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## Radion Stabilization In 5D SUGRA

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### Abstract

We present a detailed study of radion stabilization within 5D conformal SUGRA compactified on an  $S^{(1)}/Z_2$  orbifold. We use an effective 4D superfield description developed in our previous work. The effects of tree level bulk and boundary couplings, and in particular of one loop contributions and of a non perturbative correction on the radion stabilization are investigated. We find new examples of radion stabilization in non SUSY and (meta-stable) SUSY preserving Minkowski vacua.

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# 1 Introduction

The stabilization of moduli is a central problem if one is searching for vacua of string theory and their possible realization in the real world. Models with one extra dimension allow to study this and other questions in a simplified setting. They might lead to interesting insights in physics [1], e.g. concerning the generation of hierarchies, the discussion of SUSY breaking, or a new look to inflation in early cosmology. They even may be directly related to string/M-theory in an intriguing way, if Calabi-Yau and 5D compactification differ in scale. Last not least 5D supergravity is an (admittedly still technically complicated but) in principle well understood field.

For supergravity with a fifth dimension on a  $S^{(1)}/Z_2$  orbifold off-shell SUGRA (i.e. with auxiliary field components, not integrated out) certainly is a very good choice. This was pioneered in [2]. In our line of research [3–5] we rely on the work of Fujita, Kugo, Ohashi [6] on 5D superconformal gravity including also vector and hypermultiplets and using compensators of that type. We have reformulated important parts of this work<sup>4</sup> in terms of a 4D superfield formalism (see also [8] and subsequent developments [9]) including vector and chiral multiplets and 4D supergravity. In a recent paper we [5] succeeded to integrate out a multiplier real superfield  $W_y$  contained in the even part of the 5D Weyl multiplet. Thus we obtained a superspace action where the radion field  $e_y^5$  is now contained only in chiral superfields. Depending on the problem one wants to attack one can gauge fix the superconformal symmetries already in the 5D theory or only in the effective 4D theory which turns out to be still 4D superconformal invariant in this version.

There have been numerous studies of moduli (radion) stabilization (see for instance [10]–[19], [5] and references therein) and various mechanisms were suggested. In this work we consider this issue within our formalism which greatly simplifies investigations. Also, new examples of moduli stabilization are presented.

The paper is organized as follows: Starting point in sect. 2 is a brief presentation of our formalism developed further in our recent work [5]. In sect. 3 the issue of moduli stabilization is discussed and examples of stabilization due to 1-loop corrections and non perturbative effects are presented. In section 4 we discuss the uplift and stabilization in SUSY breaking and SUSY preserving (meta-stable) Minkowski ( $M_4$ ) vacua. In sect. 5 we present some mass formulas and estimates. We also mention some possible phenomenological implications of considered scenarios. In sect. 6 we conclude with a short summary. Appendix A deals with calculational tools of the various 1-loop corrections to the Kähler potential.

## 2 Effective 4D Superconformal Description

We use an effective 4D superconformal description of 5D SUGRA in which case the gauge and hypermultiplet action can be built by knowing the forms of the Kähler potential  $\mathcal{K}$ , the superpotential  $W$  and the gauge kinetic function  $f_{IJ}$ . The Lagrangian couplings are

$$\mathcal{L}_D^{(4D)} = -3 \int d^4\theta e^{-\mathcal{K}/3} \phi^+ \phi , \quad (1)$$

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<sup>4</sup>What is still lacking is a superfield formulation of the odd part of the Weyl multiplet and of related covariant derivatives. This would be needed for a genuine loop calculation containing gravitational higher KK modes in super language like in [7].

$$\mathcal{L}_W^{(4D)} = \int d^2\theta \phi^3 W , \quad (2)$$

$$\mathcal{L}_V^{(4D)} = \frac{1}{4} \int d^2\theta f_{IJ}(\Sigma) \tilde{\mathcal{W}}^{\alpha I} \tilde{\mathcal{W}}_\alpha^J + \text{h.c.} \quad (3)$$

where  $\phi$  is the 4D conformal compensator, a chiral superfield.

If the theory under discussion is 5D orbifold SUGRA, then as derived in [5], at tree level we have the Kähler potential

$$\mathcal{K} = -\ln \tilde{\mathcal{N}}(\Sigma + \Sigma^\dagger) - 2 \ln \left( 1 - H^+ e^{-g_I V^I} H \right) , \quad (4)$$

where  $\Sigma$  is a modulus emerging from a 5D gauge supermultiplet and  $H$  is a chiral superfield which is part of the 5D hypermultiplet and which is charged under the 4D vector superfields  $V^I$  (here we assume that the hyperscalar manifold is the simplest one). The  $\tilde{\mathcal{W}}$  denotes the 4D gauge field-strength superfield, while  $f_{IJ}(\Sigma) = -\hat{\mathcal{N}}_{IJ}(\Sigma)$  with double derivative  $\hat{\mathcal{N}}_{IJ}$  obtained from the 3<sup>rd</sup> order norm function  $\mathcal{N}(\Sigma)$  by setting at the end the odd fields to zero (more relations between 4D and 5D objects are given below).

The scalar potential derived from (1)-(3) consists of two parts: the F-term potential is given in terms of the Kähler potential and the superpotential as

$$V_F = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{I\bar{J}} D_I W D_{\bar{J}} \bar{W} - 3|W|^2 \right) , \quad \text{with } D_I \equiv \partial_I + \mathcal{K}_I \quad (5)$$

where  $I$  runs over chiral multiplets emerging from both hyper and vector scalars ( $I = H, \Sigma$ ). It is easy to write down also the  $D$ -term potential

$$V_D = -\frac{1}{4} M_{\text{Pl}}^4 \hat{\mathcal{N}}_{IJ} D^I D^J , \quad \text{with } D^I = -\hat{\mathcal{N}}^{IJ} g_{AJ} q_i \mathcal{K}_{H_i} H_i , \quad (6)$$

where we took into account that the  $\mathcal{K}$  includes the hypermultiplets as  $\mathcal{K}(H_i^\dagger e^{-q_i g_{4I} V^I} H_i)$  (here we have restored the  $U(1)$  charge  $q_i$  of the chiral superfield  $H_i$ ). Eq. (6) does not contain contributions from FI terms arising if the compensator hypermultiplet is charged under  $U(1)_R$ .  $D$ -terms may play crucial role for uplifting to the Minkowski or de-Sitter vacua [19]. In this work we will not consider  $D$ -term potentials but instead will investigate the role of  $F$ -term potentials (of eq. (5)) in the radion stabilization.

Our investigation will focus mostly on the *vector* moduli  $\Sigma$ , as the *hyper* moduli  $H$  can be stabilized in a fairly simple way by using tree-level brane potentials. Stabilization of vector moduli is more complicated because they do not couple directly to the branes. Here indeed lies the difficulty of radion stabilization. Recall that the radion (the  $e_y^5$  component of 5D Weyl multiplet) is related to the chiral superfields  $\Sigma_5^I$

$$\Sigma_5^I = \frac{1}{2} (e_y^5 M^I + i A_y^I) + \dots \quad (7)$$

Since for this paper we consider unwarped geometries, the zero modes of the various superfields are  $y$ -independent. In the following we consider a dimensionless  $e_y^5$  and a fixed length  $R$  ( $-\pi R < y < \pi R$ ) is introduced. Note that  $R$  is not dynamical, but for practical reasons we will take it to coincide with the size of the extra-dimension (after the latter is stabilized).

Before closing this section we present relations between 5D and 4D objects and give some explanation for the notations used in the paper. The fields  $\Sigma$ ,  $H$  in eq. (4) and appearing in 4D Lagrangian densities are dimensionless and are related to the corresponding 5D components as

$$\Sigma = \kappa_5 \Sigma_5, \quad H = \kappa_5 H_5, \quad \text{with} \quad \kappa_5 = M_5^{-3/2}, \quad M_{\text{Pl}}^2 = 2\pi R M_5^3. \quad (8)$$

Moreover, we have

$$\tilde{\mathcal{N}} = \kappa_5^2 \mathcal{N}, \quad \tilde{\mathcal{W}} = \sqrt{2\pi R} \mathcal{W}, \quad \sqrt{2\pi R} g_4 = g_5. \quad (9)$$

At some places (for instance in eqs. (14), (16), (A.9), (A.21)) we also present 5D Lagrangian couplings. In order to make expressions more compact there we omit subscript 5 for the 5D objects, hoping that this will not create any confusion.

