

Gaugino condensates and D -terms from D7-branes

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ABSTRACT: We investigate, at the microscopic level, the compatibility between D -term potentials from world-volume fluxes on D7-branes and non-perturbative superpotentials arising from gaugino condensation on a different stack of D7-branes. This is motivated by attempts to construct metastable de Sitter vacua in type IIB string theory via D -term uplifts. We find a condition under which the Kähler modulus, T , of a Calabi-Yau 4-cycle gets charged under the anomalous $U(1)$ on the branes with flux. If in addition this 4-cycle is wrapped by a stack of D7-branes on which gaugino condensation takes place, the question of $U(1)$ -gauge invariance of the (T -dependent) non-perturbative superpotential arises. In this case an index theorem guarantees that strings, stretching between the two stacks, yield additional charged chiral fields which also appear in the superpotential from gaugino condensation. We check that the charges work out to make this superpotential gauge invariant, and we argue that the mechanism survives the inclusion of higher curvature corrections to the D7-brane action.

KEYWORDS: Flux compactifications, Intersecting branes models, Supersymmetric Effective Theories, Anomalies in Field and String Theories.

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1. Introduction

One of the most far-reaching consequences of the recent impressive advances in observational cosmology is the strong evidence for an accelerated expansion of the present universe, apparently driven by a small positive cosmological constant, or a similar form of “dark energy”.

Implementing a positive cosmological constant in a semi-realistic string theory set-up has been a subject of great interest during the past couple of years. In the standard low energy framework of string compactifications, one attempts to find local positive minima of the effective 4D supergravity potential in which all moduli are stabilized at sufficiently large mass scales so as to respect several other phenomenological constraints. Whereas moduli stabilization in general is already inflicted with a number of technical challenges (many of which have been overcome during the past years), the stabilization of moduli in a local de Sitter (dS) minimum requires some extra care due to a number of peculiarities of de Sitter spacetimes.

Some of these problems can already be understood from the lower-dimensional supergravity point of view and are due to the fact that there is no dS analogue of the Anti-de Sitter (AdS) Breitenlohner-Freedman bound [1] and that supersymmetry is necessarily broken in a dS background. The first property implies that, unlike for some AdS spaces, the stability of a dS vacuum does not tolerate tachyonic directions at the critical point,

whereas the lack of supersymmetry means that tachyonic directions cannot be ruled out using any supersymmetry arguments, as is e.g. possible for supersymmetric Minkowski vacua (for a recent discussion, see [2–4]).¹ These observations alone make stable dS vacua in supergravity theories appear less generic than stable Minkowski or AdS vacua.²

Besides these purely lower-dimensional issues, there are additional difficulties if one wants to obtain dS vacua from a compactification of 10D/11D string/M-theory. In fact, there exist several no-go theorems against dS compactifications [9–11], and finding a consistent dS compactification requires a violation of at least one of the premises that underlie these no-go theorems.

In [12], a strategy for the construction of (meta-)stable dS vacua in a class of type IIB compactifications was proposed, based on a combination and a careful balance of a number of different effects. The key ingredients used in [12] are background fluxes on a type IIB CY_3 orientifold (or an F-theory generalization thereof), a non-perturbative superpotential, W_{np} , due to gaugino condensation on wrapped D7-branes or Euclidean D3 instantons, as well as a contribution to the scalar potential due to anti-D3-branes ($\overline{D3}$ -branes). The rôle of the background fluxes is to generate a superpotential W_{flux} for the complex structure moduli and the dilaton, whereas the non-perturbative superpotential W_{np} leads to a non-trivial potential also for the Kähler moduli.

The gaugino condensation of pure SYM leads to a non-perturbative superpotential of the form

$$W_{np} = A(z)e^{-\alpha T^G}, \tag{1.1}$$

where T^G is the complex Kähler modulus of the 4-cycle wrapped by the D7-branes on which the gaugino condensation takes place, and $A(z)$ denotes a possible function of the complex structure moduli, the dilaton and the open string fields, which arises due to 1-loop threshold corrections to the D7-brane gauge kinetic function and/or due to world-volume fluxes and higher curvature corrections on the world volume of the D7-branes.³ If, as in the original proposal [12], one has only one Kähler modulus denoted by T , $W_{flux} + W_{np}$ can generate AdS minima in which all moduli are stabilized. These AdS minima can finally be “uplifted” by the $\overline{D3}$ -brane contribution to the potential,

$$V_{\overline{D3}} \sim \frac{\mu_3}{[\text{Re}(T)]^2}, \tag{1.2}$$

where μ_3 is the tension of the $\overline{D3}$ -brane. This may lead to a stable dS vacuum under certain conditions [15–17].

¹If a dS vacuum arises from a “gentle” uplift of a supersymmetric Minkowski vacuum, as in [2–4], one might obtain a nearby dS solution that inherits the absence of tachyonic directions from the Minkowski vacuum.

²In fact, despite the enormous body of literature from the 1980’s, stable de Sitter vacua were not found in extended gauged supergravity theories until rather recently [5–8].

³The superpotential in the case of Euclidean D3-branes has a similar form. In this text, we restrict our attention to the mechanism of gaugino condensation. For a recent discussion of Euclidean D3-branes, see [13, 14]. In writing down (1.1) we assumed that there is no light matter charged under the D7-brane gauge group. We will come back to the more general case including matter in the (anti)fundamental representation below.

In a variant of this construction, Burgess, Kallosh and Quevedo [18] suggested an alternative mechanism for the uplift of the AdS to the dS minimum that is not based on a $\overline{D3}$ -brane, but instead on world volume fluxes on D7-branes. In contrast to the $\overline{D3}$ -brane potential, the effects of the D7-brane world volume fluxes can easily be incorporated in the framework of a standard supergravity Lagrangian, which helps one to maintain technical control at all stages of the construction.⁴

In order to understand the effect of the world volume fluxes, let us denote by $D7_{\mathbb{G}}$ the D7-brane stack on which the gaugino condensation takes place and by $D7_{\mathbb{F}}$ a D7-brane (stack) with a non-trivial world-volume flux. A priori, these two stacks could be the same and/or wrap the same 4-cycle, which is the original situation described in [18]. This case was already discussed in [23] and we will, in this paper, always assume that $D7_{\mathbb{G}}$ and $D7_{\mathbb{F}}$ are different and wrap different 4-cycles. Likewise, the Kähler moduli of the 4-cycles $\Sigma^{\mathbb{F}}$ and $\Sigma^{\mathbb{G}}$ that are wrapped by $D7_{\mathbb{F}}$ and $D7_{\mathbb{G}}$ are denoted by $T^{\mathbb{F}}$ and $T^{\mathbb{G}}$, respectively, as already done in (1.1), and we use $G_{\mathbb{F}}$ and $G_{\mathbb{G}}$ for the respective gauge groups realized on these brane stacks. $G_{\mathbb{F}}$ is assumed to have an Abelian factor $U(1)_{\mathbb{F}}$. A generic four-cycle will be denoted by Σ^{α} .

Under certain topological conditions to be derived in the next section, the effect of the world volume flux on $D7_{\mathbb{F}}$ is to generate a D -term that corresponds to a gauged shift symmetry of the Kähler modulus $T^{\mathbb{G}}$,

$$T^{\mathbb{G}} \rightarrow T^{\mathbb{G}} + iq\epsilon, \tag{1.3}$$

with “charge” q and gauge parameter ϵ , so that the axion $a^{\mathbb{G}} \equiv \text{Im}(T^{\mathbb{G}})$ appears with a gauge covariant derivative

$$\mathcal{L} \sim (\partial_{\mu} a^{\mathbb{G}} + qA_{\mu}^{(\mathbb{F})})^2 \tag{1.4}$$

in the effective supergravity Lagrangian.⁵ In this expression, the vector field $A_{\mu}^{(\mathbb{F})}$ arises from the $U(1)_{\mathbb{F}}$ world volume gauge field on $D7_{\mathbb{F}}$. The gauging of (1.3) entails a non-trivial D -term potential,

$$V_D = \frac{1}{2g_{\mathbb{F}}^{-2}} \left[M_{\mathbb{P}}^2 q \partial_{T^{\mathbb{G}}} K + \sum M_{\mathbb{P}}^2 q_I |\Phi_I|^2 \right]^2, \tag{1.5}$$

where we introduced the reduced Planck mass $M_{\mathbb{P}}^2$ (which appears in the Einstein-Hilbert action as $\frac{M_{\mathbb{P}}^2}{2} \int d^4x \sqrt{-g} R$).⁶ The first term in (1.5), $q \partial_{T^{\mathbb{G}}} K$, arises from the standard expression $D \sim \eta^{T^{\mathbb{G}}} \partial_{T^{\mathbb{G}}} K$ for a D -term with Kähler potential K and a constant Killing vector $\eta^{T^{\mathbb{G}}} = iq$ corresponding to the gauged shift symmetry (1.3). The fields Φ_I in (1.5) denote other possibly present fields which transform linearly under $U(1)_{\mathbb{F}}$ with charges q_I and which, for simplicity, are assumed to have a canonical Kähler potential. The tree-level gauge kinetic function yields $g_{\mathbb{F}}^{-2} \sim \text{Re}(T^{\mathbb{F}}) + \dots$, where the dots are flux-dependent

⁴Similar uplifting methods using non-perturbative potentials in the context of the heterotic string were discussed in [19–22].

⁵To anticipate the result, a necessary condition for the appearance of a non-vanishing D -term involving $T^{\mathbb{G}}$ is that the 4-cycles wrapped by $D7_{\mathbb{G}}$ and $D7_{\mathbb{F}}$ intersect over a 2-cycle on which the world-volume flux is non-trivial, cf. (2.29) and (2.47).

⁶For the appearance of the $M_{\mathbb{P}}^2$ -factors, compare with formulas (4.6) and (5.15) in [24]. In general, there will be other Kähler moduli charged under $U(1)_{\mathbb{F}}$, cf. eq. (2.46), which we did not display in (1.5).

and/or curvature corrections involving the dilaton (see eq. (2.65)). In some simple cases, the tree-level Kähler potential for $T^{\mathbb{G}}$ is of the form $K(T^{\mathbb{G}} + \bar{T}^{\mathbb{G}}) \sim \ln(T^{\mathbb{G}} + \bar{T}^{\mathbb{G}})$. To make contact with [18], we assume this simple Kähler potential for the moment, neglect the above-mentioned (or any other) corrections to the gauge kinetic function and specialize to the case of one Kähler modulus, denoted by T . Furthermore, assuming that any additional matter field Φ_I has a vanishing vacuum expectation value (vev), one finds that the above D -term is simply proportional to $(\text{Re}(T))^{-3}$ and can be used to uplift an AdS minimum, just as the $\overline{D3}$ potential (1.2) [18].

The question as to whether the vevs of the Φ_I can really be assumed to be zero is discussed e.g. in [25, 18, 26] and is model-dependent. As was argued in [26], the Φ_I are in fact often non-zero at the minimum, and the real question is instead whether the Φ_I can acquire vevs that lead to a vanishing D -term or whether $V_D \neq 0$ at the minimum.⁷ In any case, as stressed in [28, 29], one should keep in mind that a D -term potential can only be used to uplift a *non*-supersymmetric AdS F -Term vacuum.

