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On Supersymmetry Breaking in Intersecting Brane Models*

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ABSTRACT

We discuss a framework to analyze the transmission of supersymmetry breaking in models of intersecting D-branes. Generically, different intersections preserve different fractions of an extended bulk supersymmetry, thus breaking supersymmetry completely but in a non-local way. We analyze this mechanism in a 5D toy model where two brane intersections, which are separated in the fifth dimension, break different halves of an extended $\mathcal{N} = 2$ supersymmetry. The sector of the theory on one brane intersection feels the breakdown of the residual $\mathcal{N} = 1$ supersymmetry only through two-loop interactions involving a coupling to fields from the other intersection. We compute the diagrams that contribute to scalar masses on one intersection and find that the masses are proportional to the compactification scale up to logarithmic corrections. We also compute the three-loop diagrams relevant to the Casimir energy between the two intersections and find a repulsive Casimir force.

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1 Introduction

Supersymmetry breaking in brane-world models has several very attractive features. The possibility of breaking supersymmetry on a distant brane offers a geometric realization of the idea of hidden sectors. Supersymmetry can be completely broken in a non-local way by partially breaking supersymmetry on different branes in such a way that each brane preserves a different fraction of the extended bulk supersymmetry, e.g., [1, 2, 3, 4]. This mechanism is naturally realized in many realistic D-brane models of string theory.

There are two approaches to embed the Standard Model of particle physics in string theory. The traditional way to achieve a phenomenologically viable model is to compactify the heterotic string on an orbifold or a Calabi-Yau manifold. In the strong coupling limit this amounts to compactifying M-theory on an orbifold or a G_2 manifold. Since the discovery of D-branes, it has become clear that type I and II string theory also provide a very promising framework for string model building. Gauge and matter fields are localized to the world-volume of the D-branes, giving rise to a nice geometrical picture of the Standard Model embedding.

Our aim is to understand the mechanism of supersymmetry breaking in D-brane models where different stacks of D-branes break different fractions of the extended bulk supersymmetry. It is crucial to realize that a D-brane breaks half of the bulk supersymmetry by its mere presence. The fields of the effective world-volume theory only fill multiplets of the smaller supersymmetry algebra. Supersymmetry can be completely broken either by adding anti-D-branes, which break the half of supersymmetry that is preserved by the D-branes, or by considering configurations of intersecting D-branes where different intersections preserve different fractions of the bulk supersymmetry. Such a scenario has been called pseudo-supersymmetry [5].

In this brief review, we are interested in the radiatively generated mass-splittings for the scalar matter fields of the following 5D toy model. Two 3-branes[1] are located at $x^5 = 0$ and $x^5 = \pi R$ in $M^{3,1} \times S^1$. There is an $\mathcal{N}_4 = 2$ vector corresponding to the gauge symmetry G in the bulk and $\mathcal{N}_4 = 1$ chiral multiplets charged under the gauge symmetry are confined to the 3-branes. The chiral multiplets from the two 3-branes couple to different bulk gauginos and thus preserve different halves of the bulk supersymmetry. Brane scalar masses are generated through Feynman diagrams involving a loop of fields from the distant brane. Such diagrams arise at the two-loop level. The explicit computation shows that the expected quadratic cutoff dependence is regulated by the finite brane separation. Only logarithmic divergences arise. They are due to wave-function renormalization of the brane fields. More precisely, we find that the brane scalar mass squared is positive and, for R much larger than the inverse cutoff scale $m_{\rm S}^{-1}$, proportional to $(2\pi R)^{-2} \ln(2\pi R\, m_{\rm S})$. Similarly, we find that the Casimir energy depends quartically on the inverse of the brane separation but only logarithmically on the cutoff scale. Thus supersymmetry breaking is soft in this class of models.

[1] In the intersecting D-brane model we have in mind, these are really $(3+1)$-dimensional intersections of higher dimensional branes. For simplicity, we will call these intersections '3-branes'. We also assume that five of the ten dimensions of type II string theory have been compactified on an orbifold of the size of the string scale, which is chosen such that it breaks the supersymmetry on the D-branes down to $\mathcal{N} = 2$.

