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Oh, wait, O8 de Sitter may be unstable!

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ABSTRACT: We analyze the stability of four-dimensional de Sitter vacua constructed by compactifying massive Type IIA supergravity in the presence of two O8 sources [1]. When embedded in String Theory the first source has a clear interpretation as an O8₋ plane, but the second one could correspond to either an O8₊ plane or to an O8₋ plane with 16 D8-branes on top. We find that this latter solution has a tachyonic instability, corresponding to the D8 branes moving away from the O8₋ plane. We comment on the possible ways of distinguishing between these sources.

KEYWORDS: Superstring Vacua, D-branes, Flux compactifications

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1 Introduction

One of the most challenging open problems in String Theory is determining whether it has metastable de Sitter vacua. An efficient way to construct such vacua is to use an effective four-dimensional theory that incorporates String-Theory ingredients [2]. However, the embedding of these ingredients in the full String Theory and their interactions therein are nontrivial, and have opened a rich debate about the validity of these constructions [3–12]. In parallel to these top-down investigations, there also exist bottom-up arguments and conjectures suggesting that metastable String-Theory compactifications with a positive cosmological constant cannot be constructed [13–16].

A way to bypass the complications inherent to the construction of de Sitter solutions using effective four-dimensional theories and to verify the validity of these conjectures is to work directly in ten (or eleven) dimensions. Upon restricting to classical contributions to the ten-dimensional stress-energy tensor, the negative-energy sources required to evade no-go theorems [17] can be realized in String Theory as orientifold planes. As a first approximation, one can consider the orientifold planes to be smeared along (some of) the internal directions. Even if, by definition, orientifold planes sit at fixed loci of some involution, this approximation may be justified and serve as a guide to the construction of more complete solutions [18–21].

Solutions with explicitly localized orientifold sources are on much firmer physical ground, but constructing them is more challenging since it requires solving the field equations point-wise on the compactification manifold, and not just the integrated (averaged) version. Furthermore, localized orientifold planes source singular supergravity solutions¹

¹See [22, section 2] for a recent review of orientifold singularities in supergravity.

and even when the leading-order divergence of the supergravity fields match those expected for the orientifold plane, there could still be subleading divergences that may signal deeper problems.²

Hence, *any* supergravity solution with singularities coming from localized orientifold sources has to be understood as a good approximation of a would-be corresponding full-fledged string-theory solution only away from these singular loci. When supersymmetry is present there exist other methods to assess the validity of the supergravity approximations, but for de Sitter solutions these methods do not work.

In [1] and in [22], Córdova, Tomasiello and one of the authors have constructed de Sitter backgrounds with localized O8 and O8-O6 sources respectively, which have the properties discussed above. We are going to refer to them as CDLT₁ and CDLT₂.

The purpose of this paper is to analyze the solutions of [1] in String Theory (CDLT₁), and to see whether they suffer from any instabilities. Despite the relative simplicity of the CDLT backgrounds, looking for tachyons in the full Kaluza-Klein spectrum is a prohibitive task, mostly because of the presence of a warp factor with a non-trivial profile.³

As we are going to review in section 2, in this class of solutions the metric takes the form

$$ds_{\text{CDLT}_1}^2 = e^{2W} ds_{\text{dS}_4}^2 + e^{-2W} (dz^2 + e^{2U_2} ds_{M_2}^2 + e^{2U_3} ds_{M_3}^2), \quad (1.1)$$

where the warp factors W and U_i (as well as the dilaton) depend on the coordinate z parametrizing a circle and M_i are compact i -dimensional Einstein spaces; at least one between M_2 and M_3 must have negative scalar curvature. There is a \mathbb{Z}_2 symmetry acting on the circle with the two fixed loci at $z = \{0, z_0\}$, where two O8 sources are located. More precisely, a source with the charges of an O8₊ plane is at $z = 0$ and a source with the charges of an O8₋ plane is at the other fixed point (see figure 1).⁴ The existence of this de Sitter solution crucially depends on the presence of sources with the charges of these O8_± planes, whose negative tension violates the standard energy conditions that are incompatible with a de Sitter solution [17].

As remarked above, near the orientifold sources the supergravity approximation breaks down. However, at the source with O8₊ (positive) mass and charge, the singularity is relatively mild: approaching it, all the physical quantities remain finite, with only a finite discontinuity in their derivative. Hence, this singularity is of the same type one encounters near normal D8 branes.

This is not the situation for the O8₋ source, which has negative mass and charge. Very near this source the solution is both strongly coupled and strongly curved. As we will review below, the region where the supergravity approximation is not valid anymore can

²This happens for example with anti-D3 branes [5].

³Computations of the full KK spectrum of warped non-supersymmetric AdS compactifications have only been feasible recently and in very specific situations by applying the techniques of exceptional field theory [23–27]. We therefore look for instabilities in the open string sector, which can be revealed using probe D-branes.

⁴We use the standard convention where an O8₋, with negative charge and tension, is the standard orientifold giving rise to an SO(2n) gauge group on a stack of n D8-branes on top of it, while O8₊, with positive charge and tension equal to that of eight D8-branes, gives rise to an Sp(2n) gauge group.

be made parametrically small but never made to disappear completely. Furthermore, as one approaches the $O8_-$ plane, the subleading behavior of the supergravity fields⁵ deviates from the flat-space one [1, 22, 29], and it is unclear whether they can be trusted at the supergravity level. Besides these conceptual issues, there are also technical issues with the numerical solution near the $O8_-$ source, which we will discuss in section 3.1. Hence, our investigation will mostly steer away from the $O8_-$ source, and focus on the $O8$ source with positive charge and mass, where supergravity is more trustworthy.

We find that at the supergravity level the non-supersymmetric $CDLT_1$ background with mobile D8-branes suffers from an instability that corresponds to the emission of a D8-brane from this $O8$ source. This instability is similar to the brane-jet instability [30] of non-supersymmetric AdS_4 solutions [31, 32], except that the brane does not come out from behind a horizon but rather from a supergravity singularity. An analogous mechanism is also responsible for the decay of an infinite class of non-supersymmetric AdS_7 solutions, in which case D6-brane sources polarize into D8-branes, destabilizing the backgrounds [33].

The key question raised by our investigation is whether there is any physics that may prevent such a D8-brane from emerging from the $O8$ source with positive charge and mass. To address this question it is important to point out that there are two possibilities for the string-theory object corresponding to this source.

The first possibility is that it corresponds to an $O8_-$ plane with 16 D8-branes on top. If this possibility is realized, the $CDLT_1$ background can be thought as just the compactified version of the very familiar configuration of two $O8_-$ planes with 16 D8-branes, that one obtains by T-dualizing Type-I String Theory. If so, the instability we found corresponds simply to one of the D8-branes moving away from the $O8$ source and this indicates that the $CDLT_1$ solution is unstable when embedded in String Theory.⁶

The second possibility is that $O8$ source with positive charge and mass corresponds to an $O8_+$ plane. If this possibility is realized the $CDLT_1$ background can be thought of as the compactified version of the $O8_+-O8_-$ configuration that one obtains by T-dualizing the Dabholkar-Park background [36–38]. Furthermore, in this realization, it appears rather unlikely that the $O8_+$ plane could emit a D8 brane, especially because there is no String-Theory object in which it could decay, at least perturbatively.

