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Twisted moduli and supersymmetry breaking

Stephen F. King and David A.J. Rayner

Department of Physics and Astronomy, University of Southampton
Southampton, SO17 1BJ, U.K.
E-mail: king@southampton.ac.uk, drayner@hep.phys.soton.ac.uk

Abstract: We consider twisted moduli contributions to supersymmetry breaking in effective type I string constructions involving intersecting $D5_i$ and D9-branes using Goldstino angles to parametrise the supersymmetry breaking. It is well known that twisted moduli enter at tree-level into the gauge kinetic functions, and can provide new sources of gaugino mass if they develop F-term vacuum expectation values. It is generally assumed that string states which are sequestered from the twisted moduli receive a zero soft mass in the twisted modulus domination limit, however the standard form of Kähler potential does not reproduce this expectation. We therefore propose a new form of the Kähler potential which is consistent at leading order with the sequestered form proposed by Randall and Sundrum, and show that it leads to exponentially suppressed sequestered soft masses. Including the effects of Green-Schwarz mixing, we write down the soft scalar masses and trilinears arising from a type I string construction involving intersecting $D5_i$ and D9-branes in the presence of untwisted and twisted moduli. If the squarks and sleptons are identified with sequestered states then in the twisted moduli dominated limit this corresponds to gaugino mediated supersymmetry breaking, and we discuss two different scenarios for this. The general results will be useful for phenomenological studies involving a combination of gravity and gaugino mediated SUSY breaking due to the dilaton, untwisted and twisted moduli contributions, and enable the soft masses to be studied as a function of the different compactification radii.

Keywords: D-branes, Supersymmetry Breaking, Supergravity Models, Supersymmetric Standard Model.

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Contents

1. Introduction

Superstring theories offer the only consistent method for unifying the four fundamental forces of Nature within a single framework. Following the discovery of string dualities [1] and Dirichlet-branes [2], heterotic strings no longer provide the only theory in which to embed the (minimal supersymmetric) Standard Model (MSSM) and type I (and II) models have all been considered. The heterotic and type I models share common features, but differ in phenomenologically important ways. For instance, both scenarios contain the dilaton and (twisted and untwisted) moduli fields that are related to the geometry of the compactified space and appear in the low-energy four-dimensional effective SUGRA theory [3]. However, in type I models the fundamental string scale M_* is no longer fixed at the grand unification M_X or Planck scales M_P, and in principle can be as low as the TeV scale with the lower bound determined by phenomenology. An important difference between heterotic and type

I is the rôle played by twisted moduli fields — closed string states that are *trapped* at fixed points in the underlying manifold due to the action of orbifold compactification. Consider the gauge kinetic function f_α that appears in the SUGRA lagrangian [3]. In weakly coupled heterotic string theory, the string coupling constant is uniquely determined by the dilaton $g_s^{-2} \sim \mathrm{Re}(S)$, and:

$$f_\alpha = k_\alpha S + \Delta_{\text{1-loop}}(T_i) \tag{1.1}$$

where k_α is the Kac-Moody level of the gauge factor[1] and $\Delta_{\text{1-loop}}(T_i)$ arises from 1-loop string threshold corrections [4].[2] In contrast, type I models have gauge kinetic functions that depend on the dilaton S for 9-branes *and* the moduli fields T_i for 5_i-branes, giving rise to different gauge couplings on different branes. In addition the gauge kinetic functions have a tree-level dependence on the twisted moduli, and this gives rise to different gauge couplings even within a particular D-brane sector. The tree-level dependence on the twisted moduli fields Y^k from the k^{th} twisted sector (within the world-volume of a given D-brane sector) are given by:

$$f_\alpha^9 = S + \sum_k \frac{s_{\alpha,k}}{4\pi} \sum_q Y^{k,q} \tag{1.2}$$

$$f_\beta^{5_i} = T_i + \sum_k \frac{s'_{\beta,k}}{4\pi} \sum_{p_i} Y^{k,p_i} \tag{1.3}$$

where the gauge coupling is found by the relation:

$$\mathrm{Re}\, f_\alpha = \frac{4\pi}{g_\alpha^2} \tag{1.4}$$

and $q(p_i)$ label the fixed points within the $9(5_i)$-brane and $s_{\alpha,k}, s'_{\beta,k}$ are calculable model-dependent coefficients. Thus twisted moduli tend to induce non-universal gauge couplings even for gauge groups living on a common brane sector. The twisted moduli also play an important rôle in the cancellation of gauge and gravitational anomalies in type I models — like the dilaton in heterotic string theory — through a generalised four-dimensional Green-Schwarz mechanism [6] that mixes twisted and untwisted moduli together.

The Kähler potential and superpotential can also receive non-perturbative contributions, and in the absence of a complete model one may adopt a phenomenologically-motivated parametrisation in order to make progress [7, 8]. The relative contributions to the overall SUSY breaking F-term vacuum expectation value (vev) from different fields can be parametrised in terms of *Goldstino* angles. In such an approach one can derive the soft parameters in terms of the Goldstino angles, and examine various limits in which the dilaton or moduli fields dominate. As envisaged by the originators of the approach, it may also be used to investigate the contributions to SUSY breaking from twisted moduli in effective type I theories, in addition to the usual dilaton and untwisted moduli fields [9]. However the analyses that have been done so far have only considered the explicit situation

[1]In the MSSM, $k_{\mathrm{SU}(3)_C} = k_{\mathrm{SU}(2)_L} = \frac{3}{5} k_{\mathrm{U}(1)_Y} = 1$.

[2]In contrast, there is a very different situation for the strongly coupled case (from M-theory) where f_α receives comparable contributions at tree-level from the dilaton and untwisted moduli fields $f \sim S + T$ [5].

where the gauge group and matter fields arise from a stack of D9-branes, and thus share the same world-volume as all of the twisted moduli fields. It is one of the purposes of this paper to extend the scope of such analyses to include more general set-ups involving intersecting $D5_i$ and D9-branes. In so doing we encounter a difficulty that is not present in the case of a single D9-brane set-up, namely the problem of sequestered states which do not share the same world volume as the twisted moduli, and we show how this problem may be successfully resolved.

In this paper, then, we shall consider twisted moduli contributions to SUSY breaking in effective type I string constructions based on a general set-up involving intersecting $D5_i$ and D9-branes, using Goldstino angles to parametrise the SUSY breaking. It is well known that the F-term vevs of the twisted moduli fields provide a new source of gaugino masses [10]. It is also generally assumed that states that do not live in the same world-volume should receive zero soft mass contributions in the twisted moduli dominated limit, which offers a possible string realisation of gaugino mediated SUSY breaking ($\tilde{g}MSB$) [10, 11]. However we show that the standard form of Kähler potential is not consistent with this physical requirement. We therefore propose a new form of the Kähler potential which is consistent at leading order with the sequestered form proposed by Randall and Sundrum [12], and which leads to exponentially suppressed sequestered soft masses, in agreement with physical expectations. Including the effects of Green-Schwarz mixing we then write down soft scalar and trilinear masses arising from a general string construction involving intersecting $D5_i$ and D9-branes in the presence of untwisted and twisted moduli. We show how the results may be applied to $\tilde{g}MSB$ and discuss two explicit scenarios for this. The general results will be useful for phenomenological studies involving a combination of gravity and gaugino mediated SUSY breaking due to the dilaton, untwisted and twisted moduli contributions, and enable the soft masses to be studied as a function of the finite compactification radii.

The layout of the remainder of the paper is as follows. In section 2 we discuss effective type I string theories in the presence of twisted moduli, point out the difficulty with the sequestered soft masses using the standard Kähler potential, and propose a new sequestered form of Kähler potential which solves the problem. In section 3 we generalise our results to include Green-Schwarz mixing, then in section 4 we write down the resulting soft scalar and trilinear masses that arise in general string constructions involving intersecting D5 and D9 branes, in the presence of twisted and untwisted moduli contributions to SUSY breaking. In section 5 we discuss gaugino mediated SUSY breaking as a simple example, and point out that our results enable gravity mediated corrections to gaugino mediation to be studied as a function of the compactification scale. For completeness we include Appendices on Supergravity basics.

2. Effective type I string theory and twisted moduli

2.1 Kähler potentials

In this section we will introduce a generic type I string construction involving intersecting $D5_i$-branes embedded within D9-branes, where coincident D-branes give rise to gauge groups localised within the world-volume of the corresponding D-brane. Chiral charged

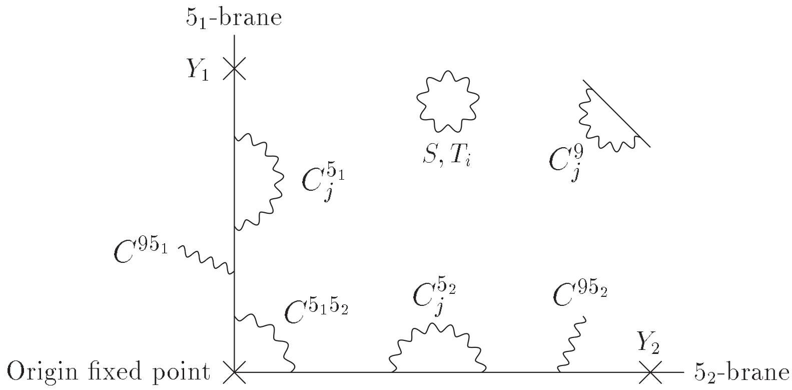

Figure 1: A generic type I string construction involving two sets of perpendicular D5-branes embedded within a D9-brane, where the D5-brane world-volumes intersect at the origin. Charged chiral fields appear as open strings with both ends attached to the same D-brane $C_j^{5_i}$ and C_j^9, or different branes $C^{5_1 5_2}$ and C^{95_i}. Closed strings (S, T_i) can live in the full 10d space, although orbifolding leads to closed strings (twisted moduli Y_k) localised at 4d fixed points within the D5_i-brane world-volume.

matter fields appear as open-strings with their ends attached to D-branes. Chan-Paton factors at the string ends carry the gauge quantum numbers under the attached gauge group. This type of construction will lead to two distinct types of matter field — $C_j^{5_i}$ and C_j^9 are open strings with both ends attached to the same D5(9)-brane, while $C^{5_i 5_j}$ and C^{95_i} have their ends attached to different D-branes and the string tension forces the inverse length of the strings to become of order the string scale M_*. The $C^{5_i 5_j}$ states become localised at the 4d intersection point between the two D5-branes, while the C^{95_i} states have one end attached anywhere along the 5_i-brane world-volume. The spectrum also contains closed strings that correspond to the gravity multiplet and dilaton (S) and moduli fields (T_i). Notice that this construction is entirely general and is T-dual to alternative scenarios involving D7- and D3-branes. A construction involving two sets of intersecting branes within a D9-brane is shown in figure 1, but our analysis can be extended for a full set of three perpendicular intersecting branes and the open/closed string states that result.

We can now exploit the string duality between 10d SO(32) heterotic theory and 10d type I theory to derive the 4d Kähler potential $K(S, T_i, C_a)$ for the dilaton, (untwisted) moduli and charged chiral fields that arise in the low energy supergravity description of the model with two sets of intersecting D5-branes embedded within a D9-brane as shown in figure 1. Ignoring the twisted moduli for the moment, the result is [8]:

$$\begin{aligned} K = & -\ln\left(S+\bar{S}-|C_1^{5_1}|^2-|C_2^{5_2}|^2\right)-\ln\left(T_1+\bar{T}_1-|C_1^9|^2-|C_3^{5_2}|^2\right)- \\ & -\ln\left(T_2+\bar{T}_2-|C_2^9|^2-|C_3^{5_1}|^2\right)-\ln\left(T_3+\bar{T}_3-|C_3^9|^2-|C_2^{5_1}|^2-|C_1^{5_2}|^2\right)+ \\ & +\frac{|C^{5_1 5_2}|^2}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}}+\frac{|C^{95_1}|^2}{(T_2+\bar{T}_2)^{1/2}(T_3+\bar{T}_3)^{1/2}}+\frac{|C^{95_2}|^2}{(T_1+\bar{T}_1)^{1/2}(T_3+\bar{T}_3)^{1/2}} \end{aligned} \tag{2.1}$$

The results can easily be extended to include a third 5_3-brane.