These results are particularly interesting in view of the recent D-brane constructions which represent embeddings of the standard model in string theory [1, 4]. The knowledge of the precise expression for the mass splittings and their dependence on the interbrane distance is is an important first step towards a phenomenological analysis of those models. Although many of the explicit models have additional features, the toy model of this article captures their main supersymmetry breaking mechanism. The quantitative results of this article are directly applicable to D-brane models if the string scale is much larger than the compactification scale since in this limit all excited string states as well as the states corresponding to strings stretching between the distant D-branes are much heavier than the Kaluza-Klein excitations and can therefore be neglected.

2 Determination of the Lagrangian

We start by determining the effective Lagrangian for the pseudo-supersymmetric model in $D = 4$ [5]. It is of the form

$$\mathcal{L} = \mathcal{L}_{\rm bulk} + \mathcal{L}^{(1)} + \mathcal{L}^{(2)}. \tag{2.1}$$

The different sectors have manifest $\mathcal{N} = 2$, $\mathcal{N} = 1$, $\mathcal{N} = 1'$ supersymmetry, respectively, where $\mathcal{N} = 1$ denotes the first half and $\mathcal{N} = 1'$ denotes the second half of the bulk supersymmetry. Thus, supersymmetry is explicitly and completely broken, but the remaining supersymmetry on each of the two 3-branes (or $(3+1)$-dimensional brane intersections) still protects the brane scalar masses from radiative corrections up to 1-loop.

2.1 Non-linearly realized supersymmetry

Although the supersymmetry breaking is explicit in the effective theory, the second supersymmetry is non-linearly realized in $\mathcal{L}^{(1)}$, and the first supersymmetry is non-linearly realized in $\mathcal{L}^{(2)}$. The non-linear invariance arises because the Lagrangian contains appropriate couplings to two Goldstino superfields [6] (see also [7]).

To see how this works, first consider chiral superfields $\Phi_{\rm b}$ on the first brane. Their interactions can be rendered invariant under the second supersymmetry by adding appropriate couplings to a Goldstino superfield Λ_g, which transforms as (see, e.g., [8])

$$\delta^{(2)}\Lambda_g = \frac{1}{k}\xi^{(2)} - ik\,(\Lambda_g\sigma^m\bar{\xi}^{(2)} - \xi^{(2)}\sigma^m\bar{\Lambda}_g)\partial_m\Lambda_g. \tag{2.2}$$

Here, $k^{-1/2}$ is the supersymmetry breaking scale.

Now, define the shifted superfields

$$\begin{aligned}
\tilde{\Phi}_{\rm b}(x,\theta,\bar{\theta}) &= \Phi_{\rm b}(x^m - ik^2\,\Lambda_g\sigma^m\bar{\Lambda}_g,\theta,\bar{\theta}),\\
\tilde{\Lambda}_g(x,\theta,\bar{\theta}) &= \Lambda(x^m - ik^2\,\Lambda_g\sigma^m\bar{\Lambda}_g,\theta,\bar{\theta}).
\end{aligned}$$

and assign to $\Phi_{\rm b}$ the transformation law

$$\delta^{(2)}\tilde{\Phi}_{\rm b} = -2ik\,\tilde{\Lambda}_g\sigma^m\bar{\xi}^{(2)}\partial_m\tilde{\Phi}_{\rm b}. \tag{2.3}$$

Then,

$$S = \int d^4x \left[\int d^2\theta d^2\bar{\theta}\, \hat{E}\, \Phi_{\rm b}^\dagger \Phi_{\rm b} + \int d^2\theta\, E_L\, \mathcal{P}(\tilde{\Phi}_{\rm b}) + \int d^2\bar{\theta}\, E_R\, \mathcal{P}(\tilde{\Phi}_{\rm b}^\dagger) \right] \tag{2.4}$$

is invariant under both supersymmetries. The densities $\hat{E}$, E_L, E_R in this expression are given by

$$\begin{aligned} \hat{E} &= 1 + \frac{k^2}{8} \bar{D}^2 \bar{\Lambda}_g^2 + \frac{k^2}{8} D^2 \Lambda_g^2 + \mathcal{O}(k^4), \\ E_L &= 1 + \frac{k^2}{4} \bar{D}^2 \bar{\Lambda}_g^2 + \mathcal{O}(k^4), \qquad E_R \;=\; 1 + \frac{k^2}{4} D^2 \Lambda_g^2 + \mathcal{O}(k^4). \end{aligned}$$

