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Unscreening the Gaugino Mass with Chiral Messengers

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Abstract

Gaugino screening, the absence of next-to-leading order corrections to gaugino masses, is a generic feature of gauge mediation models of supersymmetry breaking. We show that in a specific class of models, known as semi-direct gauge mediation, it is possible to avoid gaugino screening by allowing for a chiral messenger sector. Messengers then acquire a mass at some scale, for instance by higgsing or by some auxiliary strong coupling dynamics. We implement this idea in a simple model which we work out explicitly.

1 Introduction

In the framework of gauge mediation of supersymmetry (SUSY) breaking [1], gaugino mass screening refers to the fact that next-to-leading order radiative corrections to the gaugino mass are absent [2]. When the visible gaugino mass vanishes at leading order, as in semi-direct gauge mediation (SDGM) [3, 4], the screening implies that the gaugino mass will only arise at next-to-next-to-leading order and hence be severely suppressed with respect to sfermions masses. While such strong hierarchy would fit into a split supersymmetry scenario [5], one might wonder whether there are ways to avoid gaugino screening in gauge mediation, in the first place.

The object of this note is to present a SDGM set up where the phenomenon of gaugino mass screening does not take place. SDGM is a class of gauge mediation models where the messengers interact with the hidden sector only through (non-SM) gauge interactions but, unlike direct gauge mediation, they do not participate to the hidden sector supersymmetry breaking dynamics. A pictorial representation of SDGM is reported in figure 1.

