odd). They considered conical internal spaces of topology

$$\mathbb{R}^+ \times M^2 \times M^{6-p} \,, \tag{2.58}$$

where the field content is, and have been placed, according to

$$\begin{aligned} F_{8-p}: &\quad M^2 \times M^{6-p} \,, \\ F_{6-p}: &\quad M^{6-p} \,, \\ H_3: &\quad \mathbb{R}^+ \times M^2 \,. \end{aligned} \tag{2.59}$$

The two spaces M^2 and M^{6-p} are chosen by [HK01] as whatever spaces that satisfies the supersymmetry conditions. Since the internal space is flat and $9 - p$ dimensional, the external space-time is

$$\text{Mink}_{p+1} \,, \tag{2.60}$$

whose dimension is implied since space-filling Dp-branes are used.

Similar generalisations to various dimensions will be considered later, which is why [HK01] is mentioned here.
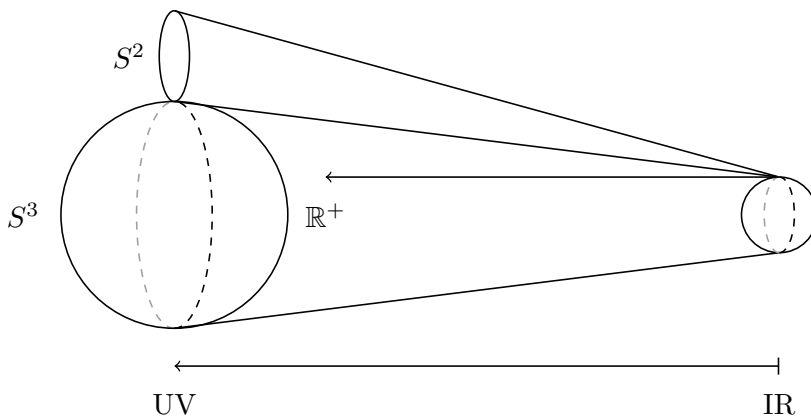
Klebanov-Strassler [KS00]

Figure 2.2. The [KS00] conifold, where the S^3 has a non-zero radius in the IR.

As was mentioned earlier, the [KT00] solution has a warp-factor singularity. The [KS00] solution is the successful *deformation* of the conifold that resolves this singularity. This is called the *deformed conifold*, not to be confused with the *resolved conifold* of [PZT00]. Both the deformed and resolved conifold can be described by adding a constant to equation (2.57) according to

$$\sum_{\alpha=1}^{4} z_\alpha^2 = \epsilon^2 \,, \tag{2.61}$$

where ϵ is related to the radius of the finite submanifold at the bottom of the conifold. For the deformed conifold the S^3 has a finite size at the bottom, see Figure 2.2, while the resolved conifold has a finite sized S^2. This deformation needs no new field content and the topology of the internal space remain the same. However it introduces new components to the fields. The internal metric are usually described by a radial coordinate τ and five one-forms $g^i|_{i=1,2,3,4,5}$ describing the $T^{1,1}$. The position of the fields and their components are in the [KS00] solution

$$
\begin{aligned}
F_5: &\quad g^{12345} \,, \\
F_3: &\quad \mathrm{d}\tau \wedge (g^{13} + g^{24}) + g^{125} + g^{345} \,, \\
H_3: &\quad \mathrm{d}\tau \wedge (g^{12} + g^{34}) + g^5 \wedge (g^{13} + g^{24}) \,,
\end{aligned}
\tag{2.62}
$$

up to coefficient τ-dependent functions. The $g^{1,2}$ describes the S^2 and $g^{3,4,5}$ the S^3 of the base. See for example [BG13] which describes an Ansatz that interpolates between the [KT00] and [KS00] Ansätze.

The D3-brane that was present at the singularity at the bottom of the [KT00] conifold is now gone, together with the singularity. This solutionis now supported only by ISD flux.
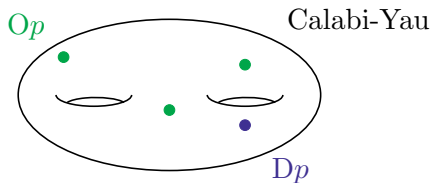
Giddings-Kachru-Polchinski [GKP02]

Figure 2.3. A representation of the [GKP02] solution.

All the solutions introduced above are non-compact. There is a simple reason as to why they cannot be made compact and this comes from

the so-called tadpole condition. Considering a D3-brane, the charge of this brane sources the Bianchi identity corresponding to F_5, that is

$$\mathrm{d}F_5 = H_3 \wedge F_3 + Q_{\mathrm{D3}}\delta_6(\mathrm{D3})\,. \qquad (2.63)$$

The charge of the brane is represented by $Q_{\mathrm{D3}} = |Q_{\mathrm{D3}}|$. For the solutions considered above, the fluxes obey the ISD condition (2.56) which means that the flux charge density can be written as

$$H_3 \wedge F_3 = \mathrm{e}^{-\phi}|H_3|^2 \star_6 1\,, \qquad (2.64)$$

that is, a positive quantity. On a compact internal manifold, equation (2.63) can be integrated to give an inconsistency

$$0 \overset{!}{=} \int_{M^6} \mathrm{e}^{-\phi}|H_3|^2 \star_6 1 + Q_{\mathrm{D3}} > 0\,. \qquad (2.65)$$

Hence D3-branes surrounded by ISD flux can not be placed on a compact manifold, nor can ISD fluxby itself support a compact manifold. Therefore all of the solutions have to be non-compact.

The non-compact solutions described above are all supersymmetric because the ISD flux and the D3-branes are mutually BPS, meaning that they can preserve the same supersymmetries. There is however other objects that are also mutually BPS with the ISD flux and the D3-branes. One such object is the three dimensional orientifold plane, O3-plane. As mentioned in Section 2.2, the O3-plane have opposite charge and tension, compared to the D3-brane. That is, $Q_{\mathrm{O3}} = -Q_{\mathrm{D3}}$. This makes it possible to substitute the D3-brane for a O3-plane and make the space compact.

This is the [GKP02] solution. It consists of the same field content as previously considered. The fluxes F_3 and H_3 obeys the ISD condition (2.56), and the source is a O3-plane. Of course it is not that strict, but any such solution allows for a number of D3-branes as well, as long as there is net O3-plane charges, see Figure 2.3. The internal space is then a compact Calabi-Yau. The exact placement of the fields in this manifold is not specified and the solution is only implicitly stated as one differential equation that solves all equations of motion and satisfies all Bianchi identities.

The BPS property of the [GKP02] solution can be seen from the following expression

$$\tilde{\nabla}\left(\mathrm{e}^{4A} - \alpha\right) = \mathrm{e}^{2A+\phi}|iG_3 - \star_6 G_3|^2 + \mathrm{e}^{-6A}|\partial(\mathrm{e}^{4A} - \alpha)|^2\,, \qquad (2.66)$$

where each term is trivially satisfied. This type of expression is expected for supersymmetry, however does not imply it. The form G_3 can have two complexity types, $(2,1)$ and $(0,3)$. Only for $(2,1)$ is supersymmetry present for the [GKP02] Ansatz. See [BJVRV12], by the present

author (article not included in this thesis), on how supersymmetry can be restored for the $(0, 3)$ complexity type case.

Kachru-Kallosh-Linde-Trivedi [KKLT03]

Up to this point some well-known supergravity solutions have been presented. The desired step now is to construct some phenomenologically relevant solutions. As mentioned in the introduction the purpose here is to break supersymmetry, keep the internal space compact, achieve a positive cosmological constant, and have the solution meta-stable, i.e. the solution is able to persist through times of the order of the age of our universe.

A solution capturing these effects is [KKLT03], and can be described as a three-step procedure. To start off, lets present the steps and then consider the details of each step.

1. Take a solution that is a no-scale compactification to Minkowski, such as the [GKP02] solution.
2. Add small non-perturbative terms which gives a negative contribution to the energy and hence makes the space Anti-de Sitter (AdS).
3. To make the energy positive and break supersymmetry, add an anti-D3-brane in a highly warped region.

The first step starts out with a nice compactification. However, [GKP02] does not have all moduli stabilised. Moduli are fields that represents the deformation of the internal space and should be fixed. Much of the discussion of moduli stabilisation is left out of this thesis. For more information on the moduli stabilisation issue for these solutions see the original papers [GKP02] and [KKLT03]. [GKP02] leaves the Kähler moduli unfixed so not only does the potential energy need to be lifted, but the Kähler moduli should also be stabilised.

The second step takes care of the Kähler moduli. The added non-perturbative effects depends on these moduli and the resulting extrema will have all moduli fixed. There is also a negative contribution to the potential energy from the non-perturbative fluxes, because they are added in such a way that the extrema is still supersymmetric. So Kähler moduli has been fixed, but supersymmetry remains and the potential energy corresponds to AdS.

For the last step, anti-D3-branes are added. These branes give a large positive contribution to the potential energy and would give de Sitter. Because of the smallness of the observed cosmological constant, one wants the positive contribution of the anti-D3-branes to overcome the negative contributions of the non-perturbative effects only slightly. This can be achieved by placing the anti-D3-branes in a very warped region, because their contribution to the potential energy is scaled down

by warping. [KKLT03] suggests that this can be done by adding the anti-D3-branes to the tip of the [KS00] conifold.
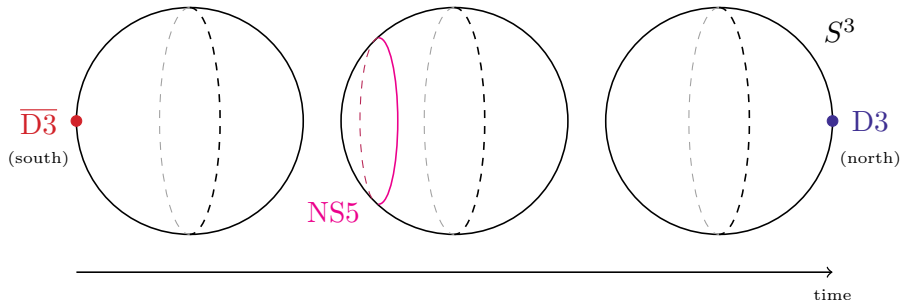
Kachru-Pearson-Verlinde [KPV02]

Figure 2.4. The [KPV02] process where anti-D3-branes decay to D3-branes through the Myers-effect, here showing the S-dual NS5-channel.

Similar to [KKLT03], [KPV02] consider anti-D3-branes on the [KS00] conifold. Assuming that the full backreaction of this problem will not affect the background significantly – such as fluxes – [KPV02] studies the physics of the (probe) anti-D3-branes living at the S^3 at the bottom of the conifold.

More precisely they consider the possibility of the anti-D3-branes polarising into a D5-brane[5] that wraps an azimuthal S^2 inside the S^3. This phenomena takes place through the so-called Myers-effect of [Mye99]. This polarised brane then have a possibility to traverse the remaining direction of the sphere and move from one pole to the other. Starting from say the south pole of the S^3 the polarised brane still have the same properties of the anti-D3-branes. As it gets closer to the north pole the effective charge changes and what remains are D3-branes, see Figure 2.4. This means that if the polarised branes manage to reach the north pole the uplifting properties of [KKLT03] would be lost.

The calculation of [KPV02] derives a potential that governs the movement of the polarised brane. This potential has a barrier that prevents the polarised brane to traverse the sphere and end up on the north pole. That is, the effective potential has a minimum close to the south pole where the polarised brane will reside. [KPV02] also estimates the possible tunnelling through this barrier and concludes that this effect is very small and the decay time would be much larger than the age of our present universe.

---

[5]Since type IIB is self-dual under S-duality, where NS5-branes are interchanged with D5-branes, [KPV02] actually considers the NS5-channel.
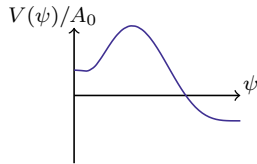
Figure 2.5. The [KPV02] meta-stable potential for the NS5-brane.