While this D -term uplifting mechanism looks quite compelling, it was already pointed out in [18] and further stressed in [26, 29–32] that there seems to be a tension between the necessity of the gauging of the shift symmetry (1.3) and the fact that the sum $W_{\text{flux}} + W_{\text{np}}$ of the superpotentials appears non-invariant under this shift. In fact, there is another seemingly non-invariant term,

$$\text{Im}(T^{\mathbb{G}})\text{tr}[F^{\mathbb{G}} \wedge F^{\mathbb{G}}], \tag{1.6}$$

where $F^{\mathbb{G}}$ denotes the Yang-Mills field of the gauge group $G_{\mathbb{G}}$, whose gauginos condense, and the prefactor arises from the tree-level relation $f^{\mathbb{G}} \sim T^{\mathbb{G}} + \dots$ for the corresponding gauge kinetic function. Both non-invariances can be cured by the same mechanism if bifundamental matter fields with non-vanishing mixed anomalies of the type $U(1)_{\mathbb{F}} - [G_{\mathbb{G}}]^2$ are present [26]. If this is the case, the mixed anomaly can cancel the classical non-invariance of (1.6) via the Green-Schwarz mechanism, and the non-perturbative superpotential (1.1) goes over to the Affleck-Dine-Seiberg (ADS) superpotential

$$W_{\text{ADS}} = \left(\frac{e^{-8\pi^2 f^{\mathbb{G}}}}{\det(M)} \right)^{\frac{1}{N_{\mathbb{G}} - N_{\mathbb{F}}}} = A(z) \left(\frac{e^{-8\pi^2 T^{\mathbb{G}}}}{\det(M)} \right)^{\frac{1}{N_{\mathbb{G}} - N_{\mathbb{F}}}}, \tag{1.7}$$

which now contains an additional factor involving the meson determinant [33, 34]

$$\det(M_j^i) \equiv \det(\tilde{\Phi}^{ia} \Phi_{ja}), \tag{1.8}$$

where i, j denotes the flavor and a, b the color index. The “quarks” and “anti-quarks” $\Phi_{jb}, \tilde{\Phi}^{ia}$ are also charged under the $U(1)_{\mathbb{F}}$ vector field $A_{\mu}^{(\mathbb{F})}$ and hence form part of the charged fields Φ_I in eq. (1.5). Obviously, the net sum of the $U(1)_{\mathbb{F}}$ charges of the “quarks” and “anti-quarks” equals the total $U(1)_{\mathbb{F}}$ charge of the determinant $\det(M_j^i)$. Thus, if this total charge is non-zero, the non-invariance of the $e^{-\alpha T^{\mathbb{G}}}$ -factor in W_{np} can in principle be compensated by the meson determinant, if the charges work out correctly. In this

⁷In [27], by contrast, the Φ_I play an important role in the uplifting, as they were assumed to give the dominant contribution to the (F -term) potential in the vacuum.

case, the mixed triangle anomaly is also non-vanishing, and the Green-Schwarz mechanism can operate to cancel the non-invariance of (1.6), so that the whole mechanism appears automatically self-consistent.

This idea was put forward and discussed at a field theoretical level in [26, 32] (see also [25] for an earlier, related discussion). Here we investigate more concretely a microscopic implementation of the idea in a D-brane setup, where the charged matter arises naturally via open strings stretched between the $D7_{\mathbb{G}}$ and $D7_{\mathbb{F}}$ stacks. We verify in particular that the flux-induced charge of the gauged axionic modulus $a^{\mathbb{G}}$ and the number of chiral matter fields (charged under both, $G_{\mathbb{G}}$ and the anomalous $U(1)_{\mathbb{F}}$) is precisely such that the non-perturbative superpotential becomes gauge invariant under $U(1)_{\mathbb{F}}$ transformations. In this way we confirm the field theoretic ideas of [26, 32] within a D-brane setup.

This is a further step in illuminating the proposal of [18]. Our analysis focuses on the issue of the gauge invariance of the non-perturbative superpotential (1.7) when the shift symmetry of the Kähler modulus $T^{\mathbb{G}}$ is gauged. More work is required to find a phenomenologically viable global D-brane model and some of the open issues will be briefly mentioned later on. Some of the presented ideas are already scattered in one or the other form in the literature, but we believe that our composition and elaboration of the facts relevant for embedding the proposal of [18] in a concrete D7-brane model could be helpful as a stepping stone for building a more complete model.

The organization of this paper is as follows: In section 2, we derive the D -term due to a world volume flux on a D7-brane ($D7_{\mathbb{F}}$) wrapped around a 4-cycle in a general Calabi-Yau using the Dirac-Born-Infeld action. We observe that a Kähler modulus gets charged only if the 4-cycle whose volume the Kähler modulus measures and the 4-cycle wrapped by the D7-brane intersect over a 2-cycle on which the world-volume flux is non-trivial. If a 4-cycle whose Kähler modulus is charged, is also wrapped by some D7-brane(s) (cf. the $D7_{\mathbb{G}}$ stack discussed before), additional chiral bifundamental matter naturally arises via open strings, exactly as needed for gauge invariance.⁸ Variants of the reduction method we apply appeared in different contexts before, e.g. in [35–38, 29]. For instance, the D -term potential resulting from world-volume fluxes on the D7-branes was previously derived in [23], however, with a completely different method, i.e. by analyzing the fermionic terms on the world-volume of the D7-branes. In our derivation we determine the D -term potential directly from the bosonic DBI-action and keep track of all factors explicitly, which allows us to read off the charge of the Kähler modulus $T^{\mathbb{G}}$ under the anomalous $U(1)_{\mathbb{F}}$. This is a prerequisite to verify the gauge invariance of (1.7) in section 3. Another ingredient in this verification is the number of chiral (anti)fundamentals of $G_{\mathbb{G}}$, charged under $U(1)_{\mathbb{F}}$. This number can be obtained using the index of the relevant Dirac operator, as we will elaborate on in section 3. Furthermore, in section 2.4 we investigate the effect of higher curvature corrections to the D7-brane action. We find that the gauge coupling potentially gets a correction depending on the dilaton field, but the charge of the Kähler moduli is not modified. Finally, section 4 summarizes our results and concludes.

⁸A related observation was made in [23], restricting though to chiral matter arising at the intersection of $D7_{\mathbb{F}}$ and its orientifold image. This is not the relevant case if $D7_{\mathbb{G}}$ and $D7_{\mathbb{F}}$ are different.

Appendix A illustrates the derivation of the D -term potential with the example of the $\mathbb{Z}_2 \times \mathbb{Z}_2$ toroidal orientifold in IIB. Appendix B gives a further alternative derivation of the D -term potential and appendix C contains the derivation of the imaginary part of the gauge kinetic function in the D7-brane compactification, including higher curvature effects.

2. D7-branes in IIB on a Calabi-Yau three-fold

In the introduction, we sketched the mechanism proposed in [26] that reconciles the presence of a gaugino condensation potential with the D -term uplifting mechanism suggested in [18]. In this section, we focus on the D -term part of this set-up by studying concrete models with D7-branes and by deriving in a more precise manner the resulting D -term potential in a KK reduction along similar lines as [35, 37, 23, 39].

2.1 The setup

We are interested in type IIB orientifolds with $\mathcal{N} = 1$ supersymmetry including D7- and D3-branes. This can be obtained if the internal Calabi-Yau manifold admits a holomorphic involution σ that acts on the Kähler form J and the $(3,0)$ -form Ω of the Calabi-Yau according to

$$\sigma J = J \quad , \quad \sigma \Omega = -\Omega \quad . \quad (2.1)$$

Modding out by

$$\mathcal{O} = (-1)^{F_L} \Omega_p \sigma \quad , \quad (2.2)$$

where F_L is the left moving fermion number and Ω_p is the worldsheet parity, leads to O3- and O7-planes at the fixed loci of σ , cf. [40]. This situation was also investigated in [41, 23]. Note that the cohomology classes $H^{(p,q)}$ split into σ eigenspaces $H_+^{(p,q)}$ and $H_-^{(p,q)}$. We do not specify the Calabi-Yau further but note that it could for example be taken to be a (blown up) toroidal orbifold $\tilde{Y} = T^6/\Gamma$, with Γ a discrete group of the kind \mathbb{Z}_N or $\mathbb{Z}_N \times \mathbb{Z}_M$. A systematic analysis of this class of models has been started in [42, 16, 17, 43]. The Calabi-Yau after modding out the involution \mathcal{O} will be denoted by $Y = \tilde{Y}/\mathcal{O}$ in the following.

In this section, we first compute the contribution of a single D7-brane with world volume flux to the D -term potential, and later take into account the contribution of other D7/D3-branes (including the mirror branes) and the O7/O3-planes. Just as in the introduction, we will use the subscript \mathbb{F} to denote all quantities pertaining to this D7-brane with flux. The stack on which gaugino condensation takes place will likewise be denoted by a subscript \mathbb{G} ; however, it will not enter until the end of section 2.2.

Let us therefore consider a single D7-brane, which we assume to be part of the \mathbb{F} stack. To lowest order, its contribution to the effective action is governed by the Dirac-Born-Infeld (DBI) and Chern-Simons (CS) action [44]. The DBI action for a D7-brane in the string frame is given by

$$S_{\text{DBI}} = -\mu_7 \int_{\mathcal{W}} d^8 \xi e^{-\phi} \sqrt{-\det(\iota^* g + \iota^* B + 2\pi \alpha' F)} \quad , \quad (2.3)$$

where \mathcal{W} is the eight-dimensional D7-brane world-volume which splits into the four-dimensional spacetime part, \mathcal{M}_4 , and a four-dimensional internal part on a 4-cycle, $\Sigma^{\mathbb{F}}$: $\mathcal{W} = \mathcal{M}_4 \times \Sigma^{\mathbb{F}}$. Furthermore, ϕ denotes the ten-dimensional dilaton, μ_7 is the D7-brane tension,⁹

$$\mu_7 = (2\pi)^{-7} \alpha'^{-4}, \tag{2.4}$$

ι^*g and ι^*B , respectively, denote the pullbacks of the ten-dimensional metric and the NSNS 2-form to \mathcal{W} , and F is the field-strength of the D7-brane gauge field. The CS action is given by

$$S_{\text{CS}} = -\mu_7 \int_{\mathcal{W}} \sum_p \iota^* C_p \wedge e^{\iota^*B + 2\pi\alpha'F}, \tag{2.5}$$

where ι^*C_p denotes the pullback of the respective RR form.

The BPS calibration conditions for D-branes wrapping Calabi-Yau cycles were derived in [45].¹⁰ In detail, for a D7-brane, the wrapped 4-cycle $\Sigma^{\mathbb{F}}$ needs to be a divisor and has to be holomorphically embedded into the Calabi-Yau. This can be equivalently expressed as [45]:

$$\frac{1}{2} (\iota^*J + i\mathcal{F}) \wedge (\iota^*J + i\mathcal{F}) = e^{i\theta} \sqrt{\frac{\det(g_{\Sigma^{\mathbb{F}}} + \mathcal{F})}{\det(g_{\Sigma^{\mathbb{F}})}}} \text{Vol}_{\Sigma^{\mathbb{F}}}, \tag{2.6}$$

$$\mathcal{F}^{2,0} = \mathcal{F}^{0,2} = 0, \tag{2.7}$$

where θ is an a priori arbitrary phase, $\text{Vol}_{\Sigma^{\mathbb{F}}}$ denotes the volume form on $\Sigma^{\mathbb{F}}$, and we have used¹¹

$$g_{\Sigma^{\mathbb{F}}} = (\iota^*g)|_{\Sigma^{\mathbb{F}}} \tag{2.8}$$

$$\mathcal{F} = (\iota^*B + 2\pi\alpha'F)|_{\Sigma^{\mathbb{F}}}. \tag{2.9}$$

Moreover, ι^*J is the pullback of the Kähler form which we can expand into harmonic forms $\iota^*\omega_\alpha$ of $H^{(1,1)}(\Sigma^{\mathbb{F}})$:

$$\iota^*J = v^\alpha \iota^*\omega_\alpha, \tag{2.10}$$

where ω_α form a basis of $H^{(1,1)}(Y, \mathbb{Z})$.

In writing down (2.3), (2.5) and (2.6), we neglected higher derivative corrections involving the Riemann tensor, cf. [47–52]. For the moment we just assume that they do not contribute, but we will come back to them in section 2.4.

In the following, we assume, for simplicity, that the negative σ eigenspaces $H_-^{(q,2-q)}(\tilde{Y})$ vanish:

$$H_-^{(q,2-q)}(\tilde{Y}) = 0. \tag{2.11}$$

⁹Note that this value for the tension is the appropriate one for T-duals of type I, where the physics is locally oriented [44]. We use the conventions of [44] throughout the paper.

