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AdS_4 flux vacua in type II superstrings and their domain-wall solutions

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Abstract: We investigate the emergence of supersymmetric negative-vacuum-energy ground states in four dimensions. First, we rely on the analysis of the effective superpotential, which depends on the background fluxes of the internal manifold, or equivalently has its origin in the underlying gauged supergravity. Four-dimensional, supersymmetric anti-de Sitter vacua with all moduli stabilized appear when appropriate Ramond and Neveu-Schwarz fluxes are introduced in IIA. Geometric fluxes are not necessary. Then the whole setup is analyzed from the perspective of the sources, namely D/NS-branes or Kaluza-Klein monopoles. Orientifold planes are also required for tadpole cancellation. The solutions found in four dimensions correspond to domain walls interpolating between AdS_4 and flat spacetime. The various consistency conditions (equations of motion, Bianchi identities and tadpole cancellation conditions) are always satisfied, albeit with source terms. We also speculate on the possibility of assigning (formal) entropies to AdS_4 flux vacua via the corresponding dual brane systems.

Keywords: D-branes, Superstring Vacua, Supersymmetric Effective Theories, Flux compactifications.

Contents

1. Introduction

The closed/open string correspondence between flux vacua and D-brane solutions has led to remarkable new insights on the structure of gauge and gravity theories. Anti-de Sitter (AdS) geometries play a crucial role in this context since they appear on one side of the AdS/CFT duality [1−3]. The latter was originally formulated between AdS_5 extended supergravity and $\mathcal{N} = 4$ SU(N) super Yang-Mills on the boundary of AdS_5. It is based on the observation that the type IIB, $AdS_5 \times S^5$ ground state with non-vanishing Ramond 5-form flux arises as the near-horizon geometry of a stack of D3-branes with open strings attached. Another example with maximal supersymmetry is $\mathcal{N} = 8$ supergravity on $AdS_{4(7)} \times S^{7(4)}$ in eleven-dimensional supergravity with non-vanishing 4-form flux. Again, these AdS background spaces arise as near-horizon geometries of M2(5)-branes.

In fact, the holographic correspondence relies on the possibility of viewing N units of 5-form flux in ten-dimensional supergravity as originating from a stack of N D3-branes. This interplay between flux backgrounds and branes is richer than the simple relation that exists between a field and the sources that create it and they will play a central role in the present note. Indeed, the sources at hand are truly dynamical objects that bring their own degrees of freedom and, depending on the regime, the latter may become important ingredients.

Anti-de Sitter geometries arise also in string backgrounds with fewer supersymmetries, such as $\mathrm{AdS}_5 \times T^{1,1}$, which arises from placing a stack of D3-branes near the six-dimensional conifold singularity [4]. This $\mathcal{N} = 2$ supersymmetric background is dual to a conformal $\mathcal{N} = 1$, $\mathrm{SU}(N) \times \mathrm{SU}(N)$ gauge theory with bifundamental matter fields. The geometric transitions [5] are another nice example of the flux/brane correspondence where one trades fluxes through internal cycles for their dual sources made of wrapped D-branes. In this way an intriguing relation emerges between the confining phase in $\mathcal{N} = 1$ effective gauge theories and their microscopic $\mathrm{SU}(N)$ descriptions. Moreover, string backgrounds of the form $\mathrm{AdS}_3 \times S^3 \times T^4(K3)$ can be constructed by conformal field theory techniques and play an important role in the context of dualities between three-dimensional gravity and two-dimensional boundary CFT's [6, 7].

The investigation of the flux/brane correspondence in conjunction with the generated anti-de Sitter backgrounds is also motivated by the possibility of comparing the macroscopic and microscopic entropies of supersymmetric black holes [8, 9]. Supersymmetric black holes in $\mathcal{N} = 2$ type IIA/B supergravity [10] arise as solutions with non-vanishing electric/magnetic charges of $\mathrm{U}(1)$ Ramond gauge fields in four dimensions. Their near-horizon geometry is given by the space $\mathrm{AdS}_2 \times S^2$. The corresponding thermodynamic entropies are determined by the area of the horizon plus corrections from higher-curvature terms in the effective action [11 − 13]. The same $\mathrm{AdS}_2 \times S^2$ geometry is obtained by adopting the brane-source picture.

The sources are in this case intersecting D-branes, point-like in four dimensions, and hence with the spatial part of their world-volumes entirely inside the internal six-dimensional part of space: D3-branes wrapped around internal Calabi-Yau (CY) 3-cycles in type IIB, or D4-branes wrapped around 4-cycles plus D0-branes in type IIA. Counting the open-string excitations living on the wrapped branes, the resulting microscopic entropy matches up the macroscopic energy of the black holes in the effective supergravity theory, as demonstrated in several explicit CY examples. These observations are far reaching and have led to many subsequent developments: wave functions of certain supersymmetric $\mathrm{AdS}_2 \times S^2$ flux vacua [14], attractor formalism and generalized non-supersymmetric black holes (for a recent review of black holes in string theory see [15]).

Clearly, much more important and phenomenologically interesting than two-dimensional AdS_2 flux vacua is to consider four-dimensional AdS_4 supergravity solutions, which will be the subject of this paper. This raises a threefold problem: (i) provide stable supersymmetric AdS_4 flux vacua within supergravity theory, (ii) identify the corresponding brane pictures and (iii) compute the entropy by counting the microscopic string states of the dual branes. In this paper we try to take the first two steps in this program by deriving

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supersymmetric, intersecting brane configurations generating four-dimensional AdS_4 flux vacua.

Our framework will be perturbative type II supergravity. We will in particular concentrate on type IIA theory with non-vanishing Ramond, Neveu-Schwarz, and possibly geometrical fluxes that can in principle generate AdS_4 ground states without using further non-perturbative effects. Despite the power of gauged-supergravity tools that can be used here, this exercise is not simple because the requirements of unbroken supersymmetry, with negative vacuum energy, and without runaway or flat directions are not easy to fulfill simultaneously. Type IIB solutions with similar properties can also be obtained. All this is the content of section 2.

The next step consists in describing the above AdS_4 flux vacua as near-horizon geometries of appropriately distributed intersecting D2/D4/D6/D8/NS5-branes and Kaluza-Klein monopoles. The examination of these, as well as of certain T-dual IIB systems, will be undertaken in section 4. Except for exotic non-geometric setups briefly described in section 2.3, full stabilization is hard to achieve perturbatively in type IIB CY compactifications because of the Khler-moduli independence of the superpotential. Nevertheless, we will introduce a new and interesting type IIB model by appropriately distributing D3/D5/D7/NS5/KK-branes/monopoles. Kaluza-Klein monopoles are sources for geometrical fluxes induced by the Scherk-Schwarz mechanism [16].[1] They introduce torsion and thereby Khler-moduli dependence. In the present case all moduli are stabilized around an AdS_4 vacuum, which appears as the near-horizon geometry of the above intersecting extended sources. The corresponding internal six-dimensional geometry is now a *nilmanifold*.

The general treatment of the source-branes, including the subtle issue of Bianchi identities and tadpole cancellation conditions is discussed in generality in section 3. It is made clear there that some of the AdS_4 vacua derived from a four-dimensional effective superpotential, *do not* admit a ten-dimensional supergravity interpretation with *localized* brane/orientifold sources. Indeed, as we will see in section 2.3, from the point-of-view of the four-dimensional effective superpotential there exist supersymmetric AdS_4 flux vacua *without* metric fluxes. However, such compactifications do not exist in the context of ten-dimensional supergravity [18], unless one considers the addition of *smeared* brane and/or orientifold sources [19]. The latter are precisely the ingredients needed for evading the necessity of metric fluxes. We note in passing that the smearing can be thought of as an artifact of the supergravity approximation, which necessarily ignores the massive Kaluza-Klein modes.

Let us finally mention that from the four-dimensional perspective, the branes act as two-branes, i.e. *domain walls*. This is now perfectly consistent, as opposed to Minkowski vacua, since the geometry is not flat but interpolates between AdS_4 and asymptotically flat space. Related considerations, including domain-wall solutions in the context of (gauged) supergravity, have previously appeared in

The domain-wall picture is very useful, since the tension of the domain walls can be

[1]From a string-theory viewpoint, the Scherk-Schwarz mechanism is implemented as a freely acting orbifold [17].

viewed as the origin of the effective $\mathcal{N}=1$ superpotential. Specifically, in this context one has to solve the supersymmetry conditions and equations of motion for the domain wall coupled to the various scalar fields. We expect that in such a solution the scalar fields would be non-singular in spacetime, since in the corresponding flux vacua they are all fixed to finite values. We therefore believe that for the supersymmetric AdS_4 domain-wall solutions there should exist an attractor mechanism that determines the values of the scalar fields at the AdS_4 horizon — in analogy to the attractor mechanism for the supersymmetric $\mathrm{AdS}_2 \times S^2$ black holes in $\mathcal{N}=2$ supergravity. To explicitly prove this statement is beyond the scope of this work.

2. AdS_4 vacua from flux compactifications

In the following we will exhibit type IIA/B models with Ramond, Neveu-Schwarz or geometrical fluxes that possess supersymmetric AdS_4 ground states. We will focus on perturbative (tree-level) contributions and use the effective superpotential description in four dimensions, omitting possible contributions to the scalar potential by D-terms.

2.1 Searching for minima

In a general $\mathcal{N}=1$ supergravity, the superpotential $W(\phi)$ is a function of chiral superfields ϕ_i. The corresponding scalar potential takes the standard form

$$V = \mathrm{e}^K \left(|D_i W|^2 - 3|W|^2\right) , \tag{2.1}$$

where the F-terms are defined as

$$F_i = \mathrm{e}^{K/2} D_i W = \mathrm{e}^{K/2} \left(\partial_{\phi_i} W + W \partial_{\phi_i} K\right) \tag{2.2}$$

with K being the Khler potential.

Our aim is to find supersymmetric extrema of V. We must therefore impose

$$F_i(\phi_{\min}) = 0 \ \forall i \, . \tag{2.3}$$

Anti-de Sitter vacua require negative vacuum energy. Equations (2.1) and (2.3) thus lead to the following requirement:

$$W(\phi_{\min}) \neq 0 \, . \tag{2.4}$$

Generically, conditions (2.3) and (2.4) are not easy to satisfy. One might assume e.g. $\partial_{\phi_i} K|_{\min} \neq 0$ for the main moduli fields. Using (2.3) and (2.4), this translates into $\partial_{\phi_i} W|_{\min} \neq 0$, which in turn shows that W must depend on all main moduli fields ϕ_i. By this reasoning, one would therefore conclude that supersymmetry together with negative vacuum energy require full moduli dependence of W and imply the stabilization of *all* moduli.[2] However, the assumption on non vanishing $\partial_{\phi_i} K|_{\min}$ does not hold generically (except for the seven main moduli considered in this paper) and supersymmetric AdS_4 do not systematically ensure all-moduli stabilization. This issue is more subtle and deserves therefore a careful analysis, which is one of our motivations here. Fortunately, this is somewhat easier that the corresponding search for Minkowskian vacua with (un)broken supersymmetry and with all moduli fixed.

[2] Strictly speaking, the real parts of all moduli are stabilized.

2.2 Fluxes in type IIA/B and mirror map

Type IIB We will start without geometrical fluxes; then the tree-level 3-form flux superpotential in type IIB on a Calabi-Yau 3-fold X is of the standard form [43–47]. It gets two kinds of contributions, namely from Ramond and Neveu-Schwarz 3-form fluxes:

$$\begin{aligned} W_{\mathrm{IIB}} &= W_H + W_F = \int_X \Omega \wedge \left(F_3^{\mathrm{R}} + S H_3^{\mathrm{NS}}\right) \\ &= e_0 + i e_i U_i + i m_0 F_0(U) + m_i F_i(U) \\ &\quad + iS(a_0 + i a_i U_i + i b_0 F_0(U) + b_i F_i(U))\,. \end{aligned} \tag{2.5}$$

Here Ω is the holomorphic 3-form on the CY space, and F_3^{R} (H_3^{NS}) is the Ramond (Neveu-Schwarz) 3-form field strength field. The U-dependent function $F(U) \equiv F_0(U)$ is the holomorphic prepotential and the $F_i(U)$ are its first derivatives. The e_I, m_I comprise the Ramond 3-form fluxes, whereas the a_I, b_I correspond to the Neveu-Schwarz 3-form fluxes ($I = 0, \ldots, h^{2,1}$). The superpotential W depends on the complex-structure moduli fields U_i ($i = 1, \ldots, h^{2,1}$) and on the dilaton S, whereas it is independent of the Kähler moduli T_m ($m = 1, \ldots, h^{1,1}$).

In type IIB the fluxes generate a C_4 tadpole given by

$$N_{\mathrm{flux}} = \int H_3 \wedge F_3 = \sum_{I=0}^{h^{2,1}} a_I m_I + b_I e_I\,. \tag{2.6}$$

This flux number is equivalent to the Ramomd charge of N_{flux} D3-branes, and has to be cancelled by the orientifold O3-planes and an appropriate number of D3-branes.

To be more precise, type II compactification on CY threefolds results in $\mathcal{N} = 2$ supergravity in four dimensions. One may then introduce orientifolds which reduce the four-dimensional supersymmetry to $\mathcal{N} = 1$. The orientifold involution introduces a grading on the spaces of harmonic forms of the CY, under which each cohomology group H^n splits into even/odd subspaces with respect to the involution: $H^n = H^n_+ \oplus H^n_-$. Accordingly, the dimension of the Kähler moduli space of type IIA/IIB is given by $h^{1,1}_{\mp}$; similarly the dimension of the complex-structure moduli space is truncated to $h^3_{\pm} = 1 + h^{2,1}$. With this understanding, we will omit the $\pm$ subscript on the various CY cohomology groups. We refer the reader to [48] for a detailed discussion of CY orientifold compactifications.

According to our previous discussion (section 2.1), for geometrical CY spaces with $h^{1,1} > 0$, since $\partial_{T_m} W = 0$, it follows that supersymmetry is not compatible with negative vacuum energy. However it is compatible with the flat no-scale models with (un)broken supersymmetry. This holds under the assumption that $\partial_{\phi_i} K \neq 0$, which is certainly fulfilled for the Khler moduli. Supersymmetric AdS_4 vacua in type IIB compactifications with all moduli stabilized seem therefore unlikely, unless one adds a non-perturbative superpotential that depends on the Kähler moduli T_m.