The observation of [KPV02] is that there seems to be a meta-stable state, but only for small p/M, where p is the amount of anti-D3-brane charges and M is the amount of background fluxes. In their notation

$$M = \frac{1}{4\pi} \int_A F_3, \quad p = \bar{N}_{\overline{\text{D3}}}, \tag{2.67}$$

where A represents the cycle where the F_3 flux is placed, in general it has all the [KS00] legs, but considering the bottom of the conifold $A = S^3$. Their estimation is that a meta-stable state is present for $p/M \lesssim 8\%$, see Figure 2.5.

Polchinski-Strassler [PS00]
Several of the solutions presented here have applications in the AdS/CFT duality which will not be mentioned much in this thesis. For this section it suffice to say that the supergravity background $\text{AdS}_5 \times S^5$ is dual to a stack of N coincident D3-branes under AdS/CFT. A certain perturbation of the three-form fluxes, which corresponds to a mass deformation in the CFT, gives rise to a naked singularity [GPPZ00].

In [PS00] it is argued how this singularity can be resolved via the so-called Myers-effect [Mye99]. As mentioned previously, when the [KPV02] calculation was discussed, the Myers-effect describes the polarisation of Dp branes into D$(p + 2)$ branes. [PS00] describes how polarised D5-branes could stabilise themselves and how they shield off the singular region, leaving a regular solution.

The [PS00] approach is summarised in for example [BGKM14] which itself will be briefly reviewed in Chapter 4. The important aspect of [PS00] which is relevant for the discussion in Chapter 4 is how they describe the possibility of certain polarisation channels to resolve singularities.

Gubser [Gub00]
The important subject of Part II is singularities and which type of singularities that should be considered as allowed, as well as how they can be resolved. Singularities can develop for various reasons in classical

theories. The subject of [Gub00] is to construct a *necessary* condition for singularities, in order for them to be allowed.

The idea is to consider the same system that presents a singularity, but at finite (Hawking) temperature. That is, an event horizon is inserted to surround the singular region. This is done through adding a blackening factor, as was introduced in Section 2.2. If the singularity remains hidden behind the horizon, it should be accepted. This is a criterion that will be used later.

Cvetič-Gibbons-Lü-Pope [CGLP03], [CGLP02]

In the same way as [KS00] is a smooth solution with dissolved D3-brane charge into flux in a non-compact Calabi-Yau, [CGLP03] is a smooth solution with M2-brane charge dissolved in flux. The geometry of the [CGLP03] solution is a conical Calabi-Yau with a base that is an S^4 fibre over an S^3, slightly more complicated than the conifold with $T^{1,1}$ base. The base of this conifold is a part of a classification of so-called Sasaki-Einstein manifolds, and is usually denoted $V_{5,2}$.[6]

It has several similarities to the [KS00] geometry. At the bottom of the conifold, the S^4 base of the fibration remains at non-zero size, while the S^3 fibre vanish. The flux placed in the [CGLP03] has the four-form flux $G_4|_{\text{non–field-strength part}} = F_4$ of 11D supergravity with a self-dual (SD) relation

$$F_4 = \star_8 F_4 \,. \tag{2.68}$$

Self-duality here is the equivalent of ISD in the case of [KS00], and it has the same charge as M2-branes. The base of the conifold is parametrised by seven one-forms $\{\tilde{\sigma}_i\}_{i=1,2,3}$, $\{\sigma_i\}_{i=1,2,3}$ and ν. The F_4 has four legs

$$F_4 \sim \nu \wedge \sigma_{123} + \mathrm{d}\tau \wedge \tilde{\sigma}_{123} + \epsilon^{ijk}\nu \wedge \sigma_i \wedge \tilde{\sigma}_{jk} + \epsilon^{ijk}\mathrm{d}\tau \wedge \sigma_{ij} \wedge \tilde{\sigma}_k \,, \tag{2.69}$$

up to coefficient factors dependent on τ. The notation here is taken from [KP11] which will be reviewed later. The first leg is proportional to the volume form of the S^4 that remains at $\tau = 0$

$$\nu \wedge \sigma_{123} \sim \star_{S^4} 1 \,. \tag{2.70}$$

In another paper [CGLP02], by the same authors, the same type of setup is considered but this time on the space denoted $\mathbb{A}_8$. This space has a vanishing tip which is a property that will be used later – anti-M2-brane on this tip can be fully localised.

Klebanov-Pufu [KP11]

Since it is possible to create a meta-stable state of anti-D3-branes in the [KS00] background, where the decay channel is via an NS5-brane as in

---

[6]So is the base of the [KS00] solution where $T^{1,1} = V_{4,2}$.

[KPV02], perhaps the same works for anti-M2-branes in the [CGLP03] background. Indeed this is what [KP11] finds.

The geometry and the objects available are slightly different in 11D supergravity. The decay channel for the anti-M2-brane is through an M5-brane that wraps an azimuthal S^3 inside the S^4 at the bottom of the conifold. The same events then transpires, as for the [KPV02] calculation. What is found by [KP11] is that for p/M as small as 0.054 the effective potential for the M5-brane has a barrier with a meta-stable minima, see Figure 2.6.
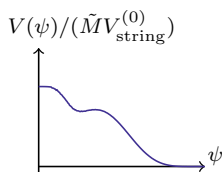

Figure 2.6. The [KP11] meta-stable potential for the M5-brane. Here drawn for $p/M = 0.05 < 0.05989$.

Part II:
Developments

3. T-duality and new solutions

This chapter will discuss Paper I, with the focus on the BPS solutions of that paper. There are additional results contained in that paper that are postponed for Chapter 4.

3.1 T-duality

In the early days of string theory, it was believed that there were several independent string theories, five to be exact. Further down the line, relations between them started to emerge. With the discovery of M-theory they were reinterpreted as all being a part of one grand theory. M/string-theory can be seen as a web that relates them through so-called dualities.

The string theories discussed in this thesis are the type II theories. It has been noted that when type IIA was compactified on a circle or radius R_A, it gave the same compactification scenario as type IIB would, compactified on a circle of radius R_B, provided

$$R_A = \frac{\alpha'}{R_B} \, .\tag{3.1}$$

For the purpose of this thesis, the following information about T-duality is needed for the results of Paper I, and also for some of the results presented in Chapter 4.

It is possible to perform T-duality in different directions. Depending on the direction used, the resulting dual solution will be different. The two examples that will be important here is T-duality along RR-sector fields or along NSNS-sector fields.

A T-duality along an RR-field would result in new RR-fields according to

$$F_{a_1 \dots a_q} \to F_{a_1 \dots a_{q-1}} \, .\tag{3.2}$$

The source giving rise to a F_q flux is of $p = 8 - q$ spatial dimensions, and after the T-duality operation it is now a $p + 1$ dimensional source. However, if the source was localised in the original solution, it will now be smeared along the T-duality direction. Hence, to be able to localise the source one has to go through the complete analysis with localised sources again.

T-dualities along the NSNS three-form flux H_3, will create new types of "fluxes". By dualising one direction of the NSNS three-form one gets

$$H_{ijk} \to f^i{}_{jk} \,, \tag{3.3}$$

where $f^i{}_{jk}$ is usually called a *metric*-flux. One might wonder whether additional T-dualities along the NSNS three-form is a sensible thing to do, and this happens to be an interesting field of its own. The complete T-duality chain is commonly denoted

$$H_{ijk} \to f^i{}_{jk} \to Q^{ij}{}_k \to R^{ijk} \,, \tag{3.4}$$

where Q and R are called *non-geometric* fluxes. They are called non-geometric because they lack a complete geometric understanding in terms of ten-dimensional supergravity. The subject of non-geometric fluxes is not relevant for this thesis and will not be discussed further. The interested reader might want to consider the review [AMN13], and references within, which considers their interpretation as coming from Double Field Theory. This review also touches upon the subject of non-geometric fluxes having an origin in string theory compactifications (as opposed to supergravity compactifications) or non-commuting ($Q^{ij} \sim [x^i, x^j]$) and non-associative ($R^{ijk} \sim [x^i, x^j, x^k]$) internal coordinates.

3.2 New solutions

In Paper I several new solutions are presented. These are all T-dual, and they are considered both fully smeared and fully localised. This serves as good examples for what information that smeared solutions can provide and how it helps to find, for solutions that are BPS, the corresponding localised solutions.

3.2.1 BPS on Ricci-flat internal space

The solutions that will be presented here are Op-plane solutions. The field-strength corresponding to an Op-plane is the F_{8-p} RR-flux, and should therefore be included. Internally, the space will be made compact and hence due to a sort of generalised Gauß's law no net charge can remain on the internal manifold. So, in order to cancel the internal charge introduced by the Op-plane, charged fluxes have to be included

$$\int \mathrm{d}F_{8-p} = 0 = \int H_3 \wedge F_{6-p} - \int Q_{\mathrm{O}p} \delta_{9-p}(\mathrm{O}p) \,, \tag{3.5}$$

as also explained in the discussion in Section 2.2.2.

In addition to this, the geometry is a warped external space and a conformal internal space

$$ds_{10}^2 = e^{2A}d\tilde{s}_{p+1}^2 + e^{2B}d\tilde{s}_{9-p}^2 \,. \tag{3.6}$$

Note also that the dilaton is in general a dynamical field that is included aswell.

Smeared solutions

As was described in Section 2.2.1 the smearing procedure constitutes removing the field-strength, warping, the conformal factor and making the dilaton constant. What remains is then simply

$$F_{6-p}, \; H_3, \; \phi = \phi_0 = \log g_s \,, \tag{3.7}$$

and of course the charge and tension of the Op-plane. In Paper I these are referred to as μ_p, but here they will be denoted $Q_{Op} = T_{Op}$ to make them distinguishable to charge and tension of other objects that will be used in subsequent sections.

To find the solutions for the smeared planes one have to make sure that all form and dilaton equations of motion, together with the Einstein equation, are satisfied. In the present setup, the equations become algebraic equations and are solved by two conditions

$$
\begin{aligned}
H_3 &= e^{(p+1)\phi_0/4} \star_{9-p} F_{6-p} \,, \\
Q_{Op} &= e^{(p+1)\phi_0/4} |F_{6-p}|^2 \,,
\end{aligned}
\tag{3.8}
$$

where these are Minkowski solutions on Ricci flat internal manifolds

$$
\begin{aligned}
R_{\mu\nu} &= 0 \,, \\
R_{ij} &= 0 \,.
\end{aligned}
\tag{3.9}
$$

The duality condition in (3.8) is for $p = 3$ often called the imaginary self-dual (ISD) condition. This is since the H_3 and F_3 can be combined into a complex three-form

$$G_3 = H_3 + SF_3 \,, \tag{3.10}$$

where $S = C_0 + ie^\phi$, and C_0 is the potential for the RR flux F_1 (put to zero here), which then is ISD

$$\star_6 G_3 = iG_3 \,. \tag{3.11}$$

These T-dual solutions give rise to similar duality relations between the fluxes.