Unfortunately supergravity is unable to distinguish between the two interpretations of the orientifold source with positive charge. Both the $O8_+$ plane and the $O8_-$ plane with 16D8 branes have the same charge and tension and all the bulk supergravity fields behave identically when approaching the two singularities. These two configurations do appear distinguishable in the full String Theory, taking into account also the open string sector. It would be interesting to understand if there is any way in which this difference could manifest itself in the low-energy supergravity description.

⁵At higher codimension these could be divergent, much like it happens near anti-D3 brane singularities [28].

⁶Such instabilities exist also in related models, as mentioned in [1, Footnote 6]. Similar related models with pairs of O_p -planes and mobile D_p -branes have also been studied from an effective field theory point of view [34]. Some of these models display other forms of perturbative instabilities [35].

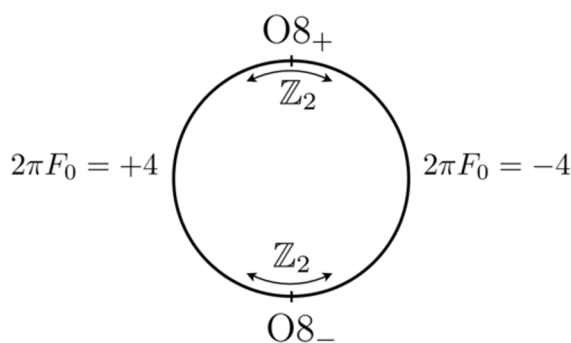


Figure 1. A picture taken from [1] describing the internal topology of the CDLT₁ solutions. Here $l_s = 1$.

The paper is organized as follows. In section 2 we review the CDLT₁ solution. In section 3 we discuss the behavior of a D8 probes in the CDLT₁ background: in particular, we provide the bosonic and fermionic action describing the fluctuations of a stack of k probes and we comment on its stability when an O8 source is approached. We conclude in section 4 summarizing the results and suggesting some candidate ways to distinguish between the two possible O8 singularities.

2 Review of CDLT₁ solutions

In this section we review the main properties of the CDLT₁ solution [1] which only contain O8 sources. More details, including the full set of equations of motion and an analysis of boundary conditions can be found in [22].

The metric of the most general CDLT₁ solution is:

$$ds^2_{\text{CDLT}_1} = e^{2W} ds^2_{\text{dS}_4} + e^{-2W} (dz^2 + e^{2U_2} ds^2_{M_2} + e^{2U_3} ds^2_{M_3}), \tag{2.1}$$

where $M_2(M_3)$ is an Einstein manifold of dimension two(three) and z is a coordinate parameterizing a circle $z \in [0, 2z_0]$ with $2z_0 \sim 0$. All the functions depend only on the circle coordinate. The solution has a Romans mass F_0 and four-form RR field strength

$$F_0 = \frac{\pm 4}{2\pi l_s}, \quad F_4 = f_4 e^{-6W+3U_3-2U_2} dz \wedge \text{vol}_{M_3}. \tag{2.2}$$

Here the two values of F_0 correspond to the two halves of the circle (see figure 1), on which there is a \mathbb{Z}_2 involution acting as the antipodal identification $z \sim -z$. The two fixed points at $z = 0$ and $z = z_0$, are two O8-sources.

The integrated Bianchi identities require the total charge in the internal space to vanish, and thus the two orientifold sources have opposite charge. Similarly, the existence of a compact solution requires also the total tension to vanish, and so the charges of these two objects are those of an O8₊ plane and an O8₋ plane. Evaluating the equations of motion on top of the source with an O8₊ plane gives a constraint on the signs of the curvatures of the various Einstein factors in (2.1). In particular a positive four-dimensional

cosmological constant (dS) is allowed only if at least one of the internal Einstein spaces has negative curvature.

There is a discrete rescaling that can be done on the solution (2.1)–(2.2). Flux quantization requires that the integral of F_4 be related to an integer, N :

$$\frac{1}{(2\pi l_s)^3} \int_{S^1 \times M_3} F_4 = N. \tag{2.3}$$

Starting from a solution where $(2\pi l_s)^{-3} \int_{S^1 \times M_3} F_4 = 1$, it is possible to generate other solutions with quantized flux (2.3) by rescaling⁷

$$g_{\mu\nu} \rightarrow N^{1/2} g_{\mu\nu}, \quad e^\phi \rightarrow N^{-1/4} e^\phi. \tag{2.4}$$

When the total flux, N , becomes large, the curvature becomes parametrically small, and the solution is weakly coupled. This rescaling acts on the warp factors in the metric (2.1) as

$$e^{2W} \rightarrow N^{1/2} e^{2W}, \quad e^{2U_i} \rightarrow N e^{2U_i}, \quad z \rightarrow N^{1/2} z \tag{2.5}$$

increasing the size of the circle. The solutions to the full Einstein equations were obtained numerically, but analytic expansion around the limiting points are available.⁸

As discussed in the Introduction, near the orientifold sources the solution is generically singular. Near the source with positive charge the singularity is mild, similar to the one encounters when approaching D8 branes. However, since the plane with O8₋ charge has negative mass and negative charge, the metric and the dilaton behave as

$$ds_{O8_-}^2 \approx \frac{1}{\sqrt{r}} \left(ds_{dS_4}^2 + ds_{M_2}^2 + ds_{M_3}^2 \right) + \sqrt{r} dr^2, \quad e^\phi \approx r^{-5/4}, \tag{2.6}$$

where $r \equiv z_0 - z$. This behavior is the same as near a source with negative D8 charge and tension (such as an O8₋ plane) in flat space or in AdS [40–44]. This can be understood as the limit $a \rightarrow 0$ of the harmonic function $e^{-4W} \equiv H = a + b|r|$, which solves the Einstein equations for an O8₋ plane with flat worldvolume. The discontinuity of the first derivative of this function at $r = 0$ is sourced by a δ -function source. In our more complicated solution the warping is not a simple harmonic function anymore, but the same δ -function is responsible for the same asymptotic behavior near $r \sim 0$ [22, section 4.4].

It is worth noticing that there exists a simple subclass obtained by tasking the two internal Einstein spaces with to have the same (negative) Einstein constant. In these solutions the warp factors are equal $U_2 = U_3 \equiv U_5$ and the four-form RR field strength F_4 vanishes. The 5 dimensional internal space is Einstein, and the metric of the form

$$ds_{10}^2 = e^{2W} ds_{dS_4}^2 + e^{-2W} (dz^2 + e^{2U_5} ds_{M_5}^2). \tag{2.7}$$

3 Probing the CDLT₁ backgrounds with D8-branes

In this section we investigate the physics of probe D8-branes in the background (2.1).

⁷The rescaling constant, c , in [1] is related to N via $N = e^{4c}$.

⁸One can construct analytic solutions valid everywhere as a formal expansion in Λ [39].