### 3 Moduli Stabilization

We start our discussion with a Kähler potential including gauge moduli only. At tree level, due to the no-scale nature of the  $\Sigma$ , which means that at tree-level  $\mathcal{K}^{\Sigma\Sigma} \mathcal{K}_\Sigma \mathcal{K}_{\bar{\Sigma}} = 3$ , and with moduli independent superpotential  $W = \text{constant}$ , from (5) we see that a vector modulus is a flat direction and  $V_F = 0$ . Inclusion of moduli dependent tree level superpotentials (consistent with gauge symmetries) can lead to moduli stabilization however in an AdS vacuum in simple minded cases [5]. The vacuum uplifting is an important issue and different mechanisms can be applied. Later in this paper we will use and discuss some of them.

As far as the hypermultiplet  $H$  moduli ('hypers') are concerned, they can be stabilized in a simpler way. The brane localized superpotential  $W(H)$  can fix the value of  $H$ . Apart from this, hyper moduli generate positive definite contribution to the potential

$$V_F(H) = M_{\text{Pl}}^4 e^{\mathcal{K}} \mathcal{K}^{H\bar{H}} D_H W D_{\bar{H}} \bar{W}, \quad (10)$$

which can also serve for a self consistent uplift as we will discuss in more detail in section 4.

#### 3.1 Stabilization and uplift by 1-loop corrections

As originally proposed in [11, 12], one-loop corrections to the Kähler potential, due to bulk multiplets, can lead to moduli stabilization in  $M_4$  vacuum. A more detailed discussion of this kind of stabilization mechanism was then presented in refs. [13, 14] (see [15] for a discussion of two-loop effects), for the one modulus (i.e. the radion) case. Here we will present the supergravity embedding of these models, generalizing them to the many moduli case, and consider *new* one-loop corrections to  $\mathcal{K}(\Sigma, \Sigma^+)$ . The new corrections will be crucial to remove the flatness of the potential in the axionic directions, as they break the continuous *shift symmetry* present at tree-level.

There are loop corrections to the Kähler potential which may be either moduli-independent or moduli-dependent. Moduli dependence can appear in a combination which measures the size of the extra dimension, the radion. Moduli also may appear in different combinations. The latter happens when some multiplets are charged under some bulk abelian gauge symmetries (that is, some isometries of the scalar manifolds are gauged by vector multiplets). These two kinds of corrections correspond to the contributions of massless and massive bulk multiplets of rigid supersymmetry of refs. [12, 14].

In our studies we will include moduli-independent corrections coming in general from bulk Weyl, vector and hyper supermultiplets. For the generation of moduli dependent corrections we will introduce bulk hypermultiplets. As we will see shortly, each type of correction will be important for moduli stabilization in Minkowski vacuum. Let's consider these 1-loop corrections to the Kähler potential in more detail.

### Moduli-independent one-loop contributions

We consider first the contributions of the gravitational sector, of abelian vector multiplets in the unbroken phase, and of uncharged hypermultiplets. These lead to a 1-loop correction to the low-energy D-term Lagrangian of the form,

$$\Delta\mathcal{L}_V = - \int d^4\theta \frac{\alpha}{R^3 \mathbb{W}_y^2} + (\text{higher powers in a super-der. expansion}), \quad (11)$$

where  $\alpha = (-2 - N_V + N_H) \frac{\zeta(3)}{2(2\pi)^5}$  depends on the number  $N_V$  of "massless" vector multiplets and on the number  $N_H$  of "massless" hypermultiplets. Note that (11) is a 5D Lagrangian coupling. Integrating out  $\mathbb{W}_y$  [5] one finds (still in the 5D theory) an effective 4D one-loop Kähler potential

$$\mathcal{K}(\Sigma, \Sigma^+) = - \ln \left( \tilde{\mathcal{N}}(\Sigma + \Sigma^+) + \Delta \right), \quad (12)$$

where

$$\Delta = \Delta_\alpha \equiv \kappa_5^2 \frac{\alpha}{R^3}. \quad (13)$$

Note that here the  $N_H$  dependent part, arising by integrating out the hypermultiplets, can be obtained from the cases presented below by taking the superpotential coupling of moduli with the hypermultiplets to be zero (setting  $g_\beta = 0$  in eq. (15)).

### Moduli-dependent one-loop contributions

There are several ways of gauging isometries of the scalar manifold and therefore there are also different possible corrections to  $\mathcal{K}(\Sigma, \Sigma^+)$ :

Let us start with a physical 5D hypermultiplet  $\mathbb{H} = (H, H^c) \sim (+, -)$  charged under the vector multiplet  $\mathbb{V} = (V, \Sigma) \sim (-, +)$ , where the parities under  $S^{(1)}/Z_2$  orbifolding are shown in brackets. The relevant coupling is [3]

$$\int d^2\theta (2H^c \partial_y H - g_\beta \Sigma (H^2 - H^c{}^2)) + \text{h.c.} . \quad (14)$$

Note again that eq. (14) is a 5D Lagrangian density (it should not be confused with the effective 4D description of (1)-(3)) and all fields and coupling is five dimensional. This form is useful as a starting point when making loop calculations. The corresponding one-loop correction due to hypermultiplets running in the loop, obtained in appendix A.1, is

$$\Delta_\beta = \kappa_5^2 \frac{\beta}{R^3} \sum_{k=1}^{\infty} \frac{1}{k^3} (e^{-k2\pi R g_\beta \Sigma} + e^{-k2\pi R g_\beta \Sigma^+}) (1 + k\pi R g_\beta (\Sigma + \Sigma^+)) , \quad (15)$$

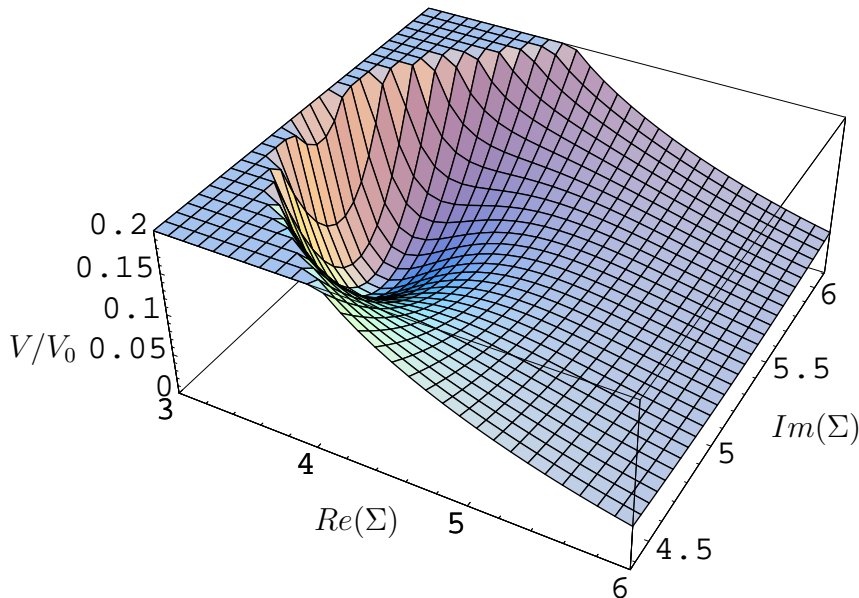


Figure 1: Effective potential as a function of  $Re(\Sigma)$  and  $Im(\Sigma)$ , for  $\bar{\alpha} = 0.437$ ,  $\bar{\beta} = 0.5$ ,  $\bar{\gamma} = 1.5$  and  $b = 0.3$ ,  $c = 1$ .  $V_0 = 8.5 \cdot 10^{-6} M_{\text{Pl}}^4 |W|^2$ .

with  $\beta = N_{H_\beta}/4(2\pi)^5$ , where  $N_{H_\beta}$  is the number of charged hypermultiplets. There are several ways of deriving eq.(15), the most elegant one using the relation of the Kähler potential of  $\mathcal{N} = 2$  supergravity to the prepotential. Another way is to make a KK reduction and then to use the standard tools in order to evaluate the 1-loop correction to the Kähler potential. These two ways are presented in appendix A.1. For our purposes the higher powers in the superderivative expansion do not play an important rôle, we will therefore discard them in the following. Another feature which we did not display is that (15) is only valid for  $Re(\Sigma) > 0$ , for  $Re(\Sigma) < 0$  we must replace  $\Sigma \rightarrow -\Sigma$ . This expression leads to new contributions to  $\mathcal{K}(\Sigma, \Sigma^+)$  in eq.(12). In the most general case  $g_\beta \Sigma$  can be replaced by a linear combination of moduli  $g_I \Sigma^I$ .