Localised solutions

To proceed in finding the localised solutions, some information can be taken from the smeared solutions. One should expect the flux relation of equation (3.8), i.e. the first line, to still hold when the dilaton is promoted to a function. Also, the properties of the Ricci tensors should be preserved, but only for the unwarped and non-conformal counterparts. Using these hints, it is very easy to spot that the solution is given by four conditions and one partial differential equation

$$
\begin{aligned}
H_3 &= \mathrm{e}^{(p+1)\phi/4} \star_{9-p} F_{6-p}\,, \\
\alpha &= (-1)^{p+1}\mathrm{e}^{(p+1)A+(p-3)\phi/4} + \alpha_0\,, \\
\phi &= \frac{4(p-3)}{7-p}A + \phi_0\,, \\
B &= -\frac{p+1}{7-p}A\,, \\
\tilde{\nabla}^2 \mathrm{e}^{\frac{16}{p-7}A} &= -\mathrm{e}^{\frac{p-1}{2}\phi_0}|F_{6-p}|^2 + \mathrm{e}^{\frac{p-3}{4}\phi_0}Q_{\mathrm{O}p}\tilde{\delta}(\mathrm{O}p)\,.
\end{aligned}
\tag{3.12}
$$

These are implicit solutions, in the sense that the internal space is not specified, but the differential equation solves all other equations. To refer back to the smeared solution, the first and last equation of (3.12) corresponds to the same in (3.8). Setting $A = 0$ reproduces the whole smeared setup. These localised solutions are again a (warped) Minkowski on a (conformally) Ricci flat internal manifold

$$
\begin{aligned}
\tilde{R}_{\mu\nu} &= 0\,, \\
\tilde{R}_{ij} &= 0\,.
\end{aligned}
\tag{3.13}
$$

3.2.2 BPS on negatively curved internal manifold

Again Op-planes will be considered, and one has to include charged flux to cancel the tadpole, as well as warping, conformal factor and a field-strength. In this case it looks a bit different since F_{8-p} will represent both the flux charge density and the field-strength. It has the following Ansatz

$$
F_{8-p} = m_{7-p} \wedge e^9 - \mathrm{e}^{-2(p+1)A-(p-3)\phi/2}\tilde{\star}_{9-p}\mathrm{d}\alpha\,,
\tag{3.14}
$$

where the one-form e^9 contains the metric flux

$$
e^9 = \mathrm{d}x^9 + \frac{1}{2}f^9{}_{ij}x^i\mathrm{d}x^j\,,
\tag{3.15}
$$

and comes from a T-duality along the NSNS three-form, i.e. $f^9{}_{ij} = H_{ij9}$. The Op-plane position and the positioning of all other flux components are according to Table 3.1.

	(external)	(internal)	
	$\mu = 0 \dots p-1$	$i = p \dots 8$	9
Op	x		x
$f^9{}_{ij}$		x	
dx^9			x
m_{7-p}		x	

Table 3.1. *Table of constituents for the compactification including metric flux. The* x *marks the respecive objects position.*

Smeared solutions
The smeared solutions are again solved by two conditions

$$
\begin{aligned}
\mathrm{d}e^9 &= (-1)^p g^{99} \mathrm{e}^{(p-3)\phi_0/4} \star_{9-p} m_{7-p}\,, \\
Q_{\mathrm{O}p} &= \mathrm{e}^{(p-3)\phi_0/4} |F_{8-p}|^2\,,
\end{aligned}
\tag{3.16}
$$

that is, a duality relation like before and a condition that solves the tadpole. Where this is again a Minkowski space-time, but in p space-time dimensions, and now on a negatively curved background

$$
\begin{aligned}
R_{\mu\nu} &= 0\,, \\
g^{ij} R_{ij} + 2 g^{9i} R_{9i} + g^{99} R_{99} &< 0\,.
\end{aligned}
\tag{3.17}
$$

Localised solutions
For the localised solutions the field-strength is included according to (3.14).

$$
\begin{aligned}
\mathrm{d}e^9 &= (-1)^p g^{99} \mathrm{e}^{(p-3)\phi/4} \star_{9-p} m_{7-p}\,, \\
\alpha &= (-1)^{p+1} \mathrm{e}^{(p+1)A+(p-3)\phi/4} + \alpha_0\,, \\
\phi &= \frac{4(p-3)}{7-p} A + \phi_0\,, \\
B &= -\frac{p+1}{7-p} A\,, \\
\tilde{\nabla}^2 \mathrm{e}^{\frac{16}{p-7}A} &= -\mathrm{e}^{(p-3)\phi_0/2}|\tilde{\hat{F}}_{8-p}|^2 + \mathrm{e}^{(p-3)\phi_0/4} Q_{\mathrm{O}p} \tilde{\delta}(\mathrm{O}p)\,.
\end{aligned}
\tag{3.18}
$$

This is warped p-dimensional Minkowski with negatively curved internal manifold

$$
\int \sqrt{g_{(10)}} R_{(10-p)}
$$

$$
= \int \sqrt{\tilde{g}} \left(-8\frac{p+1}{7-p}(\tilde{\nabla}A)^2 - \frac{1}{4}\mathrm{e}^{\frac{16}{7-p}A} \tilde{g}_{99}\tilde{g}^{ij}\tilde{g}^{kl} f^9{}_{ik} f^9{}_{jl} \right) < 0\,.
\tag{3.19}
$$

3.3 Motivation for the new solutions and summary

The complete T-duality schema that relates these solutions can be described as

$$
H \propto \star_{9-p} F_{6-p} \quad \underset{T_{F_{7-p}}}{\overset{T_{O_p}}{\rightleftarrows}} \quad H \propto \star_{10-p} F_{7-p} \quad \underset{T_{e^9}}{\overset{T_H}{\rightleftarrows}} \quad \mathrm{d}e^9 \propto \star_{9-p} \imath_9 F_{8-p} .
$$

(3.20)

Even though these solutions are just T-duals of each other, they could be useful for the study of higher dimensional space-times. This could be interesting since the dimensionality of the internal space would be smaller and hence easier to classify. For an example of applying this logic for de Sitter searches see [VR12].

The BPS expression for the Ricci flat solutions looks like

$$
\begin{aligned}
\tilde{\nabla}^2 \big(\mathrm{e}^{(p+1)A + \frac{p-3}{4}\phi} &+ (-1)^p \alpha \big) \\
&= \mathrm{e}^{\frac{(p-3)^2}{p-7}A} \mathrm{e}^{\frac{p-3}{4}\phi} \tilde{R}_{p+1} \\
&+ \mathrm{e}^{\frac{(p+1)(9-p)}{p-7}A - \frac{p-3}{4}\phi} \left| \partial \left(\mathrm{e}^{(p+1)A + \frac{p-3}{4}\phi} + (-1)^p \alpha \right) \right|^2 \\
&+ \tfrac{1}{2} \mathrm{e}^{\frac{(p+1)(p-5)A}{p-7} + \frac{3p-5}{4}\phi} \left| F_{6-p} - (-1)^p \mathrm{e}^{-\frac{p+1}{4}\phi} \star_{9-p} H \right|^2 ,
\end{aligned}
$$

(3.21)

Two conditions enter here; the duality condition for the fluxes and the relation between the field-strength potential and the warping and dilaton. Integrating this expression over the compact internal space, the total derivative on the left hand side vanish. Consequently, the two last terms lowers the curvature of the external space. Only using these constituents hence can therefore only give AdS space-times, or Minkowski when the BPS conditions are saturated. For a smeared setup, only the curvature term and the flux term would be present. Breaking the duality condition would give AdS and these solutions will be covered in the next chapter.

4. Branes in oppositely charged fluxes and their singularities

The so-called *de Sitter landscape* is a concept in which the expansion of our universe, i.e. one that could possibly be explained via a positive cosmological constant, is just one of many. Subscribers to this idea would say that there are several universes. Some of which would have similar properties to ours and others would be significantly different and even hostile to life as we know it. The construction [KKLT03], reviewed shortly in Section 2.2.2, gives credence to this point of view. Our universe would be at an elevated extremal point in this landscape and could decay into other regimes of the landscape with lower vacuum energy. The important thing about such construction is that the decay time is much larger than the age of our universe, or else it would be difficult to explain why our universe still would look like it does. This concept is called *meta-stability*, i.e. a solution is meta-stable when it has decay channels but the decay times of these are very large.

[KKLT03] describes a possible construction for de Sitter space-times, leaving several aspects of the construction for future, more detailed, investigation. One of these aspects, that will be the only one considered here, is the last step of their procedure. This step involves placing an anti-D3-brane at a highly warped region of a compact internal manifold. A brane is associated with supersymmetry projectors that partially breaks supersymmetry of the background. The background where the anti-branes are placed already have supersymmetry broken to some degree. The anti-brane is not only responsible for uplifting the vacuum energy to become positive, but also for breaking the remaining supersymmetry.

When the details of placing anti-branes in backgrounds that preserve opposite supersymmetry were investigated, it was discovered that a certain singularity develops. Some singularities are already expected. This is a classical theory involving sources, and as such the field-strength associated to a certain source would develop singular self-energy. This singularity can however be easily understood. An analogy can be made to the self-energy of the electron, which is a classical singularity resolved by the quantum theory. In contrast to this understood singularity, the newly found singularity is present in the flux charge density that surrounds the source and so far has no universal resolution nor a fully understood interpretation.

As this singularity of the flux charge density was ascertained, assumptions and approximations were used, in particluar partial smearing and perturbative expansions. A first important step is to verify that the singularity is still present beyond these assumptions. That is, perhaps the singularity is simply a result of a careless use of these tools? If the singularity remains after this step it becomes important to investigate the possible resolutions of the singularity. For example, just as with the self-energy that is resolved by the quantum theory, perhaps the supergravity (classical theory) singularity is resolved by the quantum theory (string theory). If even this step fails one have to investigate possible interpretations of the singularity.

All of these aspects will be considered in the rest of this chapter. Although only the anti-D3-brane construction was mentioned above, a similar flux charge density singularity have been observed in other systems, for anti-D6-branes and for anti-M2-branes.[1] These three systems will be covered in this chapter.

As a starting point the anti-D6-brane will be considered, the papers relevant to this setup which are included in this thesis are Papers I, II and III, as well as a particular case of Paper V. In the analysis of the singularity associated to the anti-D6-brane it will be concluded that a singularity is present beyond partial smearing and perturbative expansions. Furthermore, a short review will be presented regarding other developments that concerns the resolution of the singularity.

The next step is to consider the anti-D3-branes. The conclusions made by other authors, and their working assumptions, will be reviewed, and a summary of Paper IV will be presented. Paper IV considers a possible interpretation of the singularity. The main idea is that the singularity is present due to the fact that in all of these cases a static Ansatz is used to describe a time-dependent system which ultimately leads to the singularity. By introducing this time-dependence in an adiabatic manner, a new channel of instability arise.

A flux energy density singularity is also present in the 11D supergravity, and the status of this will be reviewed subsequently. The relevant paper included in this thesis is Paper VI.

The chapter ends with a summary of the present status of this programme, collecting all results into a final conclusion.

---

[1]It would also be possible to consider anti-D2-branes in the [CGLP01] background. However none of the articles included into this thesis considers them explicitly and hence they will also be left out of the discussion. The reader is instead referred to [CGLP01] and [GGO12].

4.1 The anti-D6-brane solution and their singularity

4.1.1 The smeared solutions

In the effort of investigating the solutions of Paper I, discussed in Chapter 3, a BPS type of relation was derived where equations of motions have been combined to form squares

$$
\tilde{\nabla}^2 \left(e^{(p+1)A + \frac{(p-3)}{4}\phi} + (-1)^p \alpha \right) = e^{\frac{(p-3)^2}{p-7}A + \frac{(p-3)}{4}\phi} \tilde{R}_{p+1}
$$
$$
+ e^{\frac{(p+1)(9-p)}{p-7}A - \frac{p-3}{4}\phi} \left| \partial \left(e^{(p+1)A + \frac{(p-3)}{4}\phi} + (-1)^p \alpha \right) \right|^2
$$
$$
+ \frac{1}{2} e^{\frac{(p+1)(p-5)}{p-7}A + \frac{3p-5}{4}\phi} \left| F_{6-p} - (-1)^p e^{-\frac{p+1}{4}\phi} \star_{9-p} H_3 \right|^2 .
$$
$$(4.1)$$

The BPS relation between the fluxes are such that both squares and the term containing the Laplacian are trivially zero. In the last term one can define a generalisation of the ISD relation (2.56) to any dimension, determined by p. Even if this relation does not imply the existence of a complex form that is ISD, from here on even this more general type of flux will be referred to as the *ISD*-relation. If one would break any of these BPS relations, one would get AdS_{p+1} solutions. Since the BPS solutions have already been covered, the interest lies in what these AdS_{p+1} solutions are.