3.1 The bosonic and fermionic action

The bosonic action of the D8 branes has a DBI and a Wess-Zumino part. From the perspective of the worldvolume D8-brane theory the interval position z_{D8} of a single D8-brane is U(1) worldvolume scalar field. Besides this scalar field, the worldvolume theory has a gauge field. For k D8-branes, the gauge group of the worldvolume theory becomes $SU(k)$ and the action is⁹

$$S_{\text{D8}}^{(B)} = -\tau_{\text{D8}} \int d^9\xi e^{-\phi} \sqrt{-\det(P[g] + \lambda \mathcal{F})} \pm \tau_{\text{D8}} \int P[C \wedge e^{\lambda \mathcal{F}}], \quad (3.1)$$

where \mathcal{F} is the worldvolume field strength and we already assumed the NSNS B -field to be vanishing. In the previous expression, $\tau_{\text{D8}} = 2\pi/(2\pi l_s)^9$ is the brane tension while $\lambda = 2\pi l_s^2$ is the worldvolume gauge coupling. Each function of z can be expanded in Taylor series around a reference point, z_0 , and each polynomial term in the series corresponds to a polynomial interaction of the worldvolume adjoint scalar field, Φ , describing the positions of the D8 branes:

$$f(z) = \sum_{i=0}^{\infty} f_{\text{D8}}^{(i)} (z - z_{\text{D8}})^i \Rightarrow \sum_{i=0}^{\infty} f_{\text{D8}}^{(i)} \lambda^i \text{Tr} \Phi^i \quad \text{with} \quad f_{\text{D8}}^{(i)} = \partial_z^i f(z)|_{z_{\text{D8}}}. \quad (3.2)$$

Neglecting for the moment the contribution of F_4 , expanding (3.1) at second order in λ gives:

$$S_{\text{D8}}^{(B)} = -\tau_{\text{D8}} \lambda^2 \text{Tr} \int d^9\xi \sqrt{-g_{(0)}} \left\{ \frac{1}{\lambda^2} (\Delta_{\text{D8}} + c_9) + \frac{1}{\lambda} (\partial_z \Delta_{\text{D8}} + f_{10}) \Phi + \right. \\ \left. + \frac{1}{2} e^{-2W} \Delta_{\text{D8}} D_\alpha \Phi D^\alpha \Phi + \frac{1}{4} \Delta_{\text{D8}} \mathcal{F}_{\alpha\beta} \mathcal{F}^{\alpha\beta} + \right. \\ \left. + \frac{1}{2} (\partial_z^2 \Delta_{\text{D8}} + \partial_z f_{10}) \Phi^2 \right\} + O(\lambda^3), \quad (3.3)$$

where we have defined

$$\Delta_{\text{D8}} \equiv e^{3U_3 + 2U_2 - W - \phi}, \quad F_{10} \equiv f_{10} dz \wedge \text{vol}_{\text{D8}} = d(c_9 \text{vol}_{\text{D8}}), \quad f_{10} \equiv F_0 \Delta_{\text{D8}} e^{\phi - W} \quad (3.4)$$

and $g_{(0)}$ is the metric on $\text{AdS}_4 \times M_5$ in absence of warping. Observe that worldvolume indices are raised and lowered with the worldvolume metric $g_{\alpha\beta}$ (including the warping factors) and $D_\alpha \equiv \partial_\alpha - i[\mathcal{A}_\alpha, \cdot]$.

It is convenient to use canonically normalized scalar fields, which we can obtain by defining:

$$\tilde{\Phi} \equiv \frac{1}{g_{\text{YM}}} \Phi, \quad g_{\text{YM}}^2 \equiv \frac{e^{-3U_3 - 2U_2 + \phi}|_{z_{\text{D8}}}}{\lambda^2 \tau_{\text{D8}}} \quad (3.5)$$

and to use a normalized worldvolume metric:

$$(\gamma_{\text{D8}})_{\alpha\beta} = \left(e^{-2W} g_{\alpha\beta} \right) \Big|_{z_{\text{D8}}}. \quad (3.6)$$

⁹The full gauge group is $U(k)$ but the Abelian $U(1)$ center-of-mass degrees of freedom decouple from the $SU(k)$ non-Abelian sector.

Performing the previous redefinitions and renaming for clarity $\tilde{\Phi} \rightarrow \Phi$, the action (3.3) becomes:

$$S_{\text{D8}}^{(B)} = -\text{Tr} \int d^9 \xi \sqrt{-\gamma_{\text{D8}}} \left\{ -\frac{e^{4W} F_{\text{D8}}}{\Delta_{\text{D8}} g_{\text{YM}}} \Phi + \frac{1}{2} D_\alpha \Phi D^\alpha \Phi + \frac{1}{4 g_{\text{YM}}^2} \mathcal{F}_{\alpha\beta} \mathcal{F}^{\alpha\beta} - \frac{m_B^2}{2} \Phi^2 \right\}, \quad (3.7)$$

where the worldvolume indices are now raised and lowered using γ_{D8} . The parameters F_{D8} and m_B can be identified as the force acting on the stack of brane and as the mass of the worldvolume bosonic field, and are given by:

$$m_B^2 = -e^{4W} \Delta_{\text{D8}}^{-1} (\partial_z^2 \Delta_{\text{D8}} + \partial_z f_{10}) \quad (3.8)$$

$$F_{\text{D8}} = -\lambda \tau_{\text{D8}} \Delta_{\text{D8}} (\partial_z \log \Delta_{\text{D8}} + F_0 e^{\phi-W}). \quad (3.9)$$

Let us stress that the effective Yang-Mills coupling g_{YM} , the boson mass and the force are finite on the whole interval in the CDLT₁ backgrounds. When the four-form is non-vanishing, one should also take into account the topological term coming from the WZ term:

$$\delta S_{\text{D8}}^{(B)} = -\tau_{\text{D8}} \lambda^2 \int \text{Tr} (C_5 \wedge \mathcal{F} \wedge \mathcal{F}), \quad (3.10)$$

where C_5 is the potential for the dual six-form $F_6 = -\star F_4 = -f_4 \text{vol}_{M_2} \wedge \text{vol}_{\text{dS}_4}$.

The fermionic action can be also computed. The details of the computation are given in appendix A. The final result is:

$$S_{\text{D8}}^{(F)} = \frac{1}{2} \int d^9 \xi \sqrt{\gamma_{\text{D8}}} \bar{\chi} \{ i \gamma^\alpha \nabla_\alpha \chi + m_F \chi + [A_\alpha, \gamma^\alpha \chi] - i g_{\text{YM}} [\Phi, \chi] \}, \quad (3.11)$$

where χ is a nine-dimensional adjoint (Majorana) spinor with mass

$$m_F = \frac{e^{2W}}{2} (\partial_z \log \Delta_{\text{D8}} + e^{\phi-W} F_0). \quad (3.12)$$

Observe that the Wess-Zumino coupling in (3.11) exactly equals the Yang-Mills coupling. In appendix A we also show that the fermionic action is not affected by the presence of the four-form flux.

3.2 Approaching the singularities

The dynamics of probe D-branes can unveil the presence of pathologies of a given background. Quite often, such pathologies manifest themselves as perturbative instabilities or non-perturbative decay channels [30, 45–49].