The problem of the Khler-moduli perturbative independence can be in principle resolved either by introducing geometrical fluxes, i.e. torsion, hence abandoning the CY structure, or in backgrounds which are non-geometrical and do not possess *at all* Kähler moduli (i.e. $h^{1,1} = 0$ in the framework of CY). These may be thought of as asymmetric

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orbifolds, or as the mirror of the Z-manifold, and have been recently discussed in a different context [49].[3] In this case there are no associated F-term conditions, and supersymmetric AdS_4 can be obtained in type IIB directly from the tree-level flux superpotential (eq. (2.5)) with S and all complex-structure moduli U_i stabilized (see section 2.3 for the detailed computation on the type IIA side). Concerning the models with torsion, a new example will be discussed in section 4.2.2.

Type IIA The moduli dependence is different in type IIA compactifications. Even in the absence of geometrical fluxes, the problem of the Khler-moduli perturbative independence does not generically occur. Situations of this type, involving only Neveu-Schwarz and Ramond fields, will be examined later. Furthermore, in most cases, the inclusion of geometrical fluxes is not an option but becomes compulsory. Indeed, in type IIA the backreaction due to the Ramond and Neveu-Schwarz fluxes is usually so strong that the underlying geometry Y is no longer a CY space, but rather a space with a certain SU(3) or SU(2) group structure. These spaces have generically $\mathrm{d}J \neq 0$, J being the Kähler form, which leads precisely to geometrical-flux contributions.[4] Those reinforce the Khler-moduli dependence of the perturbative effective superpotential.

Notice that there are nonetheless examples where not all moduli are present. The type IIA mirror of the Khler-moduli-free example cited above has no complex-structure moduli (in the CY language this corresponds to $\tilde{h}^{2,1} = 0$). Although this is not a priori necessary for the perturbative stabilization of all moduli, we will discuss this particular scenario in section 2.3 because stabilization occurs in this case thanks to the particular structure of the Khler potential.

The type IIA effective superpotential receives three kinds of contributions (see e.g. [51–54]):

$$W_{\mathrm{IIA}} = W_H + W_F + W_{\mathrm{geom}} \,. \tag{2.7}$$

The first term is due to the Neveu-Schwarz 3-form fluxes and depends on the dilaton S and the type IIA complex-structure moduli U_m ($m = 1, \ldots, \tilde{h}^{2,1}$):

$$W_H(S,U) = \int_Y \Omega_c \wedge H_3 = i\tilde{a}_0 S + i\tilde{c}_m U_m \,, \tag{2.8}$$

where in type IIA the 3-form Ω_c is defined by $\Omega_c = C_3 + i\mathrm{Re}(C\Omega)$. Second, we have the contribution from Ramond 0-, 2-, 4-, 6-form fluxes (the 0-form flux corresponds to the mass parameter $\tilde{m}_0$ in massive IIA supergravity — see also [18]):

$$\begin{aligned} W_F(T) &= \int_Y \mathrm{e}^{J_c} \wedge F^{\mathrm{R}} \\ &= \tilde{m}_0 \frac{1}{6} \int_Y (J_c \wedge J_c \wedge J_c) + \frac{1}{2} \int_Y \left(F_2^{\mathrm{R}} \wedge J_c \wedge J_c\right) + \int_Y F_4^{\mathrm{R}} \wedge J_c + \int_Y F_6^{\mathrm{R}} \\ &= i\tilde{m}_0 F_0(T) - \tilde{m}_i F_i(T) + i\tilde{e}_i T_i + \tilde{e}_0 \,. \end{aligned} \tag{2.9}$$

[3] It is fair to say that the existence of such compactifications in type IIB is still conjectural. It is inferred by T-duality from type IIA, where a precise geometric construction is available, based on the identification of all complex-structure moduli with S. This construction was also discussed some time ago in [50].

[4] These geometrical fluxes originate from the mirror transformation (fibre-wise T-duality) of the Neveu-Schwarz 3-form fluxes a_i on the type IIB side.

Here $F(T) := F_0(T)$ is the type IIA prepotential, which depends on the IIA Kähler moduli T_i ($i = 1, \ldots, \tilde{h}^{1,1}$) and $F_i(T) := \partial F_0/\partial T_i$. We use the notation J_c for the complexified Kähler metric $J_c := B + iJ$. Finally we have the contribution of the geometrical (metric) fluxes, which captures the non-Calabi-Yau property of Y:

$$W_{\text{geom}}(S,T,U) = i\int_Y \Omega_c \wedge \mathrm{d}J = -\tilde{a}_i S T_i - \tilde{d}_{im} T_i U_m \,, \tag{2.10}$$

where the metric fluxes $\tilde{a}_i$, $\tilde{d}_{im}$ parameterize the non-vanishing of $\mathrm{d}J$.

The type IIA Ramond tadpole follows from the equation of motion of the field C_7. Specifically it is of the form [54]:

$$\tilde{N}_{\text{flux}} = \int \left(C_7 \wedge \mathrm{d}F_2 + C_7 \wedge (\tilde{a}_0 H_3 + \mathrm{d}\bar{F}_2) \right) \,, \tag{2.11}$$

where $G_2 = \mathrm{d}C_1 + \tilde{a}_0 B_2 + \bar{F}_2$ and ${}^*F_2 = F_8 = \mathrm{d}C_7$. The metric fluxes $\tilde{a}_i$ contribute to $\mathrm{d}\bar{F}_2$, and one gets for non-vanishing fluxes $\tilde{a}_I$ and $\tilde{m}_I$ that

$$\tilde{N}_{\text{flux}} = \sum_{I=0}^{\tilde{h}^{1,1}} \tilde{a}_I \tilde{m}_I \,. \tag{2.12}$$

This non-vanishing flux tadpole, which corresponds to a non-vanishing D6-brane charge, must be cancelled by the orientifold O6-planes and an appropriate number of D6-branes:

$$\tilde{N}_{\text{flux}} + N_{\text{D6}} = 2N_{\text{O6}} \,. \tag{2.13}$$

The mirror map Let us now assume that the two spaces X and Y are mirror (T-)dual to each other. This implies that $h^{1,1} = \tilde{h}^{2,1}$, $h^{2,1} = \tilde{h}^{1,1}$. Moroever we identify $S^{\text{IIA}} = S^{\text{IIB}}$, $T^{\text{IIA}} = U^{\text{IIB}}$, $U^{\text{IIA}} = T^{\text{IIB}}$. Then we can see that on the various fluxes in the type IIA/B superpotential mirror symmetry acts as follows. On the one hand, all Ramond fluxes are mapped onto each other, i.e.

$$e_I = \tilde{e}_I \,, \quad m_I = \tilde{m}_I \,. \tag{2.14}$$

On the other hand, Neveu-Schwarz 3-fluxes in type IIB are generically mapped on metric fluxes in type IIA and vice versa. The precise mapping is determined by the orientation of the T^3 in X on which the T-duality is acting. We assume that the IIB Neveu-Schwarz 3-form flux a_0 remains a Neveu-Schwarz 3-form flux in type IIA, whereas the Neveu-Schwarz 3-form fluxes a_i become metric fluxes. Specifically we get

$$a_I = \tilde{a}_I \,. \tag{2.15}$$

The type IIB Neveu-Schwarz 3-form fluxes b_I as well as the type IIA Neveu-Schwarz 3-form fluxes $\tilde{c}_m$ and metric fluxes $\tilde{d}_{im}$ do also have their mirror-duals. We will not elaborate any further on the precise correspondence, but mention that tadpoles (eq. (2.12)) do also match when using the full mirror dictionary (including (2.14) and (2.15)).

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2.3 Supersymmetric AdS$_4$ vacua in type IIA and their IIB mirror duals

We now come to the crucial point of generating supersymmetric AdS$_4$ ground states in type IIA with all main moduli stabilized. One of the few instances where such vacua were constructed is provided in [51]. The present construction is somewhat different and we will comment on this later, in particular in section 3.

The situation we will consider falls in the class where the assumption $\partial_{\phi_i} K|_{\rm min} \neq 0$ holds. Therefore as emphasized in section 2.1 the superpotential must depend on all chiral fields for the vacuum energy to be negative with unbroken supersymmetry. Following [51, 53, 54] we will concentrate on the case without metric fluxes, i.e. $\tilde{a}_i = \tilde{d}_{im} = 0$. Furthermore, the 6-form fluxes as well as the 2-form fluxes can be shown to be gauge dependent and hence can be set to zero: $\tilde{e}_0 = \tilde{m}_i = 0$ [53, 54]. The fluxes $\tilde{m}_0$ and $\tilde{a}_0$ must be non-zero for W to be kept non-vanishing. Finally, we combine $\tilde{e}_i$ and $\tilde{m}_0$ as

$$\gamma_i = \tilde{m}_0 \tilde{e}_i \,. \tag{2.16}$$

Assuming a simple (toroidal) cubic prepotential $F = T_1 T_2 T_3$, the superpotential has the generic form:

$$\begin{aligned} W_{\rm IIA} &= W_F + W_H = \tilde{m}_0 \int_Y (J \wedge J \wedge J) + \int_Y F_4^{\rm R} \wedge J + \int_Y \Omega_c \wedge H_3 \\ &= i\tilde{e}_i T_i + i\tilde{m}_0 T_1 T_2 T_3 + i\tilde{a}_0 S + i\tilde{c}_m U_m \,. \end{aligned} \tag{2.17}$$

Using eq. (2.12), the D6-tadpole of corresponding fluxes is simply given by

$$\tilde{N}_{\rm flux} = \tilde{a}_0 \tilde{m}_0 \,. \tag{2.18}$$

According to eq. (2.13) this number has to be balanced by the D6-branes and the O6-planes.

We may also consider the generalization to the case where the prepotential is given by $F = \frac{1}{6} c_{ijk} T_i T_j T_k$. In CY compactifications, the c_{ijk} would be the classical triple intersection numbers and the corresponding superpotential would read:

$$W_{\rm IIA} = W_F + W_H = i\tilde{e}_i T_i + i\tilde{m}_0 \frac{1}{6} c_{ijk} T_i T_j T_k + i\tilde{a}_0 S + i\tilde{c}_m U_m \,. \tag{2.19}$$

However, in order to keep the algebra simple we will focus in the following on the toroidal prepotential with $c_{ijk} = 1$.

Before analyzing the IIA situation captured by (2.17) we would like to consider the simpler situation we have advertised in section 2.2, namely the compactification without complex-structure moduli. In this case the flux superpotential has the simple form:

$$W_{\rm IIA} = W_F + W_H = i\tilde{e}_i T_i + i\tilde{m}_0 T_1 T_2 T_3 + i\tilde{a}_0 S \,, \tag{2.20}$$

where the last term of (2.17) is now missing because the U's are absent. Furthermore, the Khler potential reads:

$$K = -\log(S + \bar{S})^4 \prod_{i=1}^{3} (T_i + \bar{T}_i). \tag{2.21}$$

It is clear from the latter that this compactification is achieved as an orbifold which effectively identifies all U's and S. The precise dependence of K on the S-field is crucial for the equations $F_S = F_{T_1} = F_{T_2} = F_{T_3} = 0$ to admit a non-trivial solution. This solution is

$$|\gamma_1|T_1 = |\gamma_2|T_2 = |\gamma_3|T_3 = \sqrt{\frac{5|\gamma_1\gamma_2\gamma_3|}{3\tilde{m}_0^2}}\,, \quad S = -\frac{8}{3\tilde{m}_0\tilde{a}_0}\gamma_i T_i\,, \tag{2.22}$$

and describes a supersymmetric AdS_4 vacuum. Note that we must have $\gamma_i < 0$ for this vacuum to exist.

Coming back to the more generic case with superpotential (2.17), and Khler potential $K = -\log(S+\bar{S})\prod_{i=1}^3(T_i+\bar{T}_i)\prod_{i=1}^3(U_i+\bar{U}_i)$, the equations $F_{\phi_i} = 0$ admit the following solution:

$$|\gamma_i|T_i = \sqrt{\frac{5|\gamma_1\gamma_2\gamma_3|}{3\tilde{m}_0^2}}\,, \quad S = -\frac{2}{3\tilde{m}_0\tilde{a}_0}\gamma_i T_i\,, \quad \tilde{c}_m U_m = -\frac{2}{3\tilde{m}_0}\gamma_i T_i\,. \tag{2.23}$$

This solution corresponds again to supersymmetric AdS_4 vacua, as shown in [54].

The above analysis deserves several comments. Contrary to the case studied in [51] where AdS_4 with moduli stabilized was obtained with all fluxes on, here no geometric fluxes are present (see (2.17)). This might seem inconsistent at first glance. Indeed, the geometric fluxes were introduced in ref. [51] as a consequence of the Jacobi identity and the tadpole conditions that the set of fluxes must satisfy. Fluxes are indeed gauging parameters since flux compactifications are gauged supergravities.

The resolution of this puzzle goes as follows. The background analyzed in [51] originated from localized branes and orientifold planes. As already mentioned in the introduction, the vacua under investigation here *do not* admit a ten-dimensional supergravity interpretation with localized brane/orientifold sources [18]. One must consider the addition of *smeared* brane and/or orientifold sources [19]. This will be made more transparent in section 3, when discussing the brane-sources that create the background at hand. It is clear however that in the presence of smeared branes, Jacobi or equivalently Bianchi identities and tadpole conditions are slightly modified and consistency is recovered despite the non-vanishing $a_0 m_0$ (see eq. (2.18)). This issue has been nicely discussed in a similar context by Villadoro and Zwirner [41].

Let us end the present section by discussing some T-dual/mirror transforms of the of the IIA models. As mentioned already, T-duality will in general transform the NS-fluxes into geometrical fluxes. We can for instance investigate within the toroidal models the T-duality transformation in the internal directions x^1 and x^2, acting as $T_1 \to 1/T_1$, $T_{2,3} \to T_{2,3}$. Then the T-dual superpotential of eq. (2.17) becomes

$$W_{\mathrm{IIA}} = \tilde{e}_1 + \tilde{e}_2 T_1 T_2 + \tilde{m}_0 T_2 T_3 + \tilde{e}_3 T_3 T_1 + \tilde{a}_0 S T_1 + \tilde{c}_m T_1 U_m\,. \tag{2.24}$$

The fluxes $\tilde{a}_0$ and $\tilde{c}_m$ become now geometrical. The corresponding AdS_4 ground states can be simply obtained by replacing T_1 by $1/T_1$ in eq. (2.23).