Before continuing, between Papers I and II it was noticed that the generalised broken ISD relation on the form

$$
F_{6-p} = (-1)^p \kappa e^{-\frac{p+1}{4}\phi} \star_{9-p} H_3 , \tag{4.2}
$$

does not capture the whole dynamics. This relation is referred to as *broken* ISD because of the inserted parameter κ. It turns out, however, that one instead should use

$$
H_3 = \lambda e^{\frac{p+1}{4}\phi} \star_{9-p} F_{6-p} , \tag{4.3}
$$

with $\lambda = \frac{1}{\kappa}$. The use of the parameter κ will from now on be dropped completely. The reason is that λ will be shown to have regular sign switches in the bulk, which means κ would have bulk singularities. Therefore the appropriate function to use is λ.

The first simplifying step towards determining these solutions is to use the smearing procedure. This will set the warp factor A and the potential α to zero and the dilaton to a constant $\phi \to \phi_0 = \log g_s$. This is according to what have been explained already in Section 2.2.1, and this reduces the equations to algebraic ones. Before leaving the discussion of the ISD condition it should be mentioned that the relation

can be broken even further. The smeared AdS_{p+1} solutions of Paper I considers an ISD relation of the form

$$H_3 = \lambda e^{\frac{p+1}{4}\phi} \star_{9-p} \left(F_{6-p} - \overline{F}_{6-p}\right) , \qquad (4.4)$$

where $\overline{F}_{6-p}$ is closed and co-closed and satisfies $\overline{F}_{6-p} \wedge H_3 = 0$. Dropping this newly introduced flux component induces a condition on the parameter λ which now have to be fixed to the value

$$\lambda = -\frac{p-1}{2} . \qquad (4.5)$$

The important result here is that this value is necessarily negative, hence it sources the Bianchi identity for the F_{8-p} field-strength with the opposite signed charge. Two other important aspects that needs to be derived are the curvature of the internal space and the tension/charge of the source. Via the trace of the internal and external components of the Einstein equation one deduce the internal curvature to be

$$R_{9-p} = -\frac{2p}{p+1} R_{p+1} , \qquad (4.6)$$

hence necessarily positive. The internal space can be split such that there is a three dimensional subspace containing H_3 and a $(6-p)$ dimensional one containing F_{p-6} flux, where the curvature of each subspace is given by

$$\begin{aligned} R_{ij}^{(3)} &= 2(p+1)(p-2)e^{(p-1)\phi_0/2}|\tilde{F}_{6-p}|^2 g_{ij}^{(3)} , \\ R_{ij}^{(6-p)} &= 2(p+1)e^{(p-1)\phi_0/2}|\tilde{F}_{6-p}|^2 g_{ij}^{6-p} . \end{aligned} \qquad (4.7)$$

Except for a two special cases; $p = 5$ where $6 - p = 1$ no positively curved space is available since the space is one-dimensional, and $p = 2$ where the $R_{ij}^{(3)}$ vanish, the internal space can then be taken to be a direct product of spheres

$$\mathcal{M}^{9-p} = S^3 \times S^{6-p} . \qquad (4.8)$$

By combining the ISD relation (4.3) and the Bianchi identity the charge is given by

$$Q_{Dp} = -\lambda e^{\frac{p+1}{4}\phi_0} |F_{6-p}|^2 . \qquad (4.9)$$

As mentioned in Section 2.2.2 the [GKP02] solution is possible to make compact since the positive ISD-flux charge is cancelled by the negative O3-plane charge. In this case the tadpole is instead solved via a cancellation of negative, imaginary anti-self-dual IASD[2] flux (IASD) and positive D-brane charge. For a better overview of these solutions Table 4.1 have been added.

---

[2]Note that only the for $p = 3$ the flux determined by (4.5) is strictly IASD. What is meant with IASD here is that this flux has the opposite signed charge dissolved in the fluxes, or in other words net IASD "components".

	ISD	IASD
Dp	[KT00] + [HK01] (non-compact)	Paper I + [Sil04] (compact)
$\overline{\text{D}p}$	(Paper I + [Sil04])* (compact)	([KT00] + [HK01])* (non-compact)
Op	Paper I + [GKP02] (compact)	(non-compact)
$\overline{\text{O}p}$	(non-compact)	(Paper I + [GKP02])* (compact)

Table 4.1. *All the solutions. Charge conjugation labelled with an asterisk. Note that some of these are explicit solutions while others are implicit (given in terms of one differential equation that solves all equations of motion), and some are smeared and some are localised.*

4.1.2 Localising the non-BPS solutions

The first indication that localisation would be problematic came from considering how the source term influences the external Ricci curvature. By examining an implicit Ansatz for the metric

$$\mathrm{d}s_{10}^2 = \mathrm{e}^{2A}\mathrm{d}\tilde{s}_4^2 + \mathrm{d}s_6^2\,, \qquad (4.10)$$

where $\mathrm{d}s_6^2$ is left unfixed but could include a conformal factor. The warp factor A and the potential α for the field-strength F_5 are both included to capture some of the localisation effects. Using this Ansatz one can relate the external curvature to only source terms

$$\mathrm{e}^{-2A}\tilde{R}_4 = -(1 \mp 1)Q_{\mathrm{D3}}\delta(\mathrm{O3/D3})\,. \qquad (4.11)$$

The negative sign refers to the O3-plane charge, which accurately reproduces the Minkowski space-time of [GKP02]. The positive sign, however, corresponds to the D3-brane charge and there is a constant negative curvature when the source is smeared $\delta(\mathrm{D3}) \to 1$. When the source is localised, the AdS$_4$ curvature instead collapses to a point. This analysis does not prove much, it should only be taken as a hint: localisation will get troublesome.

From here on only the charge conjugated solution from the previous section will be used. This is to make connection to the [KKLT03] construction where anti-brane uplifting is usually referred to as "placing anti-branes in an ISD background". The dimension will also be set to $p = 6$ such that localisation will only occur on a single manifold – S^3.

An important effect that happens upon localisation is that the profile of the flux distribution (4.3) will change such that it can vary across the internal manifold. For a stable solution to arise one might expect the flux around the source to be mutually BPS, as in Figure 4.1. Since

the global charge of the internal flux is mutually non-BPS, it should therefore get a non-trivial profile which respects the smearing procedure, that is the net amount of flux is ISD. Consequently

$$\lambda = \frac{5}{2} \quad \to \quad \lambda(\theta) \text{ such that } \quad \int \lambda(\theta) \mathrm{e}^{7\phi/4+3B} \tilde{\star}_3 \, 1 = \frac{5}{2} g_s^{7/4} \int \tilde{\star}_3 \, 1 \,. \tag{4.12}$$
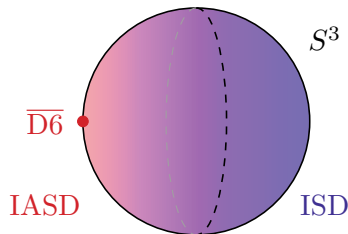

Figure 4.1. Mutually BPS flux close to the brane.

The generalisation by including λ can also be motivated by [BGH10] where is was shown that the backreaction of anti-D3-branes in the [KS00] conifold also changes the background ISD profile. This will be elaborated further upon later, see Section 4.2.1.

The statement made above, i.e. that the flux will arrange itself such that it becomes mutually BPS at the brane depicted in Figure 4.1, while desired for a stable solution, it will be shown that this is not what happens.
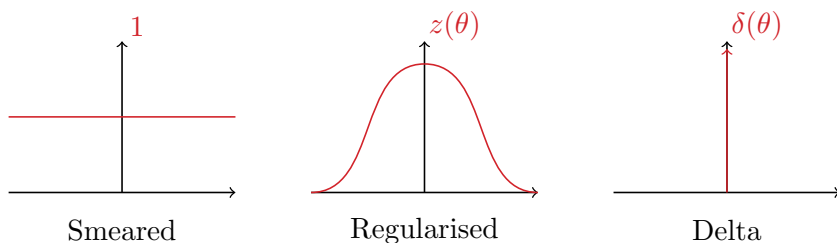
Regularised sources

Figure 4.2. The source profiles to be considered.

One attempt made in Paper II was to localise the source to a regularised source profile, see Figure 4.2. This was done by introducing a function $z(\theta)$ that relates to the source term as

$$\delta\left(\overline{\mathrm{D6}}\right) = \frac{1}{\sqrt{g^{(3)}}} \tilde{\delta}\left(\overline{\mathrm{D6}}\right) \to \mathrm{e}^{-3B} z(\theta) \,. \tag{4.13}$$

All functions were evaluated in a Taylor series at an arbitrary point on the three-sphere and one regular source profile solved all equations of motion. However, all fields except $z(\theta)$ and $B(\theta)$ turned out to be constant. In fact, the combination of z and B that appears in the source term (4.13) is also constant. This means that what has been found is just the smeared solution in terms after a change of coordinate θ. This attempt showed that there are only two possible sources – the completely localised and the smeared source.

Topological no-go

The main result of Paper II is that by a certain combination of the H_3 equation of motion and the F_2 Bianchi identity, it is possible to find a relation that restricts the behaviour of the potential function α.

The way it works is that the H_3 equation of motion provides a relation between α and flux density λ

$$\alpha = \lambda e^{7A+\frac{3}{4}\phi} + \beta \,. \tag{4.14}$$

By choosing a gauge[3] where $\beta = 0$, the dependence of λ can be completely removed from the Bianchi identity for F_2 to get an equation involving all but λ

$$\frac{1}{e^{3B}\sin^2\theta}\partial_\theta\left(e^{-7A+B-\frac{3}{2}\phi}\sin^2\theta\partial_\theta\alpha\right) = \alpha e^{-7A+\phi}F_0^2 - Q_{\mathrm{D6}}\delta(\overline{\mathrm{D6}})\,. \tag{4.15}$$

This is a quite complicated differential equation, especially since it couples to the dilaton equation of motion and the Einstein equation. However, when considering the above at a point $\theta = \vartheta$ such that $\partial_\theta\alpha|_{\theta=\vartheta} = 0$, this equation can be reduced to the following statement

$$\mathrm{sgn}\,\alpha = \mathrm{sgn}\,\alpha''\,. \tag{4.16}$$

This simple relation restricts the profile of the function α and in turn, through (4.14), the flux distribution λ. The importance of (4.14) is that it shows that

$$\mathrm{sgn}\,\alpha = \mathrm{sgn}\,\lambda\,. \tag{4.17}$$

In fact, one can also integrate the Bianchi identity and find that at the source position

$$\mathrm{sgn}\,\alpha'|_{\theta=0} = -\mathrm{sgn}\,Q_{\mathrm{D6}}\,. \tag{4.18}$$

As motivated earlier, λ was introduced in order to establish a degree of freedom that could make the flux become mutually BPS with the brane at the brane position. A anti-D6-brane in a mutually BPS flux background, $\lambda = -1$, would have the following behaviour for $\alpha = -e^{7A+\frac{3}{4}\phi}$

$$\alpha'|_{\theta=0} < 0\,, \quad \alpha(0) = 0\,, \tag{4.19}$$

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[3]This gauge parameter β is referred to as α_0 in Paper II.

as depicted in Figure 4.3. However, since the smeared solution should
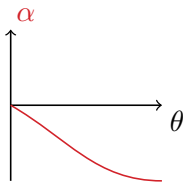

Figure 4.3. The would be BPS α profile.

originate from the integrated localised solution, α must switch sign. The reasoning is as follows: a correspondence between the smeared and localised solution means that λ must satisfy (4.12), i.e. have an integrated positive sign, and α must have the same sign as λ according to (4.14). Although this statement must be obeyed, according to (4.16) α cannot change signs in this way, see Figure 4.4.
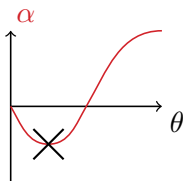

Figure 4.4. The behaviour of α satisfying the smearing procedure.