¹⁰Note that these conditions also hold under the inclusion of background fluxes [46].

¹¹In the rest of this paper, we will sometimes also use the symbol \mathcal{F} for the full expression $(\iota^*B + 2\pi\alpha'F)$, i.e., not only for the restriction to $\Sigma^{\mathbb{F}}$. It should be clear from the context which of these two meanings of \mathcal{F} is intended.

This implies that the NSNS 2-form B , which is odd under world-sheet parity and hence, as it lives in the whole bulk, needs to be expanded in elements of the above negative cohomology groups, vanishes. Otherwise, the definition of the Kähler moduli becomes rather cumbersome [23].¹² This also means that B does not contribute to $\mathcal{F} \equiv (\iota^* B + 2\pi\alpha' F)|_{\Sigma^{\mathbb{F}}}$, i.e., we assume

$$\mathcal{F} = 2\pi\alpha' F. \tag{2.12}$$

This would automatically be true if the D7-brane was sitting on top of an O7-plane, as $\Sigma^{\mathbb{F}}$ would then also be wrapped by the O7-plane, i.e. it would be σ -invariant. However, we do not assume that the D7-brane is sitting on top of the O7-plane here. We will come back to a more precise description of the D-brane setup at the beginning of section 3.

In general, the forms $\iota^* \omega_\alpha$ which arise as pullbacks of (1,1)-forms of the ambient Calabi-Yau Y , do not necessarily form a complete basis of $H^{(1,1)}(\Sigma^{\mathbb{F}})$ [23]. There might be additional (1,1)-forms that are harmonic only locally on $\Sigma^{\mathbb{F}}$, but cannot be extended to harmonic (1,1)-forms on the whole of Y , i.e. they lie in the cokernel of ι^* . We denote a basis of those (1,1)-forms by $\tilde{\omega}_a$, and the full cohomology group $H^{(1,1)}(\Sigma^{\mathbb{F}})$ is split according to

$$H^{(1,1)}(\Sigma^{\mathbb{F}}) = \iota^* H^{(1,1)}(Y) \oplus \tilde{H}^{(1,1)}(\Sigma^{\mathbb{F}}). \tag{2.13}$$

As argued in [23], the basis $\tilde{\omega}_a$ can always be chosen in such a way that

$$\int_{\Sigma^{\mathbb{F}}} \iota^* \omega_\alpha \wedge \tilde{\omega}_a = 0. \tag{2.14}$$

2.2 From DBI to D-terms

Using a product ansatz for the spacetime, the DBI action (2.3) can be rewritten as:

$$S_{\text{DBI}} = -\mu_7 \int_{\mathcal{M}_4} d^4 x e^{-\phi} \sqrt{-\det(g_{(4)})} \sqrt{\det\left(1 + 2\pi\alpha' g_{(4)}^{-1} F_{(4)}\right)} \Gamma_{\mathbb{F}}, \tag{2.15}$$

where $F_{(4)}$ denotes the four-dimensional field-strength and $g_{(4)}$ is the (string frame) metric of \mathcal{M}_4 . Further, we defined

$$\Gamma_{\mathbb{F}} = \int_{\Sigma^{\mathbb{F}}} d^4 z \sqrt{\det(g_{\Sigma^{\mathbb{F}}} + \mathcal{F})}. \tag{2.16}$$

A low energy derivative expansion of (2.15) shows that the contribution of the D7-brane to the four-dimensional scalar potential is given by

$$V_{D7} = \mu_7 e^{-\phi} \Gamma_{\mathbb{F}} (\mathcal{V} e^{-2\phi})^{-2} = \mu_7 e^{3\phi} \mathcal{V}^{-2} \Gamma_{\mathbb{F}}. \tag{2.17}$$

The factor $(\mathcal{V} e^{-2\phi})^{-2}$ appears in the potential after transforming to the four-dimensional Einstein frame ($g_{\mu\nu}^{(\text{str})} = g_{\mu\nu}^{(\text{E})} (\mathcal{V} e^{-2\phi})^{-1}$), where the (dimensionless) volume of the Calabi-Yau orientifold, measured in the ten-dimensional string frame metric, is given by¹³

$$\mathcal{V} = \frac{1}{6} (\sqrt{\alpha'})^{-6} \int_Y J \wedge J \wedge J = \frac{1}{6} \mathcal{K}_{\alpha\beta\gamma} v^\alpha v^\beta v^\gamma. \tag{2.18}$$

¹²The additional moduli arising from B in the case of $H^2(\tilde{Y}) \neq 0$ can be stabilized by D -terms [17].

¹³We use $\sqrt{\alpha'}$ as the unit of length.

Here $\mathcal{K}_{\alpha\beta\gamma}$ are the triple intersection numbers

$$\mathcal{K}_{\alpha\beta\gamma} = (\sqrt{\alpha'})^{-6} \int_Y \omega_\alpha \wedge \omega_\beta \wedge \omega_\gamma . \quad (2.19)$$

For future reference, we note that the reduced Planck mass after the Weyl transformation to the four-dimensional Einstein frame is given by

$$M_{\mathbb{P}}^2 = \frac{\alpha'^3}{\kappa_{10}^2} = 2(2\pi)^{-7} \alpha'^{-1} . \quad (2.20)$$

In addition to (2.17), the low energy expansion of (2.15) also determines the gauge coupling to be

$$g_{\mathbb{F}}^{-2} = \mu_7 (2\pi\alpha')^2 e^{-\phi} \Gamma_{\mathbb{F}} . \quad (2.21)$$

Thus both, the contribution to the potential and to the gauge coupling, are determined by $\Gamma_{\mathbb{F}}$. In order to proceed further, we therefore have to calculate (2.16). In similar situations with D6-branes at angles (or, in the T-dual picture, D9-branes with world-volume fluxes), different methods were applied in the past. In [35, 37] each brane stack is separately chosen to keep some supersymmetry intact. This allows to make use of (2.6), where a priori each brane stack might be calibrated using a different θ -phase. However, only if all branes are calibrated with the same θ as the O-planes, i.e. using $\theta = 0$ [37], an overall supersymmetry is unbroken; otherwise, a D -term potential arises.

An alternative method to treat the analogue of (2.16) in the case of D6-branes at angles in a toroidal orientifold has recently been applied in [39]. As the metric and world-volume fluxes are known explicitly in toroidal orientifolds, the integral analogous to (2.16) can be calculated explicitly and expanded around a supersymmetric vacuum, which also allows to read off the D -term potential. It turns out that this method can not only be carried over to the D7-brane case, but it is also possible to apply it to a general Calabi-Yau for which the metric and world-volume fluxes are not explicitly known. We defer that calculation to appendix B and here proceed with a variant of the method using (2.6) and (2.7) as a starting point.¹⁴

Due to (2.7) we can expand the 2-form flux \mathcal{F} into basis elements of $H^{(1,1)}(\Sigma^{\mathbb{F}})$:¹⁵

$$\mathcal{F} = f^\alpha \iota^* \omega_\alpha + \tilde{f}^a \tilde{\omega}_a . \quad (2.22)$$

To avoid any misunderstanding, let us stress that we only consider fluxes on the D7-brane wrapped on $\Sigma^{\mathbb{F}}$ and thus ι^* always refers to the pullback onto the world-volume of this brane (we did not want to overload the notation with an index \mathbb{F} on ι^* though). In particular,

¹⁴Still another method to calculate the D -term potential in the case of D7-branes was applied in [23]. The authors determined the D -term both by considering the gauging of sigma-model symmetries after turning on world-volume fluxes, as well as by analyzing the fermionic terms on the world-volume of the D7-branes. We here reproduce their result directly from a reduction of the bosonic DBI action.

¹⁵Note that our \tilde{f} is not identical to the one used in [23], because we do not expand the world-volume flux on the brane and its mirror simultaneously. Thus, in the end we still have to sum over the branes and their mirror images. The end result is, however, equivalent to the one of [23].

$f^\alpha \neq 0$ denotes the $\iota^* \omega_\alpha$ -component of the flux on $D7_{\mathbb{F}}$ and not a flux on another possibly present D7-brane wrapped on the cycle Σ^α .

Now using (2.6) in (2.16), we obtain

$$\Gamma_{\mathbb{F}} = \tilde{\Gamma}_{\mathbb{F}} e^{-i\theta} = |\tilde{\Gamma}_{\mathbb{F}}| e^{i(\tilde{\theta} - \theta)} \tag{2.23}$$

with

$$\tilde{\Gamma}_{\mathbb{F}} \equiv \frac{1}{2} \int_{\Sigma^{\mathbb{F}}} (\iota^* J \wedge \iota^* J - \mathcal{F} \wedge \mathcal{F}) + i \int_{\Sigma^{\mathbb{F}}} (\iota^* J \wedge \mathcal{F}) . \tag{2.24}$$

The phase θ is fixed since the tension of the brane should be real and positive. This gives the relation

$$\theta = \tilde{\theta} + 2n\pi, \tag{2.25}$$

where $n \in \mathbb{Z}$ and

$$\tilde{\theta} = \arctan \left(\frac{2 \int_{\Sigma^{\mathbb{F}}} \iota^* J \wedge \mathcal{F}}{\int_{\Sigma^{\mathbb{F}}} (\iota^* J \wedge \iota^* J - \mathcal{F} \wedge \mathcal{F})} \right) \tag{2.26}$$

is the phase of $\tilde{\Gamma}_{\mathbb{F}}$. Thus, we deduce that

$$\Gamma_{\mathbb{F}} = |\tilde{\Gamma}_{\mathbb{F}}| = \sqrt{\left(\frac{1}{2} \int_{\Sigma^{\mathbb{F}}} (\iota^* J \wedge \iota^* J - \mathcal{F} \wedge \mathcal{F}) \right)^2 + \left(\int_{\Sigma^{\mathbb{F}}} \iota^* J \wedge \mathcal{F} \right)^2} . \tag{2.27}$$

We can use (2.10), (2.14) and (2.22) to express $\text{Re } \tilde{\Gamma}_{\mathbb{F}}$ and $\text{Im } \tilde{\Gamma}_{\mathbb{F}}$ as

$$\begin{aligned} \text{Re } \tilde{\Gamma}_{\mathbb{F}} &= \frac{1}{2} \int_{\Sigma^{\mathbb{F}}} (\iota^* J \wedge \iota^* J - \mathcal{F} \wedge \mathcal{F}) = \left(\frac{1}{2} v^\alpha v^\beta \mathcal{K}_{\alpha\beta\mathbb{F}} - \mathfrak{f}_{\mathbb{F}} \right) (\sqrt{\alpha'})^4, \\ \text{Im } \tilde{\Gamma}_{\mathbb{F}} &= \int_{\Sigma^{\mathbb{F}}} \iota^* J \wedge \mathcal{F} = v^\alpha Q_{\alpha\mathbb{F}} (2\pi)^2 \alpha'^2 . \end{aligned} \tag{2.28}$$

Here, $\mathcal{K}_{\alpha\beta\mathbb{F}}$ is the triple intersection number (2.19) of the four-cycle $\Sigma^{\mathbb{F}}$ and the Poincaré dual 4-cycles of ω_α and ω_β , and we used the abbreviations

$$Q_{\alpha\mathbb{F}} = \alpha'^{-2} \int_{\Sigma^{\mathbb{F}}} \iota^* \omega_\alpha \wedge \frac{\mathcal{F}}{(2\pi)^2} = \alpha'^{-1} \int_{\Sigma^{\mathbb{F}}} \iota^* \omega_\alpha \wedge \frac{F}{2\pi} = (2\pi)^{-2} f^\beta \mathcal{K}_{\alpha\beta\mathbb{F}}, \tag{2.29}$$

and

$$\mathfrak{f}_{\mathbb{F}} = \frac{1}{2} \left(f^\alpha f^\beta \mathcal{K}_{\alpha\beta\mathbb{F}} + \tilde{f}^a \tilde{f}^b \mathcal{K}_{ab}^{(\mathbb{F})} \right) \tag{2.30}$$

with

$$\mathcal{K}_{ab}^{(\mathbb{F})} = \alpha'^{-2} \int_{\Sigma^{\mathbb{F}}} \tilde{\omega}_a \wedge \tilde{\omega}_b . \tag{2.31}$$