Another type of 5D coupling between a hypermultiplet and a vector multiplet with the same orbifold parities as in the previous example, is

$$\int d^2\theta \, 2H^c (\partial_y - g_\gamma \epsilon(y) \Sigma) H + \text{h.c.}, \quad (16)$$

where  $\epsilon(y) = \partial_y |y|$  is the periodic step-function. Integrating out the hypermultiplets we get the following contribution to  $\Delta$

$$\Delta_\gamma = \kappa_5^2 \frac{\gamma}{R^3} \sum_{k=1}^{\infty} \frac{1}{k^3} e^{-k\pi R g_\gamma (\Sigma + \Sigma^+)} (1 + k\pi R g_\gamma (\Sigma + \Sigma^+)) , \quad (17)$$

where  $\gamma = N_{H_\gamma}/2(2\pi)^5$  and  $N_{H_\gamma}$  is the number of hypermultiplets coupled with the modulus through the odd coupling  $g_\gamma \epsilon(y)$ . This expression can e.g. be obtained from formulas given in [13, 14] by

the 5D lifting of the 1-loop Kähler potential obtained by super-graph method. In appendix A.2 we present its derivation in detail. One crucial difference to the case discussed previously lies in the fact that the lowest KK mode is now massless before supersymmetry is broken. The other relevant difference to eq.(15) is that it doesn't depend on the axionic modulus  $Im(\Sigma)$ .

Now we are ready to study the issues of moduli stabilization. Let us consider the following setup: In addition to a constant (brane) superpotential  $W$  we include the three types of the 1-loop corrections to the Kähler potential discussed in this subsection. Furthermore we assume the radion to be the only existing modulus  $\Sigma$ . By a suitable arrangement of the parameters  $\bar{\alpha} = \kappa_5^2 \alpha / R^3 = 2\pi\alpha / (RM_{Pl})^2$ ,  $\bar{\beta} = \kappa_5^2 \beta / R^3 = 2\pi\beta / (RM_{Pl})^2$  and  $\bar{\gamma} = \kappa_5^2 \gamma / R^3 = 2\pi\gamma / (RM_{Pl})^2$ , as well as  $b = \pi R g_\beta / \kappa_5 = g_{4\beta} \pi R M_{Pl}$  and  $c = \pi R g_\gamma / \kappa_5 = g_{4\gamma} \pi R M_{Pl}$ , one can obtain a  $M_4$  (or dS) minimum for  $\Sigma$ . This is possible since the  $\Delta_\beta$  gives a correction opposite to  $\Delta_\gamma$  because the imaginary part stabilizes in such a way as to make some part of the potential negative. In Fig.1 we plot such a potential for a specific choice of parameters.  $Re(\Sigma)$  and  $Im(\Sigma)$  shown in the plot are dimensionless 4D fields (related to the 5D states according to eq. (8)). We tuned  $\alpha$  in such a way as to obtain a Minkowski minimum at  $Re(\Sigma) \simeq 3.5$ . Note that the axionic component of  $\Sigma$  is stabilized as well, in contrast to models previously proposed in the literature (see e.g. [13]).

Note that there is a class of parameters related to those of Fig. 1 for which the stabilization in the  $M_4$  vacuum is achieved. To see this, it is useful to realize that the scalar potential remains invariant (up to the scale factor) under the following rescalings

$$R \rightarrow sR, \quad (g_{4\beta}, g_{4\gamma}) \rightarrow s^{-1/3}(g_{4\beta}, g_{4\gamma}), \quad \Sigma \rightarrow s^{-2/3}\Sigma, \quad (18)$$

where the scale factor  $s$  is a real number. Note that under this transformation the scalar potential transforms as  $V \rightarrow s^2 V$  (due to the exponential  $e^{\mathcal{K}}$  in (5)). However, this change with factor  $s^2$  lives the potential at zero in the vacuum. With these rescalings the parameters of the model change as

$$(\bar{\alpha}, \bar{\beta}, \bar{\gamma}) \rightarrow s^{-2}(\bar{\alpha}, \bar{\beta}, \bar{\gamma}), \quad (b, c) \rightarrow s^{2/3}(b, c), \quad (19)$$

allowing to find the class of parameters from those corresponding to Fig. (1) and to have different values of stabilized  $\Sigma$  in the  $M_4$  vacuum.

The symmetry displayed in (18) is useful also for discussing some physical implications. With this rescaling the Kähler potential transforms as

$$\mathcal{K} \rightarrow \mathcal{K} + 2 \ln s. \quad (20)$$

Now it is easy to realize that the particle masses undergo linear rescalings. Namely, according to (53) and (50) the gravitino and canonically normalized radion component masses change as

$$m_{3/2} \rightarrow s \cdot m_{3/2}, \quad m_{\pm} \rightarrow s \cdot m_{\pm}. \quad (21)$$

Therefore, for the class of the parameters related to each other by the transformation (19) the ratio  $m_{3/2}/m_{\pm}$  is fixed

$$\frac{m_{\pm}}{m_{3/2}} = \mathcal{R}_{\pm}. \quad (22)$$

For the parameters of Fig. 1 we have  $\mathcal{R}_+ \simeq 0.17$  and  $\mathcal{R}_- \simeq 0.15$ . Thus, with  $m_{3/2} \simeq 1$  TeV we have  $m_+ \simeq 170$  GeV and  $m_- \simeq 150$  GeV for the two radion scalar component masses. More

generally, since we have a constant superpotential, for the scalar potential we have  $V \propto e^{\mathcal{K}}|W|^2$  and comparing with (53) we can expect  $m_{\pm} \sim m_{3/2}$  unless some cancellation occurs. Therefore with low energy SUSY (i.e.  $m_{3/2} \sim 1$  TeV) we expect to have radion's two real scalar modes in the 100 GeV-1 TeV range. This makes the model testable in ongoing and future collider experiments. More details concerning this issue will be discussed in sect 5.

Closing this subsection, let us mention that the conditions in (A.37) for this model can be easily satisfied with low energy SUSY breaking. This is due the fact that  $F$ -terms and the radion mass are close to the TeV scale. Therefore, all effects caused by higher super-derivative operators can be safely ignored.

### 3.2 Non-perturbative stabilization from gaugino condensation

As we have already mentioned, the flatness of the potential will be lifted either if the Kähler potential receives non cubic (with respect to  $\Sigma^I + \Sigma^{I\dagger}$ ) corrections, or the superpotential has a moduli dependent part. In this subsection we will show that the radion stabilization can be achieved by gaugino condensation. In this case the effective superpotential will have a moduli dependent (non perturbative) part. Also the Kähler potential gets a non perturbative correction. The whole effective action still can be written in a superconformal form

$$-3(e^{-\mathcal{K}/3}\phi^\dagger\phi)_D + (\phi^3W)_F + \text{h.c.} . \quad (23)$$

Discussing the gaugino condensation, we follow an effective 4D description dealing with zero modes of the relevant bulk fields. Thus, the arguments applied for a pure 4D gaugino condensation scenario will be appropriate. For this case the super and Kähler potentials include perturbative (p) and non-perturbative (np) parts [20]:

$$W = w_p + w_{np} , \quad e^{-\mathcal{K}/3} = e^{-\mathcal{K}_p/3} - ke^{-\mathcal{K}_{np}/3} , \quad (24)$$

where  $k$  is some constant. The action with (24) is valid below a scale  $\Lambda$  corresponding to energies where the gauge sector becomes strongly coupled.

For demonstrative purposes we will consider an example with one non Abelian  $SU(N)$  YM theory which is responsible for the gaugino condensation. Taking the norm function

$$\mathcal{N}(M^I) = M^3 - M\text{Tr}(M_g)^2 , \quad \text{with} \quad M_g = \frac{1}{2}\lambda^a M^a , \quad a = 1, \dots, N^2 - 1 , \quad (25)$$

the coupling of the moduli superfield  $\Sigma$  with the gauge field strength will be

$$-\frac{1}{4}(\mathcal{N}_{IJ}(\Sigma)\mathcal{W}^I\mathcal{W}^J)_F + \text{h.c.} \rightarrow \frac{1}{4}(\Sigma\mathcal{W}^a\mathcal{W}^a)_F + \text{h.c.} \quad (26)$$

The effective 4D superconformal theory with (23), (26) possesses Weyl and chiral  $U(1)$  symmetries at classical level. However, one superposition of these two  $U(1)$ s is anomalous on the quantum level. Namely, a mixed gauge-chiral anomalous term is generated and the counter term [21]

$$-2c\left(\frac{1}{4}\ln\phi\mathcal{W}^a\mathcal{W}^a\right)_F + \text{h.c} \quad (27)$$



is needed to take care of the anomaly cancellation. The coefficient  $c$  in (27) is related to the gauge group  $b$ -factor and is positive for asymptotically free theories.