This shows that the localised anti-D6-brane solution cannot have mutually BPS flux at the position of the brane. There is however one way to resolve this issue. It is possible to find a profile for α that satisfies all the conditions stated above. This is if α procures a non-zero, positive, value at the brane position, as in Figure 4.5. This does make λ diverge
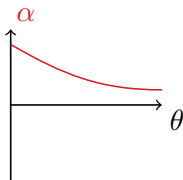

Figure 4.5. The allowed behaviour of α satisfying the smearing procedure and all topological conditions.

at the brane position, with a positive sign, since (4.14) contains $e^{7A+\frac{3}{4}\phi}$ which is expected to be zero at the position of the brane. Note that this

is an assumption and the exact boundary conditions at the brane have to be calculated explicitly to see if they indeed satisfies this assumption. However, if the assumption holds, then the flux density will diverge with a sign that is non-BPS with respect to the brane, see Figure 4.6.
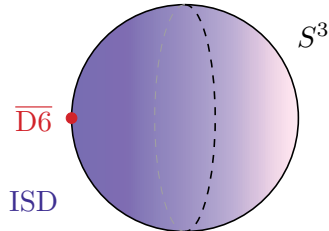

Figure 4.6. Flux distribution that satisfies the topological restrictions.

It should be noted that all of the above conclusions also hold when another anti-D6-brane is placed at the opposite side.

Confirming the boundary conditions
As just mentioned, the above reasoning are based on the assumption that a constant α at the brane position will lead to a singularity in λ. This only holds if $\mathrm{e}^{7A+\frac{3}{4}\phi}$ tends to zero at the brane position and λ diverges at the same rate.

To be able to classify all possible boundary conditions the four functions A, B, ϕ and λ are expanded in terms of θ with arbitrary powers

$$
\begin{aligned}
\mathrm{e}^{-A} &= a_0\theta^a + a_1\theta^{a+\zeta} + a_2\theta^{a+\xi} + \dots, \\
\mathrm{e}^{-2B} &= b_0\theta^b + b_1\theta^{b+\zeta} + b_2\theta^{b+\xi} + \dots, \\
\mathrm{e}^{-\frac{1}{4}\phi} &= f_0\theta^f + f_1\theta^{f+\zeta} + f_2\theta^{f+\xi} + \dots, \\
\lambda &= \lambda_0\theta^l + \lambda_1\theta^{l+\zeta} + \lambda_2\theta^{l+\xi} + \dots,
\end{aligned}
\tag{4.20}
$$

where all expansion parameters and powers[4] can take any real values. Note that one assumption is that $0 < \zeta < \xi$, and so on, such that the expansions can be consistently truncated. These expansions are then inserted into the equations of motion and all possible values for $\{a, b, f, l, \zeta, \xi\}$ that would consistently solve all equation of motion to zeroth order would be candidate boundary conditions.

Since these are small θ expansions, these expressions will only capture close to brane behaviour, such as sign of charge. The topological restrictions, which are enforced in the bulk, will have to be imposed when the

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[4]In Paper III the powers are in two cases referred to with the same symbol as the function. Some caution is needed to distinguish the functions from the expansion powers.

global properties are considered. The only goal here is to classify the boundary conditions, which is the subject of Paper III.

There are only three boundary conditions that are *physical*, in the sense that the rest of them have inconsistent sources. The physical boundary conditions are

1. No source at $\theta = 0$

$$a = b = f = l = 0 \,, \tag{4.21}$$

   implying constants at leading order.

2. The BPS boundary conditions

$$a = -\frac{1}{16}, \ b = \frac{7}{8}, \ f = -\frac{3}{16}, \ l = 0, \ \lambda_0 = -1 \,. \tag{4.22}$$

   These hence obey the BPS conditions $-16A = 16B/7 = -4\phi/3$ at least locally.

3. Boundary conditions with singular λ

$$a = -\frac{1}{16}, \ b = \frac{7}{8}, \ f = -\frac{3}{16}, \ l = -1 \,. \tag{4.23}$$

The first boundary condition tells us that one can have one pole without source, however to be able to cancel the tadpole there must be an anti-D6-brane on at least one of the poles. The second boundary condition is the BPS one, and as was shown in the previous section, these will not obey the global topological constraints. The third boundary condition is new, and is exactly what was assumed to exist. This boundary condition have a singular λ at $\theta = 0$ such that α becomes finite at the same point. This is the only boundary condition that obeys the (global) topological constraints.

Resolving the singularity

It was just shown that the flux density, when the brane is localised, displays a singularity. Whether this singularity is physical[5] or not remains as an unsolved problem. To resolve this question one can consider the criterion of Gubser [Gub00]. This is a suggested criterion that a singularity is physical and will be resolved in some way, if the singularity can be hidden behind a horizon.

This criterion has been shown to be true in several cases. As mentioned in Paper V, which deals with this issue, the Gubser criterion resolves the [GPPZ00] singularity, the [FM01] black hole, and the [KT00] singularity [KS00, Buc01, BHKPZT01, GHKT01, ABK07, Buc11].

To apply the criterion to these anti-D6-branes a horizon needs to be added that shields the branes and one should then investigate whether the singularity can be avoided. The approach will be to again write

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[5] "Physical" here means that the singularity can somehow be resolved or interpreted.

down the topological constraints and see if they are modified in such a way that the either a constant α does not imply singular λ or that $\alpha = 0$ at the brane position.

It is shown in Paper V that the topological constraints remain the same. The only difference is that the argument is now made around a point $\theta = k$, where k is the position of the horizon, instead of the origin $\theta = 0$. The important difference appears in how 4.14 changes when a blackening factor

$$\mathrm{e}^{2f} = 1 - \frac{k}{r}\,, \tag{4.24}$$

is added. The new form of the relation between α and λ becomes

$$\alpha = \lambda \mathrm{e}^{7A + \frac{3}{4}\phi + f}\,. \tag{4.25}$$

This relation shows that α must be zero at the horizon unless λ is singular. Again, the only way to satisfy the topological constraints is if α acquires a positive value. Therefore the singularity in the flux density is still present. The flux singularity has now only moved and is positioned in front of the horizon.

By working under the assumption that the Gubser criterion is a necessary one, the conclusion to be drawn here is that the singularity is non-physical. The backgrounds considered here are actually the toy backgrounds $\mathbb{R}^3 \times \mathbb{T}^{6-p}$, which have a type of conical structure. This should be a good approximation for the near-brane behaviour, when the internal space curvature is not visible.

It should be mentioned that the same conclusions were reached in [BBD13] by the means of a numerical analysis. The above argument strengthen the conclusion of their paper and gives a stronger intuition as to why the singularity is still present.

4.1.3 Parallel developments regarding anti-D6-branes

In the previous section it was shown that a singularity arises in the flux density when the anti-D6-branes are localised. However it might be possible to avoid this singularity. One method in which the singularity could be resolved is through a procedure described by [PS00]. The idea is that at a finite distance from the anti-D6-brane there resides a polarised D8-brane, using the Myers effect [Mye99]. This would resolve the singularity by cutting it off the flux density to a finite density that ends on the D8-brane.

To determine the possibility of a D8-brane, [BJKVRWZ12] derives the potential that would govern the motion of the D8-brane. This potential is a part of a Taylor series around the anti-D6-brane position, and hence only local statements can be made from this potential. The Taylor series

is made when the internal space is non-compact $\mathbb{R}^3$, see Figure 4.7, which approximates the S^3 only at short distances. The external space is also without curvature, i.e. Minkowski. What they do find when using this approach is that the polarisation potential has a certain shape without a minimum. Hence there is no (meta-)stable position for the D8-brane to reside at.
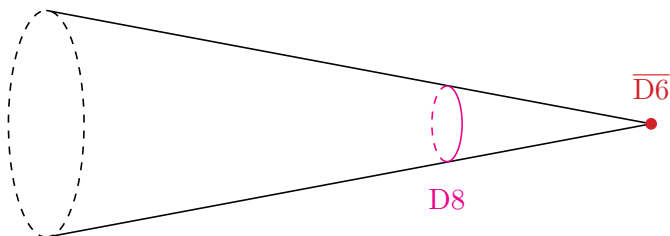

Figure 4.7. The [BJKVRWZ12] non-compact D8-polarisation.

Although, in the appendix of [BJKVRWZ12] the same calculation is performed for $\mathbb{R}^{6,1} \times \mathbb{R}^3 \to \text{AdS}_7 \times S^3$, and their conclusion is that the global properties of the solution must be known before a clear statement can be made. The non-compact analysis, however, gives an argument, since this solution is T-dual to anti-D3-branes on $\mathbb{R}^3 \times \mathbb{T}^3$ which is similar to the [KT00] and [KS00] backgrounds.

In [AFRT13] a very general analysis of the possible $AdS_7 \times M^3$ compactifications are studied. This is for the most part an implicit study where they find that only type IIA can support such geometries. In both massive and massless type IIA they verify the presence of a flux singularity in H_3, unless D8-branes are inserted to resolve them.

The difference from [BJKVRWZ12] is that [AFRT13] considers the global structure of the internal space, that is, respecting flux quantisation and tadpole cancellation. Even tough the analysis is mainly implicit they give a non-trivial numerical example in which the singularity is resolved. They place a D6- and an anti-D6-brane at each opposing pole, which are then polarised into two separate stacks of D8-branes at a finite distance from the poles. These two D8-brane stacks carry opposite D6-brane charges. At each pole the solution is BPS, see Figure 4.8, even though the supersymmetry preserved by by the two branes are opposite, the supersymmetry parameters are rotated from one pole to another and hence the system can become BPS and restore supersymmetry.

This indicates that D8-polarisations are possible, but this is numerical analysis and to determine whether more general D8-polarisations are possible, an analytic approach is needed.

The corresponding analytical answer to the the question about the possible D8-polarisations is studied in [JSZ14]. There it was analytically
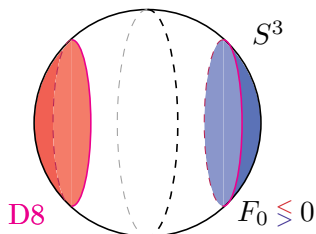
Figure 4.8. A depiction of the [AFRT13] example. The D8 stacks increase the value of F_0 integral steps from negative on the left to zero in the centre and positive on the right.

found that the external AdS curvature is responsible for the possibility to retain supersymmetry and find the meta-stable state for the polarised D8-brane. They extend the study also to more general D8-polarisations and find numerical evidence that there also exist non-supersymmetric solutions with a polarised D8-brane.

4.2 An interpretation of the anti-D3-brane singularity

Before discussing the results of Paper IV that concerns the singularity related to anti-D3-branes, there will be a short review of the development of other contributions to this area.

4.2.1 Parallel developments regarding anti-D3-branes

The singularity at a perturbative level

The anti-D3-brane flux singularity was first discovered by [MSS11]. They were considering a perturbative approach, invented by Borokhov-Gubser [BG03]. Using this approach the backreaction of the anti-D3-branes is computed as an expansion p/M, where p is the number of anti-D3-branes and M is the amount of background flux, around the supersymmetric [KS00] background.

What was observed by [MSS11] is that one of the modes that enters the expression for the fluxes have a singular mode in the first order of the expansion[6]

$$|F_3|^2 \sim |H_3|^2 \sim \frac{\mathcal{S}^2}{\tau^2} \, , \tag{4.26}$$

───────────────

[6]The coordinate τ here is the radial coordinate of the conifold. See Appendix A of [BGH10] for a dictionary between different conventions for the radial coordinate.

where $\mathcal{S}$ is a constant from the perturbative expansion that can be determined by matching UV and IR boundary conditions. In [BGH10] the singularity was confirmed, again by the same linearised perturbative approach. [BGH10] also calculates the force on a probe D3-brane that is placed at some r to be

$$F_{D3} \sim \frac{X_1}{r^5} \, , \qquad (4.27)$$

where X_1 is a parameter determined by boundary conditions. It is furthermore argued by [BGH10] that this result is the correct force by comparing to [KKLMMT03]. [KKLMMT03] investigates the a scenario of cosmic inflation from a probe anti-D3-brane falling down the [KS00] background when D3-brane charges are present at the bottom of the conifold. The radial dependence of the force of [BGH10] agrees with the result of [KKLMMT03]. One important part here is that this force is proportional to X_1, this parameter is related to $\mathcal{S}$ used by [MSS11] and the singularity that arise in the fluxes behaves as

$$|F_3|^2 \sim |H_3|^2 \sim \frac{X_1^2}{\tau^2} \, . \qquad (4.28)$$

This gives a strong argument that the singularity must be present, at least in this perturbative expansion, if an anti-D3-brane is present at the bottom of the [KS00] conifold.