Note that the part of \mathcal{F} that is only harmonic on $\Sigma^{\mathbb{F}}$, i.e. \tilde{f}^a , does not appear in $\text{Im } \tilde{\Gamma}_{\mathbb{F}}$ but only in $\text{Re } \tilde{\Gamma}_{\mathbb{F}}$. Moreover, $Q_{\alpha\mathbb{F}}$, as defined in (2.29), is integer valued because F is quantized in such a way that its integral over an arbitrary 2-cycle is a multiple of 2π , i.e.

$$\int F = 2\pi n, \quad n \in \mathbb{Z} . \tag{2.32}$$

As mentioned above, the condition for the D7-brane to preserve the same supersymmetry as the O7-plane enforces $\theta = \tilde{\theta} = 0$. Using (2.26) this translates into the condition

$$\text{Im } \tilde{\Gamma}_{\mathbb{F}} = \int_{\Sigma^{\mathbb{F}}} \iota^* J \wedge \mathcal{F} = 0 . \quad (2.33)$$

Allowing for small supersymmetry breaking, the expansion of (2.26) yields, for small $\text{Im } \tilde{\Gamma}_{\mathbb{F}}$,

$$\theta \sim \text{Im } \tilde{\Gamma}_{\mathbb{F}} . \quad (2.34)$$

Following [29], we then expand (2.27) in the limit $|\text{Im } \tilde{\Gamma}_{\mathbb{F}}| \ll |\text{Re } \tilde{\Gamma}_{\mathbb{F}}|$, and the contribution (2.17) of the D7-brane to the potential becomes

$$\begin{aligned} \mu_7 e^{3\phi} \mathcal{V}^{-2} \Gamma_{\mathbb{F}} &= \mu_7 e^{3\phi} \mathcal{V}^{-2} \text{Re } \tilde{\Gamma}_{\mathbb{F}} \sqrt{1 + \left(\frac{\text{Im } \tilde{\Gamma}_{\mathbb{F}}}{\text{Re } \tilde{\Gamma}_{\mathbb{F}}} \right)^2} \\ &\approx \mu_7 e^{3\phi} \mathcal{V}^{-2} \text{Re } \tilde{\Gamma}_{\mathbb{F}} + \frac{1}{2} \mu_7 e^{3\phi} \mathcal{V}^{-2} \frac{1}{\text{Re } \tilde{\Gamma}_{\mathbb{F}}} (\text{Im } \tilde{\Gamma}_{\mathbb{F}})^2 . \end{aligned} \quad (2.35)$$

The first term proportional to $\text{Re } \tilde{\Gamma}_{\mathbb{F}}$ is a tension term that is cancelled if one sums over all contributions of D7/D3-branes, O7/O3-planes, 3-form fluxes and possibly R^2 -terms due to RR-tadpole cancellation condition. This was discussed in similar cases for instance in [53, 37, 54]. The second term we would like to interpret as a contribution to the D -term potential. To bring it into the standard supergravity form we need an explicit formula for the D7-brane gauge coupling and the definition of the right field theoretic variables, i.e. those in which the gauge kinetic function becomes holomorphic and in which the sigma model metric is manifestly Kähler.

The gauge coupling is obtained from (2.21) using again the expansion (2.35) and keeping the first nontrivial term, which this time is the first one proportional to $\text{Re } \tilde{\Gamma}_{\mathbb{F}}$. This leads to

$$g_{\mathbb{F}}^{-2} = \mu_7 (2\pi \alpha')^2 e^{-\phi} \text{Re } \tilde{\Gamma}_{\mathbb{F}} = \mu_7 (2\pi)^2 \alpha'^4 \left(e^{-\phi} \frac{1}{2} v^\alpha v^\beta \mathcal{K}_{\alpha\beta\mathbb{F}} - e^{-\phi} f_{\mathbb{F}} \right) . \quad (2.36)$$

Comparing with (2.4), we see that the α' -factors drop out and the gauge coupling is (classically) dimensionless as it should be in four dimensions. In particular,

$$\mu_7 (2\pi)^2 \alpha'^4 = \frac{1}{(2\pi)^5} . \quad (2.37)$$

Moreover, the first term in parenthesis in (2.36) can be rewritten as

$$e^{-\phi} \frac{1}{2} v^\alpha v^\beta \mathcal{K}_{\alpha\beta\mathbb{F}} = \frac{1}{2} \hat{v}^\alpha \hat{v}^\beta \mathcal{K}_{\alpha\beta\mathbb{F}} = \partial_{\hat{v}^{\mathbb{F}}} \hat{\mathcal{V}} \equiv \hat{\mathcal{V}}^{\mathbb{F}} , \quad (2.38)$$

where

$$\hat{\mathcal{V}} = \frac{1}{6} \hat{v}^\alpha \hat{v}^\beta \hat{v}^\gamma \mathcal{K}_{\alpha\beta\gamma} = e^{-3\phi/2} \mathcal{V} \quad (2.39)$$

is the volume of the Calabi-Yau orientifold as measured in the ten-dimensional Einstein frame metric (which is related to the string frame metric by the usual dilaton factor $g_{MN}^{(E)} =$

$g_{MN}^{(\text{str})} e^{-\phi/2}$, as appropriate for ten dimensions). The corresponding 2-cycle moduli are given by

$$\hat{v}^\alpha = e^{-\phi/2} v^\alpha . \tag{2.40}$$

Up to normalization, $\hat{\mathcal{V}}^\mathbb{F}$ is the real part of the Kähler modulus $T^\mathbb{F}$ of the wrapped 4-cycle $\Sigma^\mathbb{F}$, and $e^{-\phi}$ is the real part of the dilaton, cf. for instance [55].¹⁶ Thus, choosing the normalization appropriately, (2.36) can be rewritten as

$$g_\mathbb{F}^{-2} = \text{Re } T^\mathbb{F} - \int_\mathbb{F} \text{Re } S . \tag{2.41}$$

The imaginary parts of the moduli $T^\mathbb{F}$ and S are given by RR-fields as can be seen from the reduction of the CS terms on the D7-brane (more precisely, they are given by C_0 for S and $a^\mathbb{F} = \alpha'^{-2} \int_{\Sigma^\mathbb{F}} \iota^* C^{(4)}$ for $T^\mathbb{F}$). The relative signs are determined by demanding a holomorphic gauge kinetic function (cf. appendix C for more details), and we arrive at¹⁷

$$T^\mathbb{F} = \frac{1}{(2\pi)^5} \left(\hat{\mathcal{V}}^\mathbb{F} + i a^\mathbb{F} \right) , \quad S = \frac{1}{(2\pi)^5} \left(e^{-\phi} - i C_0 \right) . \tag{2.42}$$

For general 4-cycles labelled by α , one obtains the analogous Kähler moduli

$$T^\alpha = \frac{1}{(2\pi)^5} \left(\hat{\mathcal{V}}^\alpha + i a^\alpha \right) . \tag{2.43}$$

With these moduli one verifies that

$$\hat{v}^\alpha = -\frac{2}{(2\pi)^5} \hat{\mathcal{V}}(\partial_{T^\alpha} K) , \tag{2.44}$$

where

$$K = -2 \ln \hat{\mathcal{V}} \tag{2.45}$$

is the Kähler potential for the moduli space of the Kähler moduli, and $\hat{\mathcal{V}}$ was defined in (2.39).

Before continuing, we should note that we neglected the open string moduli related to D7-brane fluctuations and Wilson-lines, which would otherwise also appear in the definitions of T^α and S . Taking them into account properly would require to consider open string 1-loop corrections in order to get holomorphic gauge kinetic functions [57].¹⁸ We thus assume that these moduli can be fixed, e.g. by the presence of fluxes [60–62].

Now we have all the ingredients to rewrite the potential arising from (2.35) in a more familiar way. As we already mentioned, tadpole cancellation implies that the first term proportional to $\text{Re } \tilde{\Gamma}$ drops out of the potential. The remaining contribution of a single D7-brane to the D -term potential is given by

$$\begin{aligned} \frac{1}{2} \mu_7 e^{3\phi} \mathcal{V}^{-2} \frac{1}{\text{Re } \tilde{\Gamma}_{\Sigma^\mathbb{F}}} (\text{Im } \tilde{\Gamma}_{\Sigma^\mathbb{F}})^2 &= \frac{1}{2g_\mathbb{F}^{-2}} \left[\frac{M_{\text{P}}^2}{4\pi^2} Q_{\alpha\mathbb{F}}(\partial_{T^\alpha} K) \right]^2 \\ &= \frac{1}{2(\text{Re } T^\mathbb{F} - \int_\mathbb{F} \text{Re } S)} \left[\frac{M_{\text{P}}^2}{4\pi^2} Q_{\alpha\mathbb{F}}(\partial_{T^\alpha} K) \right]^2 . \end{aligned} \tag{2.46}$$

¹⁶We differ from the definition in [55] by factors of i and the arbitrary normalization.

¹⁷The same relative signs were found for instance in [56], cf. their formula (8). Note that their variables differ by factors of $\pm i$ from ours.

¹⁸A closed string dual version of calculating these corrections to the gauge couplings was discussed in [58, 59].

In the first equality of (2.46) we used (2.4), (2.20), (2.28), (2.36), (2.39) and (2.44). In the last equality we wrote out explicitly the $U(1)_{\mathbb{F}}$ -gauge coupling to emphasize that it depends both on the Kähler modulus $T^{\mathbb{F}}$ and the dilaton S . The dilaton dependence is often ignored in the literature. Note that using the conventions of [44], there is no additional contribution to (2.46) from the orientifold image of the D7-brane.¹⁹

Our derivation of the D -term made use of the Abelian DBI action (2.3). If one had a stack of $D7_{\mathbb{F}}$ -branes instead, which does not lie on top of the $O7$ -planes, the gauge group would be $U(N_{\mathbb{F}})$. In that case, if the world volume flux lies in the diagonal $U(1)$ -subgroup, our calculation should carry over and the generalization of (2.46) gets an additional factor of $N_{\mathbb{F}}$. As the same is true for the $U(1)_{\mathbb{F}}$ -gauge couplings (2.36), the result becomes

$$V_D = \frac{1}{2g_{\mathbb{F}}^{-2}} \left[\frac{N_{\mathbb{F}} M_{\mathbb{P}}^2}{4\pi^2} Q_{\alpha\mathbb{F}} (\partial_{T^{\alpha}} K) + \dots \right]^2, \tag{2.47}$$

where the dots stand for the additional matter contributions, which cannot be derived from the DBI action, cf. eq. (1.5). Now (2.47) is of the familiar form of a D -term potential arising from a gauged $U(1)$ -symmetry. When compared with the expressions (1.3), (1.4) and (1.5), the coefficients $\frac{N_{\mathbb{F}}}{4\pi^2} Q_{\alpha\mathbb{F}}$ correspond, for $\alpha = \mathbb{G}$, to the constant Killing vector $-i\eta^{T^{\mathbb{G}}} = q$ of the gauged shift symmetry, at least up to a sign which can not be determined from (2.47) alone as it is quadratic in q .