Therefore, the coupling of the composite chiral superfield  $U = \langle \mathcal{W}^a \mathcal{W}^a \rangle$  with moduli and compensating superfields is given by

$$\frac{1}{4} ((\Sigma - 2c \ln \phi)U)_F + \text{h.c.} \quad (28)$$

Eq. (28) is the starting point for computing the non perturbative effective action  $\Gamma(\Sigma, U)$ . After obtaining the form of  $\Gamma(\Sigma, U)$ , one can minimize it with respect to  $U$  (determining the condensate  $U_0$ ) and then plug back the value of  $U = U_0$  in  $\Gamma(\Sigma, U)$  to derive the effective action for  $\Sigma$ -moduli. In case of a single condensate the non perturbative Kähler and superpotentials are given by [20]

$$\mathcal{K}_{np} = \frac{3}{2c}(\Sigma + \Sigma^\dagger), \quad w_{np} = \tilde{w} e^{-\frac{3\Sigma}{2c}}. \quad (29)$$

The tree level perturbative Kähler potential has the form  $\mathcal{K}_p = -3 \ln(\Sigma + \Sigma^\dagger)$ . Therefore the total Kähler potential will be

$$\mathcal{K} = -3 \ln \left( \Sigma + \Sigma^\dagger - k e^{-\frac{\Sigma + \Sigma^\dagger}{2c}} \right). \quad (30)$$

We also assume that the perturbative part of the superpotential is moduli independent  $w_p = \text{const.}$  (often this is the case at tree level). Note that a non perturbative correction to the Kähler potential in the problem of radion stabilization has not been taken into account so far in the literature. There is no physical reason to exclude the  $k$ -term from considerations. It also contributes to the kinetic coupling  $(\Sigma \Sigma^\dagger)_D$  for the moduli field. Since the  $k$  is related to a mass scale, we expect that it will have impact on the stabilized value of  $\Sigma$  moduli. With  $k \neq 0$  the eq. (30) lifts the flatness of the  $\mathcal{K}_I \mathcal{K}^{I\bar{J}} \mathcal{K}_{\bar{J}} - 3$  part. The resulting potential has the form

$$V = M_{\text{Pl}}^4 \frac{3|w_p|^2 \left| \frac{\tilde{w}}{w_p} \right|^2 (2+r)e^{-2r} + 2 \left| \frac{\tilde{w}}{w_p} \right| \cos(\Omega - 1.5\rho) e^{-0.5r} - \alpha}{(2c)^3 (r - \alpha e^{-r})^2 (1 + 2\alpha e^{-r} + \alpha r e^{-r})} e^{-r},$$

with  $r = \frac{1}{c} \text{Re}(\Sigma)$ ,  $\rho = \frac{1}{c} \text{Im}(\Sigma)$ ,  $\alpha = \frac{k}{2c}$ ,  $\Omega = \text{Arg} \left( \frac{\tilde{w}}{w_p} \right)$ . (31)

For  $\tilde{w} = 0$  the radion stabilization is impossible. Thus, the moduli dependent superpotential plays an important role. For  $\tilde{w} \neq 0$  the complex part of the modulus gets fixed as

$$\langle \rho \rangle = \frac{2}{3}(\Omega - \pi), \quad (32)$$

and also the real part is stabilized. In Fig. 2 we give plots for the cases  $k = 0$ ,  $k > 0$  and  $k < 0$ , where for the complex part we have taken the value of eq. (32). Obviously, the inclusion of the non perturbative part of the Kähler potential introduces significant changes. In all cases the vacuum energy in the minimum is negative. In fact, one can easily check that SUSY is unbroken in the minimum:  $D_\Sigma W = 0$ . The solution of the latter equation, i.e.

$$\left| \frac{\tilde{w}}{w_p} \right| (1+r) e^{-1.5r} e^{i(\Omega - 1.5\rho)} + 1 + \alpha e^{-r} = 0, \quad (33)$$

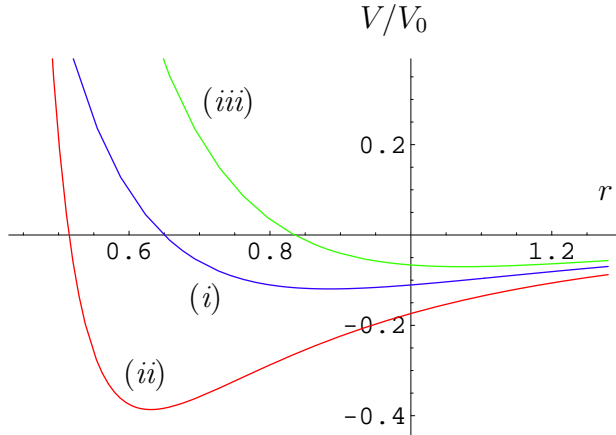


Figure 2: Moduli potential for different cases.  $V_0 = M_{\text{Pl}}^4 \frac{3|w_p|^2}{(2c)^3}$ ,  $\frac{\tilde{w}}{w_p} = 2$ . (i)  $k = 0$ ; (ii)  $\frac{k}{2c} = 0.5$ ; (iii)  $\frac{k}{2c} = -0.5$ .

coincides with the minimum of the potential in (31). Now it is clear why the vacuum ( $V = -3e^{\mathcal{K}}|W|^2$ ) is AdS. An additional contribution to the potential must be generated [10] in order to set it to zero.

We have seen that even though moduli stabilization can be achieved with the inclusion of non-perturbative effects due to bulk gaugino condensation, this mechanism still lacks a consistent *uplift* mechanism in order to obtain Minkowski or de Sitter vacua. Such possibilities will be discussed in the next section.

## 4 Uplift by Boundary Couplings and Stabilization in SUSY Minkowski Vacuum

In this section we consider the possibility of uplifting through the boundary couplings. Also the stabilization in a SUSY preserving Minkowski vacuum is discussed.

The general form of the scalar ( $F$ -term) potential is given by

$$V = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{I\bar{J}} D_I W D_{\bar{J}} \bar{W} - 3|W|^2 \right) . \quad (34)$$

Since the first term in (34) is positive definite, we see that with fixed  $W$  the minimal possible value of  $V$  is achieved if the conditions

$$D_I W = \partial_I W + \mathcal{K}_I W = 0 \quad (35)$$

are satisfied. With (35) we have indeed  $\frac{\partial V}{\partial \Sigma^I} = -3M_{\text{Pl}}^4 e^{\mathcal{K}} \bar{W} D_I W = 0$  and therefore the SUSY preserving solution (35) is an extremum. However, the vacuum is AdS ( $V = -3M_{\text{Pl}}^4 e^{\mathcal{K}} |W|^2$ ) unless the superpotential along this solution is zero (this possibility will be discussed later on).

It must be noticed that the extremum equation  $\frac{\partial V}{\partial \Sigma^I} = 0$  can have solution(s) along which (35) is not satisfied. Therefore, there is a chance of having a (meta-stable) vacuum with broken SUSY.

If in this minimum it is possible to set (fine tune)  $V = 0$ , we can have stabilization in a Minkowski vacuum.

The chiral superfields coming from 5D hypermultiplets can play a crucial role for uplifting. Consider the case with one bulk gauge modulus  $\Sigma$  and a set of 4D chiral multiplets  $H_i$  coming from bulk hypers (the  $H_i$  are even under orbifold parity). Assume that the superpotential  $W = W(\Sigma, H_i)$  is arranged in such a way that there is a minimum with broken SUSY (i.e. either  $D_\Sigma W$  or at least one  $D_{H_i} W$  is non vanishing). Then along this configuration the potential is

$$V = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{\Sigma\Sigma^\dagger} |D_\Sigma W|^2 + \mathcal{K}^{H_i H_i^\dagger} |D_{H_i} W|^2 - 3|W|^2 \right) . \quad (36)$$

If the potential (36) fixes  $\Sigma$  and  $H_i$  and if in addition

$$\mathcal{K}^{\Sigma\Sigma^\dagger} |D_\Sigma W|^2 + \mathcal{K}^{H_i H_i^\dagger} |D_{H_i} W|^2 = 3|W|^2 \quad (37)$$

is arranged, then the vacuum is  $M_4$ .