2.2 Coupling $\mathcal{N} = 2$ to $\mathcal{N} = 1$

Next, we would like to find out how to consistently couple bulk $\mathcal{N} = 2$ multiplets to brane $\mathcal{N} = 1$ multiplets. Let us decompose the $\mathcal{N} = 2$ multiplets into their $\mathcal{N} = 1$ sub-multiplets, which couple naturally to the $\mathcal{N} = 1$ fields on the brane.[2] Then, we add appropriate couplings to the Goldstino superfield Λ_g to non-linearly realize the second supersymmetry.

Example: Consider an $\mathcal{N} = 2$ vector $\mathcal{V} = (A_m, \lambda^{(1)}, \lambda^{(2)}, \phi)$ and decompose it into

$$\begin{cases} \mathcal{N} = 1 \text{ vector } V = (A_m, \lambda^{(1)}) \\ \mathcal{N} = 1 \text{ chiral } \Phi = (\phi, \lambda^{(2)}) \end{cases}$$

It is easy to see that

$$\hat{V} = V + ik\, \theta\sigma^m\bar{\theta} \left(\bar{\Lambda}_g \bar{\sigma}_m D\Phi + \Lambda_g \sigma_m \bar{D}\Phi^\dagger \right) + \mathcal{O}(k^2) \tag{2.5}$$

transforms as

$$\delta^{(2)}\hat{V} = -ik \left(\Lambda_g \sigma^m \bar{\xi}^{(2)} - \xi^{(2)} \sigma^m \bar{\Lambda}_g \right) \partial_m \hat{V}. \tag{2.6}$$

Thus,

$$S = \int d^4x\, d^2\theta d^2\bar{\theta}\, \hat{E}\, \Phi_{\rm b}^\dagger\, e^{2\hat{V}}\, \Phi_{\rm b} \tag{2.7}$$

is invariant under both supersymmetries.

2.3 Pseudo-Supersymmetry

The $\mathcal{N} = 2$ vector $\mathcal{V} = (A_m, \lambda^{(1)}, \lambda^{(2)}, \phi)$ can either be decomposed into two $\mathcal{N} = 1$ multiplets

$$\begin{aligned} V &= -\theta\sigma^m\bar{\theta} A_m + i\theta\theta\bar{\theta}\bar{\lambda}_1 - i\bar{\theta}\bar{\theta}\theta\lambda_1 + \frac{1}{2}\, \theta\theta\bar{\theta}\bar{\theta} D, \\ \Phi &= \phi + \sqrt{2}\, \theta\lambda_2 + \theta\theta F \end{aligned} \tag{2.8}$$

[2]Again, we use the term 'brane' for simplicity but really mean a (3 + 1)-dimensional intersections of higher-dimensional branes.

or into two $\mathcal{N}=1'$ multiplets

$$
\begin{aligned}
V' &= -\tilde\theta\sigma^m\bar{\tilde\theta}A_m + i\tilde\theta\tilde\theta\bar{\tilde\theta}\bar\lambda_2 - i\bar{\tilde\theta}\bar{\tilde\theta}\tilde\theta\lambda_2 - \frac{1}{2}\tilde\theta\tilde\theta\bar{\tilde\theta}\bar{\tilde\theta}D, \\
\Phi' &= \phi - \sqrt{2}\,\tilde\theta\lambda_1 + \tilde\theta\tilde\theta F^\dagger. \qquad (2.9)
\end{aligned}
$$

V couples to charged chiral fields $\Phi^{(1)}$ on the first brane, whereas V' couples to charged chiral fields $\Phi^{(2)}$ on the second brane. Note that $\Phi^{(1)}$ and $\Phi^{(2)}$ are chiral matter superfields whereas Φ and Φ' are two different ways of arranging the components of the $\mathcal{N}=1$ chiral superfield inside the $\mathcal{N}=2$ bulk vector. $\Phi^{(1)}$ and Φ transform irreducibly under the first supersymmetry whereas $\Phi^{(2)}$ and Φ' transform irreducibly under the second supersymmetry.