2.3 $U(1)$ from CS

The presence of the gauged $U(1)$ symmetry can also be inferred by looking at the part

$$-\frac{\mu_7}{2} \int_{\mathcal{W}} \iota^* C_4 \wedge \mathcal{F} \wedge \mathcal{F}, \tag{2.48}$$

of the CS action. Taking one of the \mathcal{F} in this expression to be along $\Sigma^{\mathbb{F}}$ (i.e. it is part of the world volume flux) and one as the field strength of the four-dimensional $U(1)$ -field on the D7-brane stack, and expanding

$$C_4 = C_2^{\alpha} \wedge \omega_{\alpha} + \dots, \tag{2.49}$$

gives

$$-\mu_7 \int_{\Sigma^{\mathbb{F}}} \iota^* \omega_{\alpha} \wedge \mathcal{F} \int_{\mathcal{M}_4} C_2^{\alpha} \wedge F = -\mu_7 (2\pi)^2 \alpha'^2 Q_{\alpha\mathbb{F}} \int_{\mathcal{M}_4} C_2^{\alpha} \wedge F. \tag{2.50}$$

Hence, if $Q_{\alpha\mathbb{F}} \neq 0$ the derivative of the axion a^{α} dual to the 2-forms C_2^{α} will be covariantized, just as in (1.4). From the definition of $Q_{\alpha\mathbb{F}}$ in (2.29), it is obvious that this happens if the 4-cycle Σ^{α} (which is dual to ω_{α}) and $\Sigma^{\mathbb{F}}$ intersect over a 2-cycle on which the world-volume flux is non-trivial. In particular, when $\Sigma^{\mathbb{F}}$ has self-intersections, $T^{\mathbb{F}}$ can get charged in this way, which is the case discussed in [18, 23].

¹⁹There are other conventions in the literature, where the tension and the charge is “democratically” distributed over the branes and images. In that case the contribution of a single D7-brane to the D -term potential would be half of (2.46), the other half would come from the image.

If the above requirement is fulfilled and a charged Kähler modulus T^α exists, the corresponding gauged $U(1)_\mathbb{F}$ will become anomalous if the 4-cycle Σ^α dual to ω_α is wrapped by a D7-brane (or a brane stack). More precisely, the tree level effective action contains the term

$$\int_{\mathcal{M}_4} \text{Im}(T^\alpha) \text{tr} F^\alpha \wedge F^\alpha \tag{2.51}$$

with F^α being the field strength on the branes wrapped around the 4-cycle dual to ω^α (we denote the gauge group of these branes by G_α). This seems to break the invariance under local shifts of $\text{Im}(T^\alpha)$, i.e., the gauge symmetry $U(1)_\mathbb{F}$. At the quantum level, however, there are also mixed triangle anomalies of the type $U(1)_\mathbb{F} - [G_\alpha]^2$ due to chiral bifundamental fermions from open strings stretching between the two intersecting branes (whose existence also depends on the same condition $Q_{\alpha\mathbb{F}} \neq 0$, as we elaborate on in the next section). These triangle anomalies cancel the $U(1)_\mathbb{F}$ non-invariance of (2.51) by a Green-Schwarz mechanism. It is precisely these anomalous bifundamental matter fields that also modify the naive gaugino condensation superpotential to yield the ADS superpotential (1.7), which is then naturally invariant under $U(1)_\mathbb{F}$ transformations, as required for the mechanism suggested in [26]. A more detailed analysis of the gauge invariance of the superpotential is subject of section 3.

2.4 Curvature corrections

Before moving on to that discussion, let us come back to the issue of curvature corrections. We focus here on the calibration condition (2.6). This implies

$$\text{Im} \left(e^{-i\theta} \Phi \right) = 0, \tag{2.52}$$

with

$$\Phi = \frac{1}{2} (\mathcal{F} - i t^* J) \wedge (\mathcal{F} - i t^* J). \tag{2.53}$$

For the sake of the generalization to include the higher curvature terms, the following form is more convenient:

$$\text{Im} \left(e^{-i\theta} e^{\mathcal{F} - i t^* J} \right) \Big|_{\text{top}} = 0, \tag{2.54}$$

where “top” denotes the projection of the form to its top degree on the 4-cycle, i.e. to the 4-form part.

In the analogous case of $D9$ -branes with fluxes the inclusion of the higher curvature terms (in the large volume limit) has been argued to be [63, 64]

$$\text{Im} \left(e^{-i\theta} e^{\mathcal{F} - i J} \sqrt{\hat{A}(T(Y))} \right) \Big|_{\text{top}} = 0, \tag{2.55}$$

where $\hat{A}(T(Y))$ is the A-roof genus [65]

$$\hat{A}(T(Y)) = 1 - \frac{1}{24} p_1(T(Y)) + \dots, \tag{2.56}$$

with p_1 the first Pontryagin class and $T(Y)$ the tangent bundle of the whole internal manifold. In the case at hand, i.e. for D7-branes wrapped around a 4-cycle Σ , it is plausible to conjecture the corresponding generalization

$$\text{Im} \left(e^{-i\theta} e^{\mathcal{F}-i\iota^*J} \sqrt{\frac{\hat{A}(T(\Sigma))}{\hat{A}(N(\Sigma))}} \right) \Big|_{\text{top}} = 0, \quad (2.57)$$

where $T(\Sigma)$ and $N(\Sigma)$ are the tangent- and normal bundle of Σ . We will see further evidence for this conjecture in a moment. It is straightforward to expand the bracket, leading to

$$\text{Im} \left(e^{-i\theta} \left(\frac{1}{2}(\iota^*J + i\mathcal{F}) \wedge (\iota^*J + i\mathcal{F}) + \frac{1}{48}p_1(T(\Sigma)) - \frac{1}{48}p_1(N(\Sigma)) \right) \right) = 0, \quad (2.58)$$

where we multiplied the left hand side of (2.57) by minus one in order to facilitate comparison with earlier formulas. We note that the inclusion of the higher derivative corrections amounts to a shift in $\tilde{\Gamma}_{\mathbb{F}}$ of (2.24) according to

$$\tilde{\Gamma}_{\mathbb{F}} \rightarrow \tilde{\Gamma}_{\mathbb{F}} + \frac{1}{48} \int_{\Sigma^{\mathbb{F}}} \left(p_1(T(\Sigma^{\mathbb{F}})) - p_1(N(\Sigma^{\mathbb{F}})) \right). \quad (2.59)$$

As the additional contribution is real, the imaginary part of $\tilde{\Gamma}_{\mathbb{F}}$ does not change. This does, however, not mean that the D-term potential is not changed either. In fact it is known that the Kähler potential does get corrections from higher curvature terms, at least from those arising in the bulk action [55]. These amount to a shift in the argument of the Kähler potential for the Kähler moduli, i.e.

$$K = -2 \ln \left(\hat{\mathcal{V}} + \hat{\chi} e^{-3\phi/2} \right), \quad (2.60)$$

where $\hat{\chi}$ is a constant proportional to the Euler number of the compactification space. Using this in (2.44), one would obtain instead

$$\hat{v}^\alpha = -\frac{2}{(2\pi)^5} \left(\hat{\mathcal{V}} + \hat{\chi} e^{-3\phi/2} \right) (\partial_{T^\alpha} K). \quad (2.61)$$

However, also the Einstein-Hilbert term in string frame is modified by the higher derivative corrections, leading to a term proportional to

$$e^{-2\phi} (\mathcal{V} + \hat{\chi}) R. \quad (2.62)$$

Thus, going to the four-dimensional Einstein frame now amounts to rescaling the four-dimensional metric by $e^{-2\phi} (\mathcal{V} + \hat{\chi})$, so that the left hand side of (2.46) would contain a factor of $(\hat{\mathcal{V}} + \hat{\chi} e^{-3\phi/2})^{-2}$ now instead of $e^{3\phi} \mathcal{V}^{-2} = \hat{\mathcal{V}}^{-2}$. Thus, the right hand side of (2.46) still holds, even if the Kähler potential is modified by the higher derivative corrections as in (2.60). In particular the charges of the axions are unchanged.²⁰

²⁰We should note that implicit in our argument is the assumption that the D-term and also the gauge couplings are still determined by the real and imaginary parts of $\tilde{\Gamma}_{\mathbb{F}}$. It would be nice to have a conclusive derivation of this by including the higher derivative corrections on the brane also on the right hand side of (2.6).

One can give an additional argument that these charges are not modified by repeating the analysis of the last subsection, including the known curvature terms in the CS-action. The corrected CS-action is given by [47, 49, 50]:

$$S_{\text{CS}} = -\mu_7 \int_{\mathcal{W}} \sum_p \iota^* C_p \wedge e^{\mathcal{F}} \wedge \sqrt{\frac{\hat{A}(T(\Sigma^{\mathbb{F}}))}{\hat{A}(N(\Sigma^{\mathbb{F}}))}}, \quad (2.63)$$

where, as usual, the integral picks out the 8-form part of the integrand. As the axions charged under $U(1)_{\mathbb{F}}$ originate as the duals of C_2^α , one is interested in those terms of (2.63) that contain C_4 . This, however, does not allow the curvature corrections to modify the charge of the axions, as the expansion of \hat{A} contains only forms of degree $4n$ (for $n \in \mathbb{N}$) and one also needs one power of the gauge field strength along the non-compact space-time.

In contrast to the D -term, the tension and gauge coupling are related to the real part of $\tilde{\Gamma}_{\mathbb{F}}$ and, therefore, do receive corrections from (2.59). If the tension is modified, also the tadpole condition gets a contribution from the higher curvature terms. This is the case if there are D7-branes wrapped around cycles Σ for which the integrals of $p_1(T(\Sigma))$ and/or $p_1(N(\Sigma))$ over Σ are non-vanishing, cf. a related discussion in [53].

We see that the additional contribution to the real part of $\tilde{\Gamma}_{\mathbb{F}}$ can be taken into account by adjusting the definition (2.30) to

$$\hat{f}_{\mathbb{F}} = f_{\mathbb{F}} - \frac{\alpha'^{-2}}{48} \int_{\Sigma^{\mathbb{F}}} \left(p_1(T(\Sigma^{\mathbb{F}})) - p_1(N(\Sigma^{\mathbb{F}})) \right), \quad (2.64)$$

which leads to a modification of the gauge kinetic function to

$$f^{\mathbb{F}} = T^{\mathbb{F}} - \hat{f}_{\mathbb{F}} S. \quad (2.65)$$

In appendix C, we verify that this is consistent with a reduction of the known higher curvature corrections to the CS-part of the D7-brane action (which gives an independent derivation of the corrections to the imaginary part of $f^{\mathbb{F}}$). This supports the proposal (2.57) and the method of including the higher derivative corrections on the brane by modifying $\tilde{\Gamma}_{\mathbb{F}}$, according to (2.59), in the derivation of section 2.2.

3. Anomalies and gaugino condensation on D7-branes

Whereas the previous section primarily focused on the generation of the D -terms due to the world volume fluxes on the D7-branes, the present section is devoted to the second important ingredient of the construction of [26], namely the gaugino condensation. The connecting link between these two sub-effects is provided by certain types of anomalies. We therefore start with an analysis of the relevant anomalies that might occur in our set-up. To this end, let us consider, just as in the introduction, two different stacks of D7-branes, one denoted by $D7_{\mathbb{G}}$ and the other one by $D7_{\mathbb{F}}$. On stack $D7_{\mathbb{G}}$, we assume a gauge group $G_{\mathbb{G}}$ that can undergo gaugino condensation. As generalizations of the ADS superpotential in the presence of (anti-)symmetric tensor representations are not

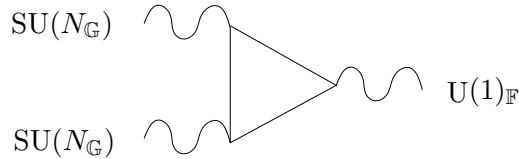


Figure 1: Mixed triangle graph.

very well understood,²¹ we assume that these representations are absent. The simplest situation in which this is the case is when the $D7_{\mathbb{G}}$ stack does not intersect the O-planes, which will henceforth be assumed. The gauge group $G_{\mathbb{G}}$ will therefore be of the form $U(N_{\mathbb{G}}) \cong SU(N_{\mathbb{G}}) \times U(1)_{\mathbb{G}}$.