This can be applied for the case with the gaugino condensation discussed in the previous section. Another interesting example which gives tree level stabilization is with linear superpotential

$$W(\Sigma) = g\Sigma . \quad (38)$$

As was discussed in [5] this superpotential arises when the compensator hypers are charged under  $\Sigma$  (the  $U(1)_R$  is gauged). Including also the superpotential  $W(H)$  for further chiral superfields, the total superpotential will be

$$W = g\Sigma + W(H) . \quad (39)$$

Restricting ourself to one  $H$ , coming from the bulk hyper, the Kähler potential is

$$\mathcal{K} = -3 \ln(\Sigma + \Sigma^\dagger) - 2 \ln(1 - H^\dagger H) . \quad (40)$$

For  $W$  and  $\mathcal{K}$  given in (39) and (40), one can show that although uplift to  $M_4$  can be obtained, not all states are stabilized. Some of them remain massless. This happens for an arbitrary form of  $W(H)$  (an interesting result). The situation can be improved if for example in  $\mathcal{K}$  the  $\Delta$  corrections (discussed in sect. 3.1) are included. However, this goes beyond the tree level analysis and is more involved. Instead, below we consider a model with a superpotential combining a gaugino condensation part and the linear one given in (38). Although separately both parts lead to the stabilization in AdS vacua, we will see that together they provide an elegant stabilization in a SUSY preserving  $M_4$  vacuum. Thus, no inclusion of the hypermultiplets is needed.

The reason we below concentrate on a stabilization in a SUSY preserving Minkowski vacuum is the following. The phenomenologically viable SUSY breaking can occur in the low energy effective 4D theory by one of the mechanisms widely discussed in the literature. Note that whatever is a SUSY breaking mechanism, it is preferable to have the gravitino mass much smaller than the mass of the moduli. In this way the latter will not be destabilized by the SUSY breaking effects. This also has other advantages especially from cosmological viewpoints [24].

## Stabilization in SUSY Minkowski vacuum

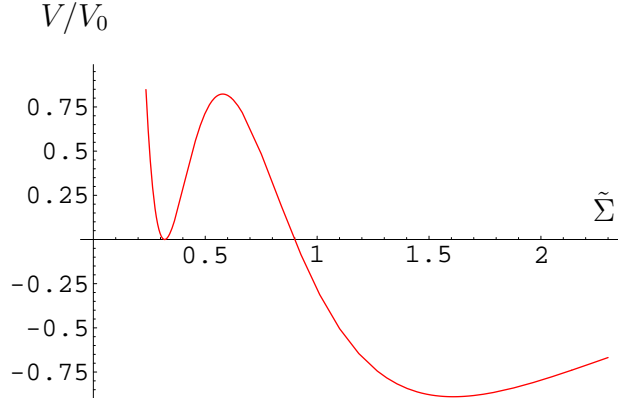


Figure 3: Moduli potential for  $\frac{g}{a} = 0.02$ ,  $b = 2.3$ .  $\tilde{\Sigma} \equiv \ln(2\text{Re}(\Sigma)/3)$ ,  $V_0 = 10^{-5}a^2M_{\text{Pl}}^4$ .

Now we discuss the issue of the stabilization in a SUSY preserving  $M_4$  vacuum. In order to have unbroken SUSY, the conditions (35) must be satisfied for all fields. With these conditions, on the other hand, the potential of (34) is negative definite unless the superpotential vanishes in the vacuum. Thus, in order to have the Minkowski vacuum we require for the superpotential

$$\langle W \rangle = 0 . \quad (41)$$

This condition in most cases can be achieved by fine tuning. Moreover, it is very likely that there is another solution with preserving SUSY but  $\langle W \rangle \neq 0$ . This means that the vacuum (41) is a local minimum. If the latter is sufficiently long lived we should not worry much about it and proceed with model building. Thus, we further invoke the conditions (35), (41) which are equivalent to

$$\langle \partial_I W \rangle = 0 , \quad \langle W \rangle = 0 , \quad I = 1, 2, \dots \quad (42)$$

Of course, one should make sure that the vacuum determined by these conditions is a minimum. Thus, second derivatives should be evaluated.

With conditions (42) we have from (34)

$$V_{I\bar{J}} \equiv \frac{\partial^2 V}{\partial \Sigma^I \partial \bar{\Sigma}^{\bar{J}}} = M_{\text{Pl}}^4 e^{\mathcal{K}} \mathcal{K}^{M\bar{N}} (\partial_M \partial_I W) (\partial_{\bar{N}} \partial_{\bar{J}} \bar{W}) , \quad V_{IJ} = V_{\bar{I}\bar{J}} = 0 . \quad (43)$$

The  $V_{I\bar{J}}$  is positive definite and one should make sure that there is no massless mode(s) in the spectrum.

In case of one modulus, we have only  $V_{\Sigma\Sigma^\dagger}$  and investigation is simplified. Consider the superpotential couplings

$$W = g\Sigma + ae^{-b\Sigma} - \frac{g}{b} \ln \frac{eab}{g} . \quad (44)$$

The first term emerges by  $U(1)_R$  gauging, while the second term can come from gaugino condensation as was discussed in the previous section. All conditions (42) are satisfied and  $\Sigma$  gets stabilized at

$$\langle \Sigma \rangle = \frac{1}{b} \ln \frac{ab}{g} . \quad (45)$$

Note that this result does not depend on the form of the Kähler potential because the latter does not appear in the conditions of eq. (42). Moreover, the corresponding second derivative  $\partial_\Sigma^2 W = gb$  is non zero. Therefore, the desired stabilization in a SUSY preserving Minkowski vacuum is obtained (imaginary part of  $\Sigma$  is also fixed). However, there is another SUSY preserving vacuum along the solution  $D_\Sigma W = 0$  (with  $W \neq 0$ ) which is AdS. The potential's profile for one possible choice of parameters (which gives  $Im(\Sigma) = 0$ ) is shown in Fig. 3. As was analyzed in early times [25] the Minkowski vacuum can be fairly long lived, i.e. meta-stable (for a recent discussion see [26]).

Before concluding this section, it is useful to present some expressions for the case of two fields. Consider one gauge modulus  $\Sigma$  and a 4D chiral superfield  $H$  (of positive parity) coming from a bulk hypermultiplet. In general, the superpotential  $W$  is a function of both superfields  $W = W(\Sigma, H)$ . For bulk modulus and hyper, the tree level Kähler potential is given by (4) and therefore  $\mathcal{K}_{\Sigma H^\dagger} = 0$ . From (43) we have

$$\begin{aligned} V_{\Sigma\Sigma^\dagger} &\equiv \frac{\partial^2 V}{\partial \Sigma \partial \Sigma^\dagger} = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{\Sigma\Sigma^\dagger} |\partial_\Sigma^2 W|^2 + \mathcal{K}^{HH^\dagger} |\partial_\Sigma \partial_H W|^2 \right) , \\ V_{\Sigma H^\dagger} &\equiv \frac{\partial^2 V}{\partial \Sigma \partial H^\dagger} = \left( \frac{\partial^2 V}{\partial \Sigma^\dagger \partial H} \right)^* = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{\Sigma\Sigma^\dagger} (\partial_\Sigma^2 W) (\partial_{\Sigma^\dagger} \partial_{H^\dagger} \bar{W}) + \mathcal{K}^{HH^\dagger} (\partial_\Sigma \partial_H W) (\partial_{H^\dagger} \bar{W}) \right) , \\ V_{HH^\dagger} &\equiv \frac{\partial^2 V}{\partial H \partial H^\dagger} = M_{\text{Pl}}^4 e^{\mathcal{K}} \left( \mathcal{K}^{\Sigma\Sigma^\dagger} |\partial_\Sigma \partial_H W|^2 + \mathcal{K}^{HH^\dagger} |\partial_H^2 W|^2 \right) , \end{aligned} \quad (46)$$

and all remaining second derivatives vanish. In order the vacuum to be a local minimum all eigenvalues of the matrix

$$\begin{pmatrix} 4V_{\Sigma\Sigma^\dagger} & 0 & V_{\Sigma H^\dagger} + V_{\Sigma^\dagger H} & i(V_{\Sigma^\dagger H} - V_{\Sigma H^\dagger}) \\ 0 & 4V_{\Sigma\Sigma^\dagger} & i(V_{\Sigma H^\dagger} - V_{\Sigma^\dagger H}) & V_{\Sigma H^\dagger} + V_{\Sigma^\dagger H} \\ V_{\Sigma H^\dagger} + V_{\Sigma^\dagger H} & i(V_{\Sigma H^\dagger} - V_{\Sigma^\dagger H}) & 4V_{HH^\dagger} & 0 \\ i(V_{\Sigma^\dagger H} - V_{\Sigma H^\dagger}) & V_{\Sigma H^\dagger} + V_{\Sigma^\dagger H} & 0 & 4V_{HH^\dagger} \end{pmatrix} , \quad (47)$$

must be positive. This condition should be required together with (42). A more detailed investigation will be required if by chance some of the eigenvalues vanish.

In the concrete case with a superpotential given by  $W = W(\Sigma) + W(H)$  all off diagonal elements in (47) vanish and the eigenvalues are positive semidefinite. Thus, the model building is extremely simplified.