The orientifold singularities of the CDLT₁ solutions can be explored using D8 branes. First, we observe that, even if the probe approaches a strongly coupled (and curved) region, the unique coupling entering the bosonic and fermionic action, g_{YM} , remains finite and can be made arbitrarily small using the rescaling symmetry (2.4)–(2.5), which acts as

$$g_{\text{YM}} \rightarrow \frac{g_{\text{YM}}}{N^{3/2}}. \quad (3.13)$$

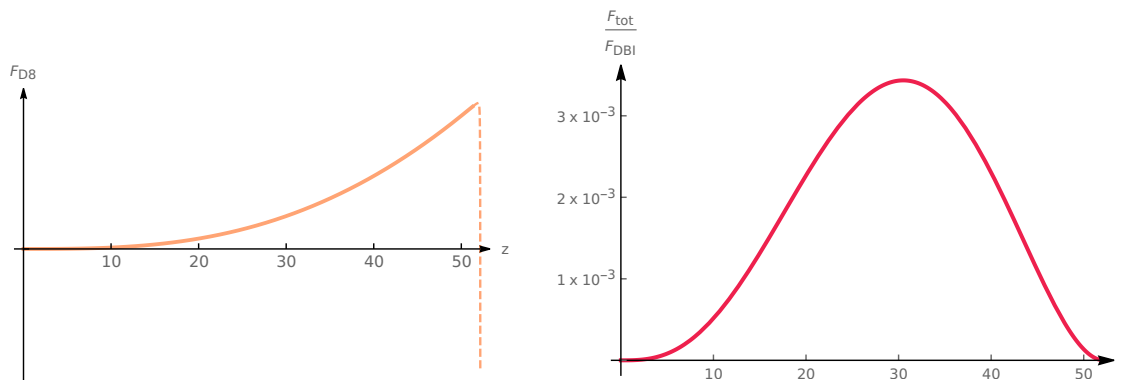


Figure 2. The left figure (orange) gives the total force on a probe D8, computed as the sum of the DBI and WZ forces. As z approaches $z_0 \sim 52$, the force drops and appears to reach a non-vanishing negative value (dashed), but we believe that this residual result is most likely a numerical artifact from the cancellation of the two diverging contributions. This can be seen from the right (red) figure, which gives the ratio between the total force and the gravitational attraction: $\frac{F_{\text{tot}}}{F_{\text{DBI}}}$. As explained in the main text, the behavior near $z \sim z_0$ has to be taken with a grain of salt, since corrections to the probe action are not negligible in the strongly-coupled region.

However, in order for the actions (3.7)–(3.11) to provide a meaningful effective description of the D8 worldvolume dynamics, the force F_{D8} should be vanishing. Evaluating (3.9) numerically for the CDLT₁ background,¹⁰ we find the force plotted in figure (2).

From the first plot in figure 2, the total force appears to have a finite value at the O8₋ singularity. This would be a qualitative difference with the common behavior observed in supersymmetric solutions with O8₋ orientifolds such as [40–44]. This discrepancy arises despite the fact that the dilaton and the metric have the right behavior at that singularity. However, while a possible orientifold source can be identified by just looking at the leading components of the warp factors and the dilaton close to the singularity, the force F_{D8} has a strong dependence on subleading components. We thus believe that this result is not entirely reliable: first, the non-trivial profile of F_{D8} is the result of a huge cancellation between the WZ and the DBI (electric and dilato-gravitational) forces acting on the probe, which diverge close to $z = z_0$ as

$$F_{\text{el, grav}} \propto \frac{1}{(z_0 - z)^2} \quad \text{for } z \rightarrow z_0 \tag{3.14}$$

where we have used (2.6). As a consequence, it is quite questionable how much the numerical evaluation of F_{D8} can be trusted in the strongly curved region, where the cancellations are very large and the numerical noise is important.

Furthermore, even if one could reach an infinite precision in evaluating the D8 action in the CDLT₁ solution, a more severe conceptual problem must be taken into account. While the divergence of the dilaton as one approaches the O8₋ could naively seem not to pose a severe problem as the couplings g_{YM} , m_B^2 , m_F in the effective action remain finite on the

¹⁰The analysis is qualitatively the same also for the particular solution (2.7) where the four-form flux is vanishing.

whole interval, we have to remember that the DBI and WZ actions are only the first terms in a more general expansion in powers of $g_s \approx e^\phi$. In the usual supergravity regime, the higher-genus contributions can be neglected and (3.1) provides a good approximation, but in a strongly coupled region such corrections cannot be considered sub-leading and become relevant. Therefore, we cannot rely on the usual effective probe action. This is exactly the situation close to the $O8_-$ singularity, and thus the full force F_{D8} should be computed using the full genus expansion that we do not have at our disposal.

There is however a silver lining. In supersymmetric solutions (such as [40–44]), one can invoke supersymmetry protection and obtain correct physics using the DBI+WZ approximation well beyond its validity regime. Furthermore, for anti-D3 branes in the KS solution [50], the loop corrections to the probe potential end up canceling to all loops, giving a flat potential despite the fact that supersymmetry is broken [51]. It is possible that this cancelation is related to the field content of the underlying theory, which is the same as that of the maximally supersymmetric theory [52]. A similar phenomenon might make the DBI+WZ approximation of the D8 brane action describe the physics correctly even in the strongly-coupling regime.

On the other hand, we have a much better control in the region close to the orientifold source with a positive charge, $z = 0$. The singularity is of exactly the (much milder) type one encounters near D8 branes: it manifests as discontinuities of the derivatives of physical quantities, the dilaton is finite and can be made arbitrarily small if $N \gg 1$. The asymptotic behavior of the total force F_{D8} close to $O8_+$ source is the following:

$$F_{D8} = \alpha z^3 + O(z^4), \tag{3.15}$$

with α a constant. The probe D8 feels no force next to the $O8_+$ source, which one can attribute to the fact that in flat space they preserve the same supersymmetries. However, the D8 branes are immediately repelled if an infinitesimal displacement occurs. Hence, the probe D8-branes are actually unstable, and their worldvolume scalar parameterizing their distance away from the $O8_+$ source is tachyonic.

As we explained in the Introduction, the key question is whether this probe D8 brane tachyon signals an instability of the background or is just an interesting tangential feature of the solution. This hinges on the interpretation of the orientifold source with positive D8 charge as an $O8_-$ plane with 16 D8 branes on top, or as an $O8_+$ plane. We know that in flat un-warped space both a configuration with an $O8_-$ plane and an $O8_+$ plane, and a configuration with two $O8_-$ planes and 16 mobile D8 branes are consistent. The former is the T-dual of the Type I' Dabholkar-Park configuration [36–38] and is dual to M-theory on the Klein bottle. The latter is the T-dual of vanilla Type I String Theory. Furthermore, in the supergravity regime it appears impossible to distinguish between an $O8_+$ and an $O8_-$ plane with 16 D8 branes on top.

Our analysis indicates that if the orientifold source with positive charge corresponds to an $O8_-$ plane with 16 D8 branes, the solution is tachyonic, and the tachyon corresponds to the D8 branes moving away from this source. Given the fact that the solution has two $O8_-$ planes, it is clear that the end point of this tachyonic instability will involve a \mathbb{Z}_2

symmetric configuration, most likely the flat-space solution where half the D8 branes are on one of the $O8_-$ planes and the other half is on the other. On the other hand, if the orientifold source with positive charge corresponds to an $O8_+$ plane, it does not appear possible for this plane to perturbatively emit D8 branes, and hence the instability of the D8 brane probes we found does not correspond to an instability if the solution.

4 Discussion

In this note we have investigated the orientifold singularities of the $CDLT_1$ de Sitter solutions using probe D8-branes. Starting from the usual DBI + WZ action, we computed the bosonic and fermionic action describing the fluctuations of a stack of k D8-branes sitting at some reference point z_{D8} . We observed that this action, despite its finiteness everywhere, becomes harder to trust when the stack of D8 branes approaches the singularity caused by vanishing warp factor near the $O8_-$ source. Accessing this region, where the action becomes strongly coupled, requires some alternative effective description that would be interesting to construct in the future.

