Brane-to-brane mediation of supersymmetry breaking in presence of D-type breaking

Thomas Grégoire
Physics department
Boston University
Boston MA, 02215

1. Introduction

As it is well known, supersymmetry is a very attractive option for physics beyond the Standard Model (SM). While it has not been discovered at LEP and the Tevatron, it will still be extensively looked for at the LHC, and many remain hopeful that it will be seen there. An important feature of supersymmetry is that it guarantees the absence of quadratic divergence in the renormalization of scalar masses, in the case of interest the Higgs mass. This property holds even if supersymmetry is broken softly, that is by operators of dimension two or three. The minimal extension of the SM into such a theory is called the minimal supersymmetric standard model (MSSM). It contains renormalizable supersymmetric interactions as well as a set of super-renormalizable, supersymmetry-breaking interactions dubbed the soft masses:

\[
L = \mathcal{L}_{\text{kin}} + \int d^2\theta \left[ y_{ij}^u Q_i U_j H_u^a + y_{ij}^d Q_i D_j H_d^a + y_{ij}^e L_i E_j H_d^a + \mu H_u^a H_d^a \right] + m_{ij}^2 q_i q_j + m_{1/2}^a \lambda^a \lambda^a + y_{ij}^u A_u q_i u_j h^u + B \mu h^u h^d + \cdots,
\]

where \( m_{ij}^2 \) are scalar soft masses squared, \( m_{1/2}^a \) gaugino masses and \( A \) trilinear scalar terms. The dots denote similar terms.

The value of these soft masses are very constrained by experiment. One constraint comes from the fact that supersymmetry has not yet been discovered, which requires the superpartners to be somewhat heavy. This typically re-introduce fine-tuning in the Higgs potential but we will not attempt to address this issue here. The most severe constraints however, come from the flavor structure of the soft masses; because observed flavor-changing neutral currents are small, the soft masses need to be either flavor-universal or aligned with the Yukawa matrices. From a low energy point of view this is not guaranteed at all, and it warrants a mechanism of supersymmetry breaking for which this happens naturally.

In general, if supersymmetry breaking occurs at high energy, even if the supersymmetry breaking transmission mechanism is flavor blind, flavor physics, responsible for generating the low energy Yukawa couplings, could leave its imprints on the soft masses and make them non-universal. Therefore a full theory of flavor would be needed to determine the flavor structure of the soft masses.

Gravity mediated supersymmetry breaking\cite{2} is of this type. In this framework, the theory is non-supersymmetric up to the highest scales of the theory. The soft masses are therefore

\footnote{Based on \cite{1}, done with Riccardo Rattazzi and Claudio Scrucca.}
sensitive to the UV, and nothing can be said of their flavor structure without a full theory of flavor.

The situation is very different if the soft masses are generated in the IR. In this case one can imagine that the flavor structure of the Standard Model Yukawa matrices is determined in the UV, when the theory is supersymmetric. When supersymmetry breaks, flavor physics has already decoupled, and a flavor blind mediation of supersymmetry would guarantee ‘flavor-safety’ at low energy. This is the case for example in gauge mediation[3], gaugino mediation[4] or anomaly mediation[5, 6]. In this talk we will concentrate on anomaly mediation.

2. Anomaly Mediation and Sequestering

Anomaly mediation of supersymmetry breaking [5, 6] is a supergravity effect that relates supersymmetry breaking in the visible sector to breaking of the conformal anomaly. The main characteristic of anomaly mediation is that while supersymmetry is broken at all scales, the soft masses at one particular scale depend only on the gauge and Yukawa coupling at that scale. Therefore soft masses are completely insensitive to flavor physics (or almost any physics for that matter) going on at a high scale, and the only flavor spurion that can enter the soft masses is given by the SM Yukawa matrices.

