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Sparticle masses from the superconformal anomaly

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Abstract: We discuss a recently proposed scenario where the sparticle masses are purely mediated by gravity through the superconformal anomaly. This scenario elegantly evades the supersymmetric flavor problem since soft masses, like the anomaly, are not directly sensitive to ultraviolet physics. However, its minimal incarnation fails by predicting tachyonic sleptons. We study the conditions for decoupling of heavy threshold effects and how these conditions are evaded. We use these results to build a realistic class of models where the non-decoupling effects of ultra-heavy vectorlike matter fields eliminate the tachyons. These models have a flavor invariant superspectrum similar to that of gauge mediated models. They, however, differ in several aspects: the gaugino masses are not unified, the colored sparticles are not much heavier than the others, the μ problem is less severe and the gravitino mass is well above the weak scale, $m_{3/2} \gtrsim 10$ TeV. We also show that in models where an R-symmetry can be gauged, the associated D-term gives rise to soft terms that are similarly insensitive to the ultraviolet.

Keywords: Supersymmetry Breaking, Supergravity Models, Supersymmetric Standard Model.

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Contents

1. Introduction

Supersymmetry is at the moment the best candidate for physics beyond the standard model (SM). It gives a natural explanation of the gauge hierarchy and its simplest realization, the minimal supersymmetric standard model (MSSM), remarkably satisfies the unification of the gauge forces. However, superparticles have not been detected yet, nor any deviation from the SM has been seen in precision experiments and flavor physics. This means that not only supersymmetry must be broken to give sparticles a mass, but also that the resulting spectrum will most likely display a clever flavor structure. The origin of supersymmetry breaking is thus crucial for phenomenology.

The simplest way to give superparticles a mass is certainly via non-renormalizable interactions in supergravity [1], as the basic ingredient, gravity, is a fact of life. Moreover gravity mediation is generic as its agent (or the underlying more fundamental interaction) couples to everything: whatever dynamics breaks supersymmetry, gravity will eventually let us know. The flipped side of the coin is that generic soft terms are not enough. They must also be very specific to avoid large flavor violations. It is possible that this problem will only be fully understood in a theory of quantum gravity such as string theory. A possibility though is that some underlying flavor symmetry solves the problem by aligning the soft terms to the fermion masses [2].

Models with gauge mediated supersymmetry breaking (GMSB) [3, 4], on the other hand, offer a calculable and flavor invariant superspectrum. These two properties motivated the recent activity on this scenario.

In this paper we focus on a particular supergravity scenario [5] in which the soft masses are generated just by the auxiliary field of the gravitational multiplet. This differs from previous ones [1] since no direct contribution arises from the hidden sector. In this sense, soft masses are *purely* mediated by gravity. This possibility was advocated in ref. [5] in theories with extra dimensions where only gravity propagates. Moreover the same results, but limited to gaugino masses and A-terms, are inevitably obtained in any model where the hidden sector breaks supersymmetry dynamically in the absence of singlets [6]. From the supergravity point of view, this idea corresponds to a no-scale form for the Kahler potential [7] and to canonical tree-level gauge kinetic terms. Both scalar and gaugino masses are zero at tree level and are purely determined by quantum effects. It is easy to show that these quantum effects are just dictated by the (super)conformal anomaly [5, 6], so that we can rightfully name this scenario anomaly mediated supersymmetry breaking (AMSB). As anomalies only depend on the low-energy effective theory, so will soft terms. Soft masses at a scale μ will be written as functions of the couplings at μ, with no additional ultraviolet (UV) dependence. This property is very interesting for the supersymmetric flavor problem. In general soft masses can pick up all sort of dangerous contributions from intermediate thresholds, as they flow to the infrared (IR). AMSB, however, provides special trajectories where all these potentially dangerous effects manage to disappear at low energy. Unfortunately the AMSB scenario cannot be applied to the MSSM, as it leads to tachyonic sleptons [5]. The goal of the present paper is to construct realistic models (without tachyons) where the original source of soft terms is the conformal anomaly. Our basic point is that an intermediate threshold can displace soft terms from the AMSB trajectory, provided that the threshold position is controlled by a light field (modulus). Using this remark we will build models where the intermediate threshold is provided by a messenger sector similar to that in GMSB model. The interesting thing is that, unlike GMSB, the present mechanism works even in the absence of tree-level mass splitting inside the messenger supermultiplets! For this reason we will call this mechanism *anti*-gauge-mediation.

This paper is organized as follows. In section 2 we review AMSB, relying mostly on a 1PI definition of the running soft terms. Section 3 focuses on decoupling and non-decoupling of heavy thresholds and on the problem with tachyon states. Sections 4 and 5 are devoted to the construction of explicit models and to a brief discussion of their phenomenology, while section 6 provides examples of solutions to the μ-problem. Section 7 is somewhat outside the main line of the paper. There we show that there is a more general family of non-trivial RG trajectories which is UV insensitive. It is interesting that this class can be determined by internal consistency in the flat limit or simply by gauging an R-symmetry in supergravity. Finally section 8 contains our conclusions.

2. Anomaly mediated supersymmetry breaking

The crucial feature of the scenario we consider is that supersymmetry breaking effects in the observable sector have a pure gravitational origin. This is equivalent to saying that the only relevant source of soft terms is the auxiliary field u of the off-shell gravitational supermultiplet for Poincaré supergravity. The natural formalism to describe these effects is the superconformal calculus formulation of supergravity [8]. In this framework one introduces a chiral superfield ϕ with Weyl weight $\lambda = 1$ playing the role of the compensating multiplet for super-Weyl transformations. With the aid of ϕ it is relatively easy to write a locally superconformal invariant lagrangean. Poincaré supergravity is then recovered by fixing the extraneous degrees of freedom by a suitable set of gauge conditions. In particular one can make the auxiliary field u to reside only in ϕ by fixing $\phi = 1 + \theta^2 u/3$. (We will later comment on different choices of super-Weyl-Kahler gauge fixing). Then the effects of $\langle u\rangle/3 = F_\phi \neq 0$ on any operator are simply determined by inserting the suitable powers of $\phi, \phi^\dagger$ that render the operator Weyl invariant. It is convenient to assign to each physical field a Weyl weight equal to its mass dimension. Then the most general Weyl invariant action has the form

$$\int d^4\theta \phi\phi^\dagger \mathcal{K}\left(\frac{\phi^{1/2}}{\phi^\dagger}D_\alpha, \frac{Q}{\phi}, \frac{W_\alpha}{\phi^{3/2}}, V\right) + \mathrm{Re}\int d^2\theta \phi^3 \mathcal{W}\left(\frac{Q}{\phi}, \frac{W_\alpha}{\phi^{3/2}}\right) , \tag{2.1}$$

where by Q and W_α we collectively indicate the matter and gauge field strength chiral superfields. For simplicity we have omitted the dependence of $\mathcal{K}$ on $\bar{D}_{\dot\alpha}$, $Q^\dagger$, etc.