Alternatively let us consider the IIB mirror transform of the superpotential (2.17), which is obtained by applying T-duality transformations in the three directions x^1, x^3

and x^5. This exchanges the IIA Khler moduli by the IIB complex-structure moduli and vice versa. In the presence of the IIA NS-fluxes $\tilde{c}_m$ as in (2.17), the type IIB mirror superpotential will necessarily contain geometrical fluxes:

$$W_{\mathrm{IIB}} = i\tilde{e}_i U_i + i\tilde{m}_0 U_1 U_2 U_3 + i\tilde{a}_0 S + i\tilde{c}_m T_m \,. \tag{2.25}$$

In this case, T-duality takes the system away from the original CY framework. However, we can consider the simpler version with no IIA complex-structure deformations and superpotential (2.20). Now, there are no geometrical fluxes on the mirror IIB side. This model was indeed motivated by type IIB, where it provides the only possibility for perturbatively stabilizing all moduli without geometrical fluxes. The dual geometry is not of the standard form, since the space has no Kähler deformation. In fact, it is a non-geometric space that does not allow for a large-radius description, but rather for CFT or Landau-Ginzburg description [49, 55].[5] Ignoring the absence of a sensible large radius limit, the mirror-dual type IIB superpotential takes the form

$$W_{\mathrm{IIB}} = \int_X \Omega \wedge \left(F_3^{\mathrm{R}} + S H_3^{\mathrm{NS}}\right) = ie_i U_i + im_0 U_1 U_2 U_3 + ia_0 S \,. \tag{2.26}$$

Minimization of the scalar potential goes as previously explained for type IIA (see eqs. (2.22)) and this type IIB superpotential leads to supersymmetric AdS_4 vacua with the U_i and S stabilized.

3. The source-brane picture

As we discussed in the previous section, supersymmetric AdS_4 can appear as ground states of type IIA flux compactifications. In the simplest case, it was enough to turn on a Neveu-Schwarz 3-form flux as well as a Ramond 4-form flux in gauged IIA supergravity with mass parameter $\tilde{m}_0$ (see eqs. (2.17) and (2.20)). Generically, however, metric fluxes are necessary. We have also stressed that Bianchi identities — or equivalently, Jacobi identities in the gauged-supergravity language — and tadpole conditions are always satisfied. This occurs sometimes in a subtle way, in particular when certain sources are present.

Our aim here is precisely to characterize the sources that generate the fluxes necessary for the compactifications under consideration. This complementary, or dual picture, gives another perspective to the emergence of AdS_4. The latter appears as near-horizon geometry of a certain distribution of intersecting/smeared branes and calibrated sources that act as domain walls, connecting AdS_4 to an asymptotically flat region. It also sheds light on the origin of the various terms that contribute to the Bianchi identities or the tadpole conditions. Finally, the brane picture is the first step towards the counting of microscopic states. From the viewpoint of four-dimensional gauged supergravity, one could presumably go further and consider the attractor equations and the macroscopic entropies. All this is outside the scope of the present paper.

[5] These spaces were discussed by Candelas et al. some time ago [50].

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3.1 The brane origin of the fluxes

As we will see explicitly in section 4, the appearance of AdS_4 as near-horizon geometry requires that all branes have two common spatial directions in non-compact four-dimensional spacetime, i.e. they have the geometry of a domain wall in four dimensions. Moreover, depending on their dimensionality, the branes will fill part of the internal space M_6. Keeping this structure in mind, let us summarize, the relations between the various fluxes and the corresponding source branes.

- For a Neveu-Schwarz 3-form flux H_3 through a 3-cycle Σ_3 inside M_6 the sources are NS5-branes wrapped around the dual 3-cycle $\tilde{\Sigma}_3$.
- In the Ramond sector we have fluxes of the Ramond field strengths F_n^{R} through some internal n-dimensional cycles Σ_n. The desired domain-wall configuration in spacetime, requires that these fluxes be generated by magnetic brane sources, namely by D$(8-n)$-branes, wrapped around internal cycles $\tilde{\Sigma}_{6-n}$ dual to Σ_n.
- For geometrical fluxes we expect to have Kaluza-Klein monopoles as sources. In fact, performing T-duality to directions orthogonal to the NS5-brane, one obtains a KK-monopole.

The fluxes are quantized and this reflects that the number of branes is not arbitrary. When M_6 is a CY, these numbers are related to $h^{2,1}$ or $h^{1,1}$.

3.2 Supersymmetric intersecting branes and calibrated sources

Let us briefly review the salient features of supersymmetric configurations of intersecting branes in supergravity, and clarify the issues of Bianchi identities and tadpole conditions in the presence of sources. For a comprehensive exposition of the subject the reader is referred to the literature [56 – 63].

Before we proceed, it will be useful to divide the spacetime coordinates into the following parts: (a) overall world-volume coordinates, which are common to all branes in the system; (b) overall transverse coordinates, which are transverse to all branes; (c) relative transverse coordinates, which are transverse to some (but not all) of the constituent branes. Supergravity solutions of intersecting branes can be built according to the “harmonic superposition rules” [56]. The solution in this case is constructed in terms of harmonic functions depending on the overall transverse coordinates alone. The systems we consider in the following will all be of this type.

In addition to the D-branes and NS5-branes of type II supergravity, the solutions we present in the following include supersymmetric (and therefore calibrated), smeared sources. Our analysis is based on the following observation, which holds under certain mild assumptions [64]: *any supersymmetric configuration with supersymmetric (and therefore calibrated), possibly smeared, sources will automatically obey the dilaton and Einstein equations (appropriately modified by the inclusion of the sources), provided the source-modified Bianchi identities and form equations of motion are also satisfied.* The analogue of this result in the absence of sources was shown for IIA in [18] and for type IIB in [65].

The cases of eleven-dimensional supergravity, heterotic, type I are examined in [66–68] respectively.

Let us consider Romans IIA supergravity [69]. The latter includes a massive parameter m which, from the point-of-view of the "democratic formulation" of IIA [29], can be thought of as conjugate (but not, in general, equal) to the field strength of a 9-form potential C_9. Schematically [21]:

$$\frac{\delta S_{\rm bulk}}{\delta m} = \star F_{10}\,, \tag{3.1}$$

where $S_{\rm bulk}$ is the action of massive IIA, and $F_{10} := {\rm d}C_9$. In all the examples we consider next, m is in fact equal to $\star F_{10}$. This has in particular the following consequences: (a) the supersymmetry variations of the fermions are as in Romans supergravity, with m thought of as a field strength F_0 — the Hodge-dual of F_{10}; (b) m need no longer be constant, i.e. F_0 need not be closed. The failure of m to be constant is parameterized by a one-form j_8 — the Poincaré dual to the world-volume of an 8-brane magnetic source:

$$ {\rm d}F_0 = j^{D8}\ . \tag{3.2}$$

More specifically, in the section 4 we present various supersymmetric bosonic solutions of IIA supergravity corresponding to intersecting branes. In particular, these configurations solve the supersymmetry equations[6]

$$\begin{aligned} 0 &= \left(\nabla_\mu + \frac{1}{4}H_\mu\mathcal{P} + \frac{{\rm e}^\Phi}{16}\sum_n F_n\Gamma_\mu\mathcal{P}_n\right)\epsilon\,,\\ 0 &= \left(\partial\Phi + \frac{1}{2}H\mathcal{P} + \frac{{\rm e}^\Phi}{8}\sum_n(-1)^n(5-n)F_n\mathcal{P}_n\right)\epsilon\ . \end{aligned}\tag{3.3}$$

In addition, all forms satisfy the Bianchi identities as well as their equations of motion — possibly modified by the inclusion of calibrated, smeared sources:

$$\begin{aligned} {\rm d}F + H\wedge F &= -Q\ j\\ {\rm d}\star F - H\wedge\star F &= Q\ \alpha(j)\,, \end{aligned}\tag{3.4}$$

where Q is the charge and j is the Poincaré dual to the source; the operator α acts by reversing the indices. In the above we have used polyform notation for F and j. The analysis of [64] then guarantees that the remaining equations (dilaton, Einstein) are automatically satisfied.[7]

We remark that the supersymmetry transformations (3.3) are not modified by the inclusion of sources. The latter does, however, entail a modification of the Bianchi identities as above. Equation (3.4) can be accounted for by adding to the bulk supergravity

[6] The transformations (3.3) are given in the democratic formulation; the reader is referred to [29, 64] for a detailed explanation of the notation.

[7] For the analysis of [64] to go through, one needs to check in addition that the mixed space/time components of the Einstein equations $E_{i0} = 0$ are satisfied. In all the examples considered in section 4 of the present paper, this is indeed the case. This readily follows from the fact that, in each example, either (a) all fields are static, all non-zero field-strength components are purely spatial, the metric is diagonal, or (b) the example is T-dual to one for which (a) holds.

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Lagrangian ($S_{\rm bulk}$) a source term schematically of the form:

$$S_{\rm source} = Q \int C \wedge \alpha(j) - T \int \Phi \wedge \alpha(j)\,, \tag{3.5}$$

so that our solutions correspond to supersymmetric stationary points of $S_{\rm bulk}+S_{\rm source}$. The first term on the right-hand side of (3.5) induces the modification to the Bianchi identities. The second term on the right-hand side accounts for the coupling of the calibrated source to the bulk graviton, and therefore induces a modification of the Einstein equation; Φ denotes the calibration form. In addition, the charge (Q) and tension (T) of the source obey a BPS condition $T = \pm Q$. The fact that supersymmetric D-branes and orientifolds can be thought of as (generalized) calibrated sources was shown in [70].

3.3 The tadpole equations

Any self-consistent system of spacetime-filling branes must obey the tadpole cancellation condition, which is a consequence of the (generalized) Gauß law. Alternatively, in supersymmetric configurations, this condition can be thought of as arising from the integrated Bianchi identities (3.4). Specifically for the D4/D8/NS5 example of section 4.2.1, we will need the D6/O6 tadpole cancellation condition:

$$\frac{1}{2\pi\sqrt{\alpha'}} \int_\Sigma F_0\; H_3 + N_{\rm D6} - 2N_{\rm O6} = 0\,, \tag{3.6}$$

where $N_{\rm D6}$, $N_{\rm O6}$ is the total number spacetime-filling D6-branes, O6-planes wrapping a three-cycle Σ in the internal space.

Note that from the point of view of this supergravity analysis there is an ambiguity in the interpretation: ignoring higher-order derivative corrections and possible world-volume excitations on the sources, only the difference $N_{\rm D6} - 2N_{\rm O6}$ can be determined — not the individual numbers $N_{\rm D6}$, $N_{\rm O6}$.

4. Specific examples

We are now ready to come to our explicit examples. To facilitate the comparison with the standard form of intersecting-brane solutions constructed according to the harmonic superposition rules, all solutions in the following subsections are presented in the string frame.

Keeping in mind the above subtleties that arise in the presence of sources, we will now consider a number of configurations of intersecting branes and (calibrated) sources in type IIA supergravity. They generate backgrounds which interpolate between flat space and AdS$_4$, the latter appearing as the near-horizon geometry. The examples we consider here fall into two classes:

- Configurations corresponding to supergravity vacua with many flat directions (S and U_m fields). In all these examples, the dilaton goes to zero at the horizon. These include:

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- The double D4 distribution, where we consider two stacks of intersecting D4-branes. The corresponding background can be thought of as a vacuum of IIA supergravity with superpotential (2.17) and vanishing $\tilde{m}_0, \tilde{a}_0, \tilde{c}_m$: $W_{\rm IIA} = i\tilde{e}_i T_i$. Not all moduli are stabilized in this case.
- Similarly, an appropriate D2/D6 distribution generates AdS$_4$, but the superpotential which provides this vacuum receives contributions from F_2 and F_6 (see eq. (2.9)): $W_{\rm IIA} = -\tilde{m}_i T_j T_k + \tilde{e}_0$. Again there are flat directions left.
- Performing a T-duality in the appropriate direction amounts to transforming T_1 to $1/T_1$ while leaving $T_{2,3}$ invariant. In this way we can map the previous configuration onto D4/D8 with superpotential $W_{\rm IIA} = \tilde{e}_0 T_1 - \tilde{m}_3 T_2 - \tilde{m}_2 T_3 - \tilde{m}_1 T_1 T_2 T_3$, i.e. (2.17) with $\tilde{a}_0 = \tilde{c}_m = 0$. Again, not all moduli can be fixed.

- Configurations where all moduli are stabilized. For these cases it is significant that whenever stabilization is complete, the values of the moduli found by minimizing the scalar potential are recovered by a careful analysis of the spacetime background fields near the horizon. In particular the dilaton approaches a finite constant in this limit. We will show explicitly how this works in the example of D4/D8/NS5 distribution.

  The examples under consideration here are the following:

  - Configuration with D4/D8/NS5 branes. This model contains in particular four stacks of intersecting NS5-branes. The background is a IIA ground state of the superpotential (2.17) with *all* terms non-vanishing: $W_{\rm IIA} = i\tilde{e}_i T_i + i\tilde{m}_0 T_1 T_2 T_3 + i\tilde{a}_0 S + i\tilde{c}_m U_m$. This allows for full moduli stabilization (eqs. (2.23)).
  - The next model is obtained by performing a T-duality along one direction (say x^1). This is a type IIB model with D3/D5/D7/NS5/KK-branes/monoples. Its superpotential, generated by F_1, F_3, F_5, H_3 and geometric fluxes, reads $W_{\rm IIB} = i(\tilde{e}_1 U_1 + \tilde{c}_2 U_2 + \tilde{c}_3 U_3 + \tilde{c}_1 T_1 + \tilde{e}_2 T_2 + \tilde{e}_3 T_3 + \tilde{m}_0 U_1 T_2 T_3 + \tilde{a}_0 S)$, and it exhibits an AdS$_4$ vacuum with all moduli stabilized. The latter appears as the near-horizon geometry of the brane/monopole configuration at hand. A further T-duality[8] allows to reach a new type IIA model generated by intersecting D2/D6/KK-branes/monopoles with superpotential $W_{\rm IIA} = \tilde{e}_1 + \tilde{e}_2 T_1 T_2 + \tilde{e}_3 T_1 T_3 + \tilde{m}_0 T_2 T_3 + \tilde{a}_0 S T_1 + \tilde{c}_m T_1 U_m$. This superpotential is indeed given in (2.17) with F_2 and F_6 (2.9) and geometric fluxes (2.10). All moduli are stabilized in this AdS$_4$ mirror vacuum.
  - Finally, we can simplify the first of the above setups by keeping only one out of four stacks of NS5-branes in the D4/D8/NS5. This still provides an AdS$_4$ near-horizon geometry, but the corresponding superpotential (2.20) has now vanishing $\tilde{c}_m$: $W_{\rm IIA} = i\tilde{e}_i T_i + i\tilde{m}_0 T_1 T_2 T_3 + i\tilde{a}_0 S$. In this model, the IIA side was free of complex-structure moduli. The type IIB dual of the brane configuration is a D5/NS5-brane distribution which is not the non-geometric type IIB vacuum

[8] In total we perform a T-duality along x^1 and x^2, which amounts to trading T_1 for $1/T_1$.

free of Khler moduli studied in sections 2.2 and 2.3. The D5/NS5 has Khler moduli and these are not stabilized.