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Expanding the Kähler potential in the lowest order in the matter fields (i.e. $(S+\bar{S}) \gg |C_1^{5_1}|^2 + |C_2^{5_2}|^2$) yields:

$$
\begin{aligned}
K =& -\ln\left(S+\bar{S}\right) - \sum_{i=1}^{3}\ln\left(T_i+\bar{T}_i\right) + \sum_{i=1}^{2}\frac{|C_i^{5_i}|^2}{(S+\bar{S})} + \frac{|C_3^{5_2}|^2}{(T_1+\bar{T}_1)} + \frac{|C_1^9|^2}{(T_1+\bar{T}_1)} + \\
&+ \frac{|C_3^{5_1}|^2}{(T_2+\bar{T}_2)} + \frac{|C_2^9|^2}{(T_2+\bar{T}_2)} + \frac{|C_2^{5_1}|^2}{(T_3+\bar{T}_3)} + \frac{|C_1^{5_2}|^2}{(T_3+\bar{T}_3)} + \frac{|C_3^9|^2}{(T_3+\bar{T}_3)} + \\
&+ \frac{|C^{5_15_2}|^2}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}} + \frac{|C^{95_1}|^2}{(T_2+\bar{T}_2)^{1/2}(T_3+\bar{T}_3)^{1/2}} + \frac{|C^{95_2}|^2}{(T_1+\bar{T}_1)^{1/2}(T_3+\bar{T}_3)^{1/2}}
\end{aligned}
\tag{2.2}
$$

Using appendix A and eq. (A.2), we can identify the individual Kähler metrics (which are diagonal $\tilde{K}_a = \tilde{K}_{\bar{a}b}\delta_{\bar{a}b}$) for each type of charged chiral field:

$$
\begin{aligned}
\tilde{K}_{C_i^{5_i}} &= \frac{1}{(S+\bar{S})} \\
\tilde{K}_{C_j^{5_i}} &= \frac{1}{(T_k+\bar{T}_k)} \qquad (i \neq j \neq k \neq i) \\
\tilde{K}_{C_i^9} &= \frac{1}{(T_i+\bar{T}_i)} \\
\tilde{K}_{C^{5_15_2}} &= \frac{1}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}} \\
\tilde{K}_{C^{95_i}} &= \frac{1}{(T_j+\bar{T}_j)^{1/2}(T_k+\bar{T}_k)^{1/2}} \qquad (i \neq j \neq k \neq i)\,.
\end{aligned}
\tag{2.3}
$$

The twisted moduli $Y^{k,q}$ also contribute to the Kähler potential, but the precise form of the contribution is strongly model-dependent. For simplicity we shall consider a single twisted modulus within each of the three D5-brane sectors, which we denote by Y^k where $k = 1, 2, 3$ labels the D5_k branes. Each of the Y^k may be regarded as a linear combination of all the twisted moduli within that D5_k brane, so that the simplified gauge kinetic function is from eq. (1.3),

$$
f_\alpha^{5_k} = T_k + \frac{s_\alpha}{4\pi}Y_k\,. \tag{2.4}
$$

Anomaly cancellation via the Green-Schwarz mechanism suggests that this contribution mixes twisted and untwisted moduli together while preserving modular invariance. We will work in terms of a general even function $\hat{K}$ with an argument:

$$
(Y_k+\bar{Y}_k) - \delta_{GS}^{kj}\ln(T_j+\bar{T}_j)\,. \tag{2.5}
$$

For simplicity, we will initially drop the Green-Schwarz term ($\delta_{GS} = 0$) and assume a very simple form [13]

$$
\hat{K}(Y_k, \bar{Y}_k) = \frac{1}{2}(Y_k+\bar{Y}_k)^2\,. \tag{2.6}
$$

Hence the tree-level Kähler potential for the closed string states is:

$$
\bar{K}(h,\bar{h}) = -\ln\left(S+\bar{S}\right) - \sum_{i=1}^{3}\ln\left(T_i+\bar{T}_i\right) + \sum_{k=1}^{3}\hat{K}(Y_k,\bar{Y}_k)\,. \tag{2.7}
$$

We will repeat our analysis in the presence of a Green-Schwarz mixing term ($\delta_{GS} \neq 0$) in section 3.

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The perturbative superpotential can be expressed in terms of the states present in the model by considering the set of renormalisable interactions that arise from the splitting and joining of open strings.

$$\begin{aligned}
W_{\rm ren} = & \, g_{5_1}\left(C_1^{5_1}C_2^{5_1}C_3^{5_1} + C_3^{5_1}C^{5_15_2}C^{5_15_2} + C_1^{5_1}C^{95_1}C^{95_1}\right) + \\
& + g_{5_2}\left(C_1^{5_2}C_2^{5_2}C_3^{5_2} + C_3^{5_2}C^{5_15_2}C^{5_15_2} + C_2^{5_2}C^{95_2}C^{95_2}\right) + g_{5_3}C^{5_15_2}C^{95_1}C^{95_2} + \\
& + g_9\left(C_1^9C_2^9C_3^9 + C_1^9C^{95_1}C^{95_1} + C_2^9C^{95_2}C^{95_2}\right),
\end{aligned} \tag{2.8}$$

where the Yukawa coupling constants (associated with fields arising from each 5-and 9-brane) are given by:

$$g_{5_i}^2 = \frac{4\pi}{\mathrm{Re}\, T_i}, \qquad g_9^2 = \frac{4\pi}{\mathrm{Re}\, S}. \tag{2.9}$$

However, the superpotential can also receive (unknown) non-perturbative contributions, e.g. from gaugino condensation,[3] that require the F-terms to be parametrised in terms of Goldstino angles.

In this general setup, we are assuming that SUSY breaking originates from the closed string sector. In the absence of a Green-Schwarz anomaly cancelling term in eq. (2.6), the Kähler metric is diagonal at leading order since $\mathrm{Re}\, S, \mathrm{Re}\, T_i \gg |C_a|^2$. Using eqs. (2.2), (2.6) and (A.8) we can write down the SUSY breaking F-term vev in terms of two Goldstino angles (θ, ϕ), where $\theta(\phi)$ describes the relative contributions from the dilaton and moduli (twisted and untwisted moduli) F-terms respectively, and we are assuming a vanishing cosmological constant V_0.

$$F_S = \sqrt{3}m_{3/2}\sin\theta e^{i\alpha_S}(K_{\bar{S}S})^{-1/2} = \sqrt{3}m_{3/2}\sin\theta e^{i\alpha_S}(S+\bar{S}) \tag{2.10}$$

$$F_{T_i} = \sqrt{3}m_{3/2}\cos\theta\sin\phi\,\Theta_i\, e^{i\alpha_i}(K_{\bar{T}_iT_i})^{-1/2} = \sqrt{3}m_{3/2}\cos\theta\sin\phi\,\Theta_i\, e^{i\alpha_i}(T_i+\bar{T}_i) \tag{2.11}$$

$$F_{Y_k} = \sqrt{3}m_{3/2}\cos\theta\cos\phi\,\Phi_k\, e^{i\alpha_{Y_k}}(K_{\bar{Y}_kY_k})^{-1/2} = \sqrt{3}m_{3/2}\cos\theta\cos\phi\,\Phi_k\, e^{i\alpha_{Y_k}} \tag{2.12}$$

where $\sum_{i=1}^3 \Theta_i^2 = 1$ and $\sum_{k=1}^3 \Phi_k^2 = 1$.

One can study three limits of phenomenological interest where different sources of SUSY breaking dominate: dilaton (S) domination where $\sin\theta = 1$; untwisted moduli (T_i) domination where $\cos\theta = \sin\phi = 1$; and twisted moduli ($Y_k$) domination where $\cos\theta = \cos\phi = 1$. In the next sub-section we shall see that there is a problem with the sequestered masses in the twisted moduli dominated limit, and then we shall show how this problem may be resolved.

2.2 Problems with the standard Kähler potential

In order to illustrate the problem let us consider the case of a *single* linear combination of twisted moduli located inside the 5_2-brane, which we denote by Y_2, corresponding to the simplified gauge kinetic function $f_\alpha^{5_2} = T_2 + s_\alpha Y_2/4\pi$. Thus we take the SUSY breaking parameter $\Phi_2 = 1$ in eq. (2.12). We regard this linear combination Y_2 to be located in the world-volume of the 5_2-brane at a distance $\mathcal{O}(R_{5_2})$ from the intersection states $C^{5_15_2}$.

[3] See [14] for a recent discussion in the context of stabilising the dilaton potential in type I string theory.

Figure 1 shows that only $C_j^{5_2}, C_j^9$ and C^{95_2} states can couple directly to the Y_2 twisted moduli, while $C_j^{5_1}, C^{95_1}$ are confined on the 5_1-brane, and $C^{5_15_2}$ is confined to the origin fixed point. We refer to the states $C_j^{5_1}, C^{95_1}$ and $C^{5_15_2}$ which are spatially separated from Y_2 as being sequestered from it. Using eqs. (2.10)- (2.12), (A.17) with the standard Kähler metric for the intersection and 5_1-brane states of eq. (2.3), the sequestered state scalar masses are found to be (still ignoring the Green-Schwarz mixing term $\delta_{GS} = 0$):

$$\begin{aligned}
m^2_{C^{5_15_2}} &= m^2_{3/2}\left[1 - \frac{3}{2}\left(\sin^2\theta + \Theta_3^2\cos^2\theta\,\sin^2\phi\right)\right] \\
m^2_{C_1^{5_1}} &= m^2_{3/2}\left[1 - 3\sin^2\theta\right] \\
m^2_{C_2^{5_1}} &= m^2_{3/2}\left[1 - 3\Theta_3^2\cos^2\theta\,\sin^2\phi\right] \\
m^2_{C_3^{5_1}} &= m^2_{3/2}\left[1 - 3\Theta_2^2\cos^2\theta\,\sin^2\phi\right] \\
m^2_{C^{95_1}} &= m^2_{3/2}\left[1 - \frac{3}{2}\cos^2\theta\,\sin^2\phi\left(\Theta_2^2 + \Theta_3^2\right)\right].
\end{aligned} \tag{2.13}$$

In the twisted moduli dominated limit where the F-term F_{Y_2} is the only contribution to the SUSY breaking ($\cos\theta = \cos\phi = 1$) the intersection state masses from eq. (2.13) are:

$$m^2_{C^{5_15_2}} = m^2_{C_j^{5_1}} = m^2_{C^{95_1}} = m^2_{3/2} \qquad (j = 1, 2, 3) \tag{2.14}$$

The soft masses in eq. (2.14) are independent of the separation between the origin and the fixed point at which the twisted moduli live. This is not what we expect. Since these states are sequestered from the twisted moduli we would expect that their soft masses be exponentially suppressed by the spatial separation between the two fixed points, as the following argument explains.

In the twisted moduli dominated limit, the situation regarding the sequestered states is physically equivalent to the gaugino mediated SUSY breaking scenario [11] as shown in figure 2. In gaugino mediation, SUSY is broken on a 4d "hidden sector brane" which is spatially separated along one (or more) extra dimensions from another parallel 4d "matter brane" where matter fields are localised. Scalar masses on the matter brane are exponentially suppressed at tree-level by the distance between the branes but are radiatively generated at one-loop via gaugino mediation. In the string theory realisation, the rôle of the hidden sector brane is played by the twisted moduli localised at a non-trivial fixed point separated from the origin fixed point which corresponds to the matter brane.

Clearly in the limit of large spatial separation between the two fixed points $r \sim \mathcal{O}(R_{5_2})$ the sequestered soft masses should be exponentially suppressed:[4]

$$m^2_{C^{5_15_2}}\,,\, m^2_{C_j^{5_1}}\,,\, m^2_{C^{95_1}} \sim e^{-M_* r} m^2_{3/2} \tag{2.15}$$

where M_* is the ultraviolet cutoff for the effective theory, which we will associate with the string scale. Obviously, a sufficiently large separation will lead to a negligibly small mass

[4] Physically, this suppression is due to integrating out the heavier modes (with masses above the cutoff M_*) that propagate between sectors with a Yukawa-like propagator.