A few comments on eq. (2.1) are in order. We have omitted all terms involving covariant derivatives acting on ϕ, $\phi^\dagger$ (curvature terms). This is because we are only concerned with the RG evolution of soft terms. By simple power counting, terms involving $D_\alpha\phi$, $D^2\phi$, etc. cannot arise from ultraviolet divergences. This fact was already stressed in ref. [9]. Indeed one could define the soft terms at a scale μ by considering the superspace 1PI action at external Euclidean momenta $p = \mu \gg m_{3/2} \sim F_\phi$ but truncated to terms with no covariant derivatives acting on ϕ. The neglected derivative terms correspond to the contribution from virtual momenta in the infrared domain $m_{3/2} \lesssim p \lesssim \mu$.[1]

One can easily check that the rule $D_\alpha \to D_\alpha \phi^{1/2}/\phi^\dagger$, applied to $W_\alpha = \bar{D}^2 D_\alpha V$ does give $W_\alpha \to W_\alpha/\phi^{3/2}$. Moreover, notice that the D'Alembertian is essentially $\Box \sim D^2\bar{D}^2$, so that its Weyl covariant version is $\Box/\phi\phi^\dagger$. This property of $\Box = -p^2$ crucially relates the soft terms to RG beta functions [5, 6]. In order to elucidate this relation we recall a few results. As discussed in ref. [9], in softly broken supersymmetry, the soft terms associated to a chiral superfield Q_i can be collected in a running superfield

[1]Notice that in component notation the 1PI is less useful to distinguish UV and IR contributions. For instance, in the softly broken Wess-Zumino model, the IR cut-off of the tadpole diagram correction to the scalar mass is the soft mass itself, not μ. In superspace language virtual momenta above and below μ contribute to the coefficient of different operators.

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wave function $\mathcal{Z}_i(\mu)$ such that

$$\ln \mathcal{Z}_i(\mu) = \ln Z_i(\mu) + \left(A_i(\mu)\theta^2 + \text{h.c.}\right) - m_i^2(\mu)\theta^2\bar{\theta}^2 \,. \tag{2.2}$$

Here, Z_i is the c-number wave function, m_i^2 is the sfermion mass, while A_i contributes to A-terms (for instance, a superpotential term $\lambda Q_1 Q_2 Q_3$ is associated to $A_\lambda = A_1 + A_2 + A_3$). Now, consider the present theory in the supersymmetric limit. The running wave functions can be defined as $Z_i(\mu) = c_i(p^2 = -\mu^2)$, where c_i is the coefficient of $Q_i Q_i^\dagger$ in the 1PI. Therefore, turning on F_ϕ simply amounts to the shift $\mu^2 \to \mu^2/\phi\phi^\dagger$

$$\mathcal{Z}_i(\mu) = Z_i\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}\right) . \tag{2.3}$$

By using $\phi = 1 + F_\phi\theta^2$, we can write eq. (2.3) as

$$\ln \mathcal{Z}_i(\mu) = \ln Z_i(\mu) - \frac{\gamma_i(\mu)}{2}\left(F_\phi\theta^2 + \text{h.c.}\right) + \frac{\dot{\gamma}_i(\mu)}{4}|F_\phi|^2\theta^2\bar{\theta}^2 \,, \tag{2.4}$$

where γ_i is the anomalous dimension and $\dot{\gamma}_i = d\gamma_i/d\ln\mu$. Comparing eq. (2.4) to eq. (2.2), we obtain

$$\begin{aligned} A_i(\mu) &= -\frac{\gamma_i(\mu)}{2} F_\phi \,, \\ m_i^2(\mu) &= -\frac{\dot{\gamma}_i(\mu)}{4}|F_\phi|^2 \,. \end{aligned} \tag{2.5}$$

We emphasize that this result is a simple consequence of the 1PI definition of the soft terms and of the Weyl-covariantization property of $\Box$. This makes it clear that it does not depend on the regulator. If we had worked in dimensional reduction (DRED) [10], we would have gotten the same result. This is because, in the bare lagrangean, Weyl invariance requires the scale μ to always appear in the combination $\mu/\sqrt{\phi\phi^\dagger}$. Again, soft terms would be obtained as in eq. (2.3), by deforming the RG flow into superspace.

For gauge fields the relevant quantity is also a real superfield, $R(\mu)$, with components [9]

$$R(\mu) = \frac{1}{g^2(\mu)} - 2\text{Re}\left(\frac{m_\lambda}{g^2}(\mu)\theta^2\right) + R_D\theta^2\bar{\theta}^2 \,, \tag{2.6}$$

where m_λ is the gaugino mass and where at lowest order in g^2

$$R_D = \frac{1}{8\pi^2}\left(T_G m_\lambda^2 - \sum_i T_i m_i^2\right) , \tag{2.7}$$

where T_G and T_i are the Dynkin indices of the adjoint and matter irreducible representations. The real superfield $R(\mu)$ corresponds to the physical 1PI coupling and is related to the holomorphic coupling S by [11, 12, 13]

$$R = F\left(S + S^\dagger - \frac{1}{8\pi^2}\sum_i T_i \ln \mathcal{Z}_i\right) = F(\tilde{R}) \,, \tag{2.8}$$

where F^{-1} can be expanded in a given scheme as

$$F^{-1}(x) = 1 - T_G \ln \frac{x}{8\pi^2} + \sum_{n>0} \frac{a_n}{x^n} \,. \tag{2.9}$$

The NSVZ scheme [11] corresponds to setting all $a_i = 0$.

Similarly to the matter case, one starts from the kinetic coefficient $1/g^2(\mu)$ in the supersymmetric limit and turns on F_ϕ by the shift $\mu^2 \to \mu^2/\phi\phi^\dagger$:

$$R(\mu) = g^{-2}\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}\right) \,. \tag{2.10}$$

Therefore, by using eq. (2.6), we have

$$m_\lambda = \frac{g^2}{2}\frac{dg^{-2}}{d\ln\mu}F_\phi = -\frac{\beta(g^2)}{2g^2}F_\phi \,. \tag{2.11}$$

It is also useful to consider the holomorphic coupling S, which, being a chiral superfield, is defined by the shift

$$S(\mu) = \frac{1}{g_h^2(\frac{\mu}{\phi})} = \frac{1}{g_h^2(\mu)} - \frac{b}{8\pi^2}F_\phi\theta^2 \,, \tag{2.12}$$

where g_h is the holomorphic coupling in the supersymmetric limit and b is the 1-loop β-function coefficient. Notice that although S depends only on ϕ, eq. (2.8) is a function of the combination $\phi\phi^\dagger$, as it should be. By using eqs. (2.8) and (2.9), it is easy to check that our shift $\mu \to \mu/\sqrt{\phi\phi^\dagger}$, satisfies eq. (2.7) at lowest order. The same conclusion can be reached by using directly the gauge coupling running at two loops in eq. (2.10). This is an important property of $\mu \to \mu/\sqrt{\phi\phi^\dagger}$. Indeed, quite generally, and regardless of supergravity, one may have asked whether a shift $\mu \to \mu/\sqrt{\phi\phi^\dagger}$, applied to the running parameters of a supersymmetric theory, does generate a RG trajectory for soft terms (i.e. whether the soft terms that are thus generated do solve their RG equations, rather than being meaningless expressions). The property we just mentioned is a crucial check, that our deformation does generate a meaningful softly broken theory. We will see in section 7 that a similar argument can be used to accept or discard a more general deformation of the supersymmetric RG flow.