To understand anomaly mediation, it is useful to use the conformal formalism of supergravity[7]. In this formalism, the full superconformal group is restored by introducing an redundant chiral superfield \( \phi \) called the conformal compensator. This field is non-dynamical and we will be mostly interested by its \( F \) component:

\[
\phi = 1 + F_\phi \theta^2
\]

Its usefulness lies in the fact that it parametrizes departure from scale invariance. If the theory is scale invariant at the classical level, \( \phi \) can be eliminated from the Lagrangian by a field redefinition. However, because in general the quantum theory is not scale invariant, \( \phi \) reappears in the quantum effective action. For example, the Lagrangian for a set of chiral superfields \( X \) and gauge superfield \( V \) is:

\[
\mathcal{L} = \int d^4\theta \phi \phi^{\dagger} X^\dagger e^V \mathcal{X} + \int d^2\theta \left[ \frac{1}{g^2} W^\alpha W^\alpha + \phi^3 W(X) \right].
\]

If the superpotential \( W(X) \) is cubic in \( X \), the compensator can be removed by a field redefinition \( X \rightarrow \phi X \). However, if this is the Lagrangian at some scale \( \Lambda \), then at the scale \( \mu \) the lagrangian will be:

\[
\int d^4\theta \mathcal{Z} \left( \frac{\mu}{\Lambda \sqrt{\phi^3 \phi}} \right) X^\dagger e^V \mathcal{X} + \int d^2\theta \left[ \frac{1}{g^2} \left( \frac{\mu}{\Lambda \phi} \right) W^\alpha W^\alpha + W(X) \right],
\]

where \( \mathcal{Z} \) includes the wave function renormalization of \( X \), and \( g^2(\mu/(\Lambda \phi)) \) is the holomorphic running gauge coupling constant. After replacing \( \phi \) with its vev, one gets the following expressions for the soft masses [5]:

\[
m_0^2 = -\frac{1}{4} \left( \frac{\partial \gamma}{\partial y} \beta_y + \frac{\partial \gamma}{\partial y} \beta_y \right) |F_\phi|^2
\]

\[
m_{1/2} = \frac{\beta_y}{g} F_\phi
\]

\[
A_{ijk} = \frac{1}{2} (\gamma_i + \gamma_j + \gamma_k) y_{ijk} F_\phi
\]

where \( \gamma \) is the anomalous dimension while \( \beta_y \) and \( \beta_y \) are respectively the beta function for the gauge and Yukawa coupling constant. We see that gaugino masses and trilinear \( A \) terms
appear at one loop, while scalar masses squared appear at two loops and are therefore of the same order. As mentioned before, these expressions are valid at any scale, and are therefore UV insensitive. There are however two obvious problems with anomaly mediation. First the sleptons have negative mass squared. Second, in general the contributions of anomaly mediation are suppressed with respect to those of gravity mediation. In the presence of a hidden sector with chiral superfields $X$ that break supersymmetry, the natural value for the $F$ component of the conformal compensator is:

$$F_\phi \sim \frac{F_X}{M_{pl}}$$

which leads to soft mass terms of order:

$$m_{\text{anomaly}} \sim \frac{g^2 F_X}{16\pi^2 M_{pl}}.$$  

However, gravity loops will generate operators that link the visible sector fields ($Q$) with the SUSY breaking sector:

$$\frac{Q^\dagger Q X^\dagger X}{M_{pl}^2}$$

which leads to gravity mediated soft mass terms of order:

$$m_{\text{gravity}} \sim \frac{F_X}{M_{pl}}.$$  

Therefore, the UV sensitive gravity mediated piece dominates over the UV insensitive anomaly-mediated one.

To remedy this problem one can separate the visible and hidden (SUSY breaking) sector across an extra dimension of space (this is often referred to as sequestering) [5]. In this setup, the visible sector lives on one brane and the hidden sector which breaks supersymmetry lives on another one while gravity (supergravity in fact) lives in the bulk. Here, we will argue that in the presence of D-type breaking in the hidden sector, this scenario can also remedy the first problem of tachyonic sleptons.

In a sequestered scenario anomaly mediation still contributes to the soft masses in the same way. Anomaly mediation is an infrared mechanism, and can be computed in the low energy 4D theory. However, the nature of the gravity mediated part is significantly modified in two ways.

First of all, the size of the contribution mediated by gravity is modified. And second, this contribution is now generated in the IR, at the scale of the radius of the extra-dimension: $1/R$.