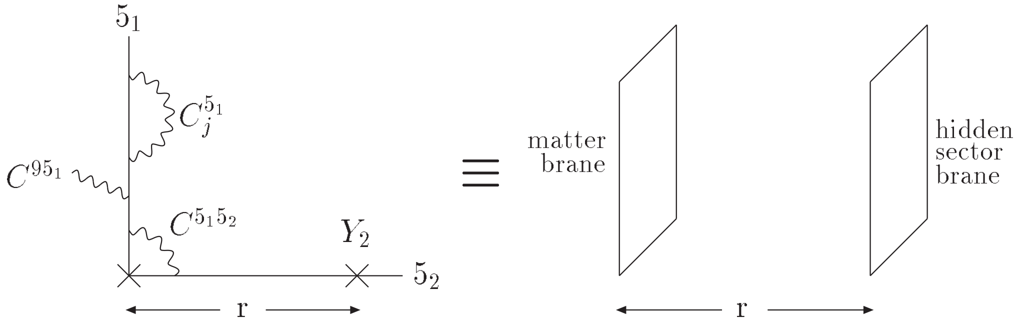

Figure 2: The intersecting D5-brane construction shares similar features with the gaugino mediated SUSY breaking model in the limit of a small compactification radius R_{5_1}. In this limit, the $C_j^{5_1}$ and C^{95_1} states are effectively localised at the origin, and these intersection states are equivalent to the matter brane while the localised twisted modulus is equivalent to the hidden sector brane where SUSY is broken. The spatial separation between the two fixed points (matter and hidden sector brane) is $r \sim \mathcal{O}(R_{5_2})$.

as in $\tilde{g}MSB$ which offers a solution to the flavour problems and suppression of flavour-changing neutral-currents. In the next sub-section we propose a new form of the Kähler potentials which give rise to the correct exponentially suppressed soft sequestered masses in eq. (2.15) rather than the result in eq. (2.14). In appendix B we consider an alternative exponential suppression factor $e^{-(M_* r)^2}$ that is attributed to non-perturbative world-sheet instanton corrections [15].

2.3 A new Kähler potential

We need to modify the intersection state Kähler potentials $\tilde{K}_{C^{5_15_2}}, \tilde{K}_{C_j^{5_1}}$ and $\tilde{K}_{C^{95_1}}$ to give the desired exponentially suppressed mass prediction in eq. (2.13) with an explicit dependence on the separation. Notice that in the limit of very small separation, we should be able to recover the previous (standard) form of eq. (2.3). To begin with we will only consider the Kähler potential for $C^{5_15_2}$ states and later generalise to the $C_j^{5_1}, C^{95_1}$ states as well.

Consider the scalar mass relation of eq. (A.17) for the intersection states $C^{5_15_2}$ in the limit of twisted moduli domination and let us determine the Kähler potential from the requirement that it be exponentially suppressed:

$$
\begin{aligned}
m^2_{C^{5_15_2}} &= m^2_{3/2} - F_{\bar{Y}_2} F_{Y_2} \partial_{\bar{Y}_2} \partial_{Y_2} \left(\ln \tilde{K}_{C^{5_15_2}} \right) \\
&= m^2_{3/2} \left[1 - \frac{3}{\tilde{K}_{C^{5_15_2}}} \partial_{\bar{Y}_2} \partial_{Y_2} \tilde{K}_{C^{5_15_2}} + \frac{3}{(\tilde{K}_{C^{5_15_2}})^2} \left(\partial_{\bar{Y}_2} \tilde{K}_{C^{5_15_2}} \right) \partial_{Y_2} \tilde{K}_{C^{5_15_2}} \right] \qquad (2.16) \\
&\equiv e^{-R_{5_2} M_*} m^2_{3/2} \qquad (2.17)
\end{aligned}
$$

where $F_{Y_2} = \sqrt{3} m_{3/2} \cos\theta \cos\phi \, e^{i\alpha_{Y_2}}$; twisted moduli domination corresponds to $\cos\theta = \cos\phi = 1$, and R_{5_2} is the compactification radius of the second complex dimension and is of order the separation between the hidden sector Y_2 moduli and the intersection states.

The 4d untwisted moduli field T_i can be decomposed into real and imaginary parts:

$$T_i = \frac{2R_i^2 M_*^2}{\lambda} + i\eta_i \tag{2.18}$$

where $R_i^2 = R_{5_i}^2$ is the compactification radius on the i^{th} torus, M_* is the string scale, λ is the 10d dilaton which is related to the fundamental (perturbative) string coupling which we can set equal to unity, and η_i is an untwisted closed string from the Ramond-Ramond sector. Hence, we can find a relationship between the real part of the untwisted T-modulus field T_2 and the compactification radius R_{5_2}:

$$R_{5_2} = \frac{\sqrt{T_2 + \bar{T}_2}}{2M_*}\,. \tag{2.19}$$

We can now solve for the Kähler potential that leads to the equivalence of eq. (2.16) and (2.17):

$$\tilde{K}_{C^{5_1 5_2}} = \exp\left[\left(1 - e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\frac{(Y_2+\bar{Y}_2)^2}{6}\right]\zeta[S,T_1,T_3] \tag{2.20}$$

where ζ is some arbitrary function of S, T_1 and/or T_3. The condition that the previous expression for the Kähler potential of eq. (2.2) is reproduced in the limit of a small compactification radius, i.e. $R_{5_2} \sim \sqrt{T_2 + \bar{T}_2} \longrightarrow 0$ fixes the function $\zeta[S,T_1,T_3]$:

$$\tilde{K}_{C^{5_1 5_2}} \longrightarrow \zeta[S,T_1,T_3] \equiv \frac{1}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}}\,. \tag{2.21}$$

Then in the limit of very large separation $R_{5_2} \sim \sqrt{T_2 + \bar{T}_2} \longrightarrow \infty$:

$$\tilde{K}_{C^{5_1 5_2}} \longrightarrow \frac{e^{(Y_2+\bar{Y}_2)^2/6}}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}} \tag{2.22}$$

In this limit, using eqs. (2.10)–(2.12), (A.17) we obtain:

$$m^2_{C^{5_1 5_2}} = m^2_{3/2}\left[1 - \frac{3}{2}\left(\sin^2\theta + \Theta_3^2\cos^2\theta\,\sin^2\phi\right) - \cos^2\theta\cos^2\phi\right] \tag{2.23}$$

which replaces the form in eq. (2.13) in the large separation limit, and which vanishes in the limit of twisted moduli domination ($\cos\theta = \cos\phi = 1$) due to the strong exponential suppression factor.

In ref. [12], Randall and Sundrum discussed the conditions under which a visible matter sector may be sequestered from a SUSY breaking hidden sector and proposed a Kähler potential that leads to vanishing scalar masses:

$$K_{\rm RS} = -3\tilde{M}_P^2 \ln\left[1 + e^{-K_{\rm vis}/3\tilde{M}_P^2} + e^{-K_{\rm hid}/3\tilde{M}_P^2}\right] \tag{2.24}$$

where $K_{\rm vis}(K_{\rm hid})$ is the separate Kähler potential for the visible (hidden) sectors, and $\tilde{M}_P = M_P/\sqrt{8\pi}$ is the reduced Planck mass. If we write the visible sector (C) and hidden sector as Y_2 Kähler potentials as:

$$K_{\rm vis} = 3\,|C|^2\,, \qquad K_{\rm hid} = \frac{3}{2}(Y_2+\bar{Y}_2)^2 \tag{2.25}$$

then the combined Kähler potential may be expanded for small $|C|^2$ and $(Y_2+\bar{Y}_2)^2$ as:

$$\begin{aligned} K_{\rm RS} &= -3\tilde{M}_P^2 \ln\left[1+e^{-|C|^2/\tilde{M}_P^2}+e^{-(Y_2+\bar{Y}_2)^2/2\tilde{M}_P^2}\right] \\ &= \frac{1}{2}(Y_2+\bar{Y}_2)^2+\mathcal{O}\left[(Y_2+\bar{Y}_2)^4\right]+|C|^2\left[1+\frac{(Y_2+\bar{Y}_2)^2}{6\tilde{M}_P^2}-\frac{(Y_2+\bar{Y}_2)^4}{72\tilde{M}_P^4}\right]+\cdots . \end{aligned} \tag{2.26}$$

The expansion of the coefficient of $|C|^2$ in eq. (2.26) is equivalent to the expansion of $e^{(Y_2+\bar{Y}_2)^2/6}$ in eq. (2.22) up to $\mathcal{O}\left[(Y_2+\bar{Y}_2)^4\right]$, where we have adopted "natural" units and set $\tilde{M}_P=1$. Therefore the numerator in eq. (2.22) is equivalent to the Randall-Sundrum sequestered form of the Kähler potential in eq. (2.26) to $\mathcal{O}\left[(Y_2+\bar{Y}_2)^4\right]$, which is sufficient for all practical purposes.

We can now write down the modified form of the tree-level Kähler potential that yields the "correct" mass for the intersection states $C^{5_15_2}$ (and similarly the 5_1-brane states $C_j^{5_1}$ and C^{95_1}) — in the limit of Y_2-domination — with an explicit dependence on the separation between the intersection point and the twisted moduli hidden sector:

$$\begin{aligned} K(S,\bar{S},T_i,\bar{T}_i,Y_2,\bar{Y}_2) =& -\ln\left(S+\bar{S}\right)-\sum_{i=1}^{3}\ln\left(T_i+\bar{T}_i\right)+\frac{1}{2}(Y_2+\bar{Y}_2)^2+ \\ &+\sum_{C_1^{5_1}}\frac{\xi(T_2,Y_2)}{(S+\bar{S})}|C_1^{5_1}|^2+\sum_{C_2^{5_1}}\frac{\xi(T_2,Y_2)}{(T_3+\bar{T}_3)}|C_2^{5_1}|^2+ \\ &+\sum_{C_3^{5_1}}\frac{\xi(T_2,Y_2)}{(T_2+\bar{T}_2)}|C_3^{5_1}|^2+\sum_{C_1^{5_2}}\frac{|C_1^{5_2}|^2}{(T_3+\bar{T}_3)}+\sum_{C_2^{5_2}}\frac{|C_2^{5_2}|^2}{(S+\bar{S})}+ \\ &+\sum_{C_3^{5_2}}\frac{|C_3^{5_2}|^2}{(T_1+\bar{T}_1)}+\sum_{C^{5_15_2}}\frac{\xi(T_2,Y_2)}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}}|C^{5_15_2}|^2+ \\ &+\sum_{C_1^9}\frac{|C_1^9|^2}{(T_1+\bar{T}_1)}+\sum_{C_2^9}\frac{|C_2^9|^2}{(T_2+\bar{T}_2)}+\sum_{C_3^9}\frac{|C_3^9|^2}{(T_3+\bar{T}_3)}+ \\ &+\sum_{C^{95_1}}\frac{\xi(T_2,Y_2)}{(T_2+\bar{T}_2)^{1/2}(T_3+\bar{T}_3)^{1/2}}|C^{95_1}|^2+\sum_{C^{95_2}}\frac{|C^{95_2}|^2}{(T_1+\bar{T}_1)^{1/2}(T_3+\bar{T}_3)^{1/2}} \end{aligned} \tag{2.27}$$

where

$$\xi(T_2,Y_2)=\exp\left[\left(1-e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\frac{(Y_2+\bar{Y}_2)^2}{6}\right]. \tag{2.28}$$

Before we repeat this analysis with the Green-Schwarz mechanism for anomaly cancellation, we will show that our result can be generalised to a construction involving three perpendicular D5-branes that all intersect at the origin. We will assume that there are three separate linear combinations of twisted moduli Y_i — one combination within each D5_i-brane world-volume — each at a distance $\mathcal{O}(R_{5_i})$ from the origin intersection states, where

$$R_{5_i}=\frac{\sqrt{T_i+\bar{T}_i}}{2M_*}. \tag{2.29}$$

We can immediately write down the form of the Kähler potential that will give the correct prediction for the masses in the Y-dominated SUSY breaking limit:

$$
\begin{aligned}
K \supset\ & \frac{\xi(T_2,Y_2)\,\xi(T_3,Y_3)}{(S+\bar{S})}|C_1^{5_1}|^2 + \frac{\xi(T_1,Y_1)\,\xi(T_3,Y_3)}{(T_3+\bar{T}_3)}|C_1^{5_2}|^2 + \\
& + \frac{\xi(T_1,Y_1)\,\xi(T_2,Y_2)}{(T_2+\bar{T}_2)}|C_1^{5_3}|^2 + \frac{\xi(T_1,Y_1)\,\xi(T_2,Y_2)\,\xi(T_3,Y_3)}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}}\,|C^{5_1 5_2}|^2 + \\
& + \frac{\xi(T_2,Y_2)\,\xi(T_3,Y_3)}{(T_2+\bar{T}_2)^{1/2}(T_3+\bar{T}_3)^{1/2}}|C^{95_1}|^2 + \frac{\xi(T_1,Y_1)\,\xi(T_3,Y_3)}{(T_1+\bar{T}_1)^{1/2}(T_3+\bar{T}_3)^{1/2}}|C^{95_2}|^2 + \cdots \,. \qquad (2.30)
\end{aligned}
$$

Notice that the $C_j^{5_i}$ and C^{95_i} states will couple directly to the twisted moduli within the same brane (Y_i), but will receive exponentially suppressed SUSY breaking contributions from the twisted moduli on different branes ($Y_{k\neq i}$). The C_j^9 states live in the full 10d space and therefore can couple to all twisted moduli.