On the other hand, D8-branes are good probes close to the orientifold source with positive D8 charge, where the singularity is milder: in particular, the dilaton can be made arbitrarily small by increasing the units of four-form flux. The position of the D8 stack is tachyonic, and this would signal an instability of the de Sitter $CDLT_1$ solution if the orientifold source with positive D8 charge had the possibility to emit D8 branes.

Since the orientifold source with positive D8 charge can correspond either to an $O8_+$ plane or to an $O8_-$ with 16 D8 branes on top, and since these two possibilities cannot be distinguished using just the supergravity solution, our work leaves open two possibilities.

- Either the $CDLT_1$ solutions are sourced by two $O8_-$ planes, one of which has 16D8 branes on top. This solution is unstable.
- Or the $CDLT_1$ solutions are sourced by an $O8_-$ and an $O8_+$ plane. This solution does not have any mobile D8 brane, and hence the probe D8 brane instability does not affect it.

There are several ways one can try to break the tie between these possibilities. The first is to understand if there is any inconsistency upon turning on a finite worldvolume cosmological constant, in either the $O8_{\pm}$ configuration or in the $2O8_- + 16$ D8 configuration. The second is to understand whether other probe branes can help distinguish between the two configurations. A preliminary exploration of the action of other types of probe branes in the $CDLT_1$ solution gives interesting physics, but unfortunately does not help us solve our conundrum. We present some of these results in appendix B

The third is to understand whether an $O8_+$ plane may be related by dualities to an $O8_-$ plane with 16 D8 branes. These two objects have the same charge, and there exist other known situations where orientifold planes with stuck D-branes are related to other orientifold planes with the same total charge. One such example is an $O3_+$ plane, which by Seiberg-type dualities can be related to an $O3_-$ plane with a stuck D3-brane [53, 54]. If such

a duality relation exists between the $O8_+$ and the $O8_-$ plane with 16 D8, the emission of a D8 brane from a $O8_- + 16$ D8 configuration would correspond to a non-perturbative tachyonic instability of the $O8_- - O8_+$ configuration, which would be fascinating to figure out.

Given the importance of orientifold planes in the construction of explicit de Sitter backgrounds of String Theory, it is crucial to understand in detail the physics of these objects. Hence another natural extension of the present work would be to investigate the richer class of CDLT₂ backgrounds, where the presence $O6_-$ planes avoids the introduction of singular $O8_-$ sources.

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A The fermionic action

The fermionic action of D8 branes in flux backgrounds can be computed evaluating the action proposed in [55] on the CDLT₁ backgrounds.

Its general form is:¹¹

$$S_{\text{Dp}}^{(F)} = \frac{\tau_{\text{Dp}}}{2} \lambda^2 \int d^{p+1} \xi e^{-\phi} \sqrt{-\det(P[g] + \mathcal{F})} \left\{ \bar{\theta}_+ \left[(M^{-1})^{\alpha\beta} \Gamma_\alpha^{(P)} \check{\nabla}_\beta - \mathcal{E}^{(1)} \right] \theta_{++} - \bar{\theta}_+ \check{\Gamma}_{\text{D8}}^{-1} \left[(M^{-1})^{\alpha\beta} \Gamma_\beta^{(P)} \Xi_\alpha - \mathcal{E}^{(2)} \right] \theta_+ \right\}. \quad (\text{A.1})$$

Let us explain the meaning of the various quantities in (A.1): first this action is already κ -fixed, in such a way that only one chiral component of the type IIA ten-dimensional Majorana spinor Θ supergravity is kept: in particular, we choose a κ -fixing such that $\theta_+ = \frac{1}{2}(1 + \Gamma_{(10)})\Theta$, with $\Gamma_{(10)}$ the ten-dimensional chirality matrix. The action is written in terms of the matrix $M_{\alpha\beta} = P[g]_{\alpha\beta} + \lambda \mathcal{F}_{\alpha\beta}$ while the pullback of the Γ -matrices is denoted by $\Gamma_\alpha^{(P)}$.

The worldvolume coordinates are labeled by the indices α, β, \dots , while the ten-dimensional space-time coordinates are labeled by the indices A, B, \dots . For a D8 brane, the other quantities in (A.1) are:

$$\check{\nabla}_\beta = \partial_\beta x^A \nabla_A, \quad \Xi_\alpha = -\frac{1}{8} e^\phi F_0 \Gamma_\alpha^{(P)}, \quad (\text{A.2})$$

$$\mathcal{E}^{(1)} = \frac{1}{2} \Gamma^A \partial_A \phi, \quad \mathcal{E}^{(2)} = -\frac{5}{8} e^\phi F_0. \quad (\text{A.3})$$

¹¹In our conventions, $\bar{\theta} = i\theta^t \Gamma_0$.

If we denote the 10d vielbein with E^A_A , the non-trivial components of the spin connection Ω are:

$$\Omega^{\alpha}_{\underline{\beta}} = \omega^{\alpha}_{\underline{\beta}}, \quad \Omega^{\alpha}_{\underline{z}} = e^W \partial_z(E^{\alpha}), \quad (\text{A.4})$$

where $\omega^{\alpha}_{\underline{\beta}}$ is nothing but the worldvolume spin connection in the absence of warping. As a consequence, $\nabla_z = \partial_z$ and:

$$\nabla_{\alpha} = \partial_{\alpha} + \frac{1}{4} \omega_{\alpha}^{\underline{\beta}\underline{\gamma}} \Gamma_{\underline{\beta}\underline{\gamma}} + \frac{1}{2} \Omega_{\alpha}^{\underline{\beta}\underline{z}} \Gamma_{\underline{\alpha}} \Gamma_{\underline{z}} = \nabla_{\alpha}^{(0)} + \frac{1}{2} e^W \partial_z(\Gamma_{\alpha}) \Gamma_{\underline{z}}, \quad (\text{A.5})$$

where $\nabla^{(0)}$ is the covariant derivative with ω as spin connection. At the lowest order in λ , we can use the following approximations:

$$M_{\alpha\beta} = g_{\alpha\beta} + O(\lambda), \quad \check{\nabla}_{\alpha} = \nabla_{\alpha} + O(\lambda), \quad (\text{A.6})$$

$$\Gamma_{\alpha}^{(P)} = \Gamma_{\alpha} + O(\lambda), \quad \Gamma_{\text{D8}}^{(0)} = \Gamma_{\underline{z}} \Gamma_{(10)} + O(\lambda). \quad (\text{A.7})$$

We can now evaluate (A.1) for a single D8 brane to the lowest order in λ , obtaining the following Abelian action:

$$S_{\text{D8}}^{(F)} = \frac{\tau_{\text{D8}}}{2} \lambda^2 \int d^9 \xi \sqrt{g^{(0)}} \Delta_{\text{D8}} \bar{\theta}_+ \left[\Gamma^{\alpha} \nabla_{\alpha} + \frac{1}{2} \left(e^W \Gamma^{\alpha} \partial_z \Gamma_{\alpha} - e^W \partial_z \phi - e^{\phi} F_0 \right) \Gamma_{\underline{z}} \right] \theta_+ \quad (\text{A.8})$$