An order of magnitude estimate for the operator that connects the visible and hidden sector and which generated by loops of supergravity in the extra-dimension is given by:

$$\frac{1}{16\pi^2} \frac{Q^\dagger Q X^\dagger X}{M_5^9 (T + T^\dagger)^4},$$

where $M_5$ is the 5D Planck mass, and $T$ is the radion superfield: $<T> = R + F_T \theta^2$. We can rewrite this contribution in term of the 4D Planck scale as:

$$\frac{1}{16\pi^2} \frac{Q^\dagger Q X^\dagger X}{M_{pl}^4} \left(\frac{1}{T + T^\dagger}\right)^2.$$  

From a low energy point of view this can be interpreted as coming from a quadratically divergent 4D graviton loop cut-off at the scale of the radius. Another operator generated by bulk gravity loops is:

$$\frac{1}{16\pi^2} \frac{Q^\dagger Q}{M_5^3 (T + T^\dagger)^3}.$$
This operator gives soft masses to the visible sector when $T$ acquire a vacuum expectation value in its $F$ component. We refer to this contribution as coming from 'radion mediation'.

From equation (12), we see that in a sequestered model, the gravity mediated contribution depends on the value of the radius. By comparing equation (9) and equation (12) keeping in mind that $F_\phi \sim F_\chi/M_{pl}$, we find that if it is of order:

$$R^2 \sim \left(\frac{16\pi^2}{g^2}\right)^2 \frac{1}{16\pi^2 M_{pl}^2},$$

the gravity mediated and anomaly mediated contributions are of the same order. Of course, for larger radius the gravity mediated contribution is smaller. The size of the radion-mediated contribution depends of the size of $F_T$, and as we will see, in our case it will be subdominant.

2.1. Gravity mediation in sequestered model

As we just saw, in a sequestered model, gravity mediated contributions to soft masses (as well as radion mediated contributions) are IR dominated, therefore calculable, and can compete with anomaly mediation. They are therefore a good candidate to solve the problem of tachyonic sleptons of anomaly mediation. We thus need to calculate 5D supergravity loops to know the exact coefficient of the operator (12) and in particular its sign. This is a non-trivial task, and has been done both using a superfield technique [8] and directly in components [9]. It was also done for a 5D ADS space using superfields in [10][see also [11]]. In the simplest cases, it was realized that both gravity mediated and radion mediated contributions to soft masses are negative. However it was pointed out by [9] (and similarly in [10] in the warped case) that the radion-mediated piece can be made positive in the presence of large kinetic term for gravity on one of the branes.

Here, we will only briefly sketch the superfield technique and quote the result. In five dimensions, supersymmetry is different than in four dimensions, because 5D spinors and therefore the SUSY generators have four components, while in 4D there are two components spinors. Therefore the particle content is enlarged, and the usual four dimensional superspace and superfields cannot be used to make the full 5D supersymmetry manifest. However, it is still possible, and useful, to us them to make a $N=1$ subgroup of the full supersymmetry manifest and write down supersymmetric bulk-brane couplings. In the case at hand, we only need a linearized version of 5D supergravity which can be written in term of superfields as [12]:

$$L = M_5^3 \int d^4\theta \left\{ \frac{1}{2} V^m K_{mn} V^n - \frac{1}{3} \Sigma^i \Sigma + \frac{2i}{3} \left( \Sigma - \Sigma^\dagger \right) \partial^m V_m - \frac{1}{2} \left[ \partial_y V_{a\dot{a}} - (\bar{D}_{a} \Psi_{a} - D_{a} \bar{\Psi}_{a}) \right]^2 + \frac{1}{4} \left[ \partial_y P_{\Sigma} - (D^a \Psi_{a} + \bar{D}_{a} \bar{\Psi}_{a}) \right]^2 - \frac{1}{2} \left[ T^\dagger (\Sigma + 2i \partial_m V^m) + \text{h.c.} \right] \right\},$$

where the real superfield $V_m$ contains the graviton, the superfield $\Psi_a$ the gravi-photon, the chiral superfield $T$ is the linearized version of $T$, $\Sigma$ the linearized version of $\phi$, $P_\Sigma$ a pre-potential for $\Sigma$, and $K_{mn}$ is a complicated two derivatives operator which we do not bother to write here (for more details see [12, 8, 10]).