We can also generalise our analysis to include more than one twisted moduli within the world-volumes of each D5_i-brane that contribute to SUSY breaking. For example, consider multiple twisted moduli Y_2^a inside the 5_2-brane at a distance $l^a\, R_{5_2} \equiv l^a\, \sqrt{T_2+\bar{T}_2}/2M_*$ from the origin fixed point where $C^{5_1 5_2}$ states are localised. Hence, the Kähler potential for these intersection states includes a sum over all twisted moduli:

$$
K_{C^{5_1 5_2}} = \sum_a \exp\left[\left(1 - e^{-l^a\, \sqrt{T_2+\bar{T}_2}/2}\right)\frac{(Y_2^a+\bar{Y}_2{}^a)^2}{6}\right]\frac{\left|C^{5_1 5_2}\right|^2}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}} \qquad (2.31)
$$

which we can easily extend for other sequestered states, and combinations of twisted moduli on different 5_i-branes.

3. Green-Schwarz mixing

In this section we will repeat our previous analysis, but with the inclusion of an anomaly cancelling Green-Schwarz term that requires mixing between the twisted and untwisted moduli fields. This mixing leads to a non-diagonal Kähler metric (at leading order) and we use a canonically normalising P-matrix in our parametrisation to define SUSY breaking F-terms as discussed in section A.2.

For simplicity, we will again assume that only a single linear combination of twisted moduli fields Y_2 (within the D5_2-brane world-volume at a distance $\mathcal{O}(R_{5_2})$ from the intersection point) contributes to the SUSY breaking. Since only the twisted moduli from the 5_2-brane contribute, it is not too unreasonable to suppose that only the T_2 untwisted modulus field participates in the anomaly cancellation. Using eqs. (2.5), (2.6) we propose that Y_2 has the following Kähler potential [13]:

$$
\hat{K}(Y_2,\bar{Y}_2) = \frac{1}{2}\left[Y_2+\bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right]^2 . \qquad (3.1)
$$

We will now calculate the Kähler metric by using eq. (2.27) with the modified twisted moduli potential:

$$K_{\bar{J}I} = \begin{pmatrix} \frac{1}{(S+\bar{S})^2} & 0 & 0 & 0 & 0 \\ 0 & \frac{1}{(T_1+\bar{T}_1)^2} & 0 & 0 & 0 \\ 0 & 0 & \frac{1}{(T_2+\bar{T}_2)^2}\left(k+\delta_{GS}^2\right) & 0 & -\frac{\delta_{GS}}{T_2+\bar{T}_2} \\ 0 & 0 & 0 & \frac{1}{(T_3+\bar{T}_3)^2} & 0 \\ 0 & 0 & -\frac{\delta_{GS}}{T_2+\bar{T}_2} & 0 & 1 \end{pmatrix} \tag{3.2}$$

where $I = S, T_i, Y_2$ and $k = 1 + \delta_{GS}\left[Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right]$. For simplicity, we assume that $|C|^2 \ll (S+\bar{S}), (T_i+\bar{T}_i)$.[5]

The Kähler metric is block-diagonal, and we can canonically normalise the metric using a (5×5) P-matrix as in [9]. Using eq. (A.8), the F-terms are:

$$F_I \equiv \begin{pmatrix} F_S \\ F_{T_1} \\ F_{T_2} \\ F_{T_3} \\ F_{Y_2} \end{pmatrix} = \sqrt{3} m_{3/2}\, P \begin{pmatrix} \sin\theta\, e^{i\alpha_S} \\ \cos\theta\, \sin\phi\, \Theta_1\, e^{i\alpha_1} \\ \cos\theta\, \sin\phi\, \Theta_2\, e^{i\alpha_2} \\ \cos\theta\, \sin\phi\, \Theta_3\, e^{i\alpha_3} \\ \cos\theta\, \cos\phi\, e^{i\alpha_{Y_2}} \end{pmatrix}, \tag{3.3}$$

where θ and ϕ are Goldstino angles, $\sum_i \Theta_i^\dagger \Theta_i = 1$ and we have included CP-phases.

Using eq. (3.2) and imposing the condition $P^\dagger K_{\bar{J}I} P = 1$ we obtain a very complicated expression for the P-matrix that can be expanded for large values of $(T_2+\bar{T}_2)$ to give:

$$P = \begin{pmatrix} S+\bar{S} & 0 & 0 & 0 & 0 \\ 0 & T_1+\bar{T}_1 & 0 & 0 & 0 \\ 0 & 0 & \frac{T_2+\bar{T}_2}{\sqrt{k}} & 0 & -\frac{\delta_{GS}}{T_2+\bar{T}_2} \\ 0 & 0 & 0 & T_3+\bar{T}_3 & 0 \\ 0 & 0 & \frac{\delta_{GS}}{\sqrt{k}} + \frac{\sqrt{k}\delta_{GS}}{(T_2+\bar{T}_2)^2} & 0 & 1 - \frac{\delta_{GS}^2}{(T_2+\bar{T}_2)^2} \end{pmatrix} + \mathcal{O}\left[\frac{1}{(T_2+\bar{T}_2)^3}\right] \tag{3.4}$$

where $k = 1 + \delta_{GS}\left[Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right]$.

From eqs. (3.3), (3.4) we find the SUSY breaking F-terms:

$$\begin{aligned} F_S &= \sqrt{3} m_{3/2} \sin\theta\, e^{i\alpha_S}\,(S+\bar{S}) \\ F_{T_1} &= \sqrt{3} m_{3/2} \cos\theta \sin\phi\, \Theta_1\, e^{i\alpha_1}\,(T_1+\bar{T}_1) \\ F_{T_2} &= \sqrt{3} m_{3/2} \cos\theta \left[\sin\phi \frac{(T_2+\bar{T}_2)}{\sqrt{k}}\, \Theta_2\, e^{i\alpha_2} - \cos\phi\, e^{i\alpha_{Y_2}}\, \frac{\delta_{GS}}{T_2+\bar{T}_2}\right] \\ F_{T_3} &= \sqrt{3} m_{3/2} \cos\theta \sin\phi\, \Theta_3\; e^{i\alpha_3}(T_3+\bar{T}_3) \\ F_{Y_2} &= \sqrt{3} m_{3/2} \cos\theta \left[\sin\phi \left(\frac{\delta_{GS}}{\sqrt{k}} + \frac{\sqrt{k}\delta_{GS}}{(T_2+\bar{T}_2)^2}\right) \Theta_2\, e^{i\alpha_2} + \cos\phi\, e^{i\alpha_{Y_2}} \left(1 - \frac{\delta_{GS}^2}{(T_2+\bar{T}_2)^2}\right)\right] \end{aligned} \tag{3.5}$$

where F_{T_2} and F_{Y_2} are expanded up to $\mathcal{O}\left[\frac{1}{(T_2+\bar{T}_2)^3}\right]$. Notice that in the limit of T_2 (or Y_2) modulus domination, *both* F_{T_2} and F_{Y_2} are non-zero. Setting $\cos\theta = \cos\phi = 1$ still

[5] Notice that in this limit the exact form of the intersection state Kähler potential $\tilde{K}_{C^{5_1 5_2}}$ is not important.

corresponds to the Y_2-domination limit, even in the presence of Green-Schwarz mixing, and we expect the intersection state masses to depend on the separation from the Y_2-fields as before.

Our previous analysis, in the absence of a Green-Schwarz mixing term, leads us to propose the following generalisation of the Kähler potential $\tilde{K}_{C^{5_1 5_2}}$ in eq. (2.20):

$$\tilde{K}_{C^{5_1 5_2}} = \frac{\exp\left[\frac{1}{6}\left(1 - e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2\right]}{(S+\bar{S})^{1/2}(T_3+\bar{T}_3)^{1/2}} \tag{3.6}$$

which leads to an exponentially suppressed intersection state mass in the limit of Y_2-domination:

$$m^2_{C^{5_1 5_2}} = e^{-\sqrt{T_2+\bar{T}_2}/2}\, m^2_{3/2} + \mathcal{O}\left[\frac{1}{(T_2+\bar{T}_2)^{3/2}}\right]. \tag{3.7}$$

Similar results apply to the $C_j^{5_1}$ and C^{95_1} states and we obtain a Kähler potential as in eq. (2.27), but with eq. (2.28) generalised to

$$\xi(T_2, Y_2) = \exp\left[\frac{1}{6}\left(1 - e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2\right]. \tag{3.8}$$

In section 5 we will consider an explicit example and analyse the soft parameters in various limits.

Notice that our comment at the end of section 2.3 about including the effects of multiple SUSY breaking twisted moduli still holds. The previous expression of eq. (2.30) is easily generalised by replacing the arguments as follows:

$$(Y_k + \bar{Y}_k) \longrightarrow (Y_k + \bar{Y}_k) - \delta_{GS}\ln(T_k + \bar{T}_k)\,. \tag{3.9}$$

4. Soft SUSY breaking parameters

We will now write down the complete list of soft scalar masses and trilinears that arise in a general string construction involving two intersecting D5-branes embedded within a D9-brane, where a single linear combination of twisted moduli Y_2 is located at a fixed point within the world-volume of one of the branes (5_2). It is straightforward to generalise these results to more twisted moduli fields, and we have explicitly discussed the case of three twisted moduli Y_i in section 2. The results presented in this section will be useful for performing more general phenomenological analyses of sparticle spectra in string theory than have been done so far in the literature.[6]

Note that the gaugino masses require knowledge of the gauge group embedding, and therefore gaugino masses are more model-dependent. We will consider a simple example in section 5.

[6] In appendix B we consider an alternative exponential suppression factor $e^{-(M_* r)^2}$ that can be attributed to non-perturbative world-sheet instanton corrections [15], and we calculate modified expressions for scalar masses and trilinears.

4.1 Scalar masses

Using eqs. (2.27), (3.5), (3.8), (A.17) we can write down the scalar masses for the non-sequestered states $C_j^{5_2}, C_j^9$ and C^{95_2} which couple directly to the twisted moduli Y_2:

$$\begin{aligned} m^2_{C_1^{5_2}} = m^2_{C_3^9} &= m^2_{3/2}\left[1 - 3\Theta_3^2\,\cos^2\theta\,\sin^2\phi\right] \\ m^2_{C_2^{5_2}} &= m^2_{3/2}\left[1 - 3\sin^2\theta\right] \\ m^2_{C_3^{5_2}} = m^2_{C_1^9} &= m^2_{3/2}\left[1 - 3\Theta_1^2\,\cos^2\theta\,\sin^2\phi\right] \\ m^2_{C_2^9} &= m^2_{3/2}\left[1 - 3\Theta_2^2\,\cos^2\theta\,\sin^2\phi\right] \\ m^2_{C^{95_2}} &= m^2_{3/2}\left[1 - \frac{3}{2}\cos^2\theta\,\sin^2\phi\left(\Theta_1^2 + \Theta_3^2\right)\right]. \end{aligned} \tag{4.1}$$

The sequestered states $C_j^{5_2}, C^{5_1 5_2}$ and C^{95_1} are spatially separated from the twisted modulus field Y_2 and have masses of the form:

$$\begin{aligned} m^2_{C^{5_1 5_2}} &= \tilde{m}^2 - \frac{3}{2}m^2_{3/2}\left(\sin^2\theta + \Theta_3^2\,\cos^2\theta\,\sin^2\phi\right) \\ m^2_{C^{95_1}} &= \tilde{m}^2 - \frac{3}{2}m^2_{3/2}\,\cos^2\theta\,\sin^2\phi\left(\frac{\Theta_2^2}{k} + \Theta_3^2\right) \\ m^2_{C_1^{5_1}} &= \tilde{m}^2 - 3m^2_{3/2}\sin^2\theta \\ m^2_{C_2^{5_1}} &= \tilde{m}^2 - 3m^2_{3/2}\Theta_3^2\cos^2\theta\sin^2\phi \\ m^2_{C_3^{5_1}} &= \tilde{m}^2 - \frac{3}{k}m^2_{3/2}\Theta_2^2\cos^2\theta\sin^2\phi \end{aligned} \tag{4.2}$$

where

$$\begin{aligned} \tilde{m}^2 = m^2_{3/2}\Bigg[& 1 - \cos^2\theta\cos^2\phi\left(1 - e^{-\sqrt{T_2+\bar{T}_2}/2}\right) - \\ & - \frac{\cos^2\theta\sin^2\phi\,\Theta_2^2\,\delta_{GS}}{k}\left(1 - e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\} + \\ & + \frac{\cos^2\theta\sin^2\phi\,\Theta_2^2\,e^{-\sqrt{T_2+\bar{T}_2}/2}}{32k}\sqrt{T_2+\bar{T}_2}\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2 \times \\ & \times\left(2 + \sqrt{T_2+\bar{T}_2}\right) - \\ & - \frac{\cos^2\theta\cos\phi\sin\phi\left(\Theta_2\,e^{i(\alpha_2-\alpha_{Y_2})} + \Theta_2^\dagger\,e^{-i(\alpha_2-\alpha_{Y_2})}\right)e^{-\sqrt{T_2+\bar{T}_2}/2}}{32\sqrt{k}\,(T_2+\bar{T}_2)} \times \\ & \times\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\} \times \\ & \times\left(8(T_2+\bar{T}_2)^{3/2} + \delta_{GS}\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right)\Bigg] + \mathcal{O}\left[\frac{\delta_{GS}}{(T_2+\bar{T}_2)^{3/2}}\right]. \end{aligned} \tag{4.3}$$