Following [56], the non-Abelian action can be obtained promoting the fermion, θ , to an adjoint (Majorana) fermion and performing the usual covariantization $\partial_{\alpha} \rightarrow \partial_{\alpha} - i[A_{\alpha}, \cdot] + O(\lambda)$, introducing the additional coupling of the form $\Gamma^{\alpha}[A_{\alpha}, \cdot]$; consistency with T-duality requires an additional companion term of the form $\Gamma_z[\Phi, \cdot]$. The final result is:

$$S_{\text{D8}}^{(F)} = \frac{\tau_{\text{D8}}}{2} \lambda^2 \int d^9 \xi \sqrt{g^{(0)}} \Delta_{\text{D8}} \bar{\theta}_+ \left\{ \Gamma^{\alpha} \nabla_{\alpha} \theta_+ + \frac{1}{2} \left(e^W \partial_z \log \Delta_{\text{D8}} + e^{\phi} F_0 \right) \Gamma_{\underline{z}} \theta_+ + \right. \\ \left. - i[A_{\alpha}, \Gamma^{\alpha} \theta_+] - i e^{-W} [\Phi, \Gamma_{\underline{z}} \theta_+] \right\} \quad (\text{A.9})$$

We can now introduce a smart choice of Γ -matrices that make manifest the 9 + 1 splitting of the ten-dimensional space:

$$\Gamma_{\underline{\alpha}} = \gamma_{\underline{\alpha}} \otimes \sigma_2, \quad \Gamma_{\underline{z}} = \mathbb{I} \otimes \sigma_1, \quad (\text{A.10})$$

where $\gamma_{\underline{\alpha}}$ are 9-dimensional purely-imaginary γ -matrices with Minkowskian signature.¹² With these conventions, the 10d chirality matrix can be written as:

$$\Gamma_{\underline{10}} = \Gamma_{\underline{0}} \dots \Gamma_{\underline{9}} = \mathbb{I} \otimes \sigma_3, \quad (\text{A.11})$$

where we used the fact that the 9-dimensional γ -matrices can be chosen such that $\gamma_{\underline{0}} \dots \gamma_{\underline{8}} = i \mathbb{I}$; in this basis. Hence, a chiral Majorana spinor can be decomposed as

$$\theta_+ = \chi \otimes \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad (\text{A.12})$$

¹²A way to construct them, for instance, is to start from 7-dimensional Euclidean γ -matrices, that can be chosen to be purely imaginary. Then, one can construct $\gamma_{\underline{i}} = \gamma_{\underline{i}}^{(7)} \otimes \sigma_3$, $\gamma_{\underline{8}} = \mathbb{I} \otimes \sigma_2$, $\gamma_{\underline{9}} = \mathbb{I} \otimes (i\sigma_1)$.

where χ is a 9d Majorana spinor. Using such decomposition of the spinors, it is straightforward to show that the following relations hold:

$$\bar{\theta}\Gamma_{\underline{\alpha}}\theta = i\bar{\chi}\gamma_{\underline{\alpha}}\chi, \quad \bar{\theta}\Gamma_{\underline{z}}\theta = \bar{\chi}\chi, \quad (\text{A.13})$$

where, as before, $\bar{\chi} = i\chi^t\gamma_0$. As we did for the bosonic action, it is best to normalize the fermions and in particular to rescale the spinor χ by a factor $\frac{1}{\lambda\sqrt{\tau_{\text{D8}}}}e^{(\phi-8W)/2}$. The canonically normalized action is then:

$$S_{\text{D8}}^{(F)} = \frac{1}{2} \int d^9\xi \sqrt{\gamma_{\text{D8}}} \bar{\chi} \{i\gamma^\alpha \nabla_\alpha \chi + m_F \chi + [A_\alpha, \gamma^\alpha \chi] - i g_{\text{YM}}[\Phi, \chi]\}. \quad (\text{A.14})$$

where the fermion mass is defined as:

$$m_F \equiv \frac{e^{2W}}{2} \left(\partial_z \log \Delta_{\text{D8}} + e^{\phi-W} F_0 \right). \quad (\text{A.15})$$

Finally, let us comment about a possible contribution coming from the four-form flux. In fact, the fermionic action in principle contains a Lagrangian term of the form:

$$\delta S_{F_4} = \frac{\tau_{\text{D8}}}{2} \lambda^2 \int d^9\xi e^{-\phi} \sqrt{-\det P[g]} \frac{e^\phi}{8} \bar{\theta}_+ \Gamma_{\underline{z}} [\Gamma^\alpha \not{F}_4 \Gamma_\alpha - \not{F}_4], \quad (\text{A.16})$$

where $\not{F}_4 = \frac{1}{4!} F_{mnpq} \Gamma^{mnpq}$. We recall that in the CDLT₁ backgrounds the four-form flux has the form (2.2), and define for simplicity $\tilde{f}_4 \equiv f_4 e^{-6W+3\lambda_3-2\lambda_2}$. It is easy to see that:

$$\begin{aligned} \bar{\theta}_+ \Gamma_{\underline{z}} [\Gamma^\alpha \not{F}_4 \Gamma_\alpha - \not{F}_4] \theta_+ &= \tilde{f}_4 \bar{\theta}_+ \Gamma_{\underline{z}} [\Gamma^\alpha \Gamma_{\underline{z}} \text{v}\not{1}_{M_3} \Gamma_\alpha - \Gamma_{\underline{z}} \text{v}\not{1}_{M_3}] \\ &= \tilde{f}_4 \bar{\theta}_+ [3\text{v}\not{1}_{M_3} - \text{v}\not{1}_{M_3}] \theta_+ = 2\bar{\theta}_+ \text{v}\not{1}_{M_3} \theta_+ \end{aligned} \quad (\text{A.17})$$

However, this combination vanishes for a ten-dimensional chiral spinor, given the symmetry properties of the Γ -matrices. This establishes that the non-trivial four-form flux does not affect our computation at the lowest order in λ .

B Probe D0-branes

Besides the action of probe D8 branes, it is instructive to see whether other probe D-branes have interesting physics in the CDLT₁ solution. The most interesting are D0 branes, whose action in this background naively contains only a DBI term: $-\tau_{\text{D0}} \int e^{-\phi} \sqrt{-g_{00}}$. However, in the presence of D8 branes (and O8 planes as well), these D0 branes come with F1 strings attached. One way to see this is from imposing tadpole cancelation in the presence of a Romans mass [57] but this does not indicate which D8 branes the F1 strings terminate on. A simpler way is to realize that in a region with no Romans mass, D0 branes have no strings attached, but passing them through n_0 D8 branes creates n_0 F1 strings via the Hanany-Witten effect [58].

Hence, the complete action describing D0 branes in our solution is:

$$S_{\text{D0}}^{F_0} = S_{\text{D0}} + F_0 S_{\text{F1}} = -\tau_{\text{D0}} \int dt e^{-\phi} \sqrt{-g_{00}} - n_0 \tau_{\text{F1}} \int_\gamma dt d\xi \sqrt{-\det P[g]}, \quad (\text{B.1})$$

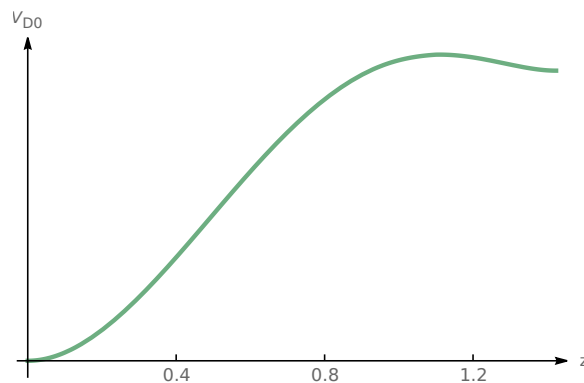


Figure 3. – The D0-brane potential, V_{D0} . It can be seen numerically that the unstable equilibrium point only exists when $F_4 \neq 0$ and that it gets closer to z_0 as N grows.