Notice that, by eq. (2.1), the soft terms associated to each operator are essentially determined by the (quantum) dimension of its coefficient. Indeed the soft masses of eqs. (2.5) and (2.11) are just determined by the RG scaling of the couplings, *i.e* by the conformal anomaly. For a scale invariant theory, like for instance $\mathcal{N} = 4$ Yang-Mills, the compensator dependence drops out in eq. (2.1) and soft terms are not generated. This is why we refer to this scenario as "anomaly mediated supersymmetry breaking".

We should recall that in previous studies [14], based on the effective string action of ref. [15], contributions to gaugino masses proportional to the β-functions have been found. However, in that case, unlike ours, the direct source of the effect is the F-term of a modulus. Moreover, prior to ref. [5], no contribution to scalar squared-masses beyond one loop was discussed.

We now spend a few words on Kahler invariance versus the physical meaning of F_ϕ.[2] In the superconformal approach the tree-level supergravity lagrangean is written as [8, 16]

$$\begin{aligned}\mathcal{L} &= -3\left[e^{-K(Q,Q^\dagger)/3}\phi\phi^\dagger\right]_D + \left[W(Q)\phi^3\right]_F &(2.13)\\ &= \frac{1}{2}\varphi\varphi^\dagger e^{-K(\tilde{q},\tilde{q}^\dagger)/3}R + \cdots, &(2.14)\end{aligned}$$

where, for simplicity, we only display the Einstein term in the component supergravity action. The two expressions in eq. (2.13) reduce, in the absence of gravitational fields, respectively to $d^2\theta d^2\bar{\theta}$ and $d^2\theta$ integrals. Before fixing the extraneous superconformal gauge freedom the compensator is a general chiral superfield $\phi = \varphi + \chi\theta + F_\phi\theta^2$. Eq. (2.13) is also invariant under Kahler transformations

$$\begin{aligned}K &\to K + f + \bar{f}\,,\\ W &\to e^{-f}W\,,\\ \phi &\to e^{f/3}\phi\,, \end{aligned}\qquad (2.15)$$

where $f = f(Q)$ is a function of the chiral matter fields. As ϕ is not a physical degree of freedom, Kahler invariance is not a true symmetry of the lagrangean. Rather it states that physical quantities depend only on the combination $G = e^K WW^\dagger$. Now, the quantity F_ϕ/φ, on which we have been focusing so far, is not left invariant when f involves hidden sector fields with non-zero F-components. However also the direct contribution from hidden fields to soft terms is not Kahler invariant, but precisely compensates the change of F_ϕ/φ, leaving physical quantities unaffected. The scenario of pure mediation by F_ϕ that we are considering then just corresponds to assuming the existence of a Kahler "gauge" where the hidden sector F-terms do not contribute to soft masses in the observable sector. We may call this the "convenient gauge". As the Kahler gauge is fixed, F_ϕ/φ is now a physical quantity. As discussed in ref. [5], in the convenient gauge K and W split as

$$\begin{aligned}W &= W(Q_h) + W(Q_o)\,,\\ e^{-K/3} &= 1 - f_h(Q_h, Q_h^\dagger) - f_o(Q_o, Q_o^\dagger)\,, \end{aligned}\qquad (2.16)$$

where Q_h and Q_o denote respectively fields in the hidden and observable sector. While the above form was motivated in ref. [5] by a scenario with extra space-time dimensions, it just corresponds to the old no-scale ansatz for the Kahler potential [7]. It is interesting that this form of K is stable under radiative corrections due to non-gravitational interactions. A reflection of this fact is that the parametrization for

[2] We thank Fabio Zwirner for raising and discussing this point.

soft terms just in terms of F_ϕ is valid at all renormalization scales. Again this has a beautiful geometric explanation in the scenario of ref. [5]. There the hidden and observable sectors live on different "branes" separated by a bulk where only gravity propagates: gauge and Yukawa interactions are relegated to the observable brane and cannot induce mixings to hidden fields.

To conclude this little detour we give the expression of F_ϕ obtained by solving its equation of motion

$$\frac{F_\phi}{\varphi} = \frac{1}{3}K_i F^i + e^{K/3}\frac{\varphi^{\dagger 2}}{\varphi}W^\dagger = \frac{1}{3}K_i F^i + e^{K/2}W^\dagger\,. \tag{2.17}$$

In the last equation we have made the Weyl gauge choice

$$\varphi = \varphi^\dagger = e^{K(\tilde{q},\tilde{q}^\dagger)/6}\,, \tag{2.18}$$

which renders the physics content of eq. (2.13) more explicit, by making the Einstein term canonical. Note that the second contribution to F_ϕ, $e^{K/2}W^\dagger$, is just the gravitino mass $m_{3/2}$, a Kahler invariant quantity. On the other hand, the first contribution, K_iF^i, not only should be evaluated in the convenient gauge but is also model dependent. In the case that supersymmetry is broken dynamically and without singlets in the hidden sector, we expect this first contribution to F_ϕ to be at least $\mathcal{O}(m_{3/2}^{3/2}/M_P^{1/2})$, so that we can neglect it [6]. In this interesting class of models the observable spectrum is just fixed by the gravitino mass, and by gauge and Yukawa couplings.

3. Decoupling and non-decoupling of high-energy thresholds and the problem of tachyonic states

In this section we will focus on the mass spectrum of AMSB models and on the effects of intermediate mass thresholds.

As shown in the previous section our soft terms at a scale μ are determined by the (super)conformal anomaly, which in turn is a property of the relevant interactions of the effective theory at that scale. So it seems that, by the way it was defined, in AMSB the UV thresholds affect soft masses only indirectly, via their effects on gauge and Yukawa couplings. Let us see this in more detail, by assuming that the theory has a threshold at $M \gg F_\phi \sim m_{3/2}$. Let us also assume that there are no singlets below M. Later we will see this is a crucial requirement. An example satisfying these conditions would be the MSSM with a vectorlike quark multiplet of mass $M \gg 1$ TeV. Now, it is easy to power count the M dependence of the terms in the low-energy effective action. Due to the absence of singlets, there can be no term with a positive power M (apart from a cosmological constant). Therefore the leading effects are logarithmic in M, and affect the kinetic terms of the light fields at the quantum level.[3] Since M is a parameter in the lagrangean, it has Weyl weight equal to zero. By Weyl invariance then the M dependence of the wave functions, eqs. (2.3) and (2.10), must take the form

$$\mathcal{Z}_i(\mu) = Z_i\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}, M\right), \qquad R(\mu) = g^{-2}\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}, M\right). \tag{3.1}$$

This implies that the soft masses, that arise from the ϕ-dependence of $\mathcal{Z}$ and R, are still determined by the features of the low-energy theory (at the scale μ); no extra contribution can come from the physics at M.