4.1 Examples with unstabilized moduli

4.1.1 D4

Let us consider the following system of intersecting D4-branes:

	ξ^0	ξ^1	ξ^2	y^1	y^2	y^3	x^1	x^2	x^3	x^4
D4	$\otimes$	$\otimes$	$\otimes$				$\otimes$	$\otimes$		
D4′	$\otimes$	$\otimes$	$\otimes$						$\otimes$	$\otimes$

We have used the symbol $\otimes$ to denote a direction along the brane; the boxes corresponding to directions perpendicular to the branes have been left blank. The two D4-branes intersect on a two-brane along the directions $\xi^\mu,~\mu=0,1,2$. The overall transverse directions (i.e. the directions transverse to both D4-branes) are denoted by $y^i,~i=1,2,3$, whereas the relative transverse directions (i.e. the directions transverse to one of the D4-branes but parallel to the other) are denoted by $x^a,~a=1,\dots,4$, and may be taken to parameterize a T^4. Note, however, that the D4-branes we consider here will be *smeared* in both relative and overall transverse directions.

According to the general procedure described in section 3.2, the following can be seen to be a supersymmetric solution of IIA supergravity in the presence of calibrated sources:

$$\begin{aligned}
&\mathrm{d}s^2_{10}=\frac{1}{\sqrt{H_1H_2}}\,\eta_{\mu\nu}\mathrm{d}\xi^\mu\mathrm{d}\xi^\nu+\sqrt{H_1H_2}\,\delta_{ij}\mathrm{d}y^i\mathrm{d}y^j+\sqrt{\frac{H_2}{H_1}}\sum_{a=1}^{2}(\mathrm{d}x^a)^2+\sqrt{\frac{H_1}{H_2}}\sum_{a=3}^{4}(\mathrm{d}x^a)^2\\
&\mathrm{e}^{-2\phi}=\sqrt{H_1H_2},\quad F_{x^3x^4y^iy^j}=-\delta^{k\ell}\varepsilon_{ijk}\partial_{y^\ell}H_1,\quad F_{x^1x^2y^iy^j}=-\delta^{k\ell}\varepsilon_{ijk}\partial_{y^\ell}H_2\,,
\end{aligned}\tag{4.1}$$

where ($y^2=\delta_{ij}y^iy^j$)

$$4\pi H_\alpha=1+\frac{c_\alpha}{y}+\frac{d_\alpha}{2y^2},\quad \alpha=1,2 \tag{4.2}$$

and c_α, d_α are positive constants. Moreover, the Killing spinor is constrained to satisfy

$$\begin{aligned}
&\Gamma_{y^1y^2y^3x^1x^2}\epsilon=-\epsilon\\
&\Gamma_{y^1y^2y^3x^3x^4}\epsilon=-\epsilon\,,
\end{aligned}\tag{4.3}$$

and therefore the system preserves one-fourth of the original supersymmetry (eight real supercharges, or $\mathcal{N}=1$ in six dimensions).

It is important to stress here that the functions H_α are *not* harmonic in the overall transverse directions, reflecting the presence of smeared sources. In order to assign a physical interpretation to the above configuration note that, as follows from (4.1),

$$\begin{aligned}
&3\partial_{[y^i}F_{y^jy^k]x^1x^2}=\varepsilon_{ijk}j_2\\
&3\partial_{[y^i}F_{y^jy^k]x^3x^4}=\varepsilon_{ijk}j_1\,,
\end{aligned}\tag{4.4}$$

where

$$
\begin{aligned}
j_\alpha := -\nabla^2 H_\alpha &= \frac{1}{4\pi y^2}\frac{\partial}{\partial y}\left(c_\alpha + \frac{d_\alpha}{y}\right) \\
&= c_\alpha \delta^3(y) - \frac{d_\alpha}{4\pi y^4}\,,
\end{aligned}
\tag{4.5}
$$

and the nabla acts in the overall transverse directions. Hence, j_1 is the *charge density* of a four-brane magnetic source along $\xi^0, \xi^1, \xi^2, x^1, x^2$, smeared along the overall transverse radial direction y. Similarly, j_2 is the charge density of a smeared four-brane magnetic source along $\xi^0, \xi^1, \xi^2, x^3, x^4$. The total charge can be read off from (4.5):

$$
Q_\alpha = c_\alpha - \frac{d_\alpha}{\varepsilon}\,,
\tag{4.6}
$$

and we have introduced a cutoff $y \geq \varepsilon$. The configuration at hand could be interpreted as consisting of smeared D4-branes with "bear" charges $c_\alpha - d_\alpha/\varepsilon$; the divergence of the charges could be interpreted as due to the smearing, and one could define finite "renormalized" charges by subtracting the infinite part.

The blowing-up of the D-brane charges and the continuous charge distributions (4.5) along the non-compact y-direction are pathological. Moreover, it should be emphasized that smearing along compact directions is qualitatively different from smearing along non-compact ones: the latter would generally invalidate a Kaluza-Klein reduction. Our point-of-view here is that this pathological behaviour reflects the fact that the present example corresponds to a setup with unstabilized moduli. In section 4.2 we will present examples corresponding to setups with stabilized moduli, which do not suffer from the pathologies mentioned above.

The near-horizon limit can be obtained by taking $y \to 0$. The ten-dimensional metric takes the form

$$
\begin{aligned}
\mathrm{d}s^2_{\mathrm{NH}} =& \frac{y^2}{\sqrt{D_1 D_2}}\,\eta_{\mu\nu}\mathrm{d}\xi^\mu \mathrm{d}\xi^\nu + \sqrt{D_1 D_2}\,\frac{\mathrm{d}y^2}{y^2} + \sqrt{D_1 D_2}\,\mathrm{d}\Omega_2^2 \\
&+ \sqrt{\frac{D_2}{D_1}}\,\sum_{a=1}^{2}(\mathrm{d}x^a)^2 + \sqrt{\frac{D_1}{D_2}}\,\sum_{a=3}^{4}(\mathrm{d}x^a)^2\,,
\end{aligned}
\tag{4.7}
$$

where $\mathrm{d}\Omega_2^2$ is the metric of the unit two-sphere in the overall transverse directions, and we have set $D_\alpha := d_\alpha/8\pi$. We can readily recognize (4.7) as the metric of $\mathrm{AdS}_4 \times S^2 \times T^4$. At $y \to \infty$ the metric asymptotes $\mathbb{R}^{1,5} \times T^4$. Finally note that, as follows from (4.1), (4.2), the string coupling $g_\mathrm{s} = \mathrm{e}^\phi$ goes to zero as $y \to 0$. This is a typical runaway behaviour. In order to stabilize the string coupling to a finite value one has to add NS5-branes, as we will see in section 4.2.

4.1.2 D2/D6

This configuration can be obtained from the previous one by applying two T-duality transformations along directions x^1, x^2. More explicitly, the branes are as follows:

	ξ^0	ξ^1	ξ^2	y^1	y^2	y^3	x^1	x^2	x^3	x^4
D2	$\otimes$	$\otimes$	$\otimes$							
D6	$\otimes$	$\otimes$	$\otimes$				$\otimes$	$\otimes$	$\otimes$	$\otimes$

The metric now reads:

$$\mathrm{d}s^2_{10} = \frac{1}{\sqrt{H_1 H_2}}\,\eta_{\mu\nu}\mathrm{d}\xi^\mu \mathrm{d}\xi^\nu + \sqrt{H_1 H_2}\;\delta_{ij}\mathrm{d}y^i\mathrm{d}y^j + \sqrt{\frac{H_1}{H_2}}\sum_{a=1}^{4}(\mathrm{d}x^a)^2\,, \tag{4.8}$$

with H_α as in (4.2). As in the previous subsection, in the near-horizon limit the metric approaches $\mathrm{AdS}_4 \times S^2 \times T^4$ and the string coupling goes to zero.

4.1.3 D4/D8

Extra T-duality along y^2, y^3 leads to the following D4/D8 brane distribution:

	ξ^0	ξ^1	ξ^2	y^1	y^2	y^3	x^1	x^2	x^3	x^4
D4	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$				
D8	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$

The discussion of this model is analogous to the previous two cases.

4.2 Examples with all moduli stabilized

4.2.1 D4/D8/NS5

We now come to the domain-wall solution of most interest, since it may be thought of as the brane-dual to the supersymmetric AdS_4 vacuum considered in section 2.3. Consider the following system of intersecting D4/D8/NS5-branes:

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
D4	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$				
D4$'$	$\otimes$	$\otimes$	$\otimes$				$\otimes$	$\otimes$		
D4$''$	$\otimes$	$\otimes$	$\otimes$						$\otimes$	$\otimes$
NS5	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$		$\otimes$	
NS5$'$	$\otimes$	$\otimes$	$\otimes$		$\otimes$			$\otimes$		$\otimes$
NS5$''$	$\otimes$	$\otimes$	$\otimes$			$\otimes$		$\otimes$	$\otimes$	
NS5$'''$	$\otimes$	$\otimes$	$\otimes$			$\otimes$	$\otimes$			$\otimes$
D8	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$

These generate a supersymmetric solution of IIA supergravity in the presence of supersymmetric (calibrated) sources. As we will see shortly, the tadpole cancellation condition induces in addition O6-planes and/or D6-branes. These brane distributions should thus be referred to as D4/D8/NS5/O6/D6, which we will avoid for simplicity, keeping only D4/D8/NS5.

Background fields and tadpole cancellation. The explicit form of the solution is given by:

$$
\begin{aligned}
\mathrm{d}s_{10}^2 =& \left\{ H^{\mathrm{D8}} \left(\prod_{\alpha=1}^{3} H_\alpha^{\mathrm{D4}} \right) \right\}^{-\frac{1}{2}} \eta_{\mu\nu} \mathrm{d}\xi^\mu \mathrm{d}\xi^\nu \\
&+ \left(\prod_{\alpha=1}^{4} H_\alpha^{\mathrm{NS5}} \right) \left\{ H^{\mathrm{D8}} \left(\prod_{\alpha=1}^{3} H_\alpha^{\mathrm{D4}} \right) \right\}^{\frac{1}{2}} \mathrm{d}y^2 \\
&+ \sqrt{\frac{H_2^{\mathrm{D4}} H_3^{\mathrm{D4}}}{H_1^{\mathrm{D4}} H^{\mathrm{D8}}}} \left\{ H_3^{\mathrm{NS5}} H_4^{\mathrm{NS5}} (\mathrm{d}x^1)^2 + H_1^{\mathrm{NS5}} H_2^{\mathrm{NS5}} (\mathrm{d}x^2)^2 \right\} \\
&+ \sqrt{\frac{H_1^{\mathrm{D4}} H_3^{\mathrm{D4}}}{H_2^{\mathrm{D4}} H^{\mathrm{D8}}}} \left\{ H_2^{\mathrm{NS5}} H_3^{\mathrm{NS5}} (\mathrm{d}x^3)^2 + H_1^{\mathrm{NS5}} H_4^{\mathrm{NS5}} (\mathrm{d}x^4)^2 \right\} \\
&+ \sqrt{\frac{H_1^{\mathrm{D4}} H_2^{\mathrm{D4}}}{H_3^{\mathrm{D4}} H^{\mathrm{D8}}}} \left\{ H_2^{\mathrm{NS5}} H_4^{\mathrm{NS5}} (\mathrm{d}x^5)^2 + H_1^{\mathrm{NS5}} H_3^{\mathrm{NS5}} (\mathrm{d}x^6)^2 \right\} ; \\
\mathrm{e}^{2\phi} =& \left(\prod_{\alpha=1}^{4} H_\alpha^{\mathrm{NS5}} \right) \left(\prod_{\alpha=1}^{3} H_\alpha^{\mathrm{D4}} \right)^{-\frac{1}{2}} \left(H^{\mathrm{D8}} \right)^{-\frac{5}{2}} ; \\
H_{x^2x^4x^6} =& -\partial_y H_1^{\mathrm{NS5}} \left(H^{\mathrm{D8}} \right)^{-1} ; \qquad H_{x^2x^3x^5} = -\partial_y H_2^{\mathrm{NS5}} \left(H^{\mathrm{D8}} \right)^{-1} ; \\
H_{x^1x^3x^6} =& -\partial_y H_3^{\mathrm{NS5}} \left(H^{\mathrm{D8}} \right)^{-1} ; \qquad H_{x^1x^4x^5} = -\partial_y H_4^{\mathrm{NS5}} \left(H^{\mathrm{D8}} \right)^{-1} ; \\
F_{x^3x^4x^5x^6} =& \partial_y H_1^{\mathrm{D4}} ; \qquad F_{x^1x^2x^5x^6} = \partial_y H_2^{\mathrm{D4}} ; \\
F_{x^1x^2x^3x^4} =& \partial_y H_3^{\mathrm{D4}} ; \qquad F = -\partial_y H^{\mathrm{D8}} \left(\prod_{\alpha=1}^{4} H_\alpha^{\mathrm{NS5}} \right)^{-1} .
\end{aligned}
\tag{4.9}
$$

The ten-dimensional Killing spinor is constrained to satisfy

$$
\begin{aligned}
\Gamma_y \epsilon &= \epsilon \\
\Gamma_{x^1x^2x^3x^4} \epsilon &= -\epsilon \\
\Gamma_{x^3x^4x^5x^6} \epsilon &= -\epsilon \\
\Gamma_{x^2x^4x^6} \epsilon &= -\Gamma_{11} \epsilon ,
\end{aligned}
\tag{4.10}
$$

therefore the system preserves one-sixteenth of the original supersymmetry (two real supercharges, or $\mathcal{N} = 1$ in three dimensions). More details can be found in appendix A.