[BGH10] also finds that the ISD relation is necessarily broken. The three legs of the three-forms, see equation (2.62), in [KS00] satisfies the ISD relation (2.56). However, when an anti-D3-brane is added, the ISD relation is broken and obtains a non-trivial profile dependent on τ. This is elaborated further in [Mas12a].

These perturbative and implicit (in terms of nested integrals) solutions, have since then been quite rigorously studied. In [BGGHM13] the nested integrals were to a large extent solved (only a few remain). The singularity was also commented upon in [Dym11] where it is argued that the singularity is only a artefact of the perturbative expansion. It is hence argued that beyond the perturbative approach, or to higher orders, new terms would cancel the singularity and make the flux densities finite at the position of the branes. It was later argued by [Mas12a] that the argument presented by [Dym11] is flawed, in the sense that [Dym11] only has an argument that would resolve one singular component of the fluxes if only one component was singular. As pointed out in [Mas12a] the flux has two singular components from the perturbative approach.

Not only do these approaches suffer from a perturbative expansion, but also partial smearing. The anti-D3-brane is only localised in the radial direction of the conifold and is smeared along the finite S^3 at the bottom. Arguments beyond both of these assumptions will be discussed next.

Beyond partial smearing and perturbative techniques

As mentioned above, the existence of the singularity is quite firmly established – under the assumptions of partial smearing and a perturbative approach. So far, in this presentation, only the anti-D6-brane singularity is established beyond these assumptions.

One argument for the existence of the singularity for the anti-D3-brane, made in [Dym11], is that it could simply be a relic of these assumptions. However, because of the well established existence of the anti-D6-brane singularity one would expect it to survive for the anti-D3-brane as well. There might be reasons not to expect this, although such opinions are absent in the literature.

The collaboration [GJZ13] discovers a new method to argue for the existence of the singularity. They derive an expression of the following form

$$\frac{8v\mathcal{V}}{d-2}\Lambda = \sum_p \left(1 + \frac{p-3}{2}c\right)\left[S_{\text{DBI}}^{(p)} + S_{\text{WZ}}^{(p)}\right] + \int \mathcal{F}(c). \qquad (4.29)$$

The first term is proportional to the unwarped curvature of the external space-time $\Lambda \sim \tilde{R}_{p+1}$. $\mathcal{V}$ and v are related to the internal and external volume respectively. On the right hand side there is a term involving the DBI and WZ actions of the sources and the last term is an expression involving the fluxes. For the full details of this expression see [GJZ13], also covered in the Ph.D. thesis [Jun13].

The importance of this expression is that they can verify the singularity of the anti-D6-branes. This is done by showing a non-perfect cancellation between the sources that gives the AdS curvature $\Lambda \sim Q_{\bar{D}6}\alpha_0$, while the $\mathcal{F}(c)$-term can be gauged to zero. α_0 being the field-strength potential at the position of the brane is hence a constant. This constant implies a singularity in the flux charge density, in the same way as it was argued in Section 4.1.2 where Paper II was covered.

It is important to note that their argument goes beyond any expansion and assumption of smearing since all statements are global.[7] The importance lies in their ability, using (4.29), to determine the presence of the singularity for [KKLT03]. That is, the full compact situation when an anti-D3-brane has been placed in the [KS00] conifold and made compact, and including non-perturbative contributions. This is a very strong argument for the existence of the singularity.

Another note to be made here is the similarity between the very well formulated argument of [GJZ13] and the not so well formulated argument presented in Section 4.1.2, covering Paper I. The equation

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[7]One way of interpreting the *smearing* procedure is that the equations considered are integrated. However [GJZ13] goes beyond this, even though they consider integrated equations only, since warping and the field-strength are appropriately accounted for.

(4.11) also displays a problem due to the imperfect cancellation of the source terms, just as for (4.29).

The singularity has also been established beyond a perturbative approach in [BGKM13a] and later strengthened in [BGKM13b]. Although these approaches involve partial smearing they still use the [KS00] background explicitly. The importance of these works is that they can verify the singularity beyond the perturbative approach.

Attempt for resolving the singularity

As was suggested by [Dym11], the singularity might be an artefact of the perturbative expansion. It was observed by [BGKM13b] that beyond the perturbative approach one singular component remains, i.e. the component of F_3 that occupies the S^3.

As mentioned in Section 4.1.3, one approach to resolve the singularity is to allow for the polarisation of a D$(p + 2)$-brane through the Myers effect [Mye99], according to the description of [PS00] which was summarised in Section 2.2.2.

In [BGKM13b] the D5-brane polarisation potential was studied. The method used was to calculate a potential for the D5-brane to determine if there exists a minimum at which the D5-brane can reside. The resulting potential is shown by [BGKM13b] to have no minimum and hence there cannot be any (meta-)stable position for the D5-brane. Hence reaching the same conclusion as for the non-compact anti-D6- to D8-brane polarisation.

Another possibility is to by hand add external curvature to the background in which the anti-D3-branes are placed. This analysis has been performed by [BG13] who finds that the singularity is not possible to resolve in this way. This is done numerically in both the [KT00] and [KS00] backgrounds. The efforts to extend their results analytically still remains.

Warping corrections and the influence of the singularity on KKLT

Often when effective theories are considered, localisation effects are left out of the calculation under the assumption that they are under control. In light of the presence of the still unresolved singularity, the [KKLT03] calculation should be redone. While revising the calculation, it should be kept in mind that localised sources induce warping and this warping is related to the singularity, even though it is present in the flux charge density [GJZ13]. This is the aim of [Jun14], where it is found that the outcome depends on a choice of boundary conditions. The singularity can show up in the effective potential, and hence problematic to interpret, or the original result of [KKLT03] can remain but only under the

assumption that this boundary condition holds. Whether this choice is reasonable still needs to be investigated further.

The existence of a CFT dual

There is another direction available arguing for a non-singular anti-D3-brane solution. This argument utilizes the AdS/CFT correspondence. The idea is that if a regular anti-D3-brane solution exists there should also exists a CFT dual to this solution. Presently no such solution has been revealed. Since none of the works included in this thesis contains any calculation or argument regarding the existence of a sensible CFT, this discussion will be kept very short. The latest developments of this approach can be found in [DM13] and [Hoo14], and the author would advice the interested reader to consider mentioned references instead of reviewing them here.

4.2.2 Resolving the anti-D3-brane singularity using time-dependence

An important observation made in the analysis of the singularity, in the case of the anti-D6-branes and the anti-D3-branes, is that the ISD components of the flux are found to be singular at the brane position. Consider the Bianchi identity for the field-strength associated with a Dp-brane

$$\mathrm{d}F_{8-p} = H_3 \wedge F_{6-p} + Q_{\mathrm{D}p}\delta_{9-p}(\mathrm{D}p)\,. \tag{4.30}$$

The ISD signed flux makes the flux charge density positive

$$\rho_{9-p}(\mathrm{ISD}) = H_3 \wedge F_{6-p} = \lambda \mathrm{e}^{(p+1)\phi/2}|F_{6-p}|^2 \star_{9-p} 1 \geq 0\,. \tag{4.31}$$

That is, it has the same signed charge as a Dp-brane. For the BPS solutions where (anti-)Dp-branes are placed in I(A)SD fluxes the net force between the flux is cancelled and a so-called *no-force condition* is achieved. While if the charge has a sign change, as in the case of an anti-Dp-brane in a ISD dominated background, there will be an attractive net force. The tension and charge signs are listed in Table 4.2. This is not an original idea, but the same type of flux–brane attraction is mentioned in [DKV04]. This leads to a simple interpretation: if there are net forces in the system the Ansatz should be time dependent. In fact, ignoring time dependence have been known to give rise to singularities in other systems, see [DGS03] and [Gre96].

It seems very plausible that in the system where a brane exists in a mutually non-BPS background the flux will be attracted to the brane. When the flux is great enough the brane will be annihilated by the surrounding flux, through the same channel as described by [KPV02].

	Tension	Charge
ISD	$+$	$+$
IASD	$+$	$-$
Dp	$+$	$+$
$\overline{\text{D}p}$	$+$	$-$

Table 4.2. *The effective tension and charge signs of fluxes and sources.*

By trusting this interpretation, the singularity would then arise because of two factors. The first being that in all cases where the singularity have been observed the Ansätze have been static. The second is that the computation, for the Ansätze where the singularity is present, are a pure closed string computation while flux-brane annihilation is an open-string process. From this interpretation one should conclude that the singularity is signalling the need of a time-dependent computation and/or a computation where flux-brane annihilation needs to be taken into account.

The idea outlined above is the motivation for Paper IV. The idea would be to revisit the [KPV02] calculation, where the flux-brane annihilation channel for the [KKLT03] construction is considered, and see if the conclusion made by [KPV02] changes as a possible time-evolution of the system is included in the analysis.

Approaching this problem with a full time-dependent Ansatz would be computationally a very difficult problem – not only would all functions acquire time dependence but also currents would be induced as the charged flux-density starts moving. Assuming that these effects are small and the fall in of flux is adiabatic, one can consider the [KPV02] computation for different times where time is only measured in the amount of flux that have fallen in. More explicitly this means that the parameter that governs the magnitude of the flux density, λ, is traced through the original calculation and appropriately included in the potential, see equation (4.32). The parameter λ is then increased to simulate the flux being attracted.

As mentioned in Section 4.2.1 the F_3 has three different components however only one component is singular. As the [KPV02] is performed at the bottom of the conifold, only the component that is occupying the S^3 at the bottom remains. This is also the component that becomes singular when the backreaction is computed.
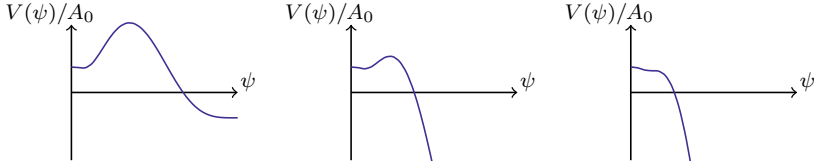
Figure 4.9. All graphs are drawn with $p/M = 3\%$ and $b_0^2 = 0.93266$. The left most picture is the [KPV02] result with $\lambda = 1$. The middle plot is after a slight build up of ISD corresponding to $\lambda = 1.3$. The right most graph corresponds to $\lambda = 1.7$ where meta-stability is lost.

The effective potential that governs the the NS5-brane movement, derived by [KPV02], acquires the following form when λ is included

$$\frac{1}{A_0} V(\psi) = \frac{1}{\pi} \sqrt{b_0^4 \sin^4 \psi + \left(\pi \frac{p}{M} - \psi + \frac{1}{2} \sin(2\psi) \right)^2} \\ - \frac{\lambda}{2\pi} (2\psi - \sin(2\psi)) \,. \tag{4.32}$$

This computation takes place at the bottom of the [KS00] conifold where only the S^3 remains. The coordinate ψ above is the direction in which the NS5-brane moves, as explained in Section 2.2.2. A_0 and b_0 are two constants, and p/M is the ratio of anti-D3-branes (p) versus the amount of background flux (M). The first term of this expression is positive and will give rise to a barrier that prevents the NS5-brane to progress over the S^3 and fully annihilate. The second term is negative, and linear in the parameter λ. This means that as λ increases, the barrier will be suppressed, see Figure 4.9.

In conclusion, singularities can arise if one ignores time dependence. The present singularity, if due to time dependence, would resolve the singularity and possibly create a new form of instability.