Couplings to a chiral superfield $Q$ on the brane are just the usual $N=1, D=4$ supergravity couplings of a chiral superfield:

$$L_i = \int d^4\theta \left[ Q^\dagger Q \left( \frac{1}{3} \Sigma^\dagger \right) \left( 1 + \frac{1}{3} \Sigma \right) + \frac{2}{3} i Q^\dagger \partial_m Q V^m - \frac{1}{6} Q^\dagger Q V^m K_{mn} V_n \right].$$
One can then calculate loops directly in superspace to show that the following operator is generated:

$$L \supset \frac{\xi(3)}{6\pi^2} \frac{Q^\dagger QX^\dagger X}{M_5^6(T + T^\dagger)^4}$$

Which, as we mentioned gives a negative mass squared for the sfermions.

3. D-Term breaking

The situation is very different if there is D-type breaking in the hidden sector. If the hidden sector has chiral superfields that are charged under some gauge group:

$$L_{\text{hidden}} = \int d^4 \theta X^\dagger e^V X + \int d^2 \theta [W_\alpha W^\alpha + W(X)],$$

then the operator generated by supergravity loops will be:

$$\frac{\xi(3)}{6\pi^2} \int d^4 \theta \frac{Q^\dagger QX^\dagger e^V X}{M_5^6(T + T^\dagger)^4},$$

and when X and V are replaced by their vev: $<X> = x + F_X \theta^2$, $<V> = \theta^2 \bar{\theta}^2 D$, we get:

$$m_0^2 = -|F_X|^2 + \frac{x^\dagger T^{(a)} D_a x}{M_5^6 R^4},$$

where $T^{(a)}$ are the hidden gauge group generators. Given the fact that $D_a = -1/2 x^\dagger T^{(a)} x$, we get:

$$m_0^2 = -|F_X|^2 + D^2/2,$$

and we see that if $D/2 > F$, the soft masses can be made positive.

There is a well known theorem (see [13]) that states that when a solution can found to the $F = 0$ equations, there is a solution to the $D = 0$ equations as well and supersymmetry is unbroken. In view of this, it is natural to ask how easy it is to have $D > F$. In particular, is it possible to have a D term that is parametrically bigger than the F term? After all, if there is a solution to the equations of motion where $F$ is small, doesn’t the aforementioned theorem guarantees that there is also a solution with small $D$? It is possible to show [1] that if there are large charges in the problem, this is in fact not the case, and it is possible to have a $D$ term parametrically larger than the $F$ term.

For example, the following superpotential of four chiral superfields:

$$W = \lambda_1 \phi_0 \left( \phi_1^{N/\sqrt{N} - 1} \right) + \lambda_2 \phi_1 \phi_{-1},$$

where the subscripts denote the charges of the different fields, results in $D^2/F^2 \sim N$. Of course, for our purpose we don’t need a parametric separation between $D$ and $F$, but this shows that it might even be generic to have $D > F$. It could also happen in models with dynamical supersymmetry breaking such as the $4 - 1$ model described in [14, 15].

4. Radius Stabilization

In order for our scenario to work, we need a radius of the right size, so that anomaly mediation and gravity mediation are comparable while radion mediation is smaller. In this section we review the radius stabilization mechanism presented in [16], and argue that it fits our needs perfectly. The setup is quite simple, and consists of two super-Yang-Mills sectors(SYM), one in
the bulk and one on the brane. These two SYM sectors become strongly coupled and generate gaugino condensations at low energies. The 4D effective superpotential can then be written as:

\[ W = \frac{1}{16\pi^2} \left( \Lambda_1^3 + \Lambda_2^3 e^{-a\Lambda_2/T} \right), \]  

where \( \Lambda_1 \) and \( \Lambda_2 \) are the strong coupling scales of respectively the boundary and bulk gauge interactions. Minimization of the potential gives the following relations between the different fields of the theory:

\[ \Lambda_2 T \simeq 3 \ln(\Lambda_2/\Lambda_1) \]  

\[ F_\phi \simeq \frac{\Lambda_1^3}{16\pi^2 M_5^2 T} \]  

\[ \frac{F_T}{T} \simeq \frac{m_{3/2}}{\Lambda_2 T}. \]  