It can be explicitly checked that (4.9) solves the supersymmetry equations as well as the equations-of-motion for the forms. In addition, the fluxes obey source-modified Bianchi identities. It then follows from the integrability theorem of [64] that the correspondingly source-modified dilaton and Einstein equations are automatically obeyed.[9] In other words, (4.9) is a *family* of solutions of type IIA supergravity, parameterized by the

[9]Strictly speaking this statement was only proven in [64] for D-branes and orientifold planes; we expect that a generalization to include NS5-brane sources should be straightforward.

functions $H_\alpha^{\rm D4}$, $H^{\rm D8}$, $H_\alpha^{\rm NS5}$. The modification of the Bianchi identities due the presence of the sources is given by

$$\begin{aligned}\partial_y F_4 &= j^{\rm D4}\\ \partial_y F_0 &= j^{\rm D8}\\ \partial_y H_3 &= j^{\rm NS5}\,,\end{aligned} \tag{4.11}$$

where the charge densities j are read off from (4.9) and the equations above. The sourceless case corresponds to setting the right-hand side to zero, $j = 0$.

We observe from (4.9) that in the "massless" IIA case (corresponding to $F_0 = 0$, $H^{\rm D8} = \text{const.}$) the F_4, H_3 fluxes satisfy their corresponding sourceless Bianchi identities if and only if the functions $H_\alpha^{\rm D4}$, $H_\alpha^{\rm NS5}$ are harmonic in the total transverse space (i.e. the y-direction). This is no longer true in the massive case, $F_0 \neq 0$. Note also that the solution has $F_2 = 0$. Consistency with the tadpole cancellation condition (3.6) then induces the presence of a non-zero charge density corresponding to O6-planes and/or D6-branes:[10]

$$j_\alpha^{\rm O6/D6} = \partial_y H^{\rm D8}\ \partial_y H_\alpha^{\rm NS5}\left(H^{\rm D8}\prod_{\beta=1}^{4} H_\beta^{\rm NS5}\right)^{-1}\,, \tag{4.12}$$

where we have used (4.9) to express F_0 and H_3 in terms of $H^{\rm D8}$, $H_\alpha^{\rm NS5}$, respectively. The induced O6-planes and/or D6-branes are along $(\xi^\mu, y, x^{1,3,5})$, $(\xi^\mu, y, x^{1,4,6})$, $(\xi^\mu, y, x^{2,4,5})$ and $(\xi^\mu, y, x^{2,3,6})$:

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
O6	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$	
O6′	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$			$\otimes$		$\otimes$
O6″	$\otimes$	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$	$\otimes$	
O6‴	$\otimes$	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$			$\otimes$

This is consistent with the supersymmetric intersection rule NS5 $\cap$ D$p(p-3)$ as well as NS5 $\cap$ D$p(p-1)$. Furthermore, the 0- and 2-brane tadpole cancellation conditions are automatically satisfied since $H_3 \wedge F_4$, $H_3 \wedge F_6$ vanish, as can be checked explicitly.

Associated with each O6 orientifold there is an internal-space involution σ_α, $\alpha \in \{1, \dots, 4\}$, reversing the parity in the directions orthogonal to the six-plane. Explicitly, $\sigma_1 : (x^2, x^4, x^6) \to (-x^2, -x^4, -x^6)$, etc. Under each of the four orientifold involutions, the background fields ϕ, g, F_0, F_4 are even, while H_3, F_2 are odd. It can readily be checked that this is indeed the case for the solution (4.9). In addition, it can be checked that the involutions leave the Killing-spinor conditions (4.10) invariant.

[10]It can be seen that the net charge of the configuration is that of orientifold six-planes. This is in agreement with the analysis of [19], which becomes relevant in the near-horizon limit as we will see in the following.

The total charges can be read off from (4.9):

$$\begin{aligned} Q^{\rm D4} &:= \int {\rm d}y \; j^{\rm D4} \\ Q^{\rm D8} &:= \int {\rm d}y \; j^{\rm D8} \\ Q^{\rm NS5} &:= \int {\rm d}y \; j^{\rm NS5} \; . \end{aligned} \tag{4.13}$$

Similarly the orientifold/D6-brane charge, $Q_\alpha^{\rm O6/D6}$, is given by the integral over the three-cycle $\Sigma_\alpha \in T^6$ which the orientifold plane/D6-brane is wrapping:

$$Q_\alpha^{\rm O6/D6} := \int_\Sigma {\rm dVol} \; j_\alpha^{\rm O6/D6} \; . \tag{4.14}$$

Note, however, that the resulting charge may in general have a y-dependence. For appropriate choices of $H^{\rm D8}$, $H_\alpha^{\rm NS5}$ it can be arranged that there is no y-dependence. This will indeed be the case in the near-horizon limit which we examine below.

The near-horizon limit. As already stressed, (4.9) corresponds to a family of solutions. There are special choices of the H-functions which lead to an $\mathrm{AdS}_4 \times T^6$ near-horizon limit. Consider for example the following:

$$\begin{aligned} H_\alpha^{\rm NS} &= \begin{cases} c_\alpha^{\rm NS} y \left\{ 1 + \frac{3}{2} \left(\frac{y}{y_0} \right)^{-\frac{5}{3}} \right\}, & y < y_0 \\ \frac{5}{2} c_\alpha^{\rm NS} y_0, & y \geq y_0 \end{cases} \\ H_\alpha^{\rm D4} &= \begin{cases} c_\alpha^{\rm D4} y \left\{ 1 - \frac{1}{2} \left(\frac{y}{y_0} \right) \right\}, & y < y_0 \\ \frac{1}{2} c_\alpha^{\rm D4} y_0, & y \geq y_0 \end{cases} \\ H^{\rm D8} &= \begin{cases} c^{\rm D8} y \left\{ 1 + \frac{3}{5} \left(\frac{y}{y_0} \right)^{-\frac{8}{3}} \right\}, & y < y_0 \\ \frac{8}{5} c^{\rm D8} y_0, & y \geq y_0 \, , \end{cases} \end{aligned} \tag{4.15}$$

where the c_α are constants ($\alpha \in \{1,2,3\}$ for D4 and $\{1,2,3,4\}$ for NS5). The leading y-dependence of the H's in the "near-horizon" $y \to 0$ limit is *uniquely* determined by the requirement that the dilaton should approach a finite constant at the horizon. The subleading behaviour is, however, largely arbitrary. We have partially constrained it by imposing that H should contain a harmonic piece, the piece linear in y, to make contact with the corresponding solution in the absence of smeared sources; additionally we have required that $\partial_y H$ should be continuous. Moreover, we have introduced an (arbitrary) point y_0 on the y-axis and imposed that $\partial_y H^{\rm D8}$, $\partial_y H^{\rm D4}$, $\partial_y H^{\rm NS5}$ vanish for $y \geq y_0$; this ensures that for $y \geq y_0$ the solution is that of flat space.

The requirements of the previous paragraph do not completely determine the H-functions, and we have made the particular choices in (4.15) for the sake of concreteness. An alternative choice with the exact same near-horizon limit as (4.15), but which ensures that $Q_\alpha^{\rm O6/D6}$ of (4.14) is everywhere y-independent,[11] is as follows. Define H^{D8} by

$$H^{\rm D8}=\begin{cases} c^{\rm D8}\exp\left\{\frac{10}{9}\frac{d_1^{\rm NS5}}{\Pi_{\alpha=1}^4 d_\alpha^{\rm NS5}}\int {\rm d}y\frac{\Pi_{\alpha=1}^4 H_\alpha^{\rm NS5}}{\partial_y H_1^{\rm NS5}}\right\}, & y\leq y_0\\ H^{\rm D8}(y_0), & y>y_0\,,\end{cases} \tag{4.16}$$

with all other H's as in (4.15), and $d^{\rm NS5}:=3/2\ y_0^{5/3}c^{\rm NS5}$. In the near-horizon limit we have $H^{\rm D8}\sim y^{-5/3}$ as before; moreover it follows that F_0H_3 is constant, leading to y-independent O6/D6 charges. The difference with (4.15) is that $\partial_y H^{\rm D8}$ is now discontinuous at $y=y_0$, signaling the presence of an eight-brane source at this point on the y-axis. Note also that $j^{O6/D6}=0$ for $y>y_0$, which implies that the O6 planes end on the eight-brane located at $y=y_0$.

We emphasize that, as can be seen from (4.15), (4.16), all H's and $\partial_y H$'s are continuous, except for $\partial_y H^{D8}$. The later has a finite jump at $y=y_0$, corresponding to a finite jump in the 'Roman's mass', as should be expected of domain wall solutions patching asymptotic regions with different 'cosmological constants'.

Anticipating the AdS$_4$ near-horizon limit, we are considering the solution only on the semi-infinite line $y>0$. One could trivially "glue" a copy of the solution on the $y<0$ side by replacing $H(y)\to H(|y|)$. This would of course result in a discontinuity of the flux at $y=0$, which could be attributed to the presence of a brane at the origin.

We now come to the crucial point of recovering the supersymmetric AdS$_4$ as the near-horizon limit of the metric, reached at $y\to 0$. The ten-dimensional metric takes the form

$$\begin{aligned}{\rm d}s^2_{\rm NH}=&\left\{d^{\rm D8}\left(\prod_{\alpha=1}^3 d_\alpha^{\rm D4}\right)\right\}^{-\frac12} y^{-\frac23}\eta_{\mu\nu}{\rm d}\xi^\mu{\rm d}\xi^\nu+\left(\prod_{\alpha=1}^4 d_\alpha^{\rm NS5}\right)\left\{d^{\rm D8}\left(\prod_{\alpha=1}^3 d_\alpha^{\rm D4}\right)\right\}^{\frac12}\frac{{\rm d}y^2}{y^2}\\ &+\sqrt{\frac{d_2^{\rm D4}d_3^{\rm D4}}{d_1^{\rm D4}d^{\rm D8}}}\left\{d_3^{\rm NS5}d_4^{\rm NS5}\ ({\rm d}x^1)^2+d_1^{\rm NS5}d_2^{\rm NS5}\ ({\rm d}x^2)^2\right\}\\ &+\sqrt{\frac{d_1^{\rm D4}d_3^{\rm D4}}{d_2^{\rm D4}d^{\rm D8}}}\left\{d_2^{\rm NS5}d_3^{\rm NS5}\ ({\rm d}x^3)^2+d_1^{\rm NS5}d_4^{\rm NS5}\ ({\rm d}x^4)^2\right\}\\ &+\sqrt{\frac{d_1^{\rm D4}d_2^{\rm D4}}{d_3^{\rm D4}d^{\rm D8}}}\left\{d_2^{\rm NS5}d_4^{\rm NS5}\ ({\rm d}x^5)^2+d_1^{\rm NS5}d_3^{\rm NS5}\ ({\rm d}x^6)^2\right\}\,,\end{aligned} \tag{4.17}$$

where we have defined $d^{\rm D4}:=c^{\rm D4}$ and $d^{\rm D8}:=3/5\ y_0^{8/3}c^{\rm D8}$, in addition to $d^{\rm NS5}:=3/2\ y_0^{5/3}c^{\rm NS5}$ defined previously. We can readily recognize (4.17) as the metric of AdS$_4\times T^6$. At $y\to\infty$ the metric asymptotes $\mathbb{R}^{1,3}\times T^6$. Note also that, as follows from (4.9), (4.15), the string coupling $g_{\rm s}={\rm e}^\phi$ approaches a finite constant in the $y\to 0$ limit:

$$ {\rm e}^\phi\longrightarrow\left(\prod_{\alpha=1}^4 d_\alpha^{\rm NS5}\right)^{\frac12}\left(\prod_{\alpha=1}^3 d_\alpha^{\rm D4}\right)^{-\frac14}\left(d^{\rm D8}\right)^{-\frac54}\,. \tag{4.18}$$

[11] Equation (4.15) leads to O6/D6 charges which are y-independent only in the near-horizon limit.

Moreover, all the flux components take on constant values:

$$
\begin{aligned}
H_{x^2x^4x^6} &= \frac{2}{3}d_1^{\mathrm{NS5}}\left(d^{\mathrm{D8}}\right)^{-1}; \quad H_{x^2x^3x^5} = \frac{2}{3}d_2^{\mathrm{NS5}}\left(d^{\mathrm{D8}}\right)^{-1} \\
H_{x^1x^3x^6} &= \frac{2}{3}d_3^{\mathrm{NS5}}\left(d^{\mathrm{D8}}\right)^{-1}; \quad H_{x^1x^4x^5} = \frac{2}{3}d_4^{\mathrm{NS5}}\left(d^{\mathrm{D8}}\right)^{-1} \\
nF_{x^3x^4x^5x^6} &= d_1^{\mathrm{D4}}; \quad F_{x^1x^2x^5x^6} = d_2^{\mathrm{D4}}; \quad F_{x^1x^2x^3x^4} = d_3^{\mathrm{D4}}; \quad F = \frac{5}{3}d^{\mathrm{D8}}\left(\prod_{\alpha=1}^{4} d_\alpha^{\mathrm{NS5}}\right)^{-1} .
\end{aligned}
\tag{4.19}
$$

As we shall see shortly (cf. eqs. (4.24), (4.25) below), this implies that all charges are regular in the near-horizon limit. In addition, from (4.11), (4.19) we see that the Poincaré duals to the D4/D8/NS5 all vanish in the limit, while the Poincaré duals to the orientifold six-planes/D6-branes become constant. In other words, *zooming in on the near-horizon region amounts to replacing the D4/D8/NS5 branes with fluxes, while adding a set of completely smeared orientifold six-planes.* This is a perfectly consistent picture from the ten-dimensional supergravity point-of-view, as we will now show. Our discussion will also establish that at the horizon there is a supersymmetry enhancement from $\mathcal{N}=1$ (two real supercharges) in three dimensions to $\mathcal{N}=1$ in four dimensions (four real supercharges).