In the limit of a large separation, $R_{5_2} \sim \sqrt{T_2 + \bar{T}_2} \to \infty$

$$\tilde{m}^2 \to m^2_{3/2}\left[1 - \cos^2\theta\cos^2\phi - \frac{\cos^2\theta\sin^2\phi\,\Theta_2^2\,\delta_{GS}}{k}\left\{Y_2 + \bar{Y}_2 - \delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right] \tag{4.4}$$

and for a small separation $R_{5_2} \sim \sqrt{T_2+\bar{T}_2} \to 0$

$$\tilde{m}^2 \to m_{3/2}^2\,. \tag{4.5}$$

Now we will consider the different limits of SUSY breaking:

- Dilaton domination ($\sin\theta = 1$):

$$\begin{aligned} m^2_{C^{5_2}_{1,3}} &= m^2_{C^{5_1}_{2,3}} = m^2_{C^9_{1,2,3}} = m^2_{C^{95_{1,2}}} = m^2_{3/2} \\ m^2_{C^{5_2}_2} &= m^2_{C^{5_1}_1} = -2m^2_{3/2} \\ m^2_{C^{5_15_2}} &= -\frac{1}{2}m^2_{3/2}\,. \end{aligned} \tag{4.6}$$

  Notice that this limit generally gives rise to tachyonic states.

- T-moduli domination ($\cos\theta = \sin\phi = 1$):

$$\begin{aligned} m^2_{C^{5_2}_1} = m^2_{C^9_3} = m^2_{3/2}\left(1-3\Theta_3^2\right), \quad & m^2_{C^{5_2}_2} = m^2_{3/2}\,, \quad m^2_{C^{5_2}_3} = m^2_{C^9_1} = m^2_{3/2}\left(1-3\Theta_1^2\right) \\ m^2_{C^9_2} &= m^2_{3/2}\left(1-3\Theta_2^2\right), \quad m^2_{C^{5_15_2}} = \tilde{m}_T^2 - \frac{3}{2}m^2_{3/2}\Theta_3^2 \\ m^2_{C^{95_2}} &= m^2_{3/2}\left[1-\frac{3}{2}\left(\Theta_1^2+\Theta_3^2\right)\right], \quad m^2_{C^{95_1}} = \tilde{m}_T^2 - \frac{3}{2}m^2_{3/2}\left(\frac{\Theta_2^2}{k}+\Theta_3^2\right) \\ m^2_{C^{5_1}_1} &= \tilde{m}_T^2\,, \quad m^2_{C^{5_1}_2} = \tilde{m}_T^2 - 3m^2_{3/2}\Theta_3^2\,, \quad m^2_{C^{5_1}_3} = \tilde{m}_T^2 - \frac{3}{k}m^2_{3/2}\Theta_2^2 \end{aligned} \tag{4.7}$$

  where

$$\begin{aligned} \tilde{m}_T^2 = m^2_{3/2}\Bigg[1 &- \frac{\Theta_2^2\delta_{GS}}{k}\left(1-e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\} + \\ &+ \frac{\Theta_2^2 e^{-\sqrt{T_2+\bar{T}_2}/2}}{32k}\sqrt{T_2+\bar{T}_2}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2 \times \\ &\times\left(2+\sqrt{T_2+\bar{T}_2}\right)\Bigg] \end{aligned}$$

- Y_2-moduli domination ($\cos\theta = \cos\phi = 1$):

$$\begin{aligned} m^2_{C^{5_2}_{1,2,3}} &= m^2_{C^9_{1,2,3}} = m^2_{C^{95_2}} = m^2_{3/2} \\ m^2_{C^{5_15_2}} &= m^2_{C^{5_1}_{1,2,3}} = m^2_{C^{95_1}} = e^{-\sqrt{T_2+\bar{T}_2}/2}\,m^2_{3/2} \end{aligned} \tag{4.8}$$

  where the 5_2-and 9-brane states couple directly to the SUSY breaking twisted moduli and are not exponentially suppressed.

Physically the twisted Y_2 moduli dominated limit, corresponds to gaugino mediated SUSY breaking, if the standard model quark and lepton states are identified with the sequestered states (see later example). The dilaton and T-moduli domination limits correspond to different examples of gravity mediated SUSY breaking. In the general case where one is not in any particular limit, SUSY breaking will have contributions from the F-terms of the dilaton and untwisted moduli as well as the twisted moduli, and then one must use the general formulae for the scalar masses in eqs. (4.1), (4.2).

4.2 Trilinears

The trilinear and Yukawa couplings arise from the superpotential, where the dominant tree-level contributions are shown in eq. (2.8) in terms of open string states. The precise structure of the Yukawa and trilinear matrices depend on the identification of these string states with MSSM fields. The leading terms are constrained by string selection rules and gauge invariance. Higher order corrections can be generated by higher-dimensional operators where powers of the model cutoff (e.g. the string or Planck scales) lead to a large suppression. We will illustrate how different identifications lead to alternative Yukawa structures in section 5, as recently discussed in ref. [16].

We will now list the dominant trilinear couplings that arise from the perturbative superpotential of eq. (2.8). Using eqs. (2.27), (3.5), (3.8), (A.19) and making the standard assumption that the Yukawa couplings Y_{abc} have no dependence on the dilaton and moduli fields ($\partial_I \ln Y_{abc} = 0$)[7] we find:

$$\begin{aligned}
A_{C_3^{5_2} C^{5_1 5_2} C^{5_1 5_2}} = & -\sqrt{3} m_{3/2} \cos\theta \times \\
& \times \Bigg[\sin\phi \frac{\Theta_2 e^{i\alpha_2}}{\sqrt{k}} + \sin\phi \frac{\Theta_2 e^{i\alpha_2}}{12\sqrt{k}} e^{-\sqrt{T_2+\bar{T}_2}/2} \sqrt{T_2+\bar{T}_2} \times \\
& \quad \times \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\}^2 - \\
& \quad - \cos\phi \frac{e^{i\alpha_{Y_2}}}{3} \left(1 + 2 e^{-\sqrt{T_2+\bar{T}_2}/2} \right) \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\} \Bigg] \qquad (4.9)
\end{aligned}$$

$$\begin{aligned}
A_{C_1^9 C^{95_1} C^{95_1}} = & -\sqrt{3} m_{3/2} \times \\
& \times \Bigg[\sin\theta e^{i\alpha_S} + \cos\theta \sin\phi \frac{\Theta_2 e^{i\alpha_2}}{12\sqrt{k}} e^{-\sqrt{T_2+\bar{T}_2}/2} \sqrt{T_2+\bar{T}_2} \times \\
& \quad \times \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\}^2 - \qquad (4.10) \\
& \quad - \cos\theta \cos\phi \frac{e^{i\alpha_{Y_2}}}{3} \left(1 + 2\, e^{-\sqrt{T_2+\bar{T}_2}/2} \right) \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\} \Bigg]
\end{aligned}$$

$$\begin{aligned}
A_{C^{5_1 5_2} C^{95_1} C^{95_2}} = & -\sqrt{3} m_{3/2} \times \\
& \times \Bigg[\frac{1}{2} \sin\theta e^{i\alpha_S} + \frac{1}{2} \cos\theta \sin\phi \left(\Theta_1\, e^{i\alpha_1} + \frac{\Theta_2 e^{i\alpha_2}}{\sqrt{k}} - \Theta_3 e^{i\alpha_3} \right) + \qquad (4.11) \\
& \quad + \cos\theta \sin\phi \frac{\Theta_2 e^{i\alpha_2}}{12\sqrt{k}} e^{-\sqrt{T_2+\bar{T}_2}/2} \sqrt{T_2+\bar{T}_2} \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\}^2 - \\
& \quad - \cos\theta \cos\phi \frac{e^{i\alpha_{Y_2}}}{3} \left(1 + 2 e^{-\sqrt{T_2+\bar{T}_2}/2} \right) \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\} \Bigg]
\end{aligned}$$

$$\begin{aligned}
A_{C_3^{5_1} C^{5_1 5_2} C^{5_1 5_2}} = & A_{C_1^{5_1} C_2^{5_1} C_3^{5_1}} = A_{C_1^{5_1} C^{95_1} C^{95_1}} \\
= & -\sqrt{3} m_{3/2} \cos\theta \Bigg[\sin\phi \Theta_1 e^{i\alpha_1} + \sin\phi \frac{\Theta_2 e^{i\alpha_2}}{8\sqrt{k}} e^{-\sqrt{T_2+\bar{T}_2}/2} \sqrt{T_2+\bar{T}_2} \times \\
& \quad \times \left\{ Y_2 + \bar{Y}_2 - \delta_{GS} \ln(T_2+\bar{T}_2) \right\}^2 - \qquad (4.12)
\end{aligned}$$

[7] Notice that by definition, the Yukawa couplings are related to the moduli fields through eq. (2.9) so this assumption is not really valid, but we make it for illustrative purposes so that our results may be compared to others in the literature.

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$$- \cos\phi\, e^{i\alpha_{Y_2}} e^{-\sqrt{T_2+\bar{T}_2}/2} \left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\Big]$$

$$\begin{aligned} A_{C_1^{5_2}C_2^{5_2}C_3^{5_2}} &= A_{C_2^{5_2}C^{95_2}C^{95_2}} \qquad (4.13)\\ &= -\sqrt{3}m_{3/2}\cos\theta\left[\sin\phi\,\frac{\Theta_2\, e^{i\alpha_2}}{\sqrt{k}} - \cos\phi\, e^{i\alpha_{Y_2}}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right]\\ A_{C_1^9C_2^9C_3^9} &= A_{C_2^9C^{95_2}C^{95_2}} \qquad (4.14)\\ &= -\sqrt{3}m_{3/2}\left[\sin\theta e^{i\alpha_S} - \cos\theta\cos\phi e^{i\alpha_{Y_2}}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right] \end{aligned}$$

Now we will consider the different limits of SUSY breaking:

- Dilaton domination ($\sin\theta = 1$):

$$\begin{aligned} A_{C_1^9C^{95_1}C^{95_1}} &= A_{C_1^9C_2^9C_3^9} = A_{C_2^9C^{95_2}C^{95_2}} = 2\,A_{C^{5_15_2}C^{95_1}C^{95_2}} = -\sqrt{3}m_{3/2}\,e^{i\alpha_S}\\ A_{C_3^{5_2}C^{5_15_2}C^{5_15_2}} &= A_{C_3^{5_1}C^{5_15_2}C^{5_15_2}} = A_{C_1^{5_1}C_2^{5_1}C_3^{5_1}} = 0\\ A_{C_1^{5_1}C^{95_1}C^{95_1}} &= A_{C_2^{5_2}C^{95_2}C^{95_2}} = A_{C_1^{5_2}C_2^{5_2}C_3^{5_2}} = 0\,. \end{aligned} \qquad (4.15)$$