We see that if \( \Lambda_2 T \gg 1 \), then we have \( F_T/T \ll F_\phi \), and radion mediation is subdominant. By further making the hypothesis that the strong coupling scales of the bulk SYM theory and of the bulk gravity are the same: \( \Lambda_2 \sim \Lambda_5 \), we get \( \alpha_5 = 1/(\Lambda_5 T)^3 \equiv 1/(16\pi^2 (M_5 T)^3) \sim 10^{-4} \), which is of the right order to make anomaly mediation and gravity mediation contributions to soft masses comparable.

5. Flavor violation

The goal of this section is to provide order of magnitude estimates for the size of flavor violations in the scenario we have presented. The first source of flavor violation comes from unknown flavor physics at a high scale. For example we could imagine that there are higher-dimensional operators suppressed by \( \Lambda_5 \) that are not flavor diagonal (in the basis where the SM Yukawa matrices have the correct values). For example an operator of the form [1]:

\[ \int d^4\theta Z_{ij} Q_i^\dagger Q_j (K_{mn} V^n)^2 \Lambda_5 \]  

where \( K_{mn} \) is a two-derivative operator could be present. If inserted in a loop the connects to the hidden sector and that is saturated at scale \( 1/R \), it would produce non-diagonal soft masses of order:

\[ \frac{\delta m_{ij}}{m_0} \sim \frac{1}{\Lambda_2^2 R^2} \sim 10^{-3}. \]  

Bounds on \( \epsilon_K \) would then require squarks of order one TeV. We could also have lighter squarks if we assume that the flavor violating factors \( Z_{ij} \) have a suppression of the order of the Cabibbo angle.

Other possible flavor violations can occur if for example our theory is embedded in a GUT [17, 18]. Since our soft masses are generated at \( 1/R \sim 10^{17}\text{GeV} > M_{\text{GUT}} \), there is additional flavor physics between the scale where we know our soft masses are diagonal, and the weak scale. This additional flavor physics is necessary because, in a GUT, leptons and quarks are embedded in a single multiplet. Therefore above the GUT scale, the Yukawa matrices are not the one of the SM. In particular, the third generation of sleptons have a large Yukawa coupling since they are unified with the top quark. In \( SO(10) \) for example, we can estimate the split between the third and the first two generations of sleptons to be of order:

\[ \frac{\delta m_3^2}{m^2} \sim \frac{15}{8\pi^2} \lambda_t^2 \ln(M_{\text{GUT}}) \sim 0.3 \left( \frac{\lambda_t}{0.8} \right)^2. \]
where $\lambda_t$ is the top Yukawa coupling at the GUT scale. This leads to an estimate for $\text{BR}(\mu \rightarrow e\gamma)$ of order:

$$5 \times 10^{-13} \left( \frac{\lambda_t}{0.8} \right)^4 \left( \frac{150 \text{GeV}}{m} \right)^4,$$

where $m$ is the average slepton mass. This is below the current experimental bound, but above the reach of planned experiment.

A similar effect is present if the small neutrino masses are explained by the see-saw mechanism, which would involve new Yukawa matrices at a high scale [19, 20]. In this case however the size of these Yukawa matrices are totally unknown, and therefore the size of this effect difficult to estimate. If the Yukawa couplings are large, and the mixing angles are of order one the effect could be big and contradict present experiments. However, if either these neutrino Yukawa are small, or if the mixing angles are CKM-like, the effect will be within the present bounds.

6. Conclusion

In this talk, we have shown that in the context of a sequestered model of supersymmetry breaking, where the hidden sector that breaks SUSY and the visible sector live on two different branes separated by a bulk where gravity propagate it is possible to get phenomenologically acceptable soft masses. It requires D-type breaking in the hidden sector in order for gravity loops to produce positive soft masses that can cure the tachyonic slepton problem of anomaly mediation. In this scenario, gaugino masses and trilinear soft terms are given by anomaly mediation, while the scalar soft masses squared get a universal positive contribution in addition to the anomaly mediated contribution.