## 5 Mass Spectrum and Phenomenological Implications

Now we present mass formulas for the stabilized modulus and also for the 4D gravitino. We consider the case of one modulus without hypers. Inclusion of the hypers and boundary chiral multiplets and a generalization of the corresponding expressions are straightforward.

The canonically normalized scalar component  $\Sigma_c = \frac{1}{\sqrt{2}}(\Sigma_c^R + i\Sigma_c^I)$  of the modulus superfield is related to  $\Sigma$  by the differential relation

$$d\Sigma_c = M_{\text{Pl}} \mathcal{K}_{\Sigma\Sigma^\dagger}^{1/2} d\Sigma . \quad (48)$$

After  $Re(\Sigma)$  and  $Im(\Sigma)$  get stabilized, in general  $\Sigma_c^R - \Sigma_c^I$  mixing occurs and the squared mass matrix is given by

$$\begin{pmatrix} V_{RR} & V_{RI} \\ V_{RI} & V_{II} \end{pmatrix}, \quad \text{with}$$

$$V_{RR} + V_{II} = \frac{2}{M_{\text{Pl}}^2 \mathcal{K}_{\Sigma\Sigma^\dagger}} V_{\Sigma\Sigma^\dagger}, \quad V_{RR}V_{II} - V_{RI}^2 = \frac{1}{M_{\text{Pl}}^4 \mathcal{K}_{\Sigma\Sigma^\dagger}^2} (V_{\Sigma\Sigma^\dagger}^2 - V_{\Sigma\Sigma} V_{\Sigma^\dagger\Sigma^\dagger}). \quad (49)$$

Here and below we omit  $\langle \dots \rangle$  symbols keeping in mind that all quantities are evaluated in the vacuum. The eigenvalues of the mass matrix in (49) are given by

$$m_{\pm}^2 = \frac{1}{M_{\text{Pl}}^2 \mathcal{K}_{\Sigma\Sigma^\dagger}} (V_{\Sigma\Sigma^\dagger} \pm |V_{\Sigma\Sigma}|). \quad (50)$$

These mass scales are phenomenologically constrained. Namely, the radion's scalar component's masses should not be less than  $\text{few} \cdot 10^{-3}$  eV [22] in order to avoid an unacceptable large deviation from the Newton's 4D gravitational potential. The source for the deviation is the exchange of light scalar modes.

In case of an unbroken SUSY in  $M_4$  vacuum (i.e.  $\partial_\Sigma W = W = 0$ ) we have from (43)  $V_{\Sigma\Sigma} = 0$  and both components of the  $\Sigma_c$  have the same mass:

$$m_{\Sigma_c}^2 = \frac{1}{M_{\text{Pl}}^2 \mathcal{K}_{\Sigma\Sigma^\dagger}} V_{\Sigma\Sigma^\dagger} = M_{\text{Pl}}^2 e^{\mathcal{K}} |\mathcal{K}^{\Sigma\Sigma^\dagger} \partial_\Sigma^2 W|^2. \quad (51)$$

Of course, for this case the fermionic partner of  $\Sigma^c$ , the modulino, has the same mass and the gravitino mass  $m_{3/2}$  vanishes. For the example presented in sect. 4 with superpotential given in eq. (44) we have

$$m_{\Sigma_c} = \frac{g}{3} \left( 2b \ln \frac{ab}{g} \right)^{1/2} M_{\text{Pl}}. \quad (52)$$

Clearly, by an appropriate selection of  $g$  (not fixed yet) it is possible to get a phenomenologically acceptable mass value.

In case of broken SUSY the gravitino obtains mass. With the  $F$ -term SUSY breaking (which we consider in this paper) in  $M_4$  vacuum we have

$$m_{3/2} = M_{\text{Pl}} |W| e^{\mathcal{K}/2}. \quad (53)$$

The masses of the radion's scalar components should be evaluated by eq. (50). The fermionic component of the superfield  $\Sigma$  is a goldstino absorbed into the 1/2-spin mode of the gravitino. (The  $F_\phi$  component of a compensator is zero and the only source here for the SUSY breaking is a non zero  $F^\Sigma$  term).

For the example presented in sect. 3.1, the value of the constant superpotential  $W$  is not fixed. Thus, this scenario allows to select  $W$  in such a way as to keep  $m_{3/2}$  in the TeV range. This is one advantage of this model over scenarios where uplifting and SUSY breaking occur by a (fixed)  $D$ -term. For the choice of parameters corresponding to Fig. 1 for the gravitino and radion's scalar component masses we have respectively

$$m_{3/2} \simeq 5 \cdot 10^{-2} M_{\text{Pl}} |W|, \quad m_+ \simeq 10^{-2} M_{\text{Pl}} |W|, \quad m_- \simeq 8 \cdot 10^{-3} M_{\text{Pl}} |W|. \quad (54)$$

With  $|W| \sim 10^{-15}$  all these states lie near the TeV scale. As was already pointed out in sect. 3.1, for this model we naturally expect  $m_{\pm} \sim m_{3/2}$ , unless due to a cancellation  $m_-$  is much smaller than  $m_{3/2}$ . This can happen with  $V_{\Sigma\Sigma^\dagger} \simeq V_{\Sigma\Sigma}$  (see eq. (50)) realized when the potential weakly depends on the axionic component of the radion, i.e. with strongly suppressed  $\bar{\beta}$ . However, in this case apart from problem of axionic component's stabilization it is difficult (if not impossible) to stabilize also the real component in the  $M_4$  vacuum. Thus, we conclude that with low energy SUSY breaking radion scalar components masses are expected to be in the 100 GeV-1 TeV energy range. This opens up a window for a collider physics with an interesting phenomenological signature [23].

## 6 Summary

In this paper we have presented several new examples of radion stabilization and found possibilities of uplifting to Minkowski vacuum which is a highly non trivial task. We have applied an effective 4D superfield description which makes investigations rather simple and provides an excellent playground for studying various ideas and for phenomenological model building. The formalism also allows us to test various mechanisms of moduli stabilization. This can for example be useful for a more detailed study of inflation together with radion stabilization, but also for a more general moduli dynamics in the early Universe. These issues will be discussed elsewhere.

### Acknowledgments

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## A One-loop corrections to the Kähler potential from bulk hypermultiplets

### A.1 Case with even coupling

First we derive a 1-loop correction to the Kähler potential due to hypermultiplets coupled with the vector modulus through even coupling [eq. (14)]. There are several ways of deriving eq.(15).

#### Calculation from prepotential

First we will take the following approach: we calculate first the prepotential  $\mathcal{F}(\Sigma)$  of  $\mathcal{N} = 2$  supergravity compactified on the  $S^1$  and then use the well-known expression (see e.g. [27])

$$\mathcal{K} = -\ln \left( -2(\mathcal{F} + \mathcal{F}^+) + (\mathcal{F}_I + \mathcal{F}_I^+)(\Sigma^I + \Sigma^{I+}) \right), \quad (\text{A.1})$$

to obtain the Kähler potential  $\mathcal{K}(\Sigma, \Sigma^+)$  from  $\mathcal{F}(\Sigma)$ . The reason one can use the  $S^1$  result also in the orbifold case is due to the fact that the relevant KK spectrum is essentially the same in both cases.

Recall that the prepotential  $\mathcal{F}(\Sigma)$  determines the holomorphic gauge couplings through its second derivative:

$$-\frac{1}{4} \int d^2\theta \mathcal{F}''(\Sigma) \mathcal{W}\mathcal{W} + \text{h.c.} \quad (\text{A.2})$$

Knowing the  $\Sigma$  dependent 1-loop correction to the gauge coupling we will be able to obtain  $\mathcal{F}$  and therefore to determine the Kähler potential. It is not difficult to find out that

$$\mathcal{F}''(\Sigma) = B \ln 2 \sinh(\pi R g_\beta \Sigma) = B \left[ \pi R g_\beta \Sigma - \sum_{k=1}^{\infty} \frac{e^{-k2\pi R g_\beta \Sigma}}{k} \right], \quad (\text{A.3})$$

with  $B = 2\pi^2 g_\beta^2 / R$ . A double integration gives

$$\mathcal{F} = B \left[ \frac{\pi R g_\beta}{6} \Sigma^3 - \sum_{k=1}^{\infty} \frac{e^{-k2\pi R g_\beta \Sigma}}{(2\pi R g)^2 k^3} \right]. \quad (\text{A.4})$$

The first term (on the r.h.s.) amounts to a renormalization of the Chern-Simons term  $\sim M^3$  in the norm-function. The second term in  $\mathcal{F}$  is the correction ( $\Delta\mathcal{F}$ ) we are searching for. It gives

$$\Delta_\beta = -2(\Delta\mathcal{F} + \Delta\mathcal{F}^+) + (\Delta\mathcal{F}' + \Delta\mathcal{F}'^+)(\Sigma + \Sigma^+). \quad (\text{A.5})$$

Using in (A.5) the second term of (A.4) we arrive at the expression given in (15). Next we present a derivation by a different method.