- T-moduli domination ($\cos\theta = \sin\phi = 1$):

$$\begin{aligned} A_{C_3^{5_2}C^{5_15_2}C^{5_15_2}} &= -\sqrt{\frac{3}{k}}\,m_{3/2}\,\Theta_2\,e^{i\alpha_2} - \hat{A}\\ A_{C_1^9C^{95_1}C^{95_1}} &= -\hat{A}\\ A_{C^{5_15_2}C^{95_1}C^{95_2}} &= -\frac{\sqrt{3}}{2}m_{3/2}\left(\Theta_1\,e^{i\alpha_1} + \frac{\Theta_2\,e^{i\alpha_2}}{\sqrt{k}} - \Theta_3\,e^{i\alpha_3}\right) - \hat{A}\\ A_{C_3^{5_1}C^{5_15_2}C^{5_15_2}} &= A_{C_1^{5_1}C_2^{5_1}C_3^{5_1}} = A_{C_1^{5_1}C^{95_1}C^{95_1}} = -\sqrt{3}m_{3/2}\,\Theta_1\,e^{i\alpha_1} - \frac{3}{2}\hat{A}\\ A_{C_1^{5_2}C_2^{5_2}C_3^{5_2}} &= A_{C_2^{5_2}C^{95_2}C^{95_2}} = -\sqrt{\frac{3}{k}}\,m_{3/2}\,\Theta_2\,e^{i\alpha_2}\\ A_{C_1^9C_2^9C_3^9} &= A_{C_2^9C^{95_2}C^{95_2}} = 0 \end{aligned} \qquad (4.16)$$

where

$$\hat{A} = m_{3/2}\frac{\Theta_2\,e^{i\alpha_2}}{4\sqrt{3k}}\,e^{-\sqrt{T_2+\bar{T}_2}/2}\,\sqrt{T_2+\bar{T}_2}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2. \qquad (4.17)$$

- Y_2-moduli domination ($\cos\theta = \cos\phi = 1$):

$$\begin{aligned} A_{C_3^{5_2}C^{5_15_2}C^{5_15_2}} &= A_{C_1^9C^{95_1}C^{95_1}} = A_{C^{5_15_2}C^{95_1}C^{95_2}} = \left(1+2\,e^{-\sqrt{T_2+\bar{T}_2}/2}\right)\tilde{A}\\ A_{C_3^{5_1}C^{5_15_2}C^{5_15_2}} &= A_{C_1^{5_1}C_2^{5_1}C_3^{5_1}} = A_{C_1^{5_1}C^{95_1}C^{95_1}} = 3e^{-\sqrt{T_2+\bar{T}_2}/2}\,\tilde{A}\\ A_{C_1^{5_2}C_2^{5_2}C_3^{5_2}} &= A_{C_2^{5_2}C^{95_2}C^{95_2}} = A_{C_1^9C_2^9C_3^9} = A_{C_2^9C^{95_2}C^{95_2}} = 3\tilde{A} \end{aligned} \qquad (4.18)$$

where

$$\tilde{A} = m_{3/2}\frac{e^{i\alpha_{Y_2}}}{\sqrt{3}}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}. \qquad (4.19)$$

5. Gaugino mediated SUSY breaking

The preceeding results are based on a general set-up of intersecting $D5_i$ and D9 branes. In order to discuss $\tilde{g}MSB$ it is sufficient to specialize to the case of just two intersecting sets of D5 branes, 5_1 and 5_2. This set-up arises for example in the explicit string constructions of [17]. We shall assume that the MSSM gauge group arises from the 5_2-brane only. This enables approximate gauge coupling unification to be achieved. The MSSM matter fields are identified as either $C^{5_1 5_2}$ or $C_j^{5_2}$ states. We assume that any $C^{5_1 5_2}$ states are gauge singlets with respect to any gauge groups on the 5_1-brane. Such a set-up may be achieved in practice by constructions involving severely asymmetric compactifications (for example $R_{5_2} \gg R_{5_1}$ [18]), where the combined gauge groups generally arise from linear combinations of groups on each set of branes. The asymmetry ensures that the dominant contributions live on the 5_2-brane, a limit we refer to as "single brane dominance".

Although the perpendicular 5_1-brane seems to be irrelevant in this scenario, in fact it plays an important rôle since the $C^{5_1 5_2}$ states are sequestered at a distance $r \sim \mathcal{O}(R_{5_2})$ from the fixed point associated with the twisted modulus Y_2. From eq. (4.2) we find the soft mass for the sequestered state to be

$$m^2_{C^{5_1 5_2}} = \tilde{m}^2 - \frac{3}{2} m^2_{3/2} \left(\sin^2\theta + \Theta_3^2 \cos^2\theta \sin^2\phi\right) \tag{5.1}$$

where $\tilde{m}^2$ is given in eq. (4.3). In the twisted moduli dominated limit ($\cos\theta = \cos\phi = 1$)

$$m^2_{C^{5_1 5_2}} = e^{-\sqrt{T_2+\bar{T}_2}/2}\, m^2_{3/2}\,. \tag{5.2}$$

If the standard model states are all identified as intersection states $C^{5_1 5_2}$ then for large radius of compactification this corresponds to gaugino mediated SUSY breaking. However the soft mass in eq. (5.1) is valid away from the twisted modulus dominated limit, and also is valid for a small compactification radius. It therefore allows more general and detailed studies of gaugino mediation to be performed, including the effects of finite radius of compactification, and the contributions from gravity mediation effects, which in type I theories correspond to the dilaton and untwisted moduli F-term vevs.

The non-sequestered soft masses are given by eq. (4.1),

$$\begin{aligned} m^2_{C_1^{5_2}} = m^2_{C_3^9} &= m^2_{3/2}\left[1 - 3\Theta_3^2 \cos^2\theta \sin^2\phi\right] \\ m^2_{C_2^{5_2}} &= m^2_{3/2}\left[1 - 3\sin^2\theta\right] \\ m^2_{C_3^{5_2}} &= m^2_{3/2}\left[1 - 3\Theta_1^2 \cos^2\theta \sin^2\phi\right]. \end{aligned} \tag{5.3}$$

The MSSM gauge groups all arise from the 5_2-brane, and using eq. (1.3) with a single linear combination of twisted moduli fields within the 5_2-brane world-volume, we find:

$$f_\alpha^{5_2} = T_2 + \frac{s_\alpha}{4\pi} Y_2 \qquad (\alpha = \mathrm{SU}(3)_C, \mathrm{SU}(2)_L, \mathrm{U}(1)_Y) \tag{5.4}$$

where s_α are model-dependent coefficients that depend on the details of the orbifold compactification. Notice that for Z_3 and Z_7 orbifolds, these coefficents are proportional to the MSSM 1-loop beta-function coefficients b_α.

We can find the gaugino masses using eqs. (1.4), (3.5), (A.18):

$$M_\alpha = \frac{\sqrt{3}m_{3/2}\, g_\alpha^2}{8\pi}\cos\theta\left[\sin\phi\,\Theta_2\, e^{i\alpha_2}\left\{\frac{T_2+\bar{T}_2}{\sqrt{k}}+\frac{s_\alpha}{4\pi}\left(\frac{\delta_{GS}}{\sqrt{k}}+\frac{\sqrt{k}\,\delta_{GS}}{(T_2+\bar{T}_2)^2}\right)\right\}-\right.$$
$$\left.-\cos\phi\, e^{i\alpha_{Y_2}}\left\{\frac{\delta_{GS}}{T_2+\bar{T}_2}-\frac{s_\alpha}{4\pi}\left(1-\frac{\delta_{GS}^2}{(T_2+\bar{T}_2)^2}\right)\right\}\right]. \tag{5.5}$$

Now consider different limits of SUSY breaking:

- Dilaton domination ($\sin\theta = 1$):

$$M_\alpha = 0 \qquad (\alpha = \mathrm{SU}(3)_C, \mathrm{SU}(2)_L, \mathrm{U}(1)_Y) \tag{5.6}$$

- T-modulus domination ($\cos\theta = \sin\phi = 1$):

$$M_\alpha = \frac{\sqrt{3}m_{3/2}\, g_\alpha^2}{8\pi}\Theta_2\, e^{i\alpha_2}\left\{\frac{T_2+\bar{T}_2}{\sqrt{k}}+\frac{s_\alpha}{4\pi}\left(\frac{\delta_{GS}}{\sqrt{k}}+\frac{\sqrt{k}\,\delta_{GS}}{(T_2+\bar{T}_2)^2}\right)\right\}. \tag{5.7}$$

- Y_2-modulus domination ($\cos\theta = \cos\phi = 1$):

$$M_\alpha = -\frac{\sqrt{3}m_{3/2}\, g_\alpha^2}{8\pi}e^{i\alpha_{Y_2}}\left\{\frac{\delta_{GS}}{T_2+\bar{T}_2}-\frac{s_\alpha}{4\pi}\left(1-\frac{\delta_{GS}^2}{(T_2+\bar{T}_2)^2}\right)\right\}. \tag{5.8}$$

5.1 Scenario A — gaugino mediated SUSY breaking for all three families

In scenario A, depicted in figure 3, the three chiral families are open string states localised at the origin fixed point, and the Higgs fields feel two extra dimensions as open string states with both ends attached to the 5_2-brane.

$$\begin{aligned} Q_i,\, L_i,\, U_i,\, D_i,\, E_i &\equiv C^{5_1 5_2} \qquad (i=1,2,3)\\ H_u,\, H_d &\equiv C_j^{5_2}\,. \end{aligned} \tag{5.9}$$

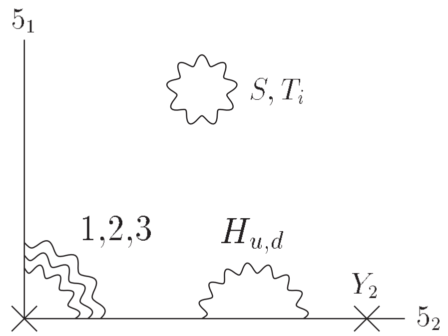

Figure 3: The allocation of charged chiral fields in scenario A which is similar to the gaugino mediated SUSY breaking model [11]. The MSSM gauge group arises from the 5_2-brane, and all three MSSM chiral families are localised at the origin, while the Higgs and MSSM gauge fields live on the 5_2-brane. The dilaton and moduli fields S, T_i live in the full 10d space and a single twisted moduli Y_2 is localised at a fixed point inside the 5_2-brane world-volume.

The Higgs states carry an extra index that plays an important rôle in constructing the perturbative superpotential from open string states. The tree-level superpotential of eq. (2.8) contains the terms:

$$W_{\rm ren} = \mathcal{O}(g_{5_2}) \left[C_1^{5_2} C_2^{5_2} C_3^{5_2} + C_3^{5_2} C^{5_1 5_2} C^{5_1 5_2} \right] . \tag{5.10}$$

In order to obtain non-zero third family Yukawa couplings (at tree-level), we can immediately see that the Higgs fields must be $C_3^{5_2}$ states. This leads to a "democratic" Yukawa texture (and trilinear matrix) where all entries are equal:

$$Y_{ij}^{ude} \sim \mathcal{O}(g_{5_2}) \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix} . \tag{5.11}$$

The democratic structure arises due to the presence of three (indistinguishable) chiral families, localised at the origin fixed point ($C^{5_1 5_2}$). However, type I compactifications do not generally lead to low-energy spectra with this property as one (or more) families generally arise with both ends attached to the same D5-brane (C^{5_i}) which is the situation in scenario B.

The squark and slepton and higgs soft masses are given by eqs. (5.1), (5.3) with the identifications in eq. (5.9). In general the squark and slepton ($C^{5_1 5_2}$) soft masses receive unsuppressed contributions from the dilaton and untwisted moduli F-term vevs, which corresponds to the string version of normal gravity mediation. In the limit of twisted moduli domination, we see that the quarks and lepton states acquire exponentially small soft scalar masses:

$$m^2_{\tilde{Q}_i} = m^2_{\tilde{U}_i} = m^2_{\tilde{D}_i} = m^2_{\tilde{L}_i} = m^2_{\tilde{E}_i} = e^{-\sqrt{T_2 + \bar{T}_2}/2}\, m^2_{3/2} \qquad (i = 1, 2, 3) \tag{5.12}$$

and the Higgs scalars obtain much larger masses due to their direct coupling with the SUSY breaking sector (twisted moduli):

$$m^2_{h_u} = m^2_{h_d} = m^2_{3/2} \,. \tag{5.13}$$

This yields the same spectrum as the gaugino mediated SUSY breaking scenario [11], where vanishingly small scalar masses (due to the separation between sectors) offers an attractive (and natural) solution to the SUSY flavour problem.[8] However, unlike $\tilde{g}MSB$, the third family trilinear $A_{33} \equiv A_{C_3^{5_2} C^{5_1 5_2} C^{5_1 5_2}}$ from eq. (4.18) is not loop suppressed and depends on the explicit function of twisted moduli $\hat{K}(Y_2, \bar{Y}_2, T_2, \bar{T}_2)$.

The general results in eqs. (5.1), (5.3) enable us to smoothly move away from the twisted moduli dominated limit (corresponding to gaugino mediated SUSY breaking) and consider the contributions of the dilaton and untwisted moduli to the soft masses (corresponding to the gravity contributions to SUSY breaking). We can also consider the effect of smoothly changing the compactification radius R_{5_2} (corresponding to varying the distance r in figure 2.)