We have argued that the radius stabilization mechanism of [16] is adequate for this scenario and quite naturally yields a radius that is of the right size for the gravity mediated and the anomaly mediated contributions to be of the same order. We have also discussed possible deviations from flavor universality and shown that they are under control.

An issue that still need to be addressed in this scenario is the generation of a $\mu$ and $B\mu$ term of the right size.

Acknowledgments

This is based on work [1] done in collaboration with Riccardo Rattazzi and Claudio Scrucca while I was at CERN. I would also like to thank the organizer of the Corfu 2005 Summer Institute for giving me the opportunity to give a talk at this nice conference.

7. References

[1] T. Gregoire, R. Rattazzi and C. A. Scrucca, Phys. Lett. B 624, 260 (2005) [arXiv:hep-ph/0505126].
[2] A. H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. 49, 970 (1982).
[3] R. Barbieri, S. Ferrara and C. A. Savoy, Phys. Lett. B 119, 343 (1982). L. J. Hall, J. D. Lykken and S. Weinberg, Phys. Rev. D 27, 2359 (1983). N. Ohta, Prog. Theor. Phys. 70, 542 (1983).
[4] H. P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B 120, 346 (1983). L. E. Ibanez, Phys. Lett. B 118, 73 (1982).
[5] N. Ohta, Prog. Theor. Phys. 70, 542 (1983).
[6] M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D 51, 1362 (1995) [arXiv:hep-ph/9408384].
[7] D. E. Kaplan, G. D. Kribs and M. Schmaltz, Phys. Rev. D 62, 035010 (2000) [arXiv:hep-ph/9911293].
[8] Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, JHEP 0001, 003 (2000) [arXiv:hep-ph/9911323].
[9] L. Randall and R. Sundrum, Nucl. Phys. B 557, 79 (1999) [arXiv:hep-th/9810155].
[10] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998) [arXiv:hep-ph/9810442].
[11] W. Siegel and S. J. J. Gates, Nucl. Phys. B 147, 77 (1979). S. J. Gates, M. T. Grisaru, M. Rocek and W. Siegel, Front. Phys. 58, 1 (1983) [arXiv:hep-th/0101200].
[12] I. L. Buchbinder, S. J. J. Gates, H. S. J. Goh, W. D. I. Linch, M. A. Luty, S. P. Ng and J. Phillips, Phys. Rev. D 70, 025008 (2004) [arXiv:hep-th/0305169].
[13] R. Rattazzi, C. A. Scrucca and A. Strumia, Nucl. Phys. B 674, 171 (2003) [arXiv:hep-th/0305184].
[10] T. Gregoire, R. Rattazzi, C. A. Scrucca, A. Strumia and E. Trincherini, Nucl. Phys. B 720, 3 (2005) [arXiv:hep-th/0411216].
[11] A. Falkowski, JHEP 0505, 073 (2005) [arXiv:hep-th/0502072].
[12] W. D. I. Linch, M. A. Luty and J. Phillips, Phys. Rev. D 68, 025008 (2003) [arXiv:hep-th/0209060].
[13] M. A. Luty and W. I. Taylor, Phys. Rev. D 53, 3399 (1996) [arXiv:hep-th/9506098].
[14] M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D 53, 2658 (1996) [arXiv:hep-ph/9507378].
[15] L. Carpenter, P. J. Fox and D. E. Kaplan, arXiv:hep-ph/0503093.
[16] M. A. Luty and R. Sundrum, Phys. Rev. D 62, 035008 (2000) [arXiv:hep-th/9910202].
[17] L. J. Hall, V. A. Kostelecky and S. Raby, Nucl. Phys. B 267, 415 (1986).
[18] R. Barbieri, L. J. Hall and A. Strumia, Nucl. Phys. B 445, 219 (1995) [arXiv:hep-ph/9501334].
[19] F. Borzumati and A. Masiero, Phys. Rev. Lett. 57, 961 (1986).
[20] A. Masiero, S. K. Vempati and O. Vives, New J. Phys. 6, 202 (2004) [arXiv:hep-ph/0407325].