The other stack, $D7_{\mathbb{F}}$, has a gauge group $G_{\mathbb{F}}$ that we require to contain at least one $U(1)$ factor denoted by $U(1)_{\mathbb{F}}$. In the generic case, when the $D7_{\mathbb{F}}$ stack does not lie on top of the O7-planes, one has the usual unitary gauge group including an Abelian $U(1)_{\mathbb{F}}$ factor. On the other hand, in the case when the $D7_{\mathbb{F}}$ stack coincides with an O7-plane, $G_{\mathbb{F}}$ becomes enhanced to a symplectic or orthogonal gauge group. This group can be broken to a unitary group with Abelian factors by switching on appropriate world volume fluxes, so this might a priori also be a valid option. However, we want to ensure at the same time that tensor representations of $SU(N_{\mathbb{G}})$ are absent and that the Kähler modulus $T^{\mathbb{G}}$ is charged under $U(1)_{\mathbb{F}}$. This is only guaranteed when the $D7_{\mathbb{F}}$ branes are *not* on top of O7-planes, which is therefore the case we assume in the rest of this paper. The reason for this is that the $D7_{\mathbb{G}}$ -stack has to intersect the $D7_{\mathbb{F}}$ stack in order for the Kähler modulus $T^{\mathbb{G}}$ to become charged under $U(1)_{\mathbb{F}}$ so as to generate a $T_{\mathbb{G}}$ -dependent D -term. But when the $D7_{\mathbb{F}}$ stack is on top of an O7-plane, the $D7_{\mathbb{G}}$ stack would then also intersect the O7-plane, contrary to our earlier assumption regarding $D7_{\mathbb{G}}$.

In this section, we are interested in the compatibility of the flux-induced D -term potential and the non-perturbative superpotential involving the Kähler modulus $T^{\mathbb{G}}$ in the case where this modulus is charged under $U(1)_{\mathbb{F}}$. As explained in the Introduction, this compatibility is equivalent to the cancellation of the mixed anomaly of the type $U(1)_{\mathbb{F}} - [SU(N_{\mathbb{G}})]^2$ by the Green-Schwarz mechanism.²² The triangle diagram of this anomaly is depicted in figure 1. Other types of anomalies will in general also appear and have to be cancelled in a complete global model, but as those other anomalies are not relevant for our discussion of gaugino condensation and D -terms, we will not consider them in this paper.

The fields running in the loop are the bifundamental fermions that arise at the intersections of stack $D7_{\mathbb{F}}$ and $D7_{\mathbb{G}}$. If $D7_{\mathbb{G}}$ intersects other branes, there might be additional fields transforming in the (anti-)fundamental representation of $SU(N_{\mathbb{G}})$. Denoting by N_F the total number of flavors of $SU(N_{\mathbb{G}})$, a non-perturbative superpotential is generated by either instantons or gaugino condensation if $N_F < N_{\mathbb{G}}$.²³

²¹See, however, [66].

²²Note that in our D-brane setup this would imply the cancellation of the mixed anomaly $U(1)_{\mathbb{F}} - U(1)_{\mathbb{G}}^2$ at the same time.

²³See, however, the example in eq. (4.43) and (4.44) of [67].

More precisely, after integrating out the non-Abelian gauge fields of $SU(N_{\mathbb{G}})$ (i.e. below the gaugino condensation scale), the term in the effective Lagrangian proportional to $\text{tr}(\langle F_{\mu\nu} \tilde{F}^{\mu\nu} \rangle)$ is

$$\left((N_{\mathbb{G}} - N_F) \text{Arg}(\langle \lambda^a \lambda^a \rangle) + \text{Arg}(\det M) + 8\pi^2 \text{Im}(T^{\mathbb{G}}) \right) \text{tr}(\langle F_{\mu\nu} \tilde{F}^{\mu\nu} \rangle). \quad (3.1)$$

The presence of the terms in (3.1) follows from the Veneziano-Yankielowicz-Taylor superpotential [68, 34, 69], i.e. they appear in

$$\int d^2\theta W_{\text{VYT}} + c.c. \quad (3.2)$$

with

$$W_{\text{VYT}} = \frac{U}{4} f^{\mathbb{G}} + \frac{U}{32\pi^2} \left((N_{\mathbb{G}} - N_F) \ln \frac{U}{\Lambda^3} + \ln \frac{\det M}{\Lambda^{2N_F}} \right), \quad (3.3)$$

where Λ is the UV cutoff scale, U is the gaugino bilinear field and we generalized (3.1) slightly by allowing for a general gauge kinetic function $f^{\mathbb{G}}$.²⁴ If one now integrates out U by its equation of motion, as for instance in [71, 67, 70], one ends up with the ADS superpotential (1.7), that we repeat here for convenience,

$$W_{\text{ADS}} = \left(\frac{e^{-8\pi^2 f^{\mathbb{G}}}}{\det M} \right)^{\frac{1}{N_{\mathbb{G}} - N_F}} = A(z) \left(\frac{e^{-8\pi^2 T^{\mathbb{G}}}}{\det M} \right)^{\frac{1}{N_{\mathbb{G}} - N_F}}. \quad (3.4)$$

The second term in (3.1), involving the meson field matrix

$$M_i^j = \tilde{\Phi}^{ja} \Phi_{ia}, \quad (3.5)$$

is necessary to ensure an anomaly free $U(1)_R$ symmetry ($\langle \lambda^a \lambda^a \rangle$ has R-charge 2, while $\det M$ has R-charge $2(N_F - N_{\mathbb{G}})$).²⁵ Here a denotes the color index and i, j the flavor index. The last term in (3.1) does not transform under $U(1)_R$.

Let us now discuss the transformation of the different terms under $U(1)_{\mathbb{F}}$. The gaugini are not charged under the anomalous $U(1)_{\mathbb{F}}$ and, therefore, do not transform. The last term does transform and, therefore, also the second one has to do so, in order to cancel the transformation of $\text{Im}(T^{\mathbb{G}})$. This is in accord with the idea that the term involving the meson fields mimics the contribution from the triangle graph in figure 1 below the confinement scale, i.e. after integrating out the gauge degrees of freedom. Given that above the confinement scale the contribution of figure 1 cancelled the transformation of $T^{\mathbb{G}}$ via the GS mechanism, it is obvious that the transformation of $\text{Arg}(\det M)$ has to perform the same job in the effective theory below the confinement scale. Indeed, it is clear from the definition (3.5) that the transformation of $\text{Arg}(\det M)$ is proportional to the sum of the $U(1)_{\mathbb{F}}$ charges of all bifundamentals charged with respect to $SU(N_{\mathbb{G}})$. If this sum vanishes, $\det M$ does not transform but at the same time the diagram in figure 1 vanishes.

²⁴For an expansion of U in components, see for example [70].

²⁵For a review of SYM, see for example [72, 73].

With our derivation of the D -term potential in section 2, we can verify the gauge invariance of (3.1) under the anomalous $U(1)_{\mathbb{F}}$ more directly. From (2.47) we read off the charge of $\text{Im} T^{\mathbb{G}}$ (up to a sign) to be

$$\frac{N_{\mathbb{F}}}{4\pi^2} Q_{\mathbb{G}\mathbb{F}}, \tag{3.6}$$

where $N_{\mathbb{F}}$ is the number of branes in the \mathbb{F} -stack (not including the mirror images) and $Q_{\mathbb{G}\mathbb{F}}$ is defined in (2.29) with $\omega_{\mathbb{G}}$ the 2-form that is Poincaré dual to the 4-cycle wrapped by the \mathbb{G} -stack. This has to be compared to the number of bifundamental fields coming from strings stretched between the \mathbb{F} - and \mathbb{G} -stacks. They each transform under $(\text{SU}(N_{\mathbb{F}}), \text{SU}(N_{\mathbb{G}}))_{U(1)_{\mathbb{F}}}$ in the representation

$$(N_{\mathbb{F}}, N_{\mathbb{G}} \oplus \bar{N}_{\mathbb{G}})_{q_{\mathbb{F}}=1}, \tag{3.7}$$

where the subscript denotes their charge under the anomalous $U(1)_{\mathbb{F}}$ (we assume that the positively charged bifundamentals are the left-handed ones appearing in the meson field (3.5); otherwise one has to take their antiparticles, which would amount to a change of sign of the $U(1)_{\mathbb{F}}$ charge). The fields in the $(N_{\mathbb{F}}, \bar{N}_{\mathbb{G}})$ -representation originate from strings stretched from $D7_{\mathbb{G}}$ to $D7_{\mathbb{F}}$, whereas the ones transforming in the $(\bar{N}_{\mathbb{F}}, N_{\mathbb{G}})$ -representation arise from strings stretched from the orientifold image of $D7_{\mathbb{G}}$ to $D7_{\mathbb{F}}$. Obviously, with only this particle content $\text{SU}(N_{\mathbb{F}})$ would be anomalous. Thus, additional fields charged under $\text{SU}(N_{\mathbb{F}})$, for instance from other brane stacks that intersect $D7_{\mathbb{F}}$, have to be present in a globally consistent model. As in the case of $D9$ -branes with fluxes, we expect that also in the case of $D7$ -branes with fluxes the number of chiral bifundamentals is given by the index of the Dirac operator on the intersection of the two $D7$ -branes and in the background of the world volume flux along this intersection. We assume that there is no world-volume flux on the $D7_{\mathbb{G}}$ -stack, otherwise the difference of the fluxes on the two stacks along their intersection locus would enter the index. This would lead to different numbers of $(N_{\mathbb{F}}, \bar{N}_{\mathbb{G}})$ - and $(\bar{N}_{\mathbb{F}}, N_{\mathbb{G}})$ -representations, which we want to avoid. Under this assumption, the number of $(N_{\mathbb{F}}, \bar{N}_{\mathbb{G}})$ -representations is given by the absolute value of [65]

$$\text{index}(\nabla) = \alpha'^{-1} \int_{\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}}} \hat{A}(T(\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}})) \wedge \text{ch}(F) = \alpha'^{-1} \int_{\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}}} \frac{F}{2\pi}, \tag{3.8}$$

where we introduced the factors α'^{-1} for dimensional reasons, \hat{A} is the A-roof genus that we already encountered in (2.56) and in the last equality we made use of the fact that the intersection locus $\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}}$ is (real) two dimensional and \hat{A} has an expansion in forms of degree $4n$ with $n \in \mathbb{N}$. Thus it can effectively be replaced by 1 in our case. We see that (3.8) exactly coincides with $Q_{\mathbb{G}\mathbb{F}}$, cf. (2.29). The number of $(N_{\mathbb{F}}, N_{\mathbb{G}})$ -representations is given by (3.8) as well. This can be argued as follows. As we said above, the fields transforming as $(N_{\mathbb{F}}, N_{\mathbb{G}})$ come from strings stretched from the images of $D7_{\mathbb{G}}$ to $D7_{\mathbb{F}}$. Thus, we have to replace $\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}}$ in (3.8) with $\Sigma^{\mathbb{F}} \cap \sigma(\Sigma^{\mathbb{G}})$, where $\sigma(\Sigma^{\mathbb{G}})$ is the image of the 4-cycle $\Sigma^{\mathbb{G}}$ under the involution σ . However,

$$\int_{\Sigma^{\mathbb{F}} \cap \sigma(\Sigma^{\mathbb{G}})} \frac{F}{2\pi} = \int_{\Sigma^{\mathbb{F}}} \iota^* \sigma^* \omega^{\mathbb{G}} \wedge \frac{F}{2\pi} = \int_{\Sigma^{\mathbb{F}}} \iota^* \omega^{\mathbb{G}} \wedge \frac{F}{2\pi} = \int_{\Sigma^{\mathbb{F}} \cap \Sigma^{\mathbb{G}}} \frac{F}{2\pi}, \tag{3.9}$$

where we used that the Poincaré-dual of $\sigma(\Sigma^{\mathbb{G}})$ is given by the σ -image (i.e. pullback) of $\omega^{\mathbb{G}}$ (denoted $\sigma^*\omega^{\mathbb{G}}$) and in the second equality we used that $\sigma^*\omega^{\mathbb{G}}$ and $\omega^{\mathbb{G}}$ have to represent the same cohomology class. Otherwise $(\sigma^*\omega^{\mathbb{G}} - \omega^{\mathbb{G}})$ would be a non-trivial element of H^2 that we assumed to be vanishing. Thus, we learn that the numbers of fundamentals and antifundamentals of $SU(N_{\mathbb{G}})$ that are charged under the anomalous $U(1)_{\mathbb{F}}$ with charge $+1$ are both given by²⁶

$$N_F = |Q_{\mathbb{G}\mathbb{F}}|N_{\mathbb{F}}. \tag{3.10}$$

Consequently, the meson determinant transforms according to

$$\det M \rightarrow e^{i\epsilon 2N_{\mathbb{F}}|Q_{\mathbb{G}\mathbb{F}}|} \det M, \tag{3.11}$$

where ϵ is the gauge parameter and the factor 2 comes from the fact that both Φ and $\tilde{\Phi}$ have charge 1 under $U(1)_{\mathbb{F}}$. Thus, the charge of $\text{Im} T^{\mathbb{G}}$ has the right value to cancel the transformation of the meson determinant.