[8] Flavour-changing neutral-current suppression places a lower limit on the size of the separation.

Notice that if we assign the Higgs fields as different 5_2-brane states — for example $H_u \equiv C_3^{5_2}$ and $H_d \equiv C_{1,2}^{5_2}$ — then it is possible to generate a μ-term in the tree-level superpotential if we add a gauge singlet that acquires a non-zero vev ($N \equiv C_1^{5_2}$ say).

$$W_{\rm ren} \supset N H_u H_d \longrightarrow \mu \sim \langle N \rangle \,. \tag{5.14}$$

5.2 Scenario B — gaugino mediated SUSY breaking for the first and second families only

In scenario B, depicted in figure 4, the third family is moved on to the 5_2-brane along with the Higgs and gauge fields.

$$\begin{aligned} Q_i,\, L_i,\, U_i,\, D_i,\, E_i &\equiv C^{5_1 5_2} \quad (i=1,2) \\ Q_3,\, L_3,\, U_3,\, D_3,\, E_3,\, H_u,\, H_d &\equiv C_j^{5_2} \end{aligned} \tag{5.15}$$

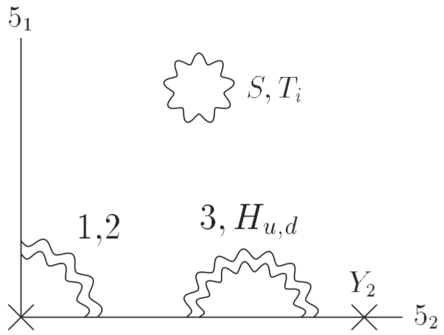

Figure 4: The allocation of charged chiral fields in scenario B which is similar to the brane mediated SUSY breaking model [18]. The MSSM gauge group arises from the 5_2-brane, and only the first two chiral families live at the origin while the third family, Higgs and gauge fields live on the 5_2-brane. The dilaton and moduli fields S, T_i live in the full 10d space and a single twisted moduli Y_2 is localised at a fixed point inside the 5_2-brane world-volume.

The separation of the third family from the first two chiral families appears frequently in type I string compactifications. Ref. [18] provides the motivation for this scenario in the limit of a vanishing 5_1-brane compactification radius. In this case, the gauge groups are dominated by their components on the 5_2-brane. Notice that an extended Pati-Salam gauge group appears instead of the MSSM group, and gauge invariance with respect to this larger symmetry prevents proton decay (and first and second family Yukawa couplings at tree-level) by forbidding R-parity violating operators.

In order to generate a third family Yukawa coupling at tree-level, the Higgs and third family singlets and doublets must carry different indices. However, we are still free to choose whether the Higgs fields allocation can give rise to first and second family Yukawa couplings at tree-level. For example, suppose that we choose the following allocations:

$$Q_3,\, L_3 \equiv C_1^{5_2} \qquad U_3,\, D_3,\, E_3 \equiv C_2^{5_2} \qquad H_u,\, H_d \equiv C_3^{5_2}. \tag{5.16}$$

We will generate block-diagonal Yukawa textures that are not consistent with experimental data:

$$Y_{ij}^{ude} \sim \mathcal{O}(g_{5_2}) \begin{pmatrix} 1 & 1 & 0 \\ 1 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}. \tag{5.17}$$

However, if we choose that $H_{u,d} \neq C_3^{5_2}$, then we generate a Yukawa texture with only a single non-zero value in the (33) entry:

$$Y_{ij}^{ude} \sim \mathcal{O}(g_{5_2}) \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \tag{5.18}$$

that is more compatible with data, as higher order corrections can generate the required structure. We also obtain a trilinear matrix with a single (33) entry:

$$\tilde{A}_{ij}^{ude} = A_{ij} Y_{ij}^{ude} \sim \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & A_{33} \end{pmatrix} \tag{5.19}$$

where $A_{33} = A_{C_1^{5_2} C_2^{5_2} C_3^{5_2}}$.

The soft masses for this scenario are very similar to scenario A, except that the third family is now a $C_2^{5_2}$ state, so is now a non-sequestered state. In the twisted moduli domination limit of eq. (4.8), the first two families receive exponentially suppressed masses as for scenario A. However, the third family and Higgs acquire large soft masses $\sim \mathcal{O}(m_{3/2})$ which may be read off from eq. (5.3).

Notice that experimental constraints from flavour-changing neutral-current data is only sensitive to the first two families, and this scenario (with a hierarchically larger third family) may not violate these constraints, thereby providing an interesting alternative solution to the flavour-changing problem.

6. Conclusions

We have considered twisted moduli contributions to supersymmetry breaking in effective type I string constructions based on intersecting D5_i and D9-branes, using the formalism of Goldstino angles and extending the scope of previous analyses which were based on a single D9-brane sector. The more general set-up allows the possibility of states which are sequestered from twisted moduli states which are located at fixed points and cannot move freely. The sequestered states should have suppressed soft mass contributions from distant twisted moduli, and this observation has been used to suggest how $\tilde{g}MSB$ might be implemented in type I string theory [10]. However, contrary to this expectation, we found that the standard form of the Kähler potential leads to non-zero soft masses for the sequestered states in the twisted moduli dominated limit. This motivated us to look for a new form of Kähler potential for the sequestered states. We have proposed a new form of the Kähler potential which is consistent at leading order with the sequestered form proposed by Randall and Sundrum [12], and which leads to exponentially suppressed sequestered soft masses. Including the effect of Green-Schwarz mixing we have written down soft scalar and trilinear masses arising from a general string construction involving intersecting D5_i and D9-branes in the presence of untwisted and twisted moduli. We have shown how the results may be applied to $\tilde{g}MSB$, and discussed two explicit scenarios for this based on two intersecting 5_1 and 5_2 brane sectors, in which the MSSM gauge group is placed on the 5_2 sector. The second scenario in which $\tilde{g}MSB$ only applies to the first two families, and the third family receives an unsuppressed soft mass was first discussed in [18].

The general results will be useful in phenomenological studies involving a combination of gravity and gaugino mediated SUSY breaking due to the dilaton, untwisted and twisted moduli contributions, and enable the soft masses to be studied as a function of the compactification radii. Previous analyses [9] have only considered the effect of twisted moduli in the case where the gauge group and matter fields live on the D9-brane, and share the same world-volume with all twisted moduli fields. However such a scenario does not give rise to localised matter fields (confined at intersection points) and in general one does not encounter states which are sequestered from twisted moduli. Hence the standard Kähler potentials used in those analyses are perfectly acceptable. By contrast our analysis opens the door for more general type I string constructions involving D9 *and* D5$_i$-branes, where potentially more realistic phenomenology and hierarchies between observables can be obtained with some or all of the matter fields sequestered from twisted moduli SUSY breaking sectors.

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A. Supergravity basics

We will now use a conventional supergravity formalism to describe the 4d effective theory that arises as the low energy limit of the type I theory. Supergravity (local SUSY) is defined in terms of a Kähler function (G) of generic chiral superfields ($\phi = h, C_a$) including the hidden sector closed strings ($h = S, T_i, Y_k$) and open string matter states ($C_a = C_j^{5_i}, C^{5_i 5_j}$):[9]

$$G(\phi, \bar{\phi}) = \frac{K(\phi, \bar{\phi})}{\tilde{M}_P^2} + \ln\left(\frac{W(\phi)}{\tilde{M}_P^3}\right) + \ln\left(\frac{W^*(\bar{\phi})}{\tilde{M}_P^3}\right). \tag{A.1}$$

The Kähler potential $K(\phi, \bar{\phi})$ is a real function of chiral superfields and may be expanded in powers of matter states C_a [3] (including non-perturbative contributions):

$$K = \bar{K}(h, \bar{h}) + \tilde{K}_{\bar{a}b}(h, \bar{h})\bar{C}_{\bar{a}}C_b + \left[\frac{1}{2}Z_{ab}(h, \bar{h})C_aC_b + \text{h.c.}\right] + \cdots \tag{A.2}$$

where $\tilde{K}_{a\bar{b}}$ is the (generally non-diagonal) matter metric and a non-zero bilinear term Z_{ab} can generate the μ-term through the Guidice-Masiero mechanism [19] subject to gauge-invariance. The superpotential $W(\phi)$ is a holomorphic function of chiral superfields that can also be expanded:

$$W = \hat{W}(h) + \frac{1}{2}\mu_{ab}(h)C_aC_b + \frac{1}{6}Y_{abc}C_aC_bC_c + \cdots. \tag{A.3}$$

Notice that it includes a trilinear Yukawa term (that will generate fermion masses) and a bilinear μ-term. However, the Kähler potential and superpotential also receive non-perturbative contributions that are often difficult to predict. To make progress we will

[9]Notice that we have included powers of the reduced Planck mass ($\tilde{M}_P$) that appear in the Kähler function to obtain the correct dimensions, although it is conventional to adopt natural units and set $\tilde{M}_P = 1$.

utilise a simple parametrisation of our ignorance of the non-perturbative sector in terms of Goldstino angles and CP-phases.

A.1 Supergravity potential

We know that in Nature SUSY must be broken and various mechanisms have been proposed. It is convenient to analyse the SUSY breaking by considering the F-part of the SUGRA scalar potential.[10] It can be expressed in terms of derivatives of the Kähler function $G(\phi, \bar{\phi})$, or equivalently in terms of the F-term auxiliary fields that can acquire non-zero vevs and trigger the SUSY breaking [3]. Using eq. (A.1) we obtain:

$$V(\phi, \bar{\phi}) = e^G \left[G_I (K^{-1})_{I\bar{J}} G_{\bar{J}} - 3\right] = F_{\bar{J}}\, K_{\bar{J}I}\, F_I - 3e^K\, |W|^2 \tag{A.4}$$

where $I, J \equiv \phi_I, \phi_J \in S, T_i, Y_k, C_a$ and

$$G_I \equiv \frac{\partial G}{\partial \phi_I} = \frac{W_I}{W} + K_I \tag{A.5}$$

$$F_I = e^{G/2} (K^{-1})_{I\bar{J}}\, G_{\bar{J}} \tag{A.6}$$

where $(K^{-1})_{I\bar{J}}$ is the inverse of the metric $K_{\bar{J}I}$, and satisfies the relation $(K^{-1})_{I\bar{J}} K_{\bar{J}L} = \delta_{IL}$. A subscript on G denotes partial differentiation, while the same subscript on F is just a label. A barred subscript on an F-term denotes its conjugate field $F_{\bar{I}} \equiv (F_I)^\dagger$. We make no distinction between upper and lower indices.

After SUSY breaking, the supersymmetric partner of the Goldstone boson (Goldstino) is *eaten* by the massless gravitino through the super-Higgs mechanism. The gravitino now has a mass given by

$$m_{3/2}^2 = e^{\langle G \rangle} = e^{\langle K \rangle}\, |\langle W \rangle|^2 = \frac{1}{3} \langle F_{\bar{J}}\, K_{\bar{J}I}\, F_I \rangle \tag{A.7}$$

and sets the overall scale of the soft parameters.

In the absence of F-term vacuum expectation values ($\langle F_I \rangle = 0 \;\forall\, \phi_I$), the locally supersymmetric vacuum is negative $V_{\rm SUSY} = -3e^G$. However if one (or more) of the auxiliary F-terms acquires a non-zero vev, the negative vacuum energy can be (partially) cancelled. This raises the exciting possibility that the vacuum energy, or rather the cosmological constant V_0, can be made vanishingly small in agreement with experimental limits. Notice that such a possibility cannot arise in *global* SUSY.

A.2 SUSY breaking F-terms

As previously mentioned, (unknown) non-perturbative contributions to the Kähler function require a parametrisation of our ignorance in terms of Goldstino angles and CP-phases that *control* the relative contributions to SUSY breaking from the various F-terms vevs. We can define a column vector of F-term vevs F in terms of a matrix P and column vector Θ

[10] We will ignore the D-term contribution to the potential that arises from the gauge sector.