An important remark. The most general form of $\mathcal{N}=1$ AdS_4 compactifications of IIA supergravity on manifolds of SU(3)-structure was given by two of the present authors in [18]. Since the group structure of T^6 is trivial, it is certainly contained in SU(3); it is therefore natural to ask whether we can recover the near-horizon solution given above, as a limiting case of the solutions of [18]. We will now show that this is indeed the case.

Setting the dilaton to zero, the solutions of [18] are given by (the reader is referred to the original paper for further details):

$$
\begin{aligned}
F_2 &= \frac{f}{18}J + \tilde{F} \\
H &= \frac{4m}{5}\mathrm{Re}(\Omega) \\
F_4 &= f\mathrm{dVol}_4 + \frac{3m}{5}J\wedge J \\
W &= -\frac{1}{5}m + \frac{i}{6}f \ .
\end{aligned}
\tag{4.20}
$$

In the above $(J,\ \Omega)$ is the SU(3) structure of the internal six-manifold and f, m are constants parameterizing the solution: f is the Freund-Rubin parameter, while m is the mass of Romans' supergravity — which can be identified with F_0 in the "democratic" formulation. The curvature of the AdS_4 is proportional to $|W|^2$. The two-form $\tilde{F}$ is the primitive part of F_2 (i.e. it is in the **8** of SU(3)) and is constrained by the Bianchi identity:

$$
\mathrm{d}\tilde{F} = \frac{1}{27}\left(\frac{108}{5}m^2 - f^2\right)\mathrm{Re}(\Omega) + j^{D6/O6}\,, \tag{4.21}
$$

where we have added a source for D6-branes/O6-planes on the right-hand side.[12] Finally, the only non-zero torsion classes of the internal manifold are given by:

$$\mathcal{W}_1^- = -\frac{4i}{9} f; \qquad \mathcal{W}_2^- = -2i\tilde{F} \ . \tag{4.22}$$

We observe that the near-horizon limit of the intersecting-brane solution given in the previous paragraph corresponds to setting $\tilde{F} = 0$ and taking the $f \to 0$ limit of (4.20), (4.21), while adding the following source term:

$$j^{D6/O6} = -\frac{4m^2}{5} \mathrm{Re}(\Omega) \ . \tag{4.23}$$

Indeed, in this limit all torsion classes of the internal manifold vanish (as they should for T^6), while the flux content exactly corresponds to (4.19).

The $\tilde{F} = 0$, $f = 0$ limit and its phenomenological implications were considered in detail by Acharya et al. in [19].

Comparison with the effective-superpotential approach. The configuration (4.9), (4.15) could be interpreted as consisting of (smeared) intersecting D4/D8/NS5-branes with charges Q^{D4}, Q^{D8}, Q^{NS5}, together with orientifold O6-planes/D6-branes. In order to make contact with the effective superpotential treatment of section 2.3, eq. (2.17), one should use the following dictionary:

$$\begin{aligned} Q_i^{\mathrm{D4}} &= \tilde{e}_i \\ Q^{\mathrm{D8}} &= \tilde{m}_0 \\ Q_1^{\mathrm{NS5}} &= \tilde{a}_0; \qquad Q_{i+1}^{\mathrm{NS5}} = \tilde{c}_i, \quad i = 1,2,3 \,, \end{aligned} \tag{4.24}$$

where we recognize the contributions of F_4, F_0 and H_3. The tadpole equation (2.18) is then reproduced for $Q_\alpha^{\mathrm{O6/D6}} = \tilde{N}_\alpha^{\mathrm{flux}}$. Moreover, substituting the near-horizon limit of (4.15) in (4.13), (4.11) and taking (4.24) into account, we obtain

$$\begin{aligned} \tilde{e}_i &= c_i^{\mathrm{D4}} \\ \tilde{m}_0 &= c^{\mathrm{D8}} \left(\Pi_{\alpha=1}^4 c_\alpha^{\mathrm{NS5}}\right)^{-1} \\ \tilde{a}_0 &= c_1^{\mathrm{NS5}} \left(c^{\mathrm{D8}}\right)^{-1}; \qquad \tilde{c}_i = c_{i+1}^{\mathrm{NS5}} \left(c^{\mathrm{D8}}\right)^{-1}, \quad i = 1,2,3 \ . \end{aligned} \tag{4.25}$$

For simplicity, here and in the remainder we ignore numerical normalizations and set $y_0 = 1$.

We would like to recover the stabilized values for the moduli, obtained in the effective-superpotential description, from the near-horizon limit of the supergravity solution discussed above. First, we would like to identify the geometric moduli corresponding to solution (4.9) with near-horizon geometry (4.17)–(4.19). The internal T^6 can be thought of as the product of three two-tori corresponding to (x^1, x^2), (x^3, x^4) and (x^5, x^6) — in

[12] The addition of this source term was first considered in [19].

accordance with the orientifold projections. There are then three Kähler moduli T_i, corresponding to the areas of the three T^2's. Explicitly:

$$T_1 = \sqrt{\frac{c_2^{\rm D4} c_3^{\rm D4}}{c_1^{\rm D4} c^{\rm D8}}} \sqrt{\Pi_{\alpha=1}^4 c_\alpha^{\rm NS5}}$$
$$T_2 = \sqrt{\frac{c_1^{\rm D4} c_3^{\rm D4}}{c_2^{\rm D4} c^{\rm D8}}} \sqrt{\Pi_{\alpha=1}^4 c_\alpha^{\rm NS5}}$$
$$T_3 = \sqrt{\frac{c_1^{\rm D4} c_2^{\rm D4}}{c_3^{\rm D4} c^{\rm D8}}} \sqrt{\Pi_{\alpha=1}^4 c_\alpha^{\rm NS5}} \, . \tag{4.26}$$

Moreover there are three complex structures τ_i, corresponding to the ratios of the radii of each of the T^2's. Explicitly:

$$\tau_1 = \sqrt{\frac{c_1^{\rm NS5} c_2^{\rm NS5}}{c_3^{\rm NS5} c_4^{\rm NS5}}}; \qquad \tau_2 = \sqrt{\frac{c_1^{\rm NS5} c_4^{\rm NS5}}{c_2^{\rm NS5} c_3^{\rm NS5}}}; \qquad \tau_3 = \sqrt{\frac{c_1^{\rm NS5} c_3^{\rm NS5}}{c_2^{\rm NS5} c_4^{\rm NS5}}} \, . \tag{4.27}$$

The τ_i's are related to the complex-structure moduli, U_i, and the dilaton, S, appearing in equation (2.19) via [54]:

$$S = {\rm e}^{-\phi_4} \frac{1}{\sqrt{\tau_1 \tau_2 \tau_3}}; \qquad U_1 = {\rm e}^{-\phi_4} \sqrt{\frac{\tau_2 \tau_3}{\tau_1}}; \qquad U_2 = {\rm e}^{-\phi_4} \sqrt{\frac{\tau_1 \tau_3}{\tau_2}}; \qquad U_3 = {\rm e}^{-\phi_4} \sqrt{\frac{\tau_1 \tau_2}{\tau_3}} \, . \tag{4.28}$$

In the above, ϕ_4 is the four-dimensional dilaton. It is given by:

$${\rm e}^{\phi_4} = {\rm e}^{\phi} \frac{1}{\sqrt{V}} \, , \tag{4.29}$$

where ϕ is the ten-dimensional dilaton appearing in (4.9), and V is the volume of the T^6. The latter can be read off in (4.9):

$$V = \left(c^{\rm D8}\right)^{-\frac{3}{2}} \left(\Pi_{\alpha=1}^4 c_\alpha^{\rm NS5}\right)^{\frac{3}{2}} \sqrt{\Pi_{\beta=1}^3 c_\beta^{\rm D4}} \, . \tag{4.30}$$

Putting (4.28), (4.29), (4.30) together we arrive at:

$$S = \sqrt{c^{\rm D8} \Pi_{\alpha=1}^3 c_\alpha^{\rm D4}} \sqrt{\frac{c_2^{\rm NS5} c_3^{\rm NS5} c_4^{\rm NS5}}{c_1^{\rm NS5}}}$$
$$U_1 = \sqrt{c^{\rm D8} \Pi_{\alpha=1}^3 c_\alpha^{\rm D4}} \sqrt{\frac{c_1^{\rm NS5} c_3^{\rm NS5} c_4^{\rm NS5}}{c_2^{\rm NS5}}}$$
$$U_2 = \sqrt{c^{\rm D8} \Pi_{\alpha=1}^3 c_\alpha^{\rm D4}} \sqrt{\frac{c_1^{\rm NS5} c_2^{\rm NS5} c_4^{\rm NS5}}{c_3^{\rm NS5}}}$$
$$U_3 = \sqrt{c^{\rm D8} \Pi_{\alpha=1}^3 c_\alpha^{\rm D4}} \sqrt{\frac{c_1^{\rm NS5} c_2^{\rm NS5} c_3^{\rm NS5}}{c_4^{\rm NS5}}} \, . \tag{4.31}$$

We are now ready to compare eqs. (4.26), (4.31) with the corresponding equation (2.23), arrived at in the effective-superpotential description. Upon taking (4.25) into account, it can be seen that indeed there is perfect agreement.

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4.2.2 D3/D5/D7/NS5/KK

Let us now perform a T-duality along x^1. As already mentioned in the introduction of section 4, this results in the following flux superpotential:

$$W_{\mathrm{IIB}} = i\tilde{e}_1 U_1 + i\tilde{c}_2 U_2 + i\tilde{c}_3 U_3 + i\tilde{a}_0 S + i\tilde{c}_1 T_1 + i\tilde{e}_2 T_2 + i\tilde{e}_3 T_3 + i\tilde{m}_0 U_1 T_2 T_3 \,, \tag{4.32}$$

where the last four terms are the geometrical-flux contributions and guarantee all-moduli stabilization around the type IIB AdS_4 vacuum, as in the type IIA mirror situation.

The corresponding brane configuration is now:

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
D3	$\otimes$	$\otimes$	$\otimes$			$\otimes$				
D5	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$	$\otimes$		
D5′	$\otimes$	$\otimes$	$\otimes$		$\otimes$				$\otimes$	$\otimes$
D7	$\otimes$	$\otimes$	$\otimes$			$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$
NS5	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$		$\otimes$	
NS5′	$\otimes$	$\otimes$	$\otimes$		$\otimes$			$\otimes$		$\otimes$
KK	$\otimes$	$\otimes$	$\otimes$		$\bullet$	$\otimes$		$\otimes$	$\otimes$	
KK′	$\otimes$	$\otimes$	$\otimes$		$\bullet$	$\otimes$	$\otimes$			$\otimes$

For simplicity, in deriving the explicit T-dualization of the components of the bulk fields we will work in the near-horizon region. After some rescalings of the coordinates of the six-torus, the original type IIA solution can be written as:

$$\begin{gathered} \mathrm{d}s^2_{10} = \mathrm{d}s^2_{\mathrm{AdS}_4} + \sum_{i=1}^{6} (\mathrm{d}x^i)^2; \quad \phi = \mathrm{const.}; \\ H_{x^2x^4x^6} = H_{x^2x^3x^5} = H_{x^1x^3x^6} = H_{x^1x^4x^5} = a; \\ F_{x^3x^4x^5x^6} = F_{x^1x^2x^5x^6} = F_{x^1x^2x^3x^4} = b; \quad F = c \,, \end{gathered}$$

where a, b, c are constants, and $b = 3/5c$. Upon T-dualising along x^1 (we follow the T-duality rules as given in [71]) we obtain the following IIB solution:

$$\begin{gathered} \mathrm{d}s^2_{10} = \mathrm{d}s^2_{\mathrm{AdS}_4} + \sum_{i=2}^{6} (\mathrm{d}x^i)^2 + (\mathrm{d}x^1 - A)^2; \quad \phi = \mathrm{const.}; \\ F_1 = c(\mathrm{d}x^1 - A); \quad F_3 = b(\mathrm{d}x^2 \wedge \mathrm{d}x^3 \wedge \mathrm{d}x^4 + \mathrm{d}x^2 \wedge \mathrm{d}x^5 \wedge \mathrm{d}x^6); \\ H_3 = a(\mathrm{d}x^2 \wedge \mathrm{d}x^4 \wedge \mathrm{d}x^6 + \mathrm{d}x^2 \wedge \mathrm{d}x^3 \wedge \mathrm{d}x^5); \quad F_{x^1x^3x^4x^5x^6} = b \,, \end{gathered} \tag{4.33}$$

where all other flux components vanish, except for the "external" component of the five-form Ramond flux $F_{\xi^0\xi^1\xi^2yx^2}$ which is determined from the one above by self-duality. These fluxes correspond precisely to the brane configuration presented in the table above.

The one-form connection A is given by:

$$A = a(x^6 \mathrm{d}x^3 + x^5 \mathrm{d}x^4) \ . \tag{4.34}$$

Hence x^1 should be thought of as the coordinate of a twisted S^1 fibre over a T^5 base parameterized by $x^2, \dots, x^6$. It follows that the geometry of the internal six-dimensional space is that of a *twisted torus*, or a *nilmanifold*. More precisely: defining left-invariant one-forms $\mathrm{e}^1 := \mathrm{d}x^1 - A$ and $\mathrm{e}^i := \mathrm{d}x^i$ for $i = 2, \dots, 6$, it is straightforward to see that the internal space can be identified with the nilmanifold 5.1 of table 4 of [72]. To make contact with the T-duality rules for NS5 branes, one may think of the internal nilmanifold as consisting of two intersecting KK monopoles (which are purely gravitational backgrounds) [73] along the directions indicated with bullets in the table above. Indeed, after T-duality, the NS5-branes are traded for Kaluza-Klein monopoles (TAUB-NUT spaces), carrying NUT-charges in the marked directions.