As an explicit example of this decoupling effect, let us consider a sector with two fields X and Y and lagrangean

$$\int d^4\theta\,(XX^\dagger + YY^\dagger) + \int d^2\theta\, Y(X^2 - M^2\phi^2)\,, \tag{3.2}$$

where by convention we have taken X and Y with Weyl weight 1. The superfield X vacuum expectation value is given by

$$\langle X\rangle = M\phi\,, \tag{3.3}$$

up to irrelevant higher-derivative terms. Both X and Y become massive and can be integrated out. After that, any low-energy wave function can depend on $\langle X\rangle$ but only through the Weyl-invariant combination $\langle X\rangle/\phi = M$. Again this does not affect the soft masses. If we were to work with component fields, we would find that the threshold contributions from the heavy fields at M combine themselves to preserve the AMSB form (eq. (2.5)) of the low-energy soft terms.

It is rather clear why the AMSB scenario is interesting for the supersymmetric flavor problem. In the SM, the three Yukawa matrices are the *only* relevant sources of flavor violation. This special property leads to a natural suppression of dangerous flavor violating processes. However this property is lost in a generic supersymmetric extension, due to the presence of squark and slepton mass matrices that are in principle new sources of flavor violation. An elegant solution to this problem is provided by models with GMSB. In these models the scale at which the three Yukawa matrices $Y_{u,d,e}$ are generated is well above the scale Λ_S at which the soft terms are induced. Therefore flavor violations in the sfermion masses are only proportional to the low-energy $Y_{u,d,e}$ and we recover natural flavor conservation. (Moreover the flavor invariant contribution due to gauge interactions dominates these masses.)

The AMSB scenario also satisfies natural flavor conservation: scalar masses are only functions of the low-energy gauge couplings and $Y_{u,d,e}$. Unlike gauge mediation, this property is rather independent on where the scale of flavor is. As we explained

[3]Higher-derivative operators scaled by M^{-n}, with $n > 0$, give a contribution to soft terms suppressed by powers of $F_\phi/M \sim m_{3/2}/M$. This class of effects can only be important when $M \sim m_{3/2}$.

above, the soft masses in AMSB models decouple from UV physics, and depend only on the IR theory. This a great virtue, but unfortunately it is also what kills this scenario by making it too (wrongly) predictive. The scalar masses are given by eq. (2.5), that leads in a pure gauge theory to

$$m_i^2(\mu) = \frac{c_i b}{8\pi^2}\alpha^2(\mu)|F_\phi|^2\,, \tag{3.4}$$

where $c_i > 0$ is the quadratic Casimir of the scalar. This contribution is therefore positive for asymptotically free gauge theories ($b > 0$) and negative for infrared free theories. In the MSSM both $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ have $b < 0$. The sleptons, whose masses are essentially determined by the $\mathrm{SU}(2)_L\times$ $\mathrm{U}(1)_Y$ gauge interactions,[4] are therefore tachyonic and the model is ruled out.

In ref. [5], in order to avoid the tachyons, it was assumed the presence of additional (non-anomaly-mediated) contributions to the scalar masses. These contributions were associated to new fields propagating in the bulk. Although these bulk contributions could lead to positive scalar masses, they certainly spoil the decoupling features of the AMSB models. These effects can be associated to terms that mix hidden with observable fields in eq. (2.16), so that we apparently go back to the standard problem: in the absence of an explicit model where the bulk effects are calculable the supersymmetric flavor problem remains open.

In the rest of the paper we will investigate the possibility of preserving the property of natural flavor conservation in AMSB models, while avoiding tachyons. For this purpose, we will reconsider the decoupling property of AMSB models. As it is often the case, the absence of light singlets turns out to be necessary in the proof of decoupling.

Our basic point is the following. Suppose that we have a superfield X such that the VEV M of its scalar component sets the threshold. If after the shift $X = M\phi + S$, the singlet S is light, it must be kept in the effective theory. In this case the Kahler potential can have a linear dependence on M:

$$\int d^4\theta \left(SS^\dagger + cM\phi S^\dagger + c^* M\phi^\dagger S\right), \tag{3.5}$$

where c is an $\mathcal{O}(1)$ number. In the presence of supersymmetry breaking these terms are important. By minimizing we get a rather large $F_S = -cMF_\phi$. (Notice that this result corresponds to the well known problem that singlets coupled to the MSSM Higgs $H_1 H_2$ can destabilize the weak scale because of their large F-terms [17].) Thus for a light singlet S, the proportionality of eq. (3.3) is violated (remember that S has zero scalar VEV by assumption). The power counting is now drastically changed. For instance, an irrelevant operator like

$$\int d^4\theta \frac{SS^\dagger}{M^2} QQ^\dagger\,, \tag{3.6}$$

[4] Yukawa coupling effects are negligible for (at least) the first two families.

leads to an unsuppressed contribution $\sim |cF_\phi|^2$ to the mass of the Q scalar. Soft terms are no longer associated to the relevant interactions of the low-energy theory and decoupling is lost. This property of light singlets is the key ingredient of the models we will construct. As we will see, in these new models, the tachyons can be eliminated by extra non-decoupling (flavor diagonal) contributions of heavy states.

Non-decoupling effects from light singlets have also been used in ref. [18] to lower the effective scale of supersymmetry breaking. The scenario in ref. [18] is different from ours as it relies on a total singlet that mixes at the Planck scale with the hidden sector fields. Also, in that scenario $m_{3/2}$ must be much below the weak scale, while in our case it is much above.

4. Anti gauge-mediated supersymmetry breaking

The first model we will present is based on the following set of extra fields: the "messenger" fields, Ψ and $\bar{\Psi}$, in vectorlike representations of the SM gauge group, and the field X, whose VEV will give mass to the messengers. The latter field has a very flat potential in order to fulfill the condition for non-decoupling that we explained above. Like in gauge mediated models we take Ψ and $\bar{\Psi}$ to fit in complete SU(5) representations in order to preserve gauge coupling unification. For instance we can take N flavors of $\mathbf{5}+\bar{\mathbf{5}}$. The superpotential of the model is given by

$$W = \lambda X \Psi \bar{\Psi} + \frac{X^n}{\Lambda^{n-3}\phi^{n-3}}\,, \tag{4.1}$$

where again all fields have Weyl weight 1 and $n > 3$. The scale Λ suppressing the non-renormalizable term could conceivably be of the order of the Planck mass, $\Lambda \sim M_P$. The form of the above superpotential could also be enforced by an (anomalous) R-symmetry. Notice that the ϕ-dependence of W is dictated by Weyl-invariance. Now, when F_ϕ is turned on, the tree-level potential along the scalar-component of X becomes

$$V(X) = n^2 \left|\frac{X^{n-1}}{\Lambda^{n-3}}\right|^2 + (n-3)\left(\frac{F_\phi X^n}{\Lambda^{n-3}} + \text{h.c.}\right). \tag{4.2}$$