Let us finally mention that the case of no bifundamentals is the one originally discussed in [12]. In that case, as well as in the case that there are only non-chiral bifundamentals, the D -term potential would be independent of $T^{\mathbb{G}}$ and, in general, a different uplift mechanism has to be envisaged.

4. Conclusions

In this paper, we studied the general compatibility of D -terms from D7-brane world volume fluxes with gaugino condensation on D7-branes. The mutual compatibility of these two features is crucial for the consistency of the proposal of [18] for obtaining (meta-)stable dS vacua in type IIB string theory via a variant of the KKLT construction [12] that does not rely on the introduction of $\overline{D3}$ -branes, which break supersymmetry explicitly.

We find that in the presence of world-volume fluxes on a D7-brane (called $D7_{\mathbb{F}}$ before), wrapped around a 4-cycle $\Sigma^{\mathbb{F}}$, any Kähler modulus T^{α} is charged if the following condition is fulfilled: The 4-cycle Σ^{α} , whose volume is measured by T^{α} , has to intersect with $\Sigma^{\mathbb{F}}$ over a 2-cycle that is threaded by non-trivial world-volume flux (on $D7_{\mathbb{F}}$).²⁷ If the cycle Σ^{α} is wrapped by a (stack of) D7-brane(s) as well (denoted by $D7_{\alpha}$), this is exactly the same condition that ensures the presence of chiral matter from strings stretching between $D7_{\mathbb{F}}$ and $D7_{\alpha}$, whose number is given by the index of the Dirac operator on the intersection 2-cycle $\Sigma^{\mathbb{F}} \cap \Sigma^{\alpha}$ in the background of the world-volume flux, cf. (3.8). In the example that the matter charged under the gauge group on $D7_{\alpha}$ transforms in the (anti)fundamental representation, we verified explicitly that the charge of T^{α} and the number of (anti)fundamentals take the right values to guarantee gauge invariance of the action via the Green-Schwarz mechanism. If the gauge group on $D7_{\alpha}$ and the matter spectrum allow for gaugino condensation (in which case we denoted $D7_{\alpha}$ as $D7_{\mathbb{G}}$ before), this implies automatically also the gauge invariance of the non-perturbative ADS superpotential (1.7), present below the

²⁶This equation assumes no other $SU(N_{\mathbb{G}})$ matter from possible additional brane stacks intersecting $D7_{\mathbb{G}}$.

²⁷With non-trivial we mean that its integral over the intersection 2-cycle $\Sigma^{\mathbb{F}} \cap \Sigma^{\alpha}$ does not vanish.

condensation-scale. Thus, our result complements the field theoretic discussion of [26, 32] by a more explicit embedding into a D-brane setup.

Furthermore, we also discussed the effects of higher curvature corrections to the D7-brane action. The gauge kinetic function possibly receives a correction that depends on the dilaton, cf. (2.65). We also argued that the charge of the axions is not modified by the higher curvature corrections, thus leaving the mechanism described in the last paragraph intact.

Our computation of the D -term potential in section 2 differs from the method used in [23] in that we determine the D -term potential directly from the dimensional reduction of the bosonic DBI-action, whereas in [23] the D -term potential is determined indirectly from the fermionic terms and the standard supergravity relations. The results, of course, agree with each other.

In our analysis, we concentrated our discussion on one flux-induced anomalous $U(1)_{\mathbb{F}}$ and one condensing gauge group $SU(N_{\mathbb{G}})$. Clearly, more work is needed to obtain a complete global model, which will be more involved. For example such a model will generically feature several anomalous $U(1)$'s and the occurrence of condensing gauge groups might be non-generic, see for instance [74]. We expect that in such a model similar mechanisms will be at work, although the structure of effective field theories with several (pseudo-) anomalous $U(1)$'s and the different types of cubic and mixed (as well as gravitational) anomalies can be quite complicated and might also involve generalized Chern-Simons-terms [75].

Another aspect that follows from the necessity of the charged matter fields Φ_I , which was also discussed in [18, 26, 32], is that the potential attains a much more complicated form, even in the case of only one Kähler modulus. Simply setting $\langle \Phi_I \rangle$ equal to zero renders W_{ADS} singular and is in general not a solution of the theory. Instead, the full scalar potential has to be minimized, also in the Φ_I -directions [26, 32], and including any tree-level contribution to the matter superpotential. Only upon including all these effects can one decide whether the vev of V_D is really non-zero or whether the matter fields relax to a vev that minimizes V_D to zero preventing it from uplifting the vacuum to a dS state. Unfortunately, the computation of tree-level matter superpotentials for intersecting D7-branes is rather complicated and the result model-dependent. Also the presence of additional B -field moduli in the case of $H^2(Y) \neq 0$ puts further constraints on the possibility of using D -term potentials for uplifting, as minimizing the D -term potential with respect to these moduli tends to cancel the uplifting D -term [23, 17]. An additional complication arises in the generic case of several Kähler moduli, because then even the matter independent part of V_D depends on both $T^{\mathbb{G}}$ and $T^{\mathbb{F}}$ (and in general also on the other Kähler moduli due to the more complicated form of the intersection numbers and hence the Kähler potential), as explained below (1.5).

When the $D7_{\mathbb{G}}$ -stack has self-intersections and/or intersects O7-planes, one encounters additional technical difficulties, as there would in general be matter fields in (anti-)symmetric tensor representations of $SU(N_{\mathbb{G}})$. A systematic analysis of the analogue of the ADS superpotential does not seem to exist in the literature for this case. For these technical reasons, we therefore restricted ourselves to the case where the $D7_{\mathbb{G}}$ -branes do not intersect the O7-planes. Thus, in order to have a non-vanishing intersection of the $D7_{\mathbb{G}}$ - and the $D7_{\mathbb{F}}$ -branes, we have to assume that also the latter are not on top of the O7-planes

(although they might intersect them). Our construction is therefore prone to F-theory corrections, which we assume to be small due to sufficient proximity of the D7-branes to the O7-planes. Taking into account F-theory and warping effects would be an interesting, but challenging problem.

Another issue of the D -term uplifting with F fluxes on D7-branes is that these are quantized and also lack the analogue of the warp factor suppression of the $\overline{D3}$ uplifting, as for D7-branes warped throats are not energetically favored.²⁸ This is sometimes criticized, as it seems to lead to a lack of sufficient tunability of the uplift potential to a small value so as to make supersymmetry breaking effects small enough. In the large volume compactifications of [76, 77], however, the D -terms are suppressed and can be very small, but in any case the charged matter contribution to the D -terms has to be taken into account. A more substantial analysis of these issues is beyond the scope of the present paper.

Some interesting topics we have not considered in this paper include Euclidean instantons and perturbative corrections to the Kähler potential. Progress on the first issue has recently been made in [13, 14]. The inclusion of perturbative corrections to the Kähler potential, like the ones discussed in [55, 78, 79], would be interesting in view of the proposal of [31] which uses D -term potentials to uplift possible AdS minima obtained from balancing effects in the F -term potential due to perturbative corrections to the Kähler potential [80, 81]. A fully consistent discussion of the uplift would need to take into account the corrections to the Kähler potential in the D -term potential as well.

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A. Example

As toy model for the derivation of the D -term potential we use the $\mathbb{Z}_2 \times \mathbb{Z}_2$ orientifold of a factorized $T_{(6)}$. The metric $g_{(6)}$ of a factorized $T_{(6)}$ has block-diagonal form $g_{(6)} = (g_1, g_2, g_3)$

²⁸However, the D7-branes might extend into the throat, like recently discussed for example in [59], which would also lead to some suppression.

with

$$g_\alpha = \begin{pmatrix} (R_1^\alpha)^2 & R_1^\alpha R_2^\alpha \cos \varphi^\alpha \\ R_1^\alpha R_2^\alpha \cos \varphi^\alpha & (R_2^\alpha)^2 \end{pmatrix}, \tag{A.1}$$

the metric of the respective sub-torus.

The geometrical moduli are defined as usual,

$$\begin{aligned} v^\alpha &:= R_1^\alpha R_2^\alpha \sin \varphi^\alpha, \\ \mathcal{U}^\alpha &:= \frac{R_2^\alpha}{R_1^\alpha} e^{i\varphi^\alpha}, \end{aligned} \tag{A.2}$$

where v^α corresponds to the 2-cycle volumes.

Now consider a D7-brane wrapped on the four cycle $\Sigma^\alpha = T_{(2)}^\beta \times T_{(2)}^\gamma$ with α, β, γ all different. Furthermore, we assume that there is a non-trivial world-volume flux present on this D7-brane. Since the metric is explicitly known, one can directly calculate $\Gamma_{\mathbb{F}}$ defined in (2.16) without invoking the supersymmetry condition (2.6) for the D-brane:²⁹

$$\Gamma_\alpha = \alpha'^2 \sqrt{(v^\beta)^2 + (f^\beta)^2} \sqrt{(v^\gamma)^2 + (f^\gamma)^2}. \tag{A.3}$$

With

$$\begin{aligned} \text{Re } \tilde{\Gamma}_\alpha &= \alpha'^2 (v^\beta v^\gamma - f^\beta f^\gamma), \\ \text{Im } \tilde{\Gamma}_\alpha &= \alpha'^2 (v^\beta f^\gamma + v^\gamma f^\beta), \end{aligned} \tag{A.4}$$

we can rewrite $\Gamma_\alpha = |\tilde{\Gamma}_\alpha|$ which can be expanded as in (2.35) for small $\text{Im } \tilde{\Gamma}_\alpha$:

$$\Gamma_\alpha \approx \text{Re } \tilde{\Gamma}_\alpha + \frac{1}{2\text{Re } \tilde{\Gamma}_\alpha} (\text{Im } \tilde{\Gamma}_\alpha)^2. \tag{A.5}$$

It is well known, that in this setup supersymmetry is preserved if the following calibration condition between the 2-cycle volumes and the amount of 2-form flux on the D7-brane world-volume is fulfilled [82]:

$$\frac{v^\beta}{v^\gamma} = -\frac{f^\beta}{f^\gamma}. \tag{A.6}$$

This condition leads to $\text{Im } \tilde{\Gamma}_\alpha = 0$. Note that, for finite values of the 2-cycle volumes and world-volume fluxes, (A.6) can only be fulfilled if f^β and f^γ differ in sign. Furthermore, in case of stabilization to a supersymmetric vacuum, this calibration condition leads to a partial fixing of the Kähler moduli [83, 84].

The field-theoretical moduli are (cf. (2.42) and (2.43))³⁰

$$\begin{aligned} \text{Re } T^\alpha &= \frac{1}{(2\pi)^5} e^{-\phi} v^\beta v^\gamma, \\ \text{Re } S &= \frac{1}{(2\pi)^5} e^{-\phi}. \end{aligned} \tag{A.7}$$

²⁹As all the (1, 1)-forms on Σ^α can be obtained by pullback from the ambient space (due to the simplicity of the torus geometry), there are no \tilde{f} fluxes appearing here.

³⁰Up to normalization, our moduli are defined as in [85].