where we already assumed that the NS potential is vanishing in the background, we denoted the temporal direction by t and we defined $n_0 \equiv 2\pi l_s F_0$. The surface γ is such that the string extends between the D0 brane and a D8 brane (or an O8 plane) sourcing the Romans mass. Let us call z_{D0} the position of the D0 brane along the interval: the tadpole cancellation requires the presence of n_0 attached strings ending on the orientifold plane located at $z = z_0$. In a configuration of minimal energy, it is natural to expect that the attached string extends only along the interval transverse to the O8 planes, so that the action (B.1) can be written as:

$$S_{D0}^{F_0} = -\tau_{D0} \int dt \left(e^{W(z_{D0}) - \phi(z_{D0})} + F_0(z_0 - z_{D0}) \right), \quad (\text{B.2})$$

where we used the fact that $\tau_{F1} = \tau_{D0}/(2\pi l_s)$ and that $\det P[g] = g_{00} g_{zz} = -1$ on the CDLT₁ background. The integrand in equation (B.2) can be thought as the interval-dependent potential V_{D0} felt by the D0 branes with (F1) strings attached. We plot this potential in figure 3. It is interesting to observe that the D0 branes are stable on both orientifold planes, and also have an unstable equilibrium point.

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References

- [1] C. Córdova, G.B. De Luca and A. Tomasiello, *Classical de Sitter solutions of 10-dimensional supergravity*, *Phys. Rev. Lett.* **122** (2019) 091601 [[arXiv:1812.04147](https://arxiv.org/abs/1812.04147)] [[INSPIRE](#)].
- [2] S. Kachru, R. Kallosh, A.D. Linde and S.P. Trivedi, *De Sitter vacua in string theory*, *Phys. Rev. D* **68** (2003) 046005 [[hep-th/0301240](https://arxiv.org/abs/hep-th/0301240)] [[INSPIRE](#)].
- [3] I. Bena, M. Graña and N. Halmagyi, *On the existence of meta-stable vacua in Klebanov-Strassler*, *JHEP* **09** (2010) 087 [[arXiv:0912.3519](https://arxiv.org/abs/0912.3519)] [[INSPIRE](#)].

- [4] I. Bena, G. Giecold, M. Graña, N. Halmagyi and S. Massai, *The backreaction of anti-D3 branes on the Klebanov-Strassler geometry*, *JHEP* **06** (2013) 060 [[arXiv:1106.6165](#)] [[INSPIRE](#)].
- [5] I. Bena, M. Graña, S. Kuperstein and S. Massai, *Giant tachyons in the landscape*, *JHEP* **02** (2015) 146 [[arXiv:1410.7776](#)] [[INSPIRE](#)].
- [6] J. Moritz, A. Retolaza and A. Westphal, *Toward de Sitter space from ten dimensions*, *Phys. Rev. D* **97** (2018) 046010 [[arXiv:1707.08678](#)] [[INSPIRE](#)].
- [7] U.H. Danielsson and T. Van Riet, *What if string theory has no de Sitter vacua?*, *Int. J. Mod. Phys. D* **27** (2018) 1830007 [[arXiv:1804.01120](#)] [[INSPIRE](#)].
- [8] I. Bena, E. Dudas, M. Graña and S. Lüst, *Uplifting runaways*, *Fortsch. Phys.* **67** (2019) 1800100 [[arXiv:1809.06861](#)] [[INSPIRE](#)].
- [9] R. Blumenhagen, D. Kläwer and L. Schlechter, *Swampland variations on a theme by KKLT*, *JHEP* **05** (2019) 152 [[arXiv:1902.07724](#)] [[INSPIRE](#)].
- [10] S. Kachru, M. Kim, L. McAllister and M. Zimet, *De Sitter vacua from ten dimensions*, [arXiv:1908.04788](#) [[INSPIRE](#)].
- [11] I. Bena, M. Graña, N. Kovensky and A. Retolaza, *Kähler moduli stabilization from ten dimensions*, *JHEP* **10** (2019) 200 [[arXiv:1908.01785](#)] [[INSPIRE](#)].
- [12] X. Gao, A. Hebecker and D. Junghans, *Control issues of KKLT*, [arXiv:2009.03914](#) [[INSPIRE](#)].
- [13] G. Obied, H. Ooguri, L. Spodyneiko and C. Vafa, *De Sitter space and the swampland*, [arXiv:1806.08362](#) [[INSPIRE](#)].
- [14] D. Andriot, *On the de Sitter swampland criterion*, *Phys. Lett. B* **785** (2018) 570 [[arXiv:1806.10999](#)] [[INSPIRE](#)].
- [15] S.K. Garg and C. Krishnan, *Bounds on slow roll and the de Sitter swampland*, *JHEP* **11** (2019) 075 [[arXiv:1807.05193](#)] [[INSPIRE](#)].
- [16] H. Ooguri, E. Palti, G. Shiu and C. Vafa, *Distance and de Sitter conjectures on the swampland*, *Phys. Lett. B* **788** (2019) 180 [[arXiv:1810.05506](#)] [[INSPIRE](#)].
- [17] J.M. Maldacena and C. Núñez, *Supergravity description of field theories on curved manifolds and a no go theorem*, *Int. J. Mod. Phys. A* **16** (2001) 822 [[hep-th/0007018](#)] [[INSPIRE](#)].
- [18] X. Dong, B. Horn, E. Silverstein and G. Torroba, *Micromanaging de Sitter holography*, *Class. Quant. Grav.* **27** (2010) 245020 [[arXiv:1005.5403](#)] [[INSPIRE](#)].
- [19] S. Baines and T. Van Riet, *Smearing orientifolds in flux compactifications can be OK*, *Class. Quant. Grav.* **37** (2020) 195015 [[arXiv:2005.09501](#)] [[INSPIRE](#)].
- [20] F. Marchesano, E. Palti, J. Quirant and A. Tomasiello, *On supersymmetric AdS₄ orientifold vacua*, *JHEP* **08** (2020) 087 [[arXiv:2003.13578](#)] [[INSPIRE](#)].
- [21] D. Junghans, *O-plane backreaction and scale separation in type IIA flux vacua*, *Fortsch. Phys.* **68** (2020) 2000040 [[arXiv:2003.06274](#)] [[INSPIRE](#)].
- [22] C. Córdova, G.B. De Luca and A. Tomasiello, *New de Sitter solutions in ten dimensions and orientifold singularities*, *JHEP* **08** (2020) 093 [[arXiv:1911.04498](#)] [[INSPIRE](#)].
- [23] E. Malek and H. Samtleben, *Kaluza-Klein spectrometry for supergravity*, *Phys. Rev. Lett.* **124** (2020) 101601 [[arXiv:1911.12640](#)] [[INSPIRE](#)].