The minimum condition for $V(X)$ is

$$\frac{X^n}{|X|^2} = -\frac{n-3}{n(n-1)} F_\phi^* \Lambda^{n-3}\,, \tag{4.3}$$

that leads to a non-zero VEV for X:

$$\langle X\rangle \sim m_{3/2}^{\frac{1}{n-2}} \Lambda^{\frac{n-3}{n-2}}\,. \tag{4.4}$$

Eq. (4.3) also fixes the value of F_X/X (the quantity relevant to soft terms)

$$\frac{F_X}{X} = \frac{nX^{*n}}{|X|^2\Lambda^{*(n-3)}} = \frac{n-3}{n-1} F_\phi\,. \tag{4.5}$$

Notice that, unlike the model of eq. (3.2), the VEV of the superfield X is no longer proportional to ϕ. Since X has weight 1, the low-energy wave functions depend on the threshold via $\tilde{X} \equiv X/\phi$, which has now a non zero F-component

$$\frac{F_{\tilde{X}}}{\tilde{X}} = -\frac{2}{n-1} F_\phi \,. \tag{4.6}$$

Then in this model the soft masses at low-energy are no longer determined by the action of $d/d\ln\mu$ on the wave functions, eqs. (2.5) and (2.11), but by

$$\begin{aligned} \frac{m_\lambda(\mu)}{g^2(\mu)} &= \frac{F_\phi}{2}\left(\frac{\partial}{\partial\ln\mu} + \frac{2}{n-1}\frac{\partial}{\partial\ln|X|}\right)\frac{1}{g^2(\mu,X)}\,, \\ A_i(\mu) &= -\frac{F_\phi}{2}\left(\frac{\partial}{\partial\ln\mu} + \frac{2}{n-1}\frac{\partial}{\partial\ln|X|}\right)\ln Z_i(\mu,X)\,, \\ m_i^2(\mu) &= -\frac{|F_\phi|^2}{4}\left(\frac{\partial}{\partial\ln\mu} + \frac{2}{n-1}\frac{\partial}{\partial\ln|X|}\right)^2 \ln Z_i(\mu,X)\,. \end{aligned} \tag{4.7}$$

From the component point of view, the soft terms in this model are determined as follows. Down to the scale X the soft terms respect the AMSB form (2.5). At this scale, the messengers add a gauge mediated contribution. This contribution does not adjust the soft terms to the AMSB trajectory of the low-energy theory. For that to happen one would need precisely $F_X/X = F_\phi$, which does not hold. Notice indeed that decoupling is recovered only in the limit $n \to \infty$, when the mass of the X excitation around the minimum becomes very large.

As another example we now consider the previous model but in the absence of the second term on the r.h.s. of eq. (4.1). In this case X is a flat direction only lifted by the effects of $F_\phi \neq 0$. The relevant term along $X \neq 0$ and $\Psi, \bar{\Psi} = 0$ is

$$\int d^4\theta Z_X\left(\sqrt{\frac{XX^\dagger}{\phi\phi^\dagger}}\right) XX^\dagger\,, \tag{4.8}$$

where $\mu^2 \to XX^\dagger/\phi\phi^\dagger$ has been taken in the wave function, since X plays now the role of IR cut-off (notice that the above equation satisfies Weyl symmetry). At leading order the potential along X is determined by the running soft mass

$$V(X) = m_X^2(|X|)|X|^2 \simeq \left|\frac{F_\phi}{16\pi^2}\right|^2 N\lambda^2(X)\left[A\lambda^2(X) - C_a g_a^2(X)\right]|X|^2\,, \tag{4.9}$$

where $A, C_a > 0$, and a sum over the gauge couplings g_a of the messengers is understood. The expression in square bracket is just the beta function of λ. It is conceivable a situation where at a large energy scale the λ^4 term dominates and $m_X^2 > 0$, while at lower values of X the gauge term becomes important leading to $m_X^2 = 0$ at some M. The X field is therefore stabilized at $\langle X\rangle \sim M$. This is just a supersymmetric version of the Coleman-Weinberg mechanism [20]. A situation like that is very easy to obtain if the gauge group has an asymptotically free factor. In the MSSM only SU(3) is weakly UV free and the presence of more than 3 flavors of messengers makes it IR free. Nonetheless it is quite possible to gauge a subgroup of the messenger flavor group obtaining a strongly UV free factor. In this case the Coleman-Weinberg stabilization of X can work very well. X could get a VEV anywhere between the weak and GUT scales. Notice that from eq. (4.9), F_X/X is a loop factor smaller that F_ϕ, $F_X/X \simeq N\lambda^2 F_\phi/16\pi^2$, and can be neglected. Therefore $F_{\tilde{X}}/\tilde{X} \simeq -F_\phi$ and the soft masses are affected by the messenger threshold. This effect is given by eqs. (4.5) and (4.7) extrapolated at $n = 3$. The essential agreement between these two limits is just due to the fact that they both correspond to classically scale invariant models in which there are no soft terms at tree level.

4.1 Superparticle spectrum

A model like the one just outlined above is very similar to GMSB models. All sources of flavor violation that are active above the scale X cannot affect the masses in a relevant way; this is because above X we are on the AMSB trajectory. Below the scale X, the soft terms are no longer on this privileged trajectory and can pick up all sort of dangerous contributions. Therefore we must assume that the scale of flavor Λ_F (below which the only sources of mixings are $Y_{u,d,e}$) is somewhat larger than X. If the MSSM unifies in a simple gauge group, then Λ_F cannot be bigger than the GUT scale $\sim 10^{16}$ GeV. Thus we will assume from now on $X \lesssim 10^{16}$ GeV. For simplicity we will also consider the case in which the VEV of X is fixed by the Coleman-Weinberg mechanism.

In order to get the sparticle spectrum from eq. (4.7), we only need to know the dependence of the gauge coupling g and the wave function Z on the scale μ and on the singlet X induced after integrating out the messengers Ψ and $\bar{\Psi}$ [19]. For a simple gauge group, these are given by

$$\alpha^{-1}(\mu, X) = \alpha^{-1}(\Lambda) + \frac{b-N}{4\pi} \ln \frac{XX^\dagger}{\Lambda^2} + + \frac{b}{4\pi} \ln \frac{\mu^2}{XX^\dagger}\,, \tag{4.10}$$

and

$$Z_i(\mu, X) = Z_i(\Lambda) \left(\frac{\alpha(\Lambda)}{\alpha(X)}\right)^{\frac{2c_i}{b-N}} \left(\frac{\alpha(X)}{\alpha(\mu)}\right)^{\frac{2c_i}{b}}\,. \tag{4.11}$$

where c_i is the quadratic Casimir. Using eqs. (4.10), (4.11) and (4.7) with $n = 3$, we obtain

$$m_\lambda(\mu) = \frac{\alpha(\mu)}{4\pi}(b-N)F_\phi\,, \tag{4.12}$$

$$A_i(\mu) = -\frac{2c_i}{4\pi}\left[\alpha(\mu) + [\alpha(X) - \alpha(\mu)]\frac{N}{b}\right] F_\phi\,, \tag{4.13}$$

$$m_i^2(\mu) = \frac{2c_i b}{(4\pi)^2}\left[\alpha^2(\mu) - \alpha^2(\mu)\frac{N}{b} + [\alpha^2(\mu) - \alpha^2(X)]\left(\frac{N^2}{b^2} - \frac{N}{b}\right)\right] |F_\phi|^2\,. \tag{4.14}$$