### Calculation in terms of $N = 1$ superfields

Since we are able to write our 5D action in terms of 4D  $N = 1$  superfields, we can take an advantage and use some results existing in 4D constructions. If the Lagrangian couplings are written in terms of 4D superfields as

$$\int d^4\theta K(\hat{\phi}, \hat{\phi}) + \int d^2\theta W_0(\phi) + \text{h.c.} + \dots \quad (\text{A.6})$$

the 1-loop correction to  $K$  (this should not be confused with  $\mathcal{K}$  used for the Kähler potential in the effective 4D superconformal description) is given by<sup>5</sup> [29]

$$\Delta K = \frac{1}{(2\pi)^4} \int \frac{d^4p}{p^2} \left[ \text{Tr} \log \hat{K} + \frac{1}{2} \text{Tr} \log \left( p^2 + \hat{K}^{-1} \hat{W}_0 \hat{K}^{-1T} \hat{W}_0 \right) - \text{Tr} \log \left( \hat{H} p^2 + \hat{S} \right) \right], \quad (\text{A.7})$$

where

$$\hat{K}_{\bar{I}J} = K_{\bar{I}J}(\hat{\phi}, \hat{\phi}), \quad (\hat{W}_0)_{IJ} = (W_0)_{IJ}(\hat{\phi}), \quad \text{etc.} \quad (\text{A.8})$$

The 5D Lagrangian of hypermultiplets coupled to the modulus by an even coupling is given by [3]

$$\mathcal{L}(H) = 2 \int d^4\theta \mathcal{W}_y (H^\dagger H + H^{c\dagger} H^c) + \int d^2\theta (2H^c \partial_y H - g\Sigma(H^2 - H^{c2})) + \text{h.c.} \quad (\text{A.9})$$

---

<sup>5</sup>One can also proceed with technics presented in [28] evaluating the log of a super determinant which is written as (A.15).



where we have set  $V = 0$ , and all objects are five dimensional. (Here and in the equations below instead of  $g_\beta$  we use  $g$  to make expressions more compact). In order to write this Lagrangian in the form of (A.6) here we perform the KK decomposition and integrate over the fifth dimension  $y$ . Thus we will be able to apply the expression (A.7) for each KK state, summing at the end over the full tower. For chiral superfields with positive and negative orbifold parities respectively we will have

$$H = \frac{1}{\sqrt{4\pi R}} H^{(0)} + \frac{1}{\sqrt{2\pi R}} \sum_{n=1}^{\infty} H^{(n)} \cos \frac{ny}{R}, \quad H^c = \frac{1}{\sqrt{2\pi R}} \sum_{n=1}^{\infty} H^{c(n)} \sin \frac{ny}{R} \quad (\text{A.10})$$

$$\begin{aligned} \int dy \mathcal{L}(H) &= \int d^4\theta \mathbb{W}_y \left( \sum_{n=0}^{\infty} H^{(n)\dagger} H^{(n)} + \sum_{n=1}^{\infty} H^{c(n)\dagger} H^{c(n)} \right) \\ &- \int d^2\theta \left( \frac{1}{R} \sum_{n=0}^{\infty} n H^{c(n)} H^{(n)} + \frac{g}{2} \Sigma \sum_{n=0}^{\infty} (H^{(n)})^2 - \frac{g}{2} \Sigma \sum_{n=1}^{\infty} (H^{c(n)})^2 \right) + \text{h.c.} \end{aligned} \quad (\text{A.11})$$

Comparing this with eqs (A.6), (A.8) we see that  $\hat{K}$  and  $\hat{W}_0$  are matrices in KK space:

$$\hat{K} = \begin{matrix} & H^{(0)} & H^{(n)} & H^{c(n)} \\ \begin{matrix} H^{(0)\dagger} \\ H^{(n)\dagger} \\ H^{c(n)\dagger} \end{matrix} & \begin{pmatrix} \mathbb{W}_y & 0 & 0 \\ 0 & \mathbb{W}_y & 0 \\ 0 & 0 & \mathbb{W}_y \end{pmatrix} \end{matrix}, \quad \hat{W} = \begin{matrix} & H^{(0)} & H^{(n)} & H^{c(n)} \\ \begin{matrix} H^{(0)} \\ H^{(n)} \\ H^{c(n)} \end{matrix} & \begin{pmatrix} -g\Sigma & 0 & 0 \\ 0 & -g\Sigma & \frac{n}{R} \\ 0 & \frac{n}{R} & g\Sigma \end{pmatrix} \end{matrix}. \quad (\text{A.12})$$

and

$$\mathcal{M}_H^2 \equiv \hat{K}^{-1} \hat{W}_0 \hat{K}^{-1T} \hat{W}_0 = \begin{pmatrix} |\omega|^2 & 0 & 0 \\ 0 & n^2 + |\omega|^2 & -n(\omega^\dagger - \omega) \\ 0 & -n(\omega - \omega^\dagger) & n^2 + |\omega|^2 \end{pmatrix} \frac{1}{R^2 \mathbb{W}_y^2}, \quad (\text{A.13})$$

where  $\omega = gR\Sigma$ . Eigenvalues of  $\mathcal{M}_H^2$  are

$$\frac{|\omega|^2}{R^2 \mathbb{W}_y^2}, \quad \frac{1}{R^2 \mathbb{W}_y^2} \left( \left( n \pm \frac{1}{2} |\omega^\dagger - \omega| \right)^2 + \frac{1}{4} (\omega + \omega^\dagger)^2 \right), \quad n = 1, 2, \dots \quad (\text{A.14})$$

Therefore,

$$\Delta K(\Sigma) = \frac{1}{32\pi^2} \int dp^2 \text{Tr} \log(p^2 + \mathcal{M}_H^2) = -\frac{1}{32\pi^2} \frac{1}{R^2 \mathbb{W}_y^2} \int \frac{dt}{t^2} \text{tr} e^{-R^2 \mathbb{W}_y^2 \mathcal{M}_H^2 t} \equiv \frac{1}{32\pi^2} \frac{1}{R^2 \mathbb{W}_y^2} \mathcal{I}. \quad (\text{A.15})$$

where we have dropped  $\Sigma$ -independent parts. We furthermore evaluate  $\mathcal{I}$ . Using Poisson resummation we have

$$\mathcal{I} = \int \frac{dt}{t^2} e^{-(\omega + \omega^\dagger)^2 t/4} \sum_{n=-\infty}^{+\infty} e^{(n + |\omega^\dagger - \omega|/2)^2 t} = \sqrt{\pi} \int \frac{dt}{t^{5/2}} e^{-(\omega + \omega^\dagger)^2 t/4} \sum_{k=-\infty}^{+\infty} e^{-\pi^2 k^2/t} e^{i\pi k |\omega^\dagger - \omega|}. \quad (\text{A.16})$$

The  $k = 0$  term gives

$$\mathcal{I}_{k=0} = \sqrt{\pi} \int \frac{dt}{t^{5/2}} e^{-(\omega + \omega^\dagger)^2 t/4} \sim |\omega + \omega^\dagger|^3 \Lambda^3 \sim (\Sigma + \Sigma^\dagger)^3 \Lambda^3, \quad (\text{A.17})$$

where  $\Lambda$  is the UV cut off ( $t_{UV} \sim 1/\Lambda^2$ ). This contribution renormalizes the tree level Kähler potential. For us contributions coming from  $k \neq 0$  terms are interesting. These contributions are finite for  $\Lambda \rightarrow \infty$ . Using the integration formula

$$\int_0^\infty \frac{dt}{t^{5/2}} e^{-at-b/t} = \frac{\sqrt{\pi}}{2b^{3/2}} (1 + 2\sqrt{ab}) e^{-2\sqrt{ab}}, \quad (\text{A.18})$$

and the abbreviation

$$2gR\hat{\Sigma} = |\omega + \omega^\dagger| + i(\omega^\dagger - \omega), \quad (\text{A.19})$$

from (A.15) and (A.16) we get finally

$$\Delta K(\Sigma) = -\frac{1}{64\pi^4} \frac{1}{R^2 \mathbb{W}_y^2} \sum_{k=1}^{\infty} \frac{1}{k^3} \left( e^{-2\pi k g R \hat{\Sigma}} + e^{-2\pi k g R \hat{\Sigma}^\dagger} \right) \left( 1 + \pi k g R (\hat{\Sigma} + \hat{\Sigma}^\dagger) \right). \quad (\text{A.20})$$

Recalling that we are dealing only with the zero mode of  $\Sigma$  we can write  $\Delta K(\Sigma) = \frac{1}{2\pi R} \int dy \Delta K(\Sigma)$  and we precisely get the expression of eq. (15) for  $\Delta_\beta$  (with taking into account that the 4D  $\Delta$  is related to 5D one as  $\Delta = \kappa_5^2 \Delta_5$ ).