The effective Lagrangian is

$$
\begin{aligned}
\mathcal{L}^{(1)} &= \int d^2\theta d^2\bar\theta\,\hat{E}^{(1)}\,\Phi^{(1)\dagger}e^{\hat V}\Phi^{(1)} + \int d^2\theta\,E_L^{(1)}\,\tilde\Lambda_g\tilde\Lambda_g \;+\; h.c. \\
\mathcal{L}_{\rm bulk} &= \int d^2\theta d^2\bar\theta\,\Phi^\dagger e^V\Phi \;+\; \frac{1}{4}\int d^2\theta\,WW + h.c. \\
\mathcal{L}^{(2)} &= \int d^2\tilde\theta d^2\bar{\tilde\theta}\,\hat{E}^{(2)}\,\Phi^{(2)\dagger}e^{\hat V'}\Phi^{(2)} + \int d^2\tilde\theta\,E_L^{(2)}\,\tilde\Lambda'_g\tilde\Lambda'_g \;+\; h.c.
\end{aligned}
$$

2.4 Generalization to D=5

Let us generalize the results of the previous subsections to $D=5$ [9]. Consider two 3-branes located at $x^5=0$ and $x^5=\pi R$ in $M^{3,1}\times S^1$, with an $\mathcal{N}_5=1$ vector multiplet in the bulk, $\mathcal{N}_4=1$ chiral multiplets on the first brane and $\mathcal{N}_4=1'$ chiral multiplets on the second brane.

The components of the $\mathcal{N}_5=1$, $D=5$ vector multiplet

$$
\begin{aligned}
&A_M, \quad \lambda^{(5)i}, \quad \phi^{(5)}, \quad X^a, \\
&M=0,\ldots,3,5,\;\; i=1,2,\;\; a=1,2,3,
\end{aligned}
$$

can be rearranged to fit into an $\mathcal{N}_4=2$, $D=4$ vector multiplet (see, e.g., [10])

$$
A_m, \quad \lambda_i, \quad \phi, \quad D, \quad F, \qquad\quad i=1,2.
$$

The precise mapping is

$$
\lambda^{(5)1} = \begin{pmatrix} \lambda_1 \\ -\bar\lambda_2 \end{pmatrix}, \quad \phi = \frac{1}{\sqrt{2}}\left(A_5 + i\,\phi^{(5)}\right), \quad D = X^3, \quad F = \frac{i}{\sqrt{2}}\left(X^1 + i\,X^2\right). \qquad (2.10)
$$

These components can either be grouped into $\mathcal{N}_4=1$ multiplets

$$
V = (A_m, \lambda_1, D), \quad \Phi = (\phi, \lambda_2, F)
$$

or into $\mathcal{N}_4 = 1'$ multiplets

$$V = (A_m, \lambda_1, -D), \quad \Phi = (\phi, -\lambda_2, F^\dagger).$$

The 5D Super-Yang-Mills Lagrangian can be written in terms of 4D superfields [11, 12]

$$\mathcal{L}_{\text{bulk}} = \frac{1}{2\, g_{(5)}^2} \,\text{tr} \left[\int d^2\theta \, W^\alpha W_\alpha + \int d^2\bar{\theta} \, \bar{W}_{\dot{\alpha}} \bar{W}^{\dot{\alpha}} + \int d^2\theta d^2\bar{\theta} \left(e^{-2V} \nabla_5 e^{2V} \right)^2 \right], \tag{2.11}$$

where

$$\nabla_5 e^{2V} = \partial_5 e^{2V} + i\sqrt{2} \left(\Phi e^{2V} - e^{2V} \Phi^\dagger \right). \tag{2.12}$$