At the scale $\mu = X$, we can see that the contributions in (4.12)–(4.14) arise from two sources. The first term is purely dictated by the superconformal anomaly. The second one is the effect of the N Ψ-$\bar{\Psi}$ threshold. This second contribution is equal in magnitude but *opposite in sign* to the usual in GMSB models. For this reason we call this scenario *anti*-GMSB. Since it is well known that the GMSB contribution to the scalar soft masses are positive, one could think that this will not help us to solve the above problem of negative scalar masses. Nevertheless, this is not the case. In *anti*-GMSB there is yet another contribution to the scalar masses: the third term in brackets in eq. (4.14) originating from the gaugino mass via RG. Therefore in order to have positive scalar masses at $\mu \sim m_W$, the third term of eq. (4.14) must overcome the others. The condition $m_i^2 > 0$ puts a lower bound on N and X. In figure 1 we plot the soft squared-masses of the left-handed squark (solid line), the right-handed down-squark (dashed-dotted), the left-handed slepton (dashed) and the right-handed slepton (dotted) normalized to the $\mathrm{SU}(2)_L$-gaugino squared-mass, as a function of $\mathrm{Log}_{10}(X/\mathrm{GeV})$. These masses are calculated at the scale ≈ 200 GeV. We have taken $N = 2, 3, 4$ pairs of Ψ and $\bar{\Psi}$ in the $\mathbf{5}$ and $\bar{\mathbf{5}}$ representation of SU(5). We see that for $X \lesssim M_{GUT} \simeq 10^{16}$ GeV, one is forced to have $N \geq 3$ in order to have positive masses. The case $N = 3$, however, implies a zero β-function for the SU(3) coupling above X, and therefore zero gluino mass. We then find that the only case that gives realistic soft masses corresponds to $N \geq 4$. As an example, let us consider the case $N = 4$ with $X = M_{GUT}$. The spectroscopy of this scenario is completely different from anyone in the literature. Neglecting the effects of Yukawa couplings, we have:

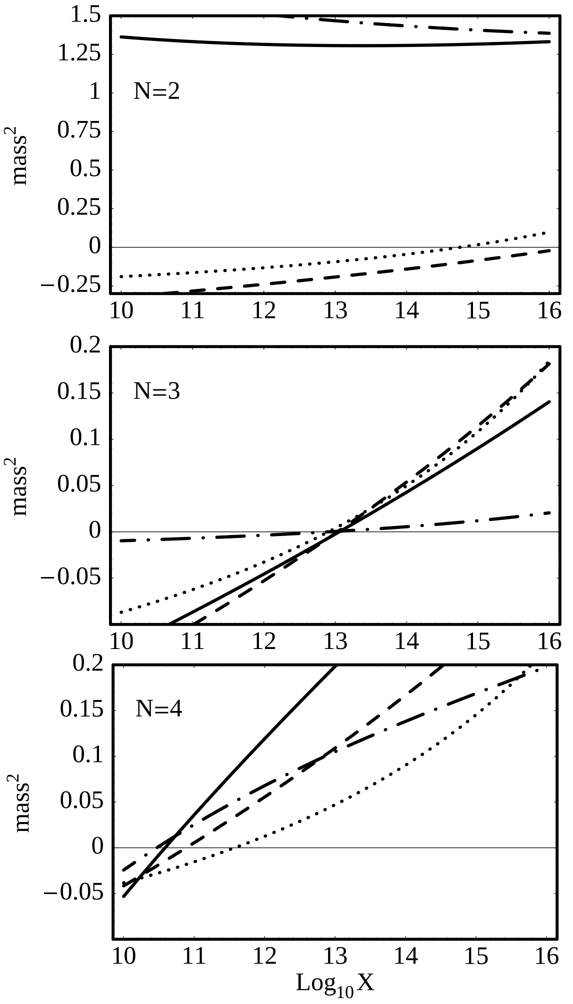

Figure 1: Soft squared-masses of the left-handed squark (solid line), the right-handed down-squark (dashed-dotted line), the left-handed slepton (dashed line) and the right-handed slepton (dotted line) normalized to the $\mathrm{SU}(2)_L$-gaugino squared-mass as a function of $\mathrm{Log}_{10}(X/\mathrm{GeV})$ for $N = 2, 3, 4$ messenger multiplets.

$$m_{\tilde{\lambda}_1} \simeq 1.1\, m_{\tilde{\lambda}_2} \simeq 1.6\, m_{\tilde{\lambda}_3} \simeq 1.6\, m_Q \simeq 1.9\, m_L \simeq 2\, m_U \simeq 2.2\, m_E \simeq 2.4\, m_D\,. \qquad (4.15)$$

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We have not specified the Higgsino mass since it depends on the μ-parameter and therefore is model dependent. From eq. (4.15) we see that the gaugino masses have an hierarchy opposite to that of GMSB. Also notice that this scenario gives some of the squarks lighter than the sleptons. The stops and the Higgs H_2 receive an extra contribution proportional to the top Yukawa coupling Y_t that is as important as eq. (4.14). For H_2 this is given by

$$\delta m^2_{H_2}(\mu) = \frac{3}{(4\pi)^2}\alpha_t(\mu)\left\{\sum_i d_i\left[-\alpha_i(\mu)+\frac{2N}{b_i}(\alpha_i(\mu)-\alpha_i(X))+\frac{F(\mu)}{E(\mu)}N\frac{\alpha_i^2(X)}{2\pi}\right]+\right.$$
$$\left.+\frac{N^2}{2\pi}\frac{G_2(\mu)}{E(\mu)}+6\alpha_t(\mu)\left[1+\frac{N}{2\pi}\frac{G_1(\mu)}{E(\mu)}\right]^2\right\}|F_\phi|^2\,, \tag{4.16}$$

where i sums over the three MSSM gauge groups, $b_i = (-6.6, -1, 3)$, $d_i = (13/15, 3, 16/3)$, $\alpha_t = Y_t^2/4\pi$ and

$$\begin{aligned}
E(\mu) &= \Pi_i\left(\frac{\alpha_i(t)}{\alpha_i(t_X)}\right)^{d_i/b_i}, \qquad t=\ln\mu\,, \qquad t_X=\ln X\,,\\
F(\mu) &= \int_{t_X}^t E(t')dt'\,,\\
G_1(\mu) &= \int_{t_X}^t E(t')\sum_i\frac{d_i}{b_i}(\alpha_i(t')-\alpha_i(t_X))dt'\,,\\
G_2(\mu) &= \int_{t_X}^t E(t')\left\{\sum_i\frac{d_i}{b_i}(\alpha_i^2(t')-\alpha_i^2(t_X))+\left[\sum_i\frac{d_i}{b_i}(\alpha_i(t')-\alpha_i(t_X))\right]^2\right\}dt'\,.
\end{aligned} \tag{4.17}$$

For the left-handed stop and right-handed stop, the top contribution is obtained by just replacing the factor 3 in front of eq. (4.16) by a factor 1 and 2 respectively. We have checked that for values of the top coupling in the region $0.7 \lesssim Y_t(X) \lesssim 0.9$ the soft mass $m^2_{H_2}$ is the only negative one. This can trigger electroweak symmetry breaking (EWSB).