(which also includes a CP-phase), where Θ has unit length and satisfies $\Theta^\dagger\Theta = 1$, and P canonically normalises the Kähler metric $P^\dagger K_{\bar{J}I} P = 1$:

$$\begin{aligned} F &= \sqrt{3}\, C\, m_{3/2}\ (P\Theta) \\ F^\dagger &= \sqrt{3}\, C\, m_{3/2}\ \left(\Theta^\dagger P^\dagger\right). \end{aligned} \tag{A.8}$$

Replacing the fields by their vevs, we can rewrite eq. (A.4) as a matrix equation:

$$\begin{aligned} \langle V\rangle \equiv V_0 &= F^\dagger\, K_{\bar{J}I}\, F - 3m_{3/2}^2 \\ &= 3C^2 m_{3/2}^2\, \Theta^\dagger\Theta \left(P^\dagger K_{\bar{J}I} P\right) - 3m_{3/2}^2 \\ &= 3m_{3/2}^2 \left(C^2 - 1\right) \end{aligned} \tag{A.9}$$

where V_0 is the cosmological constant and hence $C^2 = 1 + V_0/3m_{3/2}^2$. Therefore, choosing a vanishingly small cosmological constant sets $C = 1$.

As an example consider the case of the dilaton S and an overall moduli field T with diagonal Kähler metric. The SUGRA potential would be a "sum of squares" $V_F \sim |F_S|^2 + |F_T|^2 + \cdots - 3e^G$ and hence the P-matrix is a diagonal normalising matrix:

$$P_{I\bar{J}} = (K_{I\bar{I}})^{-1/2}\delta_{I\bar{J}} \tag{A.10}$$

In this special case we would recover the expressions of refs. [3, 7, 8]:

$$F \equiv \begin{pmatrix} F_S \\ F_T \end{pmatrix} = \sqrt{3} C\, m_{3/2} \begin{pmatrix} (K_{S\bar{S}})^{-1/2} \sin\theta\, e^{i\alpha_S} \\ (K_{T\bar{T}})^{-1/2} \cos\theta\, e^{i\alpha_T} \end{pmatrix} \tag{A.11}$$

so that dilaton(moduli) dominated SUSY breaking corresponds to $\sin\theta(\cos\theta) = 1$ respectively. However in the more general case, the potential includes terms that mix different F-terms. The action of the P-matrix is to canonically normalise the Kähler metric and maintain the validity of the parametrisation.[11]

A.3 Soft masses and trilinears

Using eqs. (A.2), (A.3) we can write down the *un-normalised* SUSY breaking masses and trilinears that arise in the soft SUGRA potential [3]:

$$V_{\text{soft}} = m_{\bar{a}b}^2 \bar{C}_{\bar{a}} C_b + \left(\frac{1}{6} A_{abc} Y_{abc} C_a C_b C_c + \text{h.c.}\right) + \cdots \tag{A.12}$$

where the Kähler metrics are in general not diagonal leading to the non-canonically normalised soft masses

$$m_{\bar{a}b}^2 = \left(m_{3/2}^2 + V_0\right) \tilde{K}_{\bar{a}b} - F_{\bar{m}} \left(\partial_{\bar{m}}\partial_n \tilde{K}_{\bar{a}b} - \partial_{\bar{m}} \tilde{K}_{\bar{a}c} (\tilde{K}^{-1})_{c\bar{d}} \partial_n \tilde{K}_{\bar{d}b}\right) F_n \tag{A.13}$$

$$\begin{aligned} A_{abc} Y_{abc} = \frac{\hat{W}^*}{|\hat{W}|}\, e^{\bar{K}/2} F_m \Big[&\bar{K}_m Y_{abc} + \partial_m Y_{abc} - \\ &- \left((\tilde{K}^{-1})_{d\bar{e}} \partial_m \tilde{K}_{\bar{e}a} Y_{dbc} + (a \leftrightarrow b) + (a \leftrightarrow c)\right)\Big] \end{aligned} \tag{A.14}$$

[11] The Kähler metric always receives off-diagonal components from the matter fields, but these are conventionally assumed to be small in comparison to the diagonal entries. However, the anomaly cancelling Green-Schwarz term mixes different fields at the same level to introduce off-diagonal components of comparable size.

where the subscript $m = h, C_a$. Notice that a non-diagonal Kähler metric for the matter states will generate a mass matrix between different fields. The physical masses and states are obtained by transforming to the canonically normalised Kähler metric,

$$\tilde{K}_{\bar{a}b}\bar{C}_{\bar{a}}C_b \longrightarrow \bar{C}'_{\bar{a}}C'_a \,. \tag{A.15}$$

The Kähler metric is canonically normalised by a transformation $\tilde{P}^\dagger\tilde{K}\tilde{P} = 1$, so that the physical canonically normalised masses m_a^2 are related to the previous non-canonical mass matrix $m_{\bar{a}b}^2$ by the relation

$$m_a^2 = \tilde{P}^\dagger m_{\bar{a}b}^2 \tilde{P} \,. \tag{A.16}$$

If the Kähler matter metric is diagonal (but not canonical) $\tilde{K}_a = \tilde{K}_{\bar{a}b}\delta_{\bar{a}b}$ then the canonically normalised scalar masses m_a^2 are simply given by

$$m_a^2 = m_{3/2}^2 - F_{\bar{J}}F_I\partial_{\bar{J}}\partial_I\left(\ln\tilde{K}_a\right) \qquad (I, J = h, C_a) \,. \tag{A.17}$$

The soft gaugino mass associated with the gauge group G_α is:

$$M_\alpha = \frac{1}{2\,\mathrm{Re}\,f_\alpha}F_I\partial_I f_\alpha \qquad (I = S, T_i, Y_k) \tag{A.18}$$

and the canonically normalised SUSY breaking trilinear term for the scalar fields $A_{abc}Y_{abc}C_aC_bC_c$ is

$$A_{abc} = F_I\left[\overline{K}_I + \partial_I\ln Y_{abc} - \partial_I\ln\left(\tilde{K}_a\tilde{K}_b\tilde{K}_c\right)\right] . \tag{A.19}$$

B. Alternative exponential suppression factor

In this appendix we consider an alternative suppression factor $e^{-(M_* r)^2}$ that is attributed to non-perturbative world-sheet instanton corrections [15], and differs from the (previous) field theory interpretation of the suppression due to propagating massive modes.

We summarize the modified scalar masses and trilinears found by repeating the earlier calculations in section 4, but with the alternative suppression factor. The non-sequestered scalar masses of eq. (4.1) remain unchanged, but eq. (4.2) becomes:

$$\begin{aligned}
m^2_{C^{5_15_2}} &= \tilde{m}^2 - \frac{3}{2}m_{3/2}^2\left(\sin^2\theta + \Theta_3^2\,\cos^2\theta\,\sin^2\phi\right) \\
m^2_{C^{95_1}} &= \tilde{m}^2 - \frac{3}{2}m_{3/2}^2\cos^2\theta\,\sin^2\phi\left(\frac{\Theta_2^2}{k} + \Theta_3^2\right) \\
m^2_{C_1^{5_1}} &= \tilde{m}^2 - 3m_{3/2}^2\sin^2\theta \\
m^2_{C_2^{5_1}} &= \tilde{m}^2 - 3m_{3/2}^2\Theta_3^2\cos^2\theta\sin^2\phi \\
m^2_{C_3^{5_1}} &= \tilde{m}^2 - \frac{3}{k}m_{3/2}^2\Theta_2^2\cos^2\theta\sin^2\phi
\end{aligned} \tag{B.1}$$

where

$$\begin{aligned}\tilde{m}^2 =\; & m_{3/2}^2\Bigg[1-\cos^2\theta\cos^2\phi\left(1-e^{-(T_2+\bar{T}_2)/4}\right)- \\ & -\frac{\cos^2\theta\sin^2\phi\,\Theta_2^2\,\delta_{GS}}{k}\left(1-e^{-(T_2+\bar{T}_2)/4}\right)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}+ \\ & +\frac{\cos^2\theta\sin^2\phi\,\Theta_2^2\,e^{-(T_2+\bar{T}_2)/4}}{32k}(T_2+\bar{T}_2)^2\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2- \\ & -\frac{\cos^2\theta\cos\phi\sin\phi\,\left(\Theta_2\,e^{i(\alpha_2-\alpha_{Y_2})}+\Theta_2^\dagger\,e^{-i(\alpha_2-\alpha_{Y_2})}\right)e^{-(T_2+\bar{T}_2)/4}}{32\sqrt{k}}\times \\ & \times\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\left(8(T_2+\bar{T}_2)+\delta_{GS}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right)\Bigg]\end{aligned}\tag{B.2}$$

which replaces eq. (4.3); and the masses are expanded up to $\mathcal{O}\left[\frac{\delta_{GS}}{T_2+\bar{T}_2}\right]$.

Similarly we can find modified expressions for the trilinears of eqs. (4.9)–(4.14).

$$\begin{aligned}A_{C_3^{5_2}C^{5_15_2}C^{5_15_2}} =\; & -\sqrt{3}m_{3/2}\,\cos\theta\times\Bigg[\sin\phi\frac{\Theta_2e^{i\alpha_2}}{\sqrt{k}}+ \\ & +\sin\phi\frac{\Theta_2e^{i\alpha_2}}{12\sqrt{k}}e^{-(T_2+\bar{T}_2)/4}(T_2+\bar{T}_2)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2- \\ & -\cos\phi\,\frac{e^{i\alpha_{Y_2}}}{3}\left(1+2\,e^{-(T_2+\bar{T}_2)/4}\right)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\Bigg]\end{aligned}\tag{B.3}$$

$$\begin{aligned}A_{C_1^{9}C^{95_1}C^{95_1}} =\; & -\sqrt{3}m_{3/2}\Bigg[\sin\theta e^{i\alpha_S}+ \\ & +\cos\theta\sin\phi\frac{\Theta_2e^{i\alpha_2}}{12\sqrt{k}}e^{-(T_2+\bar{T}_2)/4}(T_2+\bar{T}_2)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2- \\ & -\cos\theta\cos\phi\frac{e^{i\alpha_{Y_2}}}{3}\left(1+2e^{-(T_2+\bar{T}_2)/4}\right)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\Bigg]\end{aligned}\tag{B.4}$$

$$\begin{aligned}A_{C^{5_15_2}C^{95_1}C^{95_2}} =\; & -\sqrt{3}m_{3/2}\Bigg[\frac{1}{2}\sin\theta\,e^{i\alpha_S}+\frac{1}{2}\cos\theta\sin\phi\left(\Theta_1\,e^{i\alpha_1}+\frac{\Theta_2\,e^{i\alpha_2}}{\sqrt{k}}-\Theta_3\,e^{i\alpha_3}\right)+ \\ & +\cos\theta\sin\phi\frac{\Theta_2e^{i\alpha_2}}{12\sqrt{k}}e^{-(T_2+\bar{T}_2)/4}(T_2+\bar{T}_2)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2- \\ & -\cos\theta\cos\phi\,\frac{e^{i\alpha_{Y_2}}}{3}\left(1+2e^{-(T_2+\bar{T}_2)/4}\right)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\Bigg]\end{aligned}\tag{B.5}$$

$$\begin{aligned}A_{C_3^{5_1}C^{5_15_2}C^{5_15_2}} = A_{C_1^{5_1}C_2^{5_1}C_3^{5_1}} = A_{C_1^{5_1}C^{95_1}C^{95_1}} =\; & -\sqrt{3}m_{3/2}\,\cos\theta\Bigg[\sin\phi\,\Theta_1\,e^{i\alpha_1}+ \\ & +\sin\phi\frac{\Theta_2\,e^{i\alpha_2}}{8\sqrt{k}}\,e^{-(T_2+\bar{T}_2)/4}\,(T_2+\bar{T}_2)\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}^2 \\ & -\cos\phi e^{i\alpha_{Y_2}}e^{-(T_2+\bar{T}_2)/4}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\Bigg]\end{aligned}\tag{B.6}$$

$$A_{C_1^{5_2}C_2^{5_2}C_3^{5_2}} = A_{C_2^{5_2}C^{95_2}C^{95_2}}\tag{B.7}$$

$$
\begin{aligned}
&= -\sqrt{3}m_{3/2}\cos\theta\left[\sin\phi\frac{\Theta_2 e^{i\alpha_2}}{\sqrt{k}} - \cos\phi e^{i\alpha_{Y_2}}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right] \\
A_{C_1^9 C_2^9 C_3^9} &= A_{C_2^9 C^{95_2} C^{95_2}} \qquad\qquad\qquad\qquad\qquad\qquad\qquad\qquad\qquad\qquad \text{(B.8)} \\
&= -\sqrt{3}m_{3/2}\left[\sin\theta e^{i\alpha_S} - \cos\theta\cos\phi e^{i\alpha_{Y_2}}\left\{Y_2+\bar{Y}_2-\delta_{GS}\ln(T_2+\bar{T}_2)\right\}\right]
\end{aligned}
$$

where all trilinears have been expanded up to $\mathcal{O}(\frac{1}{T_2+\bar{T}_2})$.

It is now straightforward to consider these modified expressions in different limits of SUSY breaking domination as before.

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