The bulk-brane couplings are described by the Lagrangians we found in the 4D case. The complete 5D pseudo-supersymmetry Lagrangian is given by

$$\begin{aligned} \mathcal{L} &= \mathcal{L}_{\text{bulk}} + \mathcal{L}_{(1)} + \mathcal{L}_{(2)} \\ &= \frac{1}{2\, g_{(5)}^2} \,\text{tr} \left[\int d^2\theta \, W^\alpha W_\alpha + \int d^2\bar{\theta} \, \bar{W}_{\dot{\alpha}} \bar{W}^{\dot{\alpha}} + \int d^2\theta d^2\bar{\theta} \left(e^{-2V} \nabla_5 e^{2V} \right)^2 \right] \\ &\quad + \delta(x^5) \int d^2\theta d^2\bar{\theta} \, \hat{E} \, \Phi^{(1)\dagger} e^{2\hat{V}} \Phi^{(1)} \; + \; \delta(x^5 - l) \int d^2\tilde{\theta} d^2\bar{\tilde{\theta}} \, \hat{E}' \, \Phi^{(2)\dagger} e^{2\hat{V}'} \Phi^{(2)}. \end{aligned} \tag{2.13}$$

3 Computation of the soft breaking terms

Supersymmetry is explicitly broken in the effective Lagrangian but mass splittings for the supersymmetric multiplets only arise through loop corrections.

3.1 One-loop corrections to $\mathcal{L}_{\text{bulk}}$

One-loop corrections induce kinetic terms for the bulk vector and gauginos localized on the 3-branes. One finds

$$\begin{aligned} \mathcal{L}_{\text{bulk}}^{(\text{1-loop})} &= \frac{M_{\text{c}}}{2\, g_0^2} \,\text{tr} \left[\int d^2\theta \, W^\alpha W_\alpha + \int d^2\bar{\theta} \, \bar{W}_{\dot{\alpha}} \bar{W}^{\dot{\alpha}} + \int d^2\theta d^2\bar{\theta} \left(e^{-2V} \nabla_5 e^{2V} \right)^2 \right] \\ &\quad + \frac{\delta(x^5)}{2\, g_1^2} \,\text{tr} \left[\int d^2\theta \, W^\alpha W_\alpha + \int d^2\bar{\theta} \, \bar{W}_{\dot{\alpha}} \bar{W}^{\dot{\alpha}} \right] \\ &\quad + \frac{\delta(x^5 - l)}{2\, g_2^2} \,\text{tr} \left[\int d^2\tilde{\theta} \, W'^\alpha W'_\alpha + \int d^2\bar{\tilde{\theta}} \, \bar{W}'_{\dot{\alpha}} \bar{W}'^{\dot{\alpha}} \right], \end{aligned} \tag{3.1}$$

where

$$\frac{M_{\text{c}}}{g_0^2} = \frac{1}{g_{(5)}^2} + \Delta_{\text{bulk}}, \quad \frac{1}{g_i^2} = \frac{b_i}{8\pi^2} \ln \left(\frac{m}{\Lambda} \right). \tag{3.2}$$

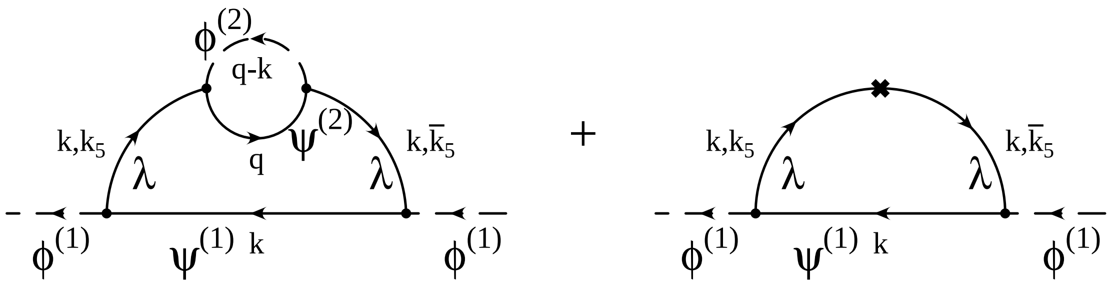

Figure 1: The leading order Feynman diagram giving rise to brane scalar masses. This diagram does not really exist in pseudo-supersymmetry but rather represents a shortcut to compute the sum of the nine two-loop diagrams that do exist. The counterterm is fixed by the condition that the one-loop corrected bulk gauge coupling constant receives no contributions from brane fields at the UV cutoff scale $m_{\rm S}$.