The μ-term and the bilinear Higgs mass term, $B\mu$, are not predicted by the model. In section 6, we will propose a way to generate them. For phenomenological purposes, however, they can be considered free parameters of the theory. Their only constraint comes from EWSB. We find that EWSB can be achieved for moderate values of μ, typically not bigger than the other soft masses. This is very convenient since it implies that the LSP in these theories can be the neutral Higgsino. This is again different from GMSB models where the LSP is the gravitino. Here the gravitino gets a tree-level mass of order F_ϕ and becomes very heavy ($\sim$ 10 TeV).

5. D-term contributions to scalar masses

One could also use the previous non-decoupling mechanism to build realistic models with extra U(1)'s at high-energies. Extra U(1)'s are well motivated as they appear

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in GUT groups of rank greater than 5, such as SO(10) or E_6. As we shall see, additional contributions to the low-energy soft masses can be obtained from the D-terms of these U(1)'s.

Let us consider two fields ψ and $\bar{\psi}$ with U(1)-charges 1 and -1. Let us also assume that the D-flat direction, $\psi = \bar{\psi} \equiv X$, is stabilized by the Coleman-Weinberg mechanism away from the origin, as in the previous section. This will break the U(1) giving a mass to the vector superfield proportional to X. We can now integrate out the vector superfield. This induces, a new spurious dependence of the low-energy wave functions on X/ϕ, and correspondingly a new contribution to the soft masses. The mass of a scalar with U(1) charge q_i will be corrected by[5]

$$\delta m_i^2 = q_i \langle D \rangle \,, \qquad\qquad \langle D \rangle = \frac{1}{2}[m^2_{\bar{\psi}}(X) - m^2_{\psi}(X)] \,, \tag{5.1}$$

where m^2_{ψ} and $m^2_{\bar{\psi}}$ are the anomaly mediated soft masses of ψ and $\bar{\psi}$, eq. (2.5).[6] This contribution to the slepton and squark masses is in general comparable to that arising from the anomaly eq. (2.5). Its exact value depend on the RG trajectory of Z_ψ and $Z_{\bar{\psi}}$ that is model dependent. For a large enough value of $\langle D \rangle$ and $q_i > 0$, the contribution (5.1) can overcome the negative contribution (2.5) and provide a realistic theory of soft masses. The contribution (5.1) is generated at the scale at which the U(1) is broken. This scale must be around the GUT scale if we do not want to spoil the gauge coupling unification of the MSSM. This does not necessary implies that the scale of flavor must be well above the GUT scale to avoid flavor violations in the soft masses. As explained in ref. [21], if q_i are flavor independent, the soft masses can be flavor diagonal at the leading order.

Gaugino masses are not affected by D-terms since the latter are R-invariant. Therefore in this scenario gaugino masses are predicted to have their anomaly mediated value eq. (2.11).

6. The μ-problem

In order to have a phenomenologically viable model, a supersymmetric Higgs mass parameter μ, of the order of the soft terms, must be generated. This is a well-known problem in GMSB theories where soft masses are induced by loop effects [22]. In our models, since soft masses are also induced at the one-loop level, one might expect the μ-problem to be equally severe. On the contrary, we will show that in these theories there is a simple way to generate the μ parameter.

[5] There can be other contributions mediated by the $U(1)$ vector multiplet at the two-loop level [19]. These contributions, however, are proportional to $g^4_{U(1)}$ and can be in principle smaller than the ones here.

[6] Notice that for $m^2_\psi = m^2_{\bar{\psi}}$ the D-term contribution eq. (5.1) is zero. Therefore ψ and $\bar{\psi}$ must have different couplings in order to generate the soft mass (5.1).

A priori, a μ mass could arise from the Weyl-invariant operator [23]

$$\int d^4\theta \frac{\phi^\dagger}{\phi} H_1 H_2 \,, \tag{6.1}$$

where $H_{1,2}$ are the two Higgs doublets of the MSSM. This operator induces both μ and $B\mu$ at tree level. While μ could be small because of a small coefficient multiplying eq. (6.1), this operator always gives $B = F_\phi$, which is much larger than all the other soft masses. Therefore this term must be forbidden. The solution we propose is to generate the μ-term by the operator

$$\int d^4\theta \, H_1 H_2 \frac{X^\dagger}{X} \tilde{Z} \left(\sqrt{\frac{XX^\dagger}{\phi\phi^\dagger}} \right) , \tag{6.2}$$

where $\tilde{Z}$ is a wave function coefficient at the renormalization scale $\mu_R \to \sqrt{XX^\dagger/\phi\phi^\dagger}$ (in this section we will denote the renormalization scale by μ_R instead of μ in order to distinguish it from the Higgs μ-parameter), and X is the light singlet with $F_X \simeq 0$ introduced in section 4.

The operator (6.2) generates a one-loop μ term and a two-loop $B\mu$ term:

$$\mu = \frac{1}{2} \frac{d\tilde{Z}}{d \ln \mu_R} \bigg|_{\mu_R = |X|} \,, \tag{6.3}$$

$$B\mu = \frac{1}{4} \frac{d^2 \tilde{Z}}{(d \ln \mu_R)^2} \bigg|_{\mu_R = |X|} \,. \tag{6.4}$$

The operator (6.2) can arise by integrating out a singlet S coupling to $H_1 H_2$ and getting its mass from X. A simple example is given by the superpotential

$$\int d^2\theta \left[\lambda S H_1 H_2 + \frac{1}{3} k S^3 + \frac{1}{2} y S^2 X \right] . \tag{6.5}$$

At one-loop, a kinetic mixing between X and S is generated

$$\int d^4\theta \tilde{Z}(\mu_R) \, S X^\dagger + \text{h.c.} \tag{6.6}$$

For a nonzero X VEV, the singlet S is massive and can be integrated out. Its equation of motion is given by

$$S \simeq -\frac{\lambda}{y} \frac{H_1 H_2}{X} \,, \tag{6.7}$$

which inserted in eq. (6.6) gives rise to eq. (6.2).