From (4.33) and the source-modified Bianchi identities $\mathrm{d}F + H \wedge F = j$ we can read off the following non-zero source components:

$$
\begin{aligned}
j^{\mathrm{O5}} &= ac(\mathrm{d}x^2 \wedge \mathrm{d}x^4 \wedge \mathrm{d}x^6 + \mathrm{d}x^2 \wedge \mathrm{d}x^3 \wedge \mathrm{d}x^5) \wedge (\mathrm{d}x^1 - A) \\
j^{\mathrm{O7}} &= ac(\mathrm{d}x^3 \wedge \mathrm{d}x^6 + \mathrm{d}x^4 \wedge \mathrm{d}x^5) \ .
\end{aligned}
\tag{4.35}
$$

Hence we have in addition O5/O7 orientifold planes along:

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
O5	$\otimes$	$\otimes$	$\otimes$	$\otimes$			$\otimes$		$\otimes$	
O5$'$	$\otimes$	$\otimes$	$\otimes$	$\otimes$				$\otimes$		$\otimes$
O7	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$	
O7$'$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$			$\otimes$

This is precisely what one would expect by applying one T-duality along x^1 to the O6-planes of the model generated by the D4/D8/NS5-branes discussed in section 4.2.1.

Performing another T-duality along x^2 brings us back to type IIA. Now all NS5-branes have been replaced by KK monopoles and the sources are D2/D6/KK-branes/monopoles distributed as

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
D2	$\otimes$	$\otimes$	$\otimes$							
D6	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$	$\otimes$	$\otimes$		
D6$'$	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$			$\otimes$	$\otimes$
D6$''$	$\otimes$	$\otimes$	$\otimes$				$\otimes$	$\otimes$	$\otimes$	$\otimes$
KK	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\bullet$	$\otimes$		$\otimes$	
KK$'$	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\bullet$		$\otimes$		$\otimes$
KK$''$	$\otimes$	$\otimes$	$\otimes$		$\bullet$	$\otimes$		$\otimes$	$\otimes$	
KK$'''$	$\otimes$	$\otimes$	$\otimes$		$\bullet$	$\otimes$	$\otimes$			$\otimes$

The superpotential receives contributions from F_2 and F_6 Ramond fields as well as from geometrical fluxes. We will not elaborate on this model here.

4.2.3 D4/D8/NS5 and D5/NS5

We would like finally to consider the situation where instead of four stacks of NS5-branes, we keep only one in the model presented in section 4.2.1:

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
D4	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$				
D4$'$	$\otimes$	$\otimes$	$\otimes$				$\otimes$	$\otimes$		
D4$''$	$\otimes$	$\otimes$	$\otimes$						$\otimes$	$\otimes$
NS5	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$		$\otimes$	
D8	$\otimes$	$\otimes$	$\otimes$		$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$	$\otimes$

This intersecting-brane distribution corresponds to the IIA superpotential eq. (2.20) with no complex-structure moduli. Its supergravity brane solution is simply obtained from (4.9), by identifying all four NS5 harmonic functions, i.e. setting all $c_\alpha^{\rm NS}$ resp. $d_\alpha^{\rm NS}$ to equal values.

The type IIB mirror brane set up of three stacks of D4-branes, one of D8-branes and one of NS5-branes is obtained by performing a T-duality in three directions (x^1, x^3, x^5):

	ξ^0	ξ^1	ξ^2	y	x^1	x^2	x^3	x^4	x^5	x^6
D5	$\otimes$	$\otimes$	$\otimes$			$\otimes$	$\otimes$		$\otimes$	
D5$'$	$\otimes$	$\otimes$	$\otimes$		$\otimes$			$\otimes$	$\otimes$	
D5$''$	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$			$\otimes$
D5$'''$	$\otimes$	$\otimes$	$\otimes$			$\otimes$		$\otimes$		$\otimes$
NS5	$\otimes$	$\otimes$	$\otimes$		$\otimes$		$\otimes$		$\otimes$	

This brane set up corresponds to the IIB superpotential (2.26), which does not depend on Kähler moduli. Here, however, the model is not the non-geometrical one discussed in sections 2.2 and 2.3, and the Khler moduli are not stabilized.

5. Conclusions

Flux compactifications have opened ways to compute macroscopically the entropy of gravitational backgrounds. The analysis of the source content of these backgrounds provides a microscopic perspective for this important issue.

With these motivations in mind, we have here explored the appearance of AdS_4 vacua in flux compactifications. The macroscopic approach is based on the effective superpotential of gauged supergravity, which is a polynomial in the chiral superfields with coefficients directly related to the flux numbers. In the absence of geometrical fluxes induced by the Scherk-Schwarz mechanism, complete moduli stabilization can be achieved in type IIA models, whereas in perturbative type IIB this is possible only in non-geometric compactifications where Khler moduli are absent. It should however be stressed that geometrical fluxes are not a problem per se; they simply enforce a generalized, torsionful structure.

An important result is the analysis of the microscopic origin of the above AdS_4 vacua. The latter appear as near-horizon geometries of brane/monopole distributions. The most interesting type IIA construction is based on a D4/D8/NS5-brane configuration, where some of the branes are smeared. This latter property plays a crucial role, and allows in particular to evade the introduction of geometrical fluxes in the macroscopic, gauged-supergravity picture. From the microscopic point-of-view all consistency conditions remain satisfied, provided the various sources and charges are carefully taken into account, and orientifold planes added appropriately.

As we saw in section 4.2.1, zooming in on the near-horizon region of the D4/D8/NS5 intersecting-brane solution amounts to replacing the D4/D8/NS5 branes with fluxes, while adding a set of completely smeared orientifold six-planes. This corresponds to a consistent ten-dimensional supergravity picture: it is contained in the class of solutions of [18] modified by the addition of smeared O6-planes. Such models were considered recently in detail by Acharya et al. [19].

A word of caution is in order: as we have seen in section 4.2.1, in the near-horizon limit the charge distributions corresponding to our brane setup are uniquely determined by the requirement that the near-horizon geometry be $\mathrm{AdS}_4 \times T^6$ and the dilaton approach a finite value. However, the subleading behaviour is largely arbitrary. It is fair to say that this behaviour should be better understood, and a physical interpretation should be provided, before a fully satisfactory picture can be said to have been given.

Upon T-dualizing type IIA models, new type IIB setups emerge. An interesting situation is reached in this way when a T-duality is performed on the above D4/D8/NS5-brane plus O6-plane system. Some of the NS5-branes become Kaluza-Klein monopoles (sources for geometric fluxes), which are purely gravitational backgrounds. The type IIB AdS_4 vacuum appears then as the near-horizon geometry of a D3/D5/D7/NS5/KK brane/monopole distribution. Going to the near-horizon region of the D3/D5/D7/NS5/KK system amounts to replacing the branes by their fluxes while adding a system of smeared O5/O7 planes.

The price for the moduli stabilization around this IIB AdS_4 vacuum is that the internal six-dimensional geometry is now a certain nilmanifold. This model is therefore a concrete realization of general ideas about flux compactifications and generalized geometries, and is an important example independently of the issues of moduli stabilization. Analyzing the emergence of spaces with generalized $G \times G$ group structure by using the tools presented in this paper, i.e. in direct contact with four-dimensional physics, deserves further attention.

Let us mention that, as already noted, there is a certain ambiguity in the interpretation of the smeared sources in the context of supergravity. One way to distinguish between the different possible interpretations would be to examine the next-to-leading order derivative corrections to our solutions; indeed, D-branes and orientifolds behave differently beyond leading order in the α' expansion [74].

Coming back to our original motivation, we would like to add that it is not straightforward to assign macroscopic entropy to an AdS_4 space, mostly because we do not know of any thermodynamical system that would support such an entropy. We could nevertheless use the dual, brane-source picture. The set of microscopic degrees of freedom certainly allows one to define an entropy, but this does not necessarily shed any light on the other

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(macroscopic) side of the duality: it merely serves as a definition. We hope to investigate these interesting directions in the future.

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A. Supersymmetry

The purpose of this appendix is to provide further details of the supersymmetry of solution (4.9) of section 4.2.1.

The calibration conditions for the branes follow from the kappa symmetry of the world-volume action, which, in the doubly-supersymmetric approach, can be thought of as a gauge-fixing of supersymmetry. In the case of the static, flat branes without world-volume excitations, considered here, the calibration conditions [70] take the following standard form (see e.g. [60, 61]):

$$\Gamma_{\xi^0\xi^1\xi^2x^{i_1}\dots x^{i_{p-2}}}\epsilon_+ = \epsilon_- \,, \tag{A.1}$$

for RR p-branes, and

$$\Gamma_{\xi^0\xi^1\xi^2x^ix^jx^k}\epsilon = \epsilon \,, \tag{A.2}$$

for type IIA NS5-branes. In (A.1) the p-brane is taken along the $\xi^0, \xi^1, \xi^2, x^{i_1}, \dots, x^{i_{p-2}}$ directions; in (A.2) the NS5-brane is along $\xi^0, \xi^1, \xi^2, x^i, x^j, x^k$. The products of gamma-matrices above are in an orthonormal frame; Γ_{11} is the chirality operator in ten dimensions, so that:

$$\epsilon_\pm = \frac{1}{2}(1 \pm \Gamma_{11})\,\epsilon \,, \tag{A.3}$$

where ϵ is the ten-dimensional Killing spinor. The latter can be taken to be constant, up to a y-dependent conformal factor, where y is the overall transverse coordinate.

Conditions (A.1) and (A.2) have to be imposed individually for each brane in the system. A priori, each condition reduces the supersymmetry by one-half, but not all conditions are necessarily independent. In the example of section 4.2.1, there are four D-branes and four NS5-branes, and therefore eight conditions following from (A.1), (A.2). However, it can easily be seen that only four of them are independent; they can be taken as in eq. (4.10).

The most straightforward way to check supersymmetry is then to simply substitute the values of all fields in (4.9) into eqs. (3.3), taking (4.10) into account. Alternatively, if the reader prefers to use the formalism of Romans' supergravity [69], the vielbein has to be rescaled from ${e_m}^a$ in the string frame to $\mathrm{e}^{-\frac{\phi}{2}}{e_m}^a$ in the Einstein frame, where ϕ is the dilaton. Moreover, the following dictionary between the democratic and the Romans' formalism (in the conventions of e.g. [18]) has to be used:

Democratic	Romans
F_0	$-2m$
F_2	$-2mB'$
F_4	$-G$
H_3	$-H$

References

[1] J.M. Maldacena, *The large-N limit of superconformal field theories and supergravity*, *Adv. Theor. Math. Phys.* **2** (1998) 231 [*Int. J. Theor. Phys.* **38** (1999) 1113] [hep-th/9711200].

[2] E. Witten, *Anti-de Sitter space and holography*, *Adv. Theor. Math. Phys.* **2** (1998) 253 [hep-th/9802150].

[3] S.S. Gubser, I.R. Klebanov and A.M. Polyakov, *Gauge theory correlators from non-critical string theory*, *Phys. Lett.* **B 428** (1998) 105 [hep-th/9802109].

[4] I.R. Klebanov and E. Witten, *Superconformal field theory on three-branes at a Calabi-Yau singularity*, *Nucl. Phys.* **B 536** (1998) 199 [hep-th/9807080].

[5] C. Vafa, *Superstrings and topological strings at large-N*, *J. Math. Phys.* **42** (2001) 2798 [hep-th/0008142].

[6] J.M. Maldacena and A. Strominger, *AdS_3 black holes and a stringy exclusion principle*, *JHEP* **12** (1998) 005 [hep-th/9804085].

[7] C. Kounnas, *Type II NS five-branes: non critical strings and their topological sectors*, *Fortschr. Phys.* **49** (2001) 599 [hep-th/0012192].

[8] A. Strominger and C. Vafa, *Microscopic origin of the Bekenstein-Hawking entropy*, *Phys. Lett.* **B 379** (1996) 99 [hep-th/9601029].

[9] C.G. Callan Jr. and J.M. Maldacena, *D-brane approach to black hole quantum mechanics*, *Nucl. Phys.* **B 472** (1996) 591 [hep-th/9602043].

[10] S. Ferrara, R. Kallosh and A. Strominger, *$N = 2$ extremal black holes*, *Phys. Rev.* **D 52** (1995) 5412 [hep-th/9508072].

[11] K. Behrndt et al., *Classical and quantum $N = 2$ supersymmetric black holes*, *Nucl. Phys.* **B 488** (1997) 236 [hep-th/9610105].

[12] J.M. Maldacena, A. Strominger and E. Witten, *Black hole entropy in M-theory*, *JHEP* **12** (1997) 002 [hep-th/9711053].

[13] G. Lopes Cardoso, B. de Wit and T. Mohaupt, *Corrections to macroscopic supersymmetric black-hole entropy*, *Phys. Lett.* **B 451** (1999) 309 [hep-th/9812082].

[14] H. Ooguri, C. Vafa and E.P. Verlinde, *Hartle-Hawking wave-function for flux compactifications*, *Lett. Math. Phys.* **74** (2005) 311 [hep-th/0502211].

JHEP09(2007)051

[15] T. Mohaupt, *Supersymmetric black holes in string theory*, *Fortschr. Phys.* **55** (2007) 519 [hep-th/0703035].

[16] J. Scherk and J.H. Schwarz, *Spontaneous breaking of supersymmetry through dimensional reduction*, *Phys. Lett.* **B 82** (1979) 60; *How to get masses from extra dimensions*, *Nucl. Phys.* **B 153** (1979) 61.