- [24] E. Malek, H. Nicolai and H. Samtleben, *Tachyonic Kaluza-Klein modes and the AdS swampland conjecture*, *JHEP* **08** (2020) 159 [[arXiv:2005.07713](#)] [[INSPIRE](#)].
- [25] E. Malek and H. Samtleben, *Kaluza-Klein spectrometry from exceptional field theory*, *Phys. Rev. D* **102** (2020) 106016 [[arXiv:2009.03347](#)] [[INSPIRE](#)].
- [26] A. Guarino, E. Malek and H. Samtleben, *Stable nonsupersymmetric anti-de Sitter vacua of massive IIA supergravity*, *Phys. Rev. Lett.* **126** (2021) 061601 [[arXiv:2011.06600](#)] [[INSPIRE](#)].
- [27] C. Eloy, *Kaluza-Klein spectrometry for AdS₃ vacua*, [arXiv:2011.11658](#) [[INSPIRE](#)].
- [28] I. Bena, M. Graña, S. Kuperstein and S. Massai, *Anti-D3 branes: singular to the bitter end*, *Phys. Rev. D* **87** (2013) 106010 [[arXiv:1206.6369](#)] [[INSPIRE](#)].
- [29] N. Cribiori and D. Junghans, *No classical (anti-)de Sitter solutions with O8-planes*, *Phys. Lett. B* **793** (2019) 54 [[arXiv:1902.08209](#)] [[INSPIRE](#)].
- [30] I. Bena, K. Pilch and N.P. Warner, *Brane-jet instabilities*, *JHEP* **10** (2020) 091 [[arXiv:2003.02851](#)] [[INSPIRE](#)].
- [31] N.P. Warner, *Some properties of the scalar potential in gauged supergravity theories*, *Nucl. Phys. B* **231** (1984) 250 [[INSPIRE](#)].
- [32] N.P. Warner, *Some new extrema of the scalar potential of gauged N = 8 supergravity*, *Phys. Lett. B* **128** (1983) 169 [[INSPIRE](#)].
- [33] F. Apruzzi, G. Bruno De Luca, A. Gnecci, G. Lo Monaco and A. Tomasiello, *On AdS₇ stability*, *JHEP* **07** (2020) 033 [[arXiv:1912.13491](#)] [[INSPIRE](#)].
- [34] E. Silverstein, *New de Sitter solutions*, plenary talk at *Strings*, June 2013.
- [35] E. Silverstein, private communications.
- [36] A. Dabholkar and J. Park, *Strings on orientifolds*, *Nucl. Phys. B* **477** (1996) 701 [[hep-th/9604178](#)] [[INSPIRE](#)].
- [37] E. Witten, *Toroidal compactification without vector structure*, *JHEP* **02** (1998) 006 [[hep-th/9712028](#)] [[INSPIRE](#)].
- [38] O. Aharony, Z. Komargodski and A. Patir, *The moduli space and M(atrrix) theory of 9d N = 1 backgrounds of M/string theory*, *JHEP* **05** (2007) 073 [[hep-th/0702195](#)] [[INSPIRE](#)].
- [39] N. Kim, *Towards an explicit construction of de Sitter solutions in classical supergravity*, *JHEP* **10** (2020) 057 [[arXiv:2004.05885](#)] [[INSPIRE](#)].
- [40] A. Brandhuber and Y. Oz, *The D4-D8 brane system and five-dimensional fixed points*, *Phys. Lett. B* **460** (1999) 307 [[hep-th/9905148](#)] [[INSPIRE](#)].
- [41] F. Apruzzi and M. Fazzi, *AdS₇/CFT₆ with orientifolds*, *JHEP* **01** (2018) 124 [[arXiv:1712.03235](#)] [[INSPIRE](#)].
- [42] G. Dibitetto, G. Lo Monaco, A. Passias, N. Petri and A. Tomasiello, *AdS₃ solutions with exceptional supersymmetry*, *Fortsch. Phys.* **66** (2018) 1800060 [[arXiv:1807.06602](#)] [[INSPIRE](#)].
- [43] A. Passias, D. Prins and A. Tomasiello, *A massive class of N = 2 AdS₄ IIA solutions*, *JHEP* **10** (2018) 071 [[arXiv:1805.03661](#)] [[INSPIRE](#)].
- [44] Y. Lozano, N.T. Macpherson, C. Núñez and A. Ramirez, *AdS₃ solutions in massive IIA with small N = (4, 0) supersymmetry*, *JHEP* **01** (2020) 129 [[arXiv:1908.09851](#)] [[INSPIRE](#)].

- [45] S.R. Coleman and F. De Luccia, *Gravitational effects on and of vacuum decay*, *Phys. Rev. D* **21** (1980) 3305 [[INSPIRE](#)].
- [46] J.M. Maldacena, J. Michelson and A. Strominger, *Anti-de Sitter fragmentation*, *JHEP* **02** (1999) 011 [[hep-th/9812073](#)] [[INSPIRE](#)].
- [47] S. Kachru, J. Pearson and H.L. Verlinde, *Brane/flux annihilation and the string dual of a nonsupersymmetric field theory*, *JHEP* **06** (2002) 021 [[hep-th/0112197](#)] [[INSPIRE](#)].
- [48] F.F. Gautason, B. Truijen and T. Van Riet, *The many faces of brane-flux annihilation*, *JHEP* **10** (2015) 152 [[arXiv:1505.00159](#)] [[INSPIRE](#)].
- [49] F. Apruzzi, G. Dibitetto and L. Tizzano, *A new 6d fixed point from holography*, *JHEP* **11** (2016) 126 [[arXiv:1603.06576](#)] [[INSPIRE](#)].
- [50] I.R. Klebanov and M.J. Strassler, *Supergravity and a confining gauge theory: duality cascades and χ SB resolution of naked singularities*, *JHEP* **08** (2000) 052 [[hep-th/0007191](#)] [[INSPIRE](#)].
- [51] I. Bena, J. Blåbäck and D. Turton, *Loop corrections to the antibrane potential*, *JHEP* **07** (2016) 132 [[arXiv:1602.05959](#)] [[INSPIRE](#)].
- [52] I. Bena, M. Graña, S. Kuperstein, P. Ntokos and M. Petrini, *D3-brane model building and the supertrace rule*, *Phys. Rev. Lett.* **116** (2016) 141601 [[arXiv:1510.07039](#)] [[INSPIRE](#)].
- [53] A.M. Uranga, *Comments on nonsupersymmetric orientifolds at strong coupling*, *JHEP* **02** (2000) 041 [[hep-th/9912145](#)] [[INSPIRE](#)].
- [54] S. Elitzur, A. Giveon, D. Kutasov and D. Tsabar, *Branes, orientifolds and chiral gauge theories*, *Nucl. Phys. B* **524** (1998) 251 [[hep-th/9801020](#)] [[INSPIRE](#)].
- [55] L. Martucci, J. Rosseel, D. Van den Bleeken and A. Van Proeyen, *Dirac actions for D-branes on backgrounds with fluxes*, *Class. Quant. Grav.* **22** (2005) 2745 [[hep-th/0504041](#)] [[INSPIRE](#)].
- [56] P. McGuirk, G. Shiu and F. Ye, *Soft branes in supersymmetry-breaking backgrounds*, *JHEP* **07** (2012) 188 [[arXiv:1206.0754](#)] [[INSPIRE](#)].
- [57] O. Bergman, M.R. Gaberdiel and G. Lifschytz, *Branes, orientifolds and the creation of elementary strings*, *Nucl. Phys. B* **509** (1998) 194 [[hep-th/9705130](#)] [[INSPIRE](#)].
- [58] A. Hanany and E. Witten, *Type IIB superstrings, BPS monopoles, and three-dimensional gauge dynamics*, *Nucl. Phys. B* **492** (1997) 152 [[hep-th/9611230](#)] [[INSPIRE](#)].