An alternative model is given by

$$\int d^2\theta S \left(\lambda H_1 H_2 + \lambda_N N^2 - \lambda_{\bar{N}} \bar{N}^2 \right) . \tag{6.8}$$

This model has a flat direction along $\lambda_N N^2 = \lambda_{\bar{N}} \bar{N}^2 \equiv X^2$. Assuming as above that X gets a VEV, the singlets get a mass. Integrating them out, one obtains eq. (6.2) with $\tilde{Z} \propto Z_N - Z_{\bar{N}}$.

We stress that in the present mechanism there is no danger of generating a large B, which was instead the case for GMSB. This is because nowhere in the visible sector there is a supermultiplet with a tree-level mass splitting. All splittings are of order $\alpha F_\phi \sim m_W$ to start with, and B cannot come out bigger than that.

7. Gauging an R-symmetry

In the first section, independent of supergravity, we could have pointed out that the shift $\mu^2 \to \mu^2/\phi\phi^\dagger$ is remarkable as it automatically defines a deformation of the RG flow of a supersymmetric theory, into the flow of a softly broken one. In this section we study under what conditions one can define a more general deformation $\mu \to \mu e^{V/3}$, where V is now a genuine vector superfield, and the factor 3 is chosen for later use. Indeed we will just need to discuss the case $V = V_D \theta^4$, since $\mu \to \mu e^{V/3}/(\phi\phi^\dagger)^{1/2}$ parametrizes the most general shift.

Let us consider first a Wess-Zumino model with superfields Q_i, $i = 1, \ldots, N$, and Yukawa couplings λ_{ijk}. To be slightly more general, we indeed define the soft terms by the formal shift

$$\mathcal{Z}_i(\mu) = e^{q_i V} Z_i\left(\mu e^{V/3}\right) \quad \longrightarrow \quad m_i^2 = -\left(q_i + \frac{\gamma_i}{3}\right) V_D \,. \tag{7.1}$$

In order to check if these soft masses do represent an RG trajectory we verify that they satisfy the evolution equations. By use of the results of ref. [9] the general RG equation at all orders has the form

$$\frac{dm_i^2}{d\ln\mu} = -\frac{\partial \gamma_i}{\partial \ln \lambda_{klm}^2} \left(m_k^2 + m_l^2 + m_m^2\right) \,, \tag{7.2}$$

where summation over k, l, m is understood. It is easy to check that eq. (7.1) solves the one above if $q_k + q_l + q_m = 0$ for any $\lambda_{klm} \neq 0$. Notice that the q's can thus be considered charges of a background gauge symmetry, see eq. (7.1). Moreover, whatever the structure of λ_{klm}, we can always define a consistent RG deformation with all q's vanishing. The situation changes when we turn to a gauge theory. In addition to the shift of eq. (7.1), the superfield gauge couplings are defined by

$$R_a = \frac{1}{g_a^2\left(\mu e^{V/3}\right)} \,. \tag{7.3}$$

We can do the same check we did for the WZ model. The Yukawa sector gives the same constraint on the q's as before. The constraint from the gauge sector is determined by eq. (2.7) for $R_a|_D$ (remember that $R|_D$ is the coefficient of a non-local

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operator [9]). By using eqs. (7.1) and (7.3), and $m_\lambda = 0$, eq. (2.7) implies

$$\left(b - \sum_i T_i\gamma_i\right)\frac{V_D}{3} = \left(\sum_i T_i(-q_i - \gamma_i/3)\right)\frac{V_D}{3}\,, \tag{7.4}$$

which, conveniently rewritten in terms of $r_i = q_i + 2/3$, is

$$N_c + \sum_i T_i(r_i - 1) = 0\,. \tag{7.5}$$

This is just the condition for an R-symmetry of charges r_i to be non anomalous. Notice also that the Yukawa constraint for r_i is also $r_i + r_j + r_k = 2$ for each $\lambda_{ijk} \neq 0$. Thus we have established that the general deformation $\mu \to \mu e^V$ can only work for theories with a non-anomalous R-symmetry. Unfortunately the MSSM does not have such a symmetry, so we cannot use this deformation to improve the AMSB scenario. A possibility would be to add the suitable matter multiplets at the weak scale [24].

It is remarkable that by the above arguments in rigid supersymmetry a gauged R symmetry has emerged. This brings us back to our natural arena: supergravity. Indeed, maybe less instructively, we could have derived this new deformation simply by gauging R in supergravity. This is done in the formalism of ref. [25], by introducing a connection e^V under which the compensator has charge $-2/3$, and matter fields have charge r_i. The lagrangean is then given by

$$\int d^4\theta\phi\phi^\dagger e^{(r_i-2/3)V}Q_iQ_i^\dagger + \int d^2\theta\phi^3 W(Q_i)\,, \tag{7.6}$$

where $W(Q_i)$ must have charge 2 to match the charge of ϕ^3.[7] Notice that when V gets a D-term we get at tree level $m_i^2 = -(r_i - 2/3)V_D = -q_iV_D$ in agreement with eq. (7.1).

8. Conclusions

Theories of supersymmetry breaking induced by the conformal anomaly [5] present several interesting properties. They are simple to realize, very predictive and the induced soft masses are independent of flavor physics. Nevertheless, they have a big drawback: they lead to tachyons (in particular, the sleptons).

Here he have reexamined these theories focusing on their ultraviolet decoupling property. We have shown that decoupling is lost whenever a singlet (a *modulus*) remains light below the high-energy threshold. These non-decoupling effects have been used to solve the tachyon problem.

[7] Indeed the usual R-symmetry is a combination of the one here defined (which commutes with supersymmetry) with the axial symmetry of the superconformal group. Both symmetries are formally broken by the VEV of the real component of the compensator.

In our first model, named anti-GMSB, the anomaly mediated masses are modified by heavy vectorlike fields in such a way that they end up being positive at low-energy (except for the Higgs). The model is still quite predictive since the spectrum depends only on the number N and the mass M of the vectorlike fields. Furthermore, the soft masses are also flavor independent if the mass of these fields is below the scale of flavor. Phenomenologically the model is quite interesting since it presents a mass spectrum different from other scenarios. For instance, for $N = 4$ the gluino is the lightest gaugino. The lightest scalar can be the down squark, but it is more typically the right-handed slepton. Depending on μ and $\tan\beta$ the LSP is either the lightest scalar or a combination of zino and neutral Higgsino. Although the μ-parameter cannot be predicted, it can be generated by the simple mechanism explained in section 6. Unlike other scenarios, the gravitino is very heavy ≈ 10 TeV since it gets its mass at tree-level.

A realistic mass spectrum can also be found in theories with extra U(1)'s broken at high-energies. After integrating out the heavy U(1) sector, non-decoupling remnants can again avoid the problem with tachyons. Flavor independence of the soft masses is guaranteed in this case if the U(1)-charges are flavor independent.

In the models considered so far supersymmetry breaking was originating from the F-term of the gravitational multiplet. In the last section we have generalized this to include a D-term as a new source of supersymmetry breaking. This requires gauging an R-symmetry and consequently enlarging the MSSM field content.

We want to conclude by emphasizing that the theories considered here make testable predictions on the sparticle spectrum. This encourages us to further explore their phenomenology at present and future experiments.

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