[17] R. Rohm, *Spontaneous supersymmetry breaking in supersymmetric string theories*, *Nucl. Phys.* **B 237** (1984) 553;
C. Kounnas and M. Porrati, *Spontaneous supersymmetry breaking in string theory*, *Nucl. Phys.* **B 310** (1988) 355;
S. Ferrara, C. Kounnas, M. Porrati and F. Zwirner, *Superstrings with spontaneously broken supersymmetry and their effective theories*, *Nucl. Phys.* **B 318** (1989) 75;
C. Kounnas and B. Rostand, *Coordinate dependent compactifications and discrete symmetries*, *Nucl. Phys.* **B 341** (1990) 641;
I. Antoniadis, *A possible new dimension at a few TeV*, *Phys. Lett.* **B 246** (1990) 377;
E. Kiritsis and C. Kounnas, *Perturbative and non-perturbative partial supersymmetry breaking: $N = 4 \to N = 2 \to N = 1$*, *Nucl. Phys.* **B 503** (1997) 117 [hep-th/9703059];
E. Kiritsis, C. Kounnas, P.M. Petropoulos and J. Rizos, *Solving the decompactification problem in string theory*, *Phys. Lett.* **B 385** (1996) 87 [hep-th/9606087];
E. Kiritsis, C. Kounnas, P.M. Petropoulos and J. Rizos, *String threshold corrections in models with spontaneously broken supersymmetry*, *Nucl. Phys.* **B 540** (1999) 87 [hep-th/9807067];
I. Antoniadis, E. Dudas and A. Sagnotti, *Supersymmetry breaking, open strings and M-theory*, *Nucl. Phys.* **B 544** (1999) 469 [hep-th/9807011]; *Brane supersymmetry breaking*, *Phys. Lett.* **B 464** (1999) 38 [hep-th/9908023];
I. Antoniadis, G. D'Appollonio, E. Dudas and A. Sagnotti, *Partial breaking of supersymmetry, open strings and M-theory*, *Nucl. Phys.* **B 553** (1999) 133 [hep-th/9812118];
I. Antoniadis, J.P. Derendinger and C. Kounnas, *Non-perturbative temperature instabilities in $N = 4$ strings*, *Nucl. Phys.* **B 551** (1999) 41 [hep-th/9902032];
C. Angelantonj, I. Antoniadis, G. D'Appollonio, E. Dudas and A. Sagnotti, *Type I vacua with brane supersymmetry breaking*, *Nucl. Phys.* **B 572** (2000) 36 [hep-th/9911081].

[18] D. Lüst and D. Tsimpis, *Supersymmetric AdS_4 compactifications of IIA supergravity*, *JHEP* **02** (2005) 027 [hep-th/0412250].

[19] B.S. Acharya, F. Benini and R. Valandro, *Fixing moduli in exact type IIA flux vacua*, *JHEP* **02** (2007) 018 [hep-th/0607223].

[20] M. Cvetič, F. Quevedo and S.J. Rey, *Stringy domain walls and target space modular invariance*, *Phys. Rev. Lett.* **67** (1991) 1836.

[21] E. Bergshoeff, M.B. Green, G. Papadopoulos and P.K. Townsend, *The IIA super-eightbrane*, hep-th/9511079;
E. Bergshoeff, M. de Roo, M.B. Green, G. Papadopoulos and P.K. Townsend, *Duality of type-II 7-branes and 8-branes*, *Nucl. Phys.* **B 470** (1996) 113 [hep-th/9601150].

[22] P.M. Cowdall, H. Lü, C.N. Pope, K.S. Stelle and P.K. Townsend, *Domain walls in massive supergravities*, *Nucl. Phys.* **B 486** (1997) 49 [hep-th/9608173].

[23] H.J. Boonstra, B. Peeters and K. Skenderis, *Brane intersections, anti-de Sitter spacetimes and dual superconformal theories*, *Nucl. Phys.* **B 533** (1998) 127 [hep-th/9803231].

[24] B. Janssen, P. Meessen and T. Ortín, *The D8-brane tied up: string and brane solutions in massive type IIA supergravity*, *Phys. Lett.* **B 453** (1999) 229 [hep-th/9901078].

[25] M. Massar and J. Troost, *$D0-D8-F1$ in massive IIA SUGRA*, *Phys. Lett.* **B 458** (1999) 283 [hep-th/9901136].

[26] K. Behrndt and M. Cvetič, *Supersymmetric domain wall world from $D=5$ simple gauged supergravity*, *Phys. Lett.* **B 475** (2000) 253 [hep-th/9909058].

[27] K. Behrndt and M. Cvetič, *Anti-de Sitter vacua of gauged supergravities with 8 supercharges*, *Phys. Rev.* **D 61** (2000) 101901 [hep-th/0001159].

[28] K. Behrndt, G. Lopes Cardoso and D. Lüst, *Curved BPS domain wall solutions in four-dimensional $N=2$ supergravity*, *Nucl. Phys.* **B 607** (2001) 391 [hep-th/0102128].

[29] E. Bergshoeff, R. Kallosh, T. Ortín, D. Roest and A. Van Proeyen, *New formulations of $D=10$ supersymmetry and $D8$ - $O8$ domain walls*, *Class. and Quant. Grav.* **18** (2001) 3359 [hep-th/0103233].

[30] G. Lopes Cardoso, G. Dall'Agata and D. Lüst, *Curved BPS domain wall solutions in five-dimensional gauged supergravity*, *JHEP* **07** (2001) 026 [hep-th/0104156].

[31] A. Ceresole, G. Dall'Agata, R. Kallosh and A. Van Proeyen, *Hypermultiplets, domain walls and supersymmetric attractors*, *Phys. Rev.* **D 64** (2001) 104006 [hep-th/0104056].

[32] R. Kallosh, S. Prokushkin and M. Shmakova, *Domain walls with strings attached*, *JHEP* **07** (2001) 023 [hep-th/0107097].

[33] G. Lopes Cardoso, G. Dall'Agata and D. Lüst, *Curved BPS domain walls and RG flow in five dimensions*, *JHEP* **03** (2002) 044 [hep-th/0201270].

[34] K. Behrndt and M. Cvetič, *Bent BPS domain walls of $D=5$ $N=2$ gauged supergravity coupled to hypermultiplets*, *Phys. Rev.* **D 65** (2002) 126007 [hep-th/0201272].

[35] K. Skenderis and M. Taylor, *Branes in AdS and pp-wave spacetimes*, *JHEP* **06** (2002) 025 [hep-th/0204054].

[36] H. Lü, C.N. Pope and J.F. Vazquez-Poritz, *From AdS black holes to supersymmetric flux-branes*, *Nucl. Phys.* **B 709** (2005) 47 [hep-th/0307001].

[37] K. Behrndt and M. Cvetič, *Supersymmetric intersecting $D6$-branes and fluxes in massive type IIA string theory*, *Nucl. Phys.* **B 676** (2004) 149 [hep-th/0308045].

[38] K. Behrndt and M. Cvetič, *General $N=1$ supersymmetric fluxes in massive type IIA string theory*, *Nucl. Phys.* **B 708** (2005) 45 [hep-th/0407263].

[39] J.P. Gauntlett, O.A.P. Mac Conamhna, T. Mateos and D. Waldram, *AdS spacetimes from wrapped $M5$ branes*, *JHEP* **11** (2006) 053 [hep-th/0605146].

[40] A. Ceresole, G. Dall'Agata, A. Giryavets, R. Kallosh and A. Linde, *Domain walls, near-BPS bubbles and probabilities in the landscape*, *Phys. Rev.* **D 74** (2006) 086010 [hep-th/0605266].

[41] G. Villadoro and F. Zwirner, *Beyond twisted tori*, *Phys. Lett.* **B 652** (2007) 118 [arXiv:0706.3049].

[42] P. Koerber and L. Martucci, *From ten to four and back again: how to generalize the geometry*, arXiv:0707.1038.

[43] S. Gukov, C. Vafa and E. Witten, *CFT's from Calabi-Yau four-folds*, *Nucl. Phys.* **B 584** (2000) 69 [*Erratum ibid.* **B 608** (2001) 477] [hep-th/9906070].

[44] T.R. Taylor and C. Vafa, *RR flux on Calabi-Yau and partial supersymmetry breaking*, *Phys. Lett.* **B 474** (2000) 130 [hep-th/9912152].

[45] P. Mayr, *On supersymmetry breaking in string theory and its realization in brane worlds*, *Nucl. Phys.* **B 593** (2001) 99 [hep-th/0003198].

[46] G. Curio, A. Klemm, D. Lüst and S. Theisen, *On the vacuum structure of type-II string compactifications on Calabi-Yau spaces with H-fluxes*, *Nucl. Phys.* **B 609** (2001) 3 [hep-th/0012213].

[47] S.B. Giddings, S. Kachru and J. Polchinski, *Hierarchies from fluxes in string compactifications*, *Phys. Rev.* **D 66** (2002) 106006 [hep-th/0105097].

[48] T.W. Grimm and J. Louis, *The effective action of $N = 1$ Calabi-Yau orientifolds*, *Nucl. Phys.* **B 699** (2004) 387 [hep-th/0403067]; *The effective action of type IIA Calabi-Yau orientifolds*, *Nucl. Phys.* **B 718** (2005) 153 [hep-th/0412277].

[49] K. Becker, M. Becker and J. Walcher, *Runaway in the Landscape*, arXiv:0706.0514.

[50] P. Candelas, E. Derrick and L. Parkes, *Generalized Calabi-Yau manifolds and the mirror of a rigid manifold*, *Nucl. Phys.* **B 407** (1993) 115 [hep-th/9304045].

[51] J.-P. Derendinger, C. Kounnas, P.M. Petropoulos and F. Zwirner, *Superpotentials in IIA compactifications with general fluxes*, *Nucl. Phys.* **B 715** (2005) 211 [hep-th/0411276].

[52] G. Villadoro and F. Zwirner, *$N = 1$ effective potential from dual type-IIA D6/O6 orientifolds with general fluxes*, *JHEP* **06** (2005) 047 [hep-th/0503169].

[53] O. DeWolfe, A. Giryavets, S. Kachru and W. Taylor, *Type IIA moduli stabilization*, *JHEP* **07** (2005) 066 [hep-th/0505160].

[54] P.G. Camara, A. Font and L.E. Ibáñez, *Fluxes, moduli fixing and MSSM-like vacua in a simple IIA orientifold*, *JHEP* **09** (2005) 013 [hep-th/0506066].

[55] K. Becker, M. Becker, C. Vafa and J. Walcher, *Moduli stabilization in non-geometric backgrounds*, *Nucl. Phys.* **B 770** (2007) 1 [hep-th/0611001].

[56] A.A. Tseytlin, *Harmonic superpositions of M-branes*, *Nucl. Phys.* **B 475** (1996) 149 [hep-th/9604035].

[57] E. Bergshoeff, M. de Roo, E. Eyras, B. Janssen and J.P. van der Schaar, *Multiple intersections of D-branes and M-branes*, *Nucl. Phys.* **B 494** (1997) 119 [hep-th/9612095].

[58] R. Argurio, F. Englert and L. Houart, *Intersection rules for p-branes*, *Phys. Lett.* **B 398** (1997) 61 [hep-th/9701042].

[59] N. Ohta, *Intersection rules for non-extreme p-branes*, *Phys. Lett.* **B 403** (1997) 218 [hep-th/9702164].

[60] J.P. Gauntlett, *Intersecting branes*, hep-th/9705011.

[61] D. Youm, *Black holes and solitons in string theory*, *Phys. Rept.* **316** (1999) 1 [hep-th/9710046].

[62] K.S. Stelle, *BPS branes in supergravity*, hep-th/9803116.

[63] H.J. Boonstra, B. Peeters and K. Skenderis, *Brane intersections, anti-de Sitter spacetimes and dual superconformal theories*, *Nucl. Phys.* **B 533** (1998) 127 [hep-th/9803231].

JHEP09(2007)051

[64] P. Koerber and D. Tsimpis, *Supersymmetric sources, integrability and generalized-structure compactifications*, arXiv:0706.1244.

[65] U. Gran, J. Gutowski, G. Papadopoulos and D. Roest, *Systematics of IIB spinorial geometry*, *Class. and Quant. Grav.* **23** (2006) 1617 [hep-th/0507087];
J.P. Gauntlett, D. Martelli, J. Sparks and D. Waldram, *Supersymmetric AdS_5 solutions of type IIB supergravity*, *Class. and Quant. Grav.* **23** (2006) 4693 [hep-th/0510125].

[66] J.P. Gauntlett and S. Pakis, *The geometry of $D = 11$ Killing spinors*, *JHEP* **04** (2003) 039 [hep-th/0212008];
I.A. Bandos, J.A. de Azcárraga, J.M. Izquierdo, M. Picón and O. Varela, *On BPS preons, generalized holonomies and $D = 11$ supergravities*, *Phys. Rev.* **D 69** (2004) 105010 [hep-th/0312266];
I. A.Bandos, J.A. de Azcárraga, M. Picón and O. Varela, *Generalized curvature and the equations of $D = 11$ supergravity*, *Phys. Lett.* **B 615** (2005) 127 [hep-th/0501007];
J. Bellorín and T. Ortín, *A note on simple applications of the Killing spinor identities*, *Phys. Lett.* **B 616** (2005) 118 [hep-th/0501246].

[67] U. Gran, P. Lohrmann and G. Papadopoulos, *The spinorial geometry of supersymmetric heterotic string backgrounds*, *JHEP* **02** (2006) 063 [hep-th/0510176].

[68] U. Gran, G. Papadopoulos, D. Roest and P. Sloane, *Geometry of all supersymmetric type I backgrounds*, hep-th/0703143.

[69] L.J. Romans, *Massive $N = 2$ a supergravity in ten dimensions*, *Phys. Lett.* **B 169** (1986) 374.

[70] G.W. Gibbons and G. Papadopoulos, *Calibrations and intersecting branes*, *Commun. Math. Phys.* **202** (1999) 593 [hep-th/9803163];
J. Gutowski and G. Papadopoulos, *AdS calibrations*, *Phys. Lett.* **B 462** (1999) 81 [hep-th/9902034]; J. Gutowski, G. Papadopoulos and P.K. Townsend, *Supersymmetry and generalized calibrations*, *Phys. Rev.* **D 60** (1999) 106006 [hep-th/9905156];
P. Koerber, *Stable D-branes, calibrations and generalized Calabi-Yau geometry*, *JHEP* **08** (2005) 099 [hep-th/0506154];
L. Martucci and P. Smyth, *Supersymmetric D-branes and calibrations on general $N = 1$ backgrounds*, *JHEP* **11** (2005) 048 [hep-th/0507099].

[71] S.F. Hassan, *T-duality, space-time spinors and RR fields in curved backgrounds*, *Nucl. Phys.* **B 568** (2000) 145 [hep-th/9907152].

[72] M. Graña, R. Minasian, M. Petrini and A. Tomasiello, *A scan for new $N = 1$ vacua on twisted tori*, *JHEP* **05** (2007) 031 [hep-th/0609124].

[73] D.J. Gross and M.J. Perry, *Magnetic monopoles in Kaluza-Klein theories*, *Nucl. Phys.* **B 226** (1983) 29.

[74] J.F. Morales, C.A. Scrucca and M. Serone, *Anomalous couplings for D-branes and O-planes*, *Nucl. Phys.* **B 552** (1999) 291 [hep-th/9812071].