The logarithmic divergences localized on the 3-branes are eliminated through standard four-dimensional renormalization. Requiring the brane-localized contributions to the bulk gauge kinetic terms to vanish at the scale of supersymmetry breaking, $m_{\rm S}$, yields

$$\frac{1}{g_i^2(\mu)} = \frac{b_i}{8\pi^2}\ln\left(\frac{\mu}{m_{\rm S}}\right). \tag{3.3}$$

3.2 Two-loop corrections to $\mathcal{L}_{(1)}$

Only two-loop (or higher loop) diagrams involving fields from the second brane contribute to masses of scalars on the first brane. These diagrams do not contain the Goldstino.

The only counterterms needed to obtain a finite result are the ones introduced to renormalize the bulk gauge coupling. As explained in [9], the sum of all relevant diagrams is equal to one fictitious diagram (plus its counterterm), which is shown in figure 1.

The explicit computation yields

$$\begin{aligned} -i\,m^2 &= (g_{(4)}\sqrt{2})^4\, C_2(r) d^2(r') \int_{k,q} \frac{i\ \mathrm{tr}\,(ik^m\bar\sigma_m ik^n\sigma_n iq^p\bar\sigma_p ik^q\sigma_q)}{(q-k)^2 k^2 (k^2-(k_5)^2) q^2 (k^2-(\bar k_5)^2)} \\ &\quad + (\text{counterterm contribution}), \end{aligned} \tag{3.4}$$

where

$$\int_{k,q} = \sum_{k_5=\frac{n}{R}} 2(-1)^n \sum_{k_5=\frac{\bar n}{R}} 2(-1)^{\bar n} \int \frac{d^4k}{(2\pi)^4}\int\frac{d^4q}{(2\pi)^4}.$$

The result is

$$m^2 = \left(\frac{g_{(4)}^2}{4\pi}\right)^2 C_2(r) d^2(r') \frac{24\,\zeta(3)}{\pi^2\,(2\pi R)^2}\left(\ln(2\pi R m_{\rm S}) - \frac{11}{6} - \frac{\zeta'(3)}{\zeta(3)} + \gamma\right). \tag{3.5}$$

For $2\pi R m_{\rm S} \gg 1$, where our approximation is valid, this leads to a positive mass squared for the tree-level massless brane scalars.

3.3 Casimir energy

The computation of the Casimir energy is very similar to the computation of the mass squared of the brane scalars. It is easy to convince oneself that Feynman diagrams without external legs involving fields from both branes are only possible at three-loop. For all diagrams that contain only fields from one brane and/or the bulk, the supersymmetric non-renormalization theorems still apply; as a consequence their contribution to the vacuum energy cancels. Again, one finds that the sum of all diagrams contributing to the vacuum energy can be obtained by computing a single fictitious diagram.

Thus, the vacuum energy is given by [9]

$$\begin{aligned} i\,E &= (g_{(4)}\sqrt{2})^4\, d^2(r) d^2(r')\, \dim(G) \int_{k,q,p} \frac{(i)^2\ \mathrm{tr}\,(ip^m \bar{\sigma}_m ik^n \sigma_n iq^p \bar{\sigma}_p ik^q \sigma_q)}{(q-k)^2 (p-k)^2 p^2 (k^2-(k_5)^2) q^2 (k^2-(\bar{k}_5)^2)} \\ &\quad + (\text{counterterm contributions}) \end{aligned} \tag{3.6}$$

The result is

$$E = \frac{4\, g_{(4)}^4}{\pi^2\, (4\pi)^4}\, d^2(r) d^2(r')\, \dim(G) \frac{\Gamma(6)\zeta(5)}{(2\pi R)^4} \left(\Big(\ln(2\pi R m_{\mathrm{S}}) + A \Big)^2 + B \right), \tag{3.7}$$

where

$$A \approx -1.679, \quad B \approx 0.203.$$

For large compactification radii, this yields a repulsive Casimir force.