4.3 Non-BPS anti-M2-brane solutions and their singularities

The two systems discussed so far are of two different types. The first type, concerning the anti-D6-branes, display a singularity when the smeared sources are localised. While the second type arises when the anti-brane is placed in oppositely charged background and backreacted, e.g. anti-D3-branes in [KS00]. The same two types are present as solutions in 11D supergravity.

In Paper VI smeared non-BPS solutions are constructed, similar to those in Paper I. Also presented in Paper VI is a first attempt to es-

tablish the singularity on similar terms as in Papers II and III, that is, when the branes are localised. This approach will be presented first in the following section.

The other system that displays flux density singularities for anti-M2-branes are the backreacted branes placed in the [CGLP03] and [CGLP02] backgrounds. Furthermore, the anti-M2-branes have been studied more extensively than any other system when it comes to an M5-brane resolution. These results will also be reviewed.

4.3.1 New non-BPS anti-M2-brane solutions and their singularities

Since in 11D supergravity there only exists one 4-form, this by it self must have components that represents the field-strength associated to the M2-brane as well as the surrounding flux. The Ansatz used in Paper VI giving this is

$$G_4 = \star_{11} F_7 + F_4 + H_4 \,. \tag{4.33}$$

Here the 7-form F_7 is the field-strength and the two 4-forms F_4 and H_4 represents the surrounding flux. These are taken to obey the following relations

$$\begin{aligned} F_7 &= \mathrm{e}^X \star_8 \mathrm{d}\alpha, \\ H_4 &= \lambda \star_8 F_4. \end{aligned} \tag{4.34}$$

These are chosen in this way to be easily compared to the type II form Ansätze, see Table 4.3.

11D field	Type II analogue
F_7	F_{8-p}
F_4	F_{6-p}
H_4	H_3

Table 4.3. *A comparison between* 11D *and type II flux Ansätze.*

Just as there exists a I(A)SD condition in the type II scenarios, this flux obeys a (anti-)self-duality ((A)SD) condition when $\lambda = \genfrac{}{}{0pt}{}{+}{(-)}1$

$$(H_4 + F_4) = \genfrac{}{}{0pt}{}{+}{(-)} \star_8 (H_4 + F_4) \,. \tag{4.35}$$

New smeared solutions
Via the same observation made in Paper I, the smeared approximation reduces the Bianchi identity to the following algebraic equation

$$0 = \lambda |F_4|^2 + Q_{\mathrm{M2}} \,. \tag{4.36}$$

To solve this, i.e. to cancel the tadpole, the flux charge density have to have the opposite sign compared to that of the brane. These smeared solutions have a negatively curved world-volume, hence can be taken to be AdS$_3$. The internal space is however more interesting. If assumed to be split in two four-dimensional parts where the fluxes are proportional to the world-volume for each part, then the internal curvature of each of these spaces are given by the internal parts of Einstein's equation

$$R_{ij}^{(H)} = \frac{1}{6}(-1 + |\lambda| + 2\lambda^2)|F_4|^2 g_{ij}^{(H)} ,$$
$$R_{ij}^{(F)} = \frac{1}{6}(2 + |\lambda| - \lambda^2)|F_4|^2 g_{ij}^{(F)} . \tag{4.37}$$

Where $R_{ij}^{(A)}$ is the Ricci tensor for one part of the internal space, with the superscript referring to the section covered by A_4 flux.

One difference to the type II setups is that here λ is a free parameter (up to flux quantisation conditions), while in type II the dilaton equation fixes λ to a specific value. The total curvature of the internal space is positive and it has to be Einsteinian. However for some values of λ either one of the two parts can be negatively curved.

As an example, in the region where both parts are positive, which is $1/2 < |\lambda| < 2$, then the internal space can be taken to be $S_{(H)}^4 \times S_{(F)}^4$. This would then be very similar to the AdS$_{p+1} \times S_{(H)}^3 \times S_{(F)}^{6-p}$ solutions of Paper I.

Singularities from localisation
From experience one would expect that these solutions would develop a singularity when the source is localised. For simplification, localisation is here only considered in one direction, along one of the H_4 legs. While this might seem a crude simplification, in no other systems have partial smearing been linked to the presence of the singularity.

Also in this system it is possible to also construct a topological restriction that forces a certain profile for α. It is in fact also the case that these singularities cannot be hidden behind a horizon, similar to what was shown in Paper V.

The argument goes through in the same way as before – the field-strength potential α is constrained by

$$\text{sgn}\, \alpha'' = \text{sgn}\, \alpha \mid_{\alpha'=0} . \tag{4.38}$$

Following the same logic as was done for the anti-D6-branes, this condition implies a singularity unless there are new boundary conditions are available that avoids this no-go condition. Furthermore, a blackening factor cannot hide the singularity. In light of [AFRT13], which manage to resolve the anti-D6-brane singularity with the use of D8-branes, it

should be possible to resolve these anti-M2-brane singularities with the use of M5-branes.

4.3.2 Parallel developments regarding anti-M2-branes

As was covered in Section 2.2.2, there are two available backgrounds that have been studied with the addition of anti-M2-branes placed at the tip. These two are [CGLP03] and [CGLP02]. As was also covered in Section 2.2.2, the [CGLP03] solution has likewise been studied for a probe anti-M2-brane and its possible meta-stable transition into a M5-brane, analogous to the [KPV02] calculation of type IIB. The backreaction of the anti-M2-branes in these two backgrounds have been studied in [BGH11, Mas12b] and [GOP13] respectively.

Singularities on $V_{5,2}$

The $V_{5,2}$ manifold used in [CGLP03, BGH11, Mas12b] has a non-vanishing S^4 at the bottom of the conifold. This causes the same technical problem that one has for placing anti-D3-branes in [KS00], i.e. the source have to be smeared over the sphere at the bottom of the conifold.[8] But as have been discussed for both the anti-D6-branes as well as the anti-D3-branes, singularities seems to still be present after full localisation.

Again utilising the [BG03] perturbation technique, implicitly solved in [BGH11] in terms of nested integrals, and analytically in [Mas12b], it was found that the bacreaction of the anti-M2-branes exhibit two singular legs

$$\mathrm{d}\tau \wedge \tilde{\sigma}_{123}, \text{ and } \epsilon^{ijk}\nu \wedge \sigma_i \wedge \tilde{\sigma}_{jk}. \tag{4.39}$$

Remember that the leg of the flux that occupies the bottom S^4 is the $\nu \wedge \sigma_{123}$ leg, which is not singular. This is in contrast to the backreaction of the anti-D3-branes, where the singular leg is the one occupying the S^3. The fact that the singular leg is occupying the S^3 was one of the motivations for the interpretations made in Paper IV, where it was argued that the meta-stability of NS5-channel vanishes. More comments on the M5-brane polarisation will be made shortly.

Singularities on $\mathbb{A}_8$

The $\mathbb{A}_8$ manifold has an advantage over $V_{5,2}$ – at the bottom of the conifold the metric is topologically flat $\mathbb{R}^8$ [CGLP02]

$$\mathrm{d}s_8^2 \sim \mathrm{d}r^2 + r^2\mathrm{d}\Omega_7. \tag{4.40}$$

Hence the anti-brane placed there only has one localisation direction, as a localised point in flat space. The aim of [GOP13] is to study

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[8]In principle it could also be localised on these $S^{4/3}$ spheres, supposedly preserving the isometries of the azimuthal $S^{3/2}$ sphere, but this is beyond current techniques.

backreaction of anti-M2-branes in this background. Compared to the $V_{5,2}$ background, $\mathbb{A}_8$ makes it easier to study the backreaction under full localisation. Again, this analysis is performed under the linearised [BG03] approach. [GOP13] verifies the presence of the singularity once more.

In parallel to [GOP13], [CGH13] performed the same calculation and also arrives at the conclusion that the singularity is present.

M5 polarisation and other problems
Analogous to [KPV02], [KP11] describes the meta-stability of placing anti-M2-branes at the bottom of the [CGLP03] background under the assumption that the background allows for an M5 polarisation channel. The very recent work of [BGKM14] extends the analysis of the anti-M2-branes significantly. First of all they show that the singularity is present beyond the perturbative expansion. Secondly they study two polarisation channels – a [KP11]-channel that governs the decay of anti-M2-branes into M2-branes and a "transversal"-channel that is suggested to cut off the singularity. The M5-brane in the [KP11]-channel wraps a S^3 inside the finite S^4, while the transversal-channel wraps the shrinking S^3.

The parameters that enter the [KP11]-channel polarisation potential are, in [BGKM14], left undetermined and do allow, for a certain range of said parameters, a minima for the M5-brane to reside at. The new thing in this approach is that close to the anti-M2-branes a local $AdS_4 \times S^7$ throat develops. Within the local region where this throat resides, a polarisation minimum do develop.

While the [KP11]-channel does seem to be allowed, the transversal-channel is absent. This channel is suggested to provide a cut off for the radial profile of the flux energy density and resolve the singularity. The absence of this channel agrees with the results for the anti-D6- and anti-D3-brane polarisations of [BJKVRWZ12] and [BGKM13b] respectively.

Even though the [KP11]-channel is present, the whole UV to IR boundary conditions have not been mapped together, and hence the presence of the singularity is not necessarily resolved. The [BGKM14] are very recent and it is too early to make a definite conclusion from these results.

Before leaving the anti-M2-branes, it should be noted that the result of Paper VI is valid for a more or less general topology. Whether the assumed structure of the internal space and the positioning of the fluxes in Paper VI actually is general enough to be applied in the $V_{5,2}$ or $\mathbb{A}_8$ geometry, has to be checked explicitly. If not, the techniques employed in Paper VI might be possible to generalise to also these cases.

4.4 Short summary

To summarise the work included in this thesis and the related papers reviewed above a short chronological map will be presented here.

1. The flux charge density singularity is discovered for anti-D3-branes in the [KS00] background by [MSS11] via the means of a perturbative expansion [BG03] in terms of p/M.
2. The singularity is verified in [BGH10] via similar means.
3. Paper I presents new smeared anti-D6-brane solutions.
4. [BGH11] shows that a singularity arise for anti-M2-branes in the [CGLP03] background, as a part of a perturbative expansion.
5. An argument is presented in [Dym11] which says that the singularity would vanish beyond the perturbative expansion.
6. The existence of an anti-D3-brane singularity is strengthened by analytically solving the nested integrals of [BGH10] in [BGGHM13].
7. A topological argument for a singularity is presented in Paper II for the localised solutions of Paper I.
8. The work of [BGH11] is strengthened by [Mas12b] through analytically solving nested integrals.
9. The topological argument of Paper II is shown to be unavoidable by characterising all boundary conditions in Paper III. This is the first results that go beyond any perturbative expansion and partial smearing.
10. An adiabatic time-dependent resolution of the anti-D3-brane singularity is presented in Paper IV. This is the first interpretation that leads to a resolution of the singularity. Furthermore it shows that the singularity can lead to a new instability.
11. [Mas12a] presents an argument against the possibility of resolving the singularity as suggested in [Dym11].
12. It is shown that the anti-D6-branes does not allow for a D8-polarisation that could have resolved the singularity by [BJKVRWZ12] in a non-compact setup.
13. By [BGKM13a] the anti-D3-brane singularity is shown to be present beyond any perturbative expansion.
14. It is shown in [BGKM13b] that the anti-D3-brane singularities cannot be resolved via a D5-brane polarisation.
15. Numerical solutions are presented in [BBD13] that show that the anti-D3-brane singularities does cannot be cloaked by a horizon.
16. A new approach is presented in [GJZ13] that goes beyond any perturbative and partial smearing techniques and shows the existence of the singularity.

17. Paper V extends the work of [BBD13] and argues analytically that the singularity does not have a string theory resolution as per the criterion of [Gub00].

18. The singularity is established for anti-M2-branes one a new background. This is done trough a perturbative expansion, however the background allows for full localisation [GOP13].