## A.2 Case with odd coupling

In this case, the relevant hypermultiplet 5D Lagrangian couplings are

$$\mathcal{L}(H) = 2 \int d^4\theta \mathbb{W}_y (H^\dagger H + H^{c\dagger} H^c) + 2 \int d^2\theta H^c (\partial_y - g\Sigma\epsilon(y)) H + \text{h.c.} \quad (\text{A.21})$$

To obtain eq.(17) we will have to resort to a different, less elegant, technique. Note that the eigenvalues of the matrix  $\mathcal{M}_H^2 \equiv \hat{K}^{-1} \hat{\mathbb{W}}_0 \hat{K}^{-1T} \hat{\mathbb{W}}_0$  coincide with the KK mass spectrum. Therefore, instead of KK decomposition we will find the mass eigenvalues by solving the 5D eigenstate equations with appropriate boundary conditions for hypers' component chiral superfields.

In the  $D$ -term part of (A.21), we use the replacement  $d^2\theta^2 = -\frac{1}{4}\bar{D}^2$  and treat the  $\mathbb{W}_y$  and  $\Sigma$  as backgrounds<sup>6</sup>:  $D^2\mathbb{W}_y = \bar{D}^2\mathbb{W}_y = D^2\Sigma = \bar{D}^2\Sigma = 0$  etc. Then from (A.21) we can write equations of motion for superfields  $H, H^c$

$$\begin{aligned} -\frac{1}{4}\mathbb{W}_y \bar{D}^2 H^\dagger - (\partial_y + g\Sigma\epsilon(y)) H^c &= 0, \\ -\frac{1}{4}\mathbb{W}_y \bar{D}^2 H^{c\dagger} + (\partial_y - g\Sigma\epsilon(y)) H &= 0. \end{aligned} \quad (\text{A.22})$$

With change of variables

$$H = e^{g\Sigma|y|} \Phi, \quad H^c = e^{-g\Sigma|y|} \Phi^c, \quad (\text{A.23})$$

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<sup>6</sup>We assume that higher super-derivative terms have negligible effect on a result. This assumption must be justified and we will discuss this issue at the end of this appendix.

using the identity  $D^2\bar{D}^2 = \bar{D}^2D^2 + 8i\bar{D}\sigma^\mu D\partial_\mu + 16\Box_4$  and making some manipulations, from (A.22) we obtain

$$\begin{aligned}\mathbb{W}_y^2 m^2 \Phi + e^{-g(\Sigma+\Sigma^\dagger)|y|} \partial_y \left( e^{g(\Sigma+\Sigma^\dagger)|y|} \partial_y \Phi \right) &= 0, \\ \mathbb{W}_y^2 m_c^2 \Phi^c + e^{g(\Sigma+\Sigma^\dagger)|y|} \partial_y \left( e^{-g(\Sigma+\Sigma^\dagger)|y|} \partial_y \Phi^c \right) &= 0,\end{aligned}\tag{A.24}$$

where we have used  $\Box_4\Phi = m^2\Phi$  and  $\Box_4\Phi^c = m_c^2\Phi^c$ . The (A.24) should be solved by boundary conditions

$$\Phi(-y) = \Phi(y), \quad \Phi^c(-y) = -\Phi^c(y), \quad H(y+2\pi R) = H(y), \quad H^c(y+2\pi R) = H^c(y), \tag{A.25}$$

which also give discrete values of  $m^2$  and  $m_c^2$ . In the interval  $0 < y < \pi R$  we have the solution

$$\Phi = e^{-g(\Sigma+\Sigma^\dagger)y/2} (Ae^{\omega y} + Be^{-\omega y}), \quad \text{with} \quad \omega = \left( \frac{g^2}{4} (\Sigma + \Sigma^\dagger)^2 - \mathbb{W}_y^2 m^2 \right)^{1/2}, \tag{A.26}$$

and for  $-\pi R < y < 0$

$$\Phi = e^{g(\Sigma+\Sigma^\dagger)y/2} (Ae^{-\omega y} + Be^{\omega y}). \tag{A.27}$$

Using (A.24) and the boundary conditions we get

$$\omega R = in. \tag{A.28}$$

Thus, for the mass eigenvalues we get

$$m_n^2 = \frac{1}{R^2 \mathbb{W}_y^2} \left( n^2 + \frac{1}{4} g^2 R^2 (\Sigma + \Sigma^\dagger)^2 \right). \tag{A.29}$$

Similarly, we obtain

$$(m_c^2)_n = \frac{1}{R^2 \mathbb{W}_y^2} \left( n^2 + \frac{1}{4} g^2 R^2 (\Sigma + \Sigma^\dagger)^2 \right). \tag{A.30}$$

Therefore, we can write

$$\Delta K(\Sigma) = -\frac{1}{32\pi^2} \frac{1}{R^2 \mathbb{W}_y^2} \int \frac{dt}{t^2} \sum_{n=-\infty}^{\infty} e^{-R^2 \mathbb{W}_y^2 m_n^2 t} = -\frac{1}{32\pi^2} \frac{1}{R^2 \mathbb{W}_y^2} \int \frac{dt}{t^2} \sum_{n=-\infty}^{\infty} e^{-[n^2 + g^2 R^2 (\Sigma + \Sigma^\dagger)^2]/4t}. \tag{A.31}$$

Making similar steps as in the previous subsection, we finally get

$$\Delta K(\Sigma) = -\frac{1}{32\pi^4} \frac{1}{R^2 \mathbb{W}_y^2} \sum_{k=1}^{\infty} \frac{1}{k^3} (1 + \pi k g R |\Sigma + \Sigma^\dagger|) e^{-\pi k g R |\Sigma + \Sigma^\dagger|}, \tag{A.32}$$

which for  $\Delta_\gamma$  gives the expression given in (17).

In our analysis we have ignored the higher superderivative terms which are powers of  $D^2\Sigma$ ,  $D^2\mathbb{W}_y$ ,  $\bar{D}^2D^2\mathbb{W}_y$  etc. In general, the solution eqs. (A.26), (A.27) receive the following type of corrections

$$\sum a_{ij} (D^2\Sigma)^i (D^2\mathbb{W}_y)^j + \sum b_i (\bar{D}^2D^2\mathbb{W}_y)^i + \dots \tag{A.33}$$

because of modification of (A.22), (A.24). One can check out that these corrections can be safely ignored if the following conditions are satisfied

$$m^2, m_c^2 \gg \frac{(\mathbb{W}_y)_D}{\mathbb{W}_y}, \quad m^2, m_c^2 \gg g_5 2\pi R |\Sigma_5| m_\Sigma^2 = g_4 2\pi R M_{\text{Pl}} |\Sigma| m_\Sigma^2. \quad (\text{A.34})$$

For  $\mathbb{W}_y$  we have [5]

$$\mathbb{W}_y = \left( \frac{\kappa_5^{-2} (\tilde{\mathcal{N}} + \Delta)}{h^\dagger h} \right)^{1/3} = M_{\text{Pl}} \frac{e^{-\mathcal{K}/3}}{(\phi^\dagger \phi)^{1/2}}, \quad (\text{A.35})$$

and therefore

$$\frac{(\mathbb{W}_y)_D}{\mathbb{W}_y} = \left| \frac{1}{3} \mathcal{K}_I F^I + \frac{1}{2} \frac{F_\phi}{\phi} \right|^2. \quad (\text{A.36})$$

For the conditions (A.34) to be satisfied it is enough to take the zero mode masses  $m_0^2$  and  $(m_c^2)_0$ . Thus, taking into account (A.29), (A.30), and (A.36), eq. (A.34) can be rewritten as

$$\frac{g_4^2}{4\mathbb{W}_y^2} (\Sigma + \Sigma^\dagger)^2 M_{\text{Pl}}^2 \gg \left| \frac{1}{3} \mathcal{K}_I F^I + \frac{1}{2} \frac{F_\phi}{\phi} \right|^2, \quad g_4 2\pi R M_{\text{Pl}} |\Sigma| m_\Sigma^2. \quad (\text{A.37})$$

Now it is easy to see that these conditions restrict the amount of SUSY breaking and the value of the radion mass. For a concrete model one can easily check out whether these conditions are met or not.

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