4 Outlook on supergravity breaking

It would be interesting to generalize the above results to situations where the bulk theory consists of pure minimal $D = 5$ supergravity. Again the scalars inside the chiral multiplets on the branes receive radiative corrections to their masses at two-loop. The components of an $\mathcal{N}_5 = 1$, $D = 5$ gravity multiplet

$$h_{MN}, \quad \psi^i_M, \quad A_M, \quad \text{auxiliary fields}, \tag{4.1}$$

can be rearranged to fit into an $\mathcal{N}_4 = 2$ gravity multiplet

$$h_{mn}, \quad \psi^i_m, \quad A_m, \quad \text{auxiliary fields}, \tag{4.2}$$

and an $\mathcal{N}_4 = 2$ vector multiplet

$$B_m, \quad \lambda^i, \quad \phi, \quad \text{auxiliary fields}. \tag{4.3}$$

The $\mathcal{N}_4 = 2$ gravity multiplet can either be split into $\mathcal{N} = 1$ multiplets

$$\begin{aligned} V_m &= \ldots + \theta\sigma^n\bar{\theta}\, h_{mn} + \bar{\theta}\bar{\theta}\,\theta\,\psi^1_m + \ldots \\ \Psi_\alpha &= \ldots + (\sigma^m\bar{\theta})_\alpha A_m + \theta\sigma^m\bar{\theta}\,\psi^2_{m\,\alpha} + \ldots \end{aligned} \tag{4.4}$$

or into $\mathcal{N}=1'$ multiplets

$$\begin{aligned} V'_m &= \ldots + \theta\sigma^n\bar\theta\, h_{mn} + \bar\theta\bar\theta\,\theta\,\psi^2_m + \ldots \\ \Psi'_\alpha &= \ldots + (\sigma^m\bar\theta)_\alpha A_m + \theta\sigma^m\bar\theta\,\psi^1_{m\,\alpha} + \ldots \end{aligned} \tag{4.5}$$

The Lagrangian for linearized $\mathcal{N}_5 = 1$, $D = 5$ supergravity can be written in terms of $\mathcal{N}_4 = 1$ superfields [13]

$$\begin{aligned} \mathcal{L} &= \\ &\int d^4\theta\Big[\frac{1}{8}V^m D^\alpha \bar D^2 D_\alpha V_m + \frac{1}{48}\left([D^\alpha, \bar D^{\dot\alpha}]V_{\alpha\dot\alpha}\right)^2 - (\partial^m V_m)^2 - \frac{1}{3}\Sigma^\dagger\Sigma + \frac{2i}{3}(\Sigma - \Sigma^\dagger)\partial^m V_m\Big] \\ &+ \int d^4\theta\Big\{[T^\dagger(\Sigma - i\partial_{\alpha\dot\alpha}V^{\dot\alpha\alpha}) + h.c.] - \frac{1}{2}[D^\alpha\hat\Psi_\alpha + \bar D_{\dot\alpha}\hat\Psi^{\dagger\dot\alpha} - \partial_5 P]^2 \\ &\qquad + [\partial_5 V_{\alpha\dot\alpha} - (\bar D_{\dot\alpha}\hat\Psi_\alpha - D_\alpha\hat\Psi^\dagger_{\dot\alpha})]^2\Big\}, \end{aligned}$$

Here, Σ is the conformal compensator multiplet for the supergravity multiplet, P is a prepotential for Σ, defined by $\Sigma = -\frac{1}{4}\bar D^2 P$, and T is the $\mathcal{N}_4 = 1$ chiral multiplet inside the $\mathcal{N}_4 = 2$ vector multiplet (4.3). Similarly, one can write down the gravitational coupling of either (4.4) or (4.5) to chiral matter multiplets on the branes. The explicit computation of the corresponding two-loop diagrams contributing to scalar masses, however, is considerably more involved and will be left for future work.

5 Conclusions

Realistic D-brane models provide an interesting supersymmetry breaking mechanism with the following features:

- supersymmetry is explicitly broken but non-linearly realized in the effective field theory
- bulk gauge fields acquire brane-localized kinetic terms at one-loop
- mass splittings for the brane supermultiplets only arise at two-loop $m^2 \sim g^4_{(4)}(2\pi R)^{-2}\ln(2\pi R m_{\rm S})$
- repulsive Casimir force at three-loop
- gravitino masses only at three-loop

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