The complex structure moduli do not need to be redefined $U^\alpha = \mathcal{U}^\alpha$. Re-expressing $\tilde{\Gamma}_\alpha$ in this basis yields:

$$\begin{aligned} \text{Re } \tilde{\Gamma}_\alpha &= (2\pi)^5 \alpha'^2 e^\phi (T_1^\alpha - f^\beta f^\gamma S_1), \\ \text{Im } \tilde{\Gamma}_\alpha &= \frac{\alpha'^2}{(2\pi)^5} \mathcal{V} e^{-\phi} \left(\frac{f^\beta}{T_1^\gamma} + \frac{f^\gamma}{T_1^\beta} \right), \end{aligned} \tag{A.8}$$

where $\mathcal{V} = v^1 v^2 v^3$ is the volume and T_1^α and S_1 denote the real parts of the respective moduli. Thus, the contribution to the D -term scalar potential due to a single D7-brane and its orientifold image, after transforming to the Einstein-frame and omitting the term that is cancelled after summing over all D-branes and O-planes, is given by (cf. (2.35))

$$\frac{1}{2} \mu_7 e^{3\phi} \mathcal{V}^{-2} \frac{1}{\text{Re } \tilde{\Gamma}_\alpha} (\text{Im } \tilde{\Gamma}_\alpha)^2 = \frac{1}{2g_\alpha^{-2}} \left(\frac{M_{\text{P}}^2}{4\pi^2} \left(\frac{(2\pi)^{-2} f^\beta}{2T_1^\gamma} + \frac{(2\pi)^{-2} f^\gamma}{2T_1^\beta} \right) \right)^2, \tag{A.9}$$

as expected from (2.47) with the tree-level Kähler potential $K = -\sum_\beta \ln(T^\beta + \bar{T}^\beta)$. In deriving (A.9) we made use of (2.20) and (2.36).

Since Σ^α is parameterized by T^α , we observe that no D -term is generated for the modulus parameterizing the 4-cycle the brane is wrapped on, but rather for the moduli of the 4-cycles intersecting the wrapped cycle as discussed before.³¹ Therefore, for a single stack of D7-branes with 2-form flux, the question of gauge invariance of the non-perturbative superpotential from gaugino condensation on that same stack of branes does not arise. If there is another stack of D7-branes on an intersecting 4-cycle, bifundamental matter will be present and the non-perturbative superpotential due to gaugino condensation on that second brane stack would be the Affleck-Dine-Seiberg superpotential which is naturally invariant under the induced shift-symmetry as explained in the main text.

B. Alternative calculation of $\Gamma_{\mathbb{F}}$

In this appendix we give an alternative derivation of (2.27) which follows in spirit more closely the method of [39] in that it does not take (2.6) as a starting point. However, we still require that the world volume flux is of type (1, 1), i.e. (2.7) holds. Furthermore, we neglect the higher derivative corrections to the DBI-action in the following.

First, we note that there is the following relation between quantities in real and complex coordinates

$$(\iota^* g)_{mn} + \mathcal{F}_{mn} = \begin{pmatrix} 0 & (\iota^* g)_{i\bar{j}} - \mathcal{F}_{i\bar{j}} \\ (\iota^* g)_{i\bar{j}} + \mathcal{F}_{i\bar{j}} & 0 \end{pmatrix}. \tag{B.1}$$

Thus

$$\begin{aligned} \det((\iota^* g)_{mn} + \mathcal{F}_{mn}) &= \det((\iota^* g)_{i\bar{j}} - \mathcal{F}_{i\bar{j}}) \det((\iota^* g)_{i\bar{j}} + \mathcal{F}_{i\bar{j}}) \\ &= \det(-i(\iota^* J)_{i\bar{j}} - \mathcal{F}_{i\bar{j}}) \det(-i(\iota^* J)_{i\bar{j}} + \mathcal{F}_{i\bar{j}}) \\ &= \det((\iota^* J)_{i\bar{j}} - i\mathcal{F}_{i\bar{j}}) \det((\iota^* J)_{i\bar{j}} + i\mathcal{F}_{i\bar{j}}), \end{aligned} \tag{B.2}$$

³¹As mentioned earlier, this changes in cases when the wrapped cycle has non-trivial self-intersections.

which implies

$$\begin{aligned}\sqrt{\det((\iota^*g)_{mn} + \mathcal{F}_{mn})} &= \sqrt{\det((\iota^*J)_{i\bar{j}} - i\mathcal{F}_{i\bar{j}}) \det((\iota^*J)_{i\bar{j}} + i\mathcal{F}_{i\bar{j}})} \\ &= \sqrt{\left(\frac{1}{2}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}((\iota^*J)_{i\bar{j}}(\iota^*J)_{k\bar{l}} - \mathcal{F}_{i\bar{j}}\mathcal{F}_{k\bar{l}})\right)^2 + \left(\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}(\iota^*J)_{i\bar{j}}\mathcal{F}_{k\bar{l}}\right)^2},\end{aligned}\tag{B.3}$$

where our epsilon-symbol takes values 0 and ± 1 .

Next, consider two harmonic $(1, 1)$ -forms $\omega^{(1)}$ and $\omega^{(2)}$ on the 4-cycle $\Sigma^{\mathbb{F}}$ that is wrapped by the D7-brane. It follows that $\omega^{(1)} \wedge \omega^{(2)} \in H^{(2,2)}(\Sigma^{\mathbb{F}})$ and

$$\begin{aligned}\star(\omega^{(1)} \wedge \omega^{(2)}) &= \star\left(\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)} dz^i \wedge d\bar{z}^{\bar{j}} \wedge dz^k \wedge d\bar{z}^{\bar{l}}\right) \\ &= \star\left(-\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)} dz^i \wedge dz^k \wedge d\bar{z}^{\bar{j}} \wedge d\bar{z}^{\bar{l}}\right) \\ &= -\frac{1}{4}\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}\sqrt{\det(\iota^*g)}^{-1} \in H^{(0,0)}(\Sigma^{\mathbb{F}}),\end{aligned}\tag{B.4}$$

where with $\sqrt{\det(\iota^*g)}$ we always mean $\sqrt{\det((\iota^*g)_{mn})} = \det((\iota^*g)_{i\bar{j}})$. As any harmonic function on a compact Riemannian manifold is constant, we infer that $\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}\sqrt{\det(\iota^*g)}^{-1}$ is constant. On the other hand, one has

$$\int_{\Sigma^{\mathbb{F}}} \omega^{(1)} \wedge \omega^{(2)} = - \int_{\Sigma^{\mathbb{F}}} \omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}d^4z = -\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}}\sqrt{\det(\iota^*g)}^{-1}\mathcal{V}_{\mathbb{F}},\tag{B.5}$$

where $\mathcal{V}_{\mathbb{F}}$ is the volume of the 4-cycle $\Sigma^{\mathbb{F}}$. Thus one concludes

$$\omega_{i\bar{j}}^{(1)}\omega_{k\bar{l}}^{(2)}\epsilon^{ik}\epsilon^{\bar{j}\bar{l}} = -\sqrt{\det(\iota^*g)}\mathcal{V}_{\mathbb{F}}^{-1} \int_{\Sigma^{\mathbb{F}}} \omega^{(1)} \wedge \omega^{(2)}.\tag{B.6}$$

Plugging this into (B.3), we derive

$$\sqrt{\det((\iota^*g) + \mathcal{F})} = \mathcal{V}_{\mathbb{F}}^{-1}\sqrt{\det(\iota^*g)}\sqrt{\frac{1}{4}\left(\int_{\Sigma^{\mathbb{F}}} \iota^*J \wedge \iota^*J - \int_{\Sigma^{\mathbb{F}}} \mathcal{F} \wedge \mathcal{F}\right)^2 + \left(\int_{\Sigma^{\mathbb{F}}} \iota^*J \wedge \mathcal{F}\right)^2}.\tag{B.7}$$

Using this in (2.16), we again arrive at

$$\Gamma_{\mathbb{F}} = \sqrt{\frac{1}{4}\left(\int_{\Sigma^{\mathbb{F}}} \iota^*J \wedge \iota^*J - \int_{\Sigma^{\mathbb{F}}} \mathcal{F} \wedge \mathcal{F}\right)^2 + \left(\int_{\Sigma^{\mathbb{F}}} \iota^*J \wedge \mathcal{F}\right)^2},\tag{B.8}$$

which coincides with (2.27).

C. Gauge kinetic function

The imaginary part of the gauge kinetic function is due to the following terms of the Chern-Simons action:

$$\begin{aligned}-\frac{1}{2}\mu_7 \int_{\mathcal{W}} \iota^*C_4 \wedge \mathcal{F} \wedge \mathcal{F}, \\ -\frac{1}{4!}\mu_7 \int_{\mathcal{W}} C_0 \wedge \mathcal{F} \wedge \mathcal{F} \wedge \mathcal{F} \wedge \mathcal{F}.\end{aligned}\tag{C.1}$$

In the first term, C_4 lives purely internal on $\Sigma^{\mathbb{F}}$ while both \mathcal{F} denote external field strengths with indices along \mathcal{M}_4 :

$$-\frac{1}{2}\mu_7(2\pi\alpha')^2 \int_{\mathcal{M}_4} \left(\int_{\Sigma^{\mathbb{F}}} \iota^* C_4 \right) F \wedge F = -\frac{1}{2}\mu_7(2\pi)^2 \alpha'^4 \int_{\mathcal{M}_4} a^{\mathbb{F}} F \wedge F, \quad (\text{C.2})$$

where we defined $a^{\mathbb{F}} = \alpha'^{-2} \int_{\Sigma^{\mathbb{F}}} \iota^* C_4$.

In the second term two \mathcal{F} are assumed to live purely on \mathcal{M}_4 while the other two \mathcal{F} are assumed to have internal indices:

$$-\frac{1}{2}\mu_7(2\pi\alpha')^2 \int_{\Sigma^{\mathbb{F}}} \mathcal{F} \wedge \mathcal{F} \int_{\mathcal{M}_4} \frac{1}{2} C_0 F \wedge F = -\frac{1}{2}\mu_7(2\pi)^2 \alpha'^4 f_{\mathbb{F}} \int_{\mathcal{M}_4} C_0 F \wedge F, \quad (\text{C.3})$$

where we expanded the 2-form flux as in (2.22) and $f_{\mathbb{F}}$ was defined in (2.30).

Thus, the imaginary part of the gauge kinetic function $\text{Im } f^{\mathbb{F}}$ is given by

$$\text{Im } f^{\mathbb{F}} = \mu_7(2\pi)^2 \alpha'^4 \left(a^{\mathbb{F}} + f_{\mathbb{F}} C_0 \right) = \frac{1}{(2\pi)^5} \left(a^{\mathbb{F}} + f_{\mathbb{F}} C_0 \right), \quad (\text{C.4})$$

where we made use of (2.37) in the second equality. Comparing with the real part (2.41) and demanding holomorphicity of the gauge kinetic functions, we arrive at (2.43).

Let us now take the curvature corrections into account. The corrections to the CS-action were already given in (2.63). Since only terms of the CS-action contribute to the imaginary part of the gauge kinetic function which possess two external field-strengths, we see that only the following additional terms arise:

$$-\frac{1}{2} \frac{1}{48} \mu_7 \int_{\mathcal{W}} C_0 \wedge \mathcal{F} \wedge \mathcal{F} \wedge \left(-p_1(T(\Sigma^{\mathbb{F}})) + p_1(N(\Sigma^{\mathbb{F}})) \right), \quad (\text{C.5})$$

where \mathcal{F} is an external field strength. Hence, similar to (C.3) we obtain

$$-\frac{1}{48} \mu_7(2\pi\alpha')^2 \int_{\Sigma^{\mathbb{F}}} \left(-p_1(T(\Sigma^{\mathbb{F}})) + p_1(N(\Sigma^{\mathbb{F}})) \right) \int_{\mathcal{M}_4} \frac{1}{2} C_0 F \wedge F. \quad (\text{C.6})$$

Thus the curvature corrected imaginary part of the gauge kinetic function is given by

$$\text{Im } f^{\mathbb{F}} = \frac{1}{(2\pi)^5} \left(a^{\mathbb{F}} + \hat{f}_{\mathbb{F}} C_0 \right), \quad (\text{C.7})$$

where $\hat{f}_{\mathbb{F}}$ was defined in (2.64).

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