19. [CGH13] describes the same system as [GOP13] and also encounters the singularity.

20. Paper VI presents new smeared (anti-)M2-brane solutions and an accompanying topological argument as to why these solutions would develop singularities and cannot be cloaked via the means of adding a horizon.

21. [AFRT13] presents a numerical example in which the anti-D6-brane singularity is resolved by having a D6-brane on the opposite side, with D8-branes stacks preserving two BPS regions.

22. In [BG13] curvature to the external space-time is added to see if the singularity can be avoided. They conclude that the singularity does not get resolved by this procedure.

23. [BGKM14] calculates the M5-brane polarisation channels and discovers a previously unknown tachyonic instability.

24. By revisiting the [KKLT03] calculation, keeping track of the singularity and taking care of warping effects, [Jun14] finds that the unresolved singularity might under a certain choice of boundary conditions not change the result.

25. The supersymmetric D8-brane polarisation potential is studied in [JSZ14], who also finds numerical evidence for non-supersymmetric polarisations.

This record shows, with very little doubt, that the singularity is present and not due to partial smearing nor some perturbative expansion. For the anti-D6-branes the singularity have been shown to be possible to resolve through the polarisation to a D8-brane. However, as noted by [AFRT13] and [JSZ14], the polarisation channel is present in these cases because of the AdS curvature. For the anti-D3-brane singularity, the external curvature available would be related to the *small*, almost vanishing, positive cosmological constant. Hence it is difficult to expect that a similar polarisation channel would exist for the anti-D3-branes for the same reasons.

One question that remains is whether the polarisation channel shown to be present in [BGKM14] is possible to resolve the singularity or not. If this polarisation do not resolve the singularity, which the result of Paper V would suggest (because the singularity fails the [Gub00] criterion), the interpretation made in Paper IV is still viable. Even if the polarisation is there and do resolve the singularity, Paper IV could still imply an

Time (arXiv)	D6	D3	M2
23 Oct 2009		[MSS11]	
17 Dec 2009		[BGH10]	
9 Sep 2010	Paper I		
9 Nov 2010			[BGH11]
8 Feb 2011		[Dym11]	
11 Feb 2011		[BGGHM13]	
24 May 2011	Paper II		
11 Oct 2011			[Mas12b]
10 Nov 2011	Paper III		
6 Feb 2012		Paper IV	
16 Feb 2012		[Mas12a]	
8 May 2012	[BJKVRWZ12]		
27 Jun 2012		[BGKM13a]	
19 Dec 2012		[BGKM13b]	
20 Dec 2012		[BBD13]	
23 Jan 2013	[GJZ13]		
29 Jan 2013	Paper V		
7 Mar 2013			[GOP13]
11 Mar 2013			[CGH13]
10 Sep 2013			Paper VI
11 Sep 2013	[AFRT13]		
4 Oct 2013		[BG13]	
10 Feb 2014			[BGKM14]
19 Feb 2014		[Jun14]	
25 Feb 2014	[JSZ14]		

Table 4.4. *This table summarises the chronological and thematic progression of the anti-brane program.*

instability of the [KKLT03] vacuum, since only a slight elevation of the flux energy density is needed for the barrier of decay to vanish.

Although no definite conclusion can be drawn at this point, these investigations have taught us new lessons about the need for caution when it comes to the study of supersymmetry breaking models.

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Svensk sammanfattning

En viktig fråga som dagens teoretiska fysik står inför är hur de fyra fundamentala krafterna skall sammanbindas inom ett teoretiskt ramverk. Det aktuella problemet handlar om att tre utav krafterna har ett gemensamt ramverk för att beskrivas medans den fjärde står för sig själv. De tre förenade krafterna kallas för *Stark*, *Elektromagnetisk* och *Svag*. Dessa tre krafter har en god beskrivning på korta avstånd och beskriver väldigt korrekt vårt mikroskopiska universum. De beskriver t.ex. hur en proton håller ihop (Stark), hur fotoner (ljuspartiklar) interagerar (Elektromagnetisk) och hur radioaktivt sönderfall sker (Svag). Det teoretiska ramverket som beskriver dessa krafter är en så kallad *kvantfältteori* med tillhörande lokala (Gauge) symmetrier, ofta kallad *Yang-Mills teori*. Ramverket som förenar dessa tre krafter kallas för *Standardmodellen*. Den fjärde kraften är *Gravitationen* som beskriver allt från planetbanor till universums kosmologiska utveckling – från Big Bang till idag samt möjliga utfall för vårt universums framtid. Det teoretiska ramverket som beskriver gravitationen är den *allmänna relativitetsteorin*. Till skillnad från de övriga tre krafterna är den allmänna relativitetsteorin inte en kvantiserad teori, vilket betyder att den inte kan användas för att beskriva gravitation på mikroskopiska avstånd. Standardmodellen beskriver alltså det lilla och den allmänna relativitetsteorin det stora i universum. Som nämndes ovan består problemet av hur dessa två beskrivningar skall sammanbindas.

En möjlig lösning till detta problem skulle kunna vara *strängteorin*. Grundidén inom strängteorin är att de minsta beståndsdelarna i vårt mikroskopiska universum inte är punkter utan strängar. Strängar är endimensionella objekt som kan vara öppna (med ändpunkter) eller slutna (utan ändpunkter). Även om strängteorin löser ett av problemen, den sammanbinder Yang-Mills teorier (öppna strängar) med den allmänna relativitetsteorin (slutna strängar), så skapar den andra. Enligt strängteorin är universum tiodimensionellt, jämfört med de fyra som vi kan observera (höjd, bredd, djup och tid). Detta innebär att antalet dimensioner bör reduceras på ett eller annat sätt. Ett annat begrepp som är viktigt för strängteori är *supersymmetri*. Supersymmetri är en symmetri mellan kraftförmedlarpartiklarna (bosoner) och materiapartiklarna (fermioner) i en teori. Även om supersymmetri finns i teorin så har dess tillämpningar för våra fundamentala partiklar, det vill säga supersymmetriska utökningar av standardmodellen, inte kunnat observerats än. Ytterligare ett problem för strängteorin är att det vanliga

grundtillståndet hos strängteoretiska modeller har en *negativ* kosmologisk konstant. Den kosmologiska konstanten beskriver universums expansion och observationer överensstämmer med en *positiv* kosmologisk konstant, det vill säga en accelererad expansion av vårt universum. Det betyder att det krävs en hel del arbete för att konstruera realistiska modeller.

En del av de problemen som vissa strängteoretiker arbetar med kan beskrivas som "hur skall antalet dimensioner reduceras?", "hur bryts supersymmetri?" och "kan strängteori beskriva den kosmologiska utveckling av vårt universum som vi observerar?". Den mest grundläggande motivationen för denna avhandling är just dessa frågor.

Ämnet för avhandlingen är att gå in mer i detalj och studera just en konstruktion som reducerar de tio dimensionerna hos strängteorin till fyra, bryter supersymmetri och beskriver en positiv kosmologisk konstant. Denna konstruktion fungerar på så sätt att de sex extra dimensionerna väljs att vara mycket små, och därför inte direkt tillgängliga för våra experiment, detta kallas *kompaktifiering*. I dessa extra dimensioner läggs sedan ett *bran* (eng. brane). Ett bran är ett objekt som har flera dimensioner. Namnet *bran* kommer från det engelska ordet för membran (eng. membrane), som är ett tvådimensionellt objekt. Ett bran kan dock ha vilken dimension som helst, ända upp till tio. Branet som används i just denna konstruktion är ett så kallat *anti*-bran. Detta betyder att det har motsatta egenskaper jämtemot övriga beståndsdelar i konstruktionen som studeras här. Det innebär att anti-branet står för supersymmetribrottet som är nödvändigt och även lyfter den kosmologiska konstanten till ett positivt värde.

Avhandlingen har därför som mål att studera denna typ av anti-bran. Detta görs dels ifrån ett perspektiv av anti-D6-bran (sex rumsdimensioner), som inte har en fenomenologiska tillämpning eftersom den slutgiltiga rumtiden är sjudimensionell. De anti-D6-bran som studeras här är helt nya system som presenteras i denna avhandling. Såväl kända anti-D3- och anti-M2-bran som nya anti-M2-bran system dyker också upp. Alla dessa resultat är från beräkningar inom en approximation som kallas för *supergravitation*, vilket är en gravitationell teori med tillhörande supersymmetri. Supergravitation är en lågenergiapproximation av strängteorin och ger möjlighet till enklare beräkningar.

De övriga beståndsdelar som är en del av dessa konstruktioner är *flöden* (kan liknas med elektromagnetiska flöden) som branen elektromagnetiskt och gravitationellt interagerar med. Dessa konstruktioner brukar kallas *flödeskompaktifieringar*. Det som diskuteras i avhandlingen är en typ av singularitet i flödestätheten. Denna singularitet uppstår i närheten av ett bran som har motsatta egenskaper jämfört med flödet som det placeras i. En stor vikt läggs vid att etablera att singulariteten faktiskt är där och inte bara är en biprodukt av de beräkningsverktyg

som används, så som störningsteorier och andra approximativa tillvägagångssätt.

Det visar sig att det råder lite tvivel, åtminstone i publicerad litteratur, om denna singularitets existens. Frågan handlar därefter om det finns sätt att undvika denna singularitet. Det vill säga: finns det något sätt att förklara varför singulariteten uppstår och går det att tolka den? Svaret kan bestämmas genom vissa etablerade test som skulle ge indikationer på att det finns sätt att upplösa singulariteten.

Det finns två sådana test. Det första handlar om att bran av en viss rumsdimension, p, kan svälla upp, eller *polariseras*, till ett högredimensionellt bran, med rumsdimension $p+2$, och omge det ursprungliga p-branet. Detta skulle då kunna begränsa värdet på det omgivande flödet och upplösa singulariteten. Det har dock visats i två fall att ingen sådan polarisering kan ske. Ett andra test är att se om det går att gömma singulariteten bakom en händelsehorisont. En händelsehorisont är ett begrepp som kommer från svarta hål, där händelsehorisonten är ett område som omger det svarta hålet och förhindrar ljus att komma ut. Om en singularitet tillåts gömmas bakom en horisont menar man att singulariteten är under kontroll. För andra singulariteter har detta visat sig vara ett fungerande test. Ett av resultaten som presenteras i denna avhandling visar dock på att det inte är möjligt att placera just denna singulariteten bakom en händelsehorisont. Båda dessa test misslyckas alltså och frågan hurvida singulariteten är under kontroll eller ej kvarstår.

Utöver detta presenteras i avhandlingen även en tolkning av singulariteten. Ett vanligt bran utsätter flödet för en attraktiv gravitationskraft och en frånstötande elektromagnetisk kraft, på ett sådant sätt att sammanlagt är kraften emellan bran och flöde frånvarande. *Anti*-branet har motsatt elektromagnetisk laddning och har därför en total attraktiv kraft på det närliggande flödet. Krafter ger upphov till acceleration och därmed rörelse i systemet, men i alla fall där denna singularitet har studerats har statiska, det vill säga tidsoberoende, beskrivningar använts. Detta medför en möjligt tolkning av singulariteten – den uppstår för att systemet inte har en möjlighet att utvecklas med tiden.

Dock så får denna tolkning oönskade konsekvenser. Man kan beräkna sönderfallet av anti-bran och det visar sig vara beroende av storleken på det omgivande flödet – större flöde leder till snabbare sönderfall. Allteftersom anti-branet drar till sig flöde från omgivningen så sönderfaller det, möjligen till den grad att det så småningom inte finns några anti-bran kvar. Detta skulle betyda att supersymmetrin återupprättas och den kosmologiska konstanten återigen blir negativ.

Att beskriva en positiv kosmologisk konstant i strängteorin är ett väldigt intressant och betydelsefullt problem. De arbeten som inkluderas i denna avhandling ger inte bara nya insikter om svårigheterna för att

bygga realistiska modeller utan även mer förståelse för hur olika objekt inom strängteori beter sig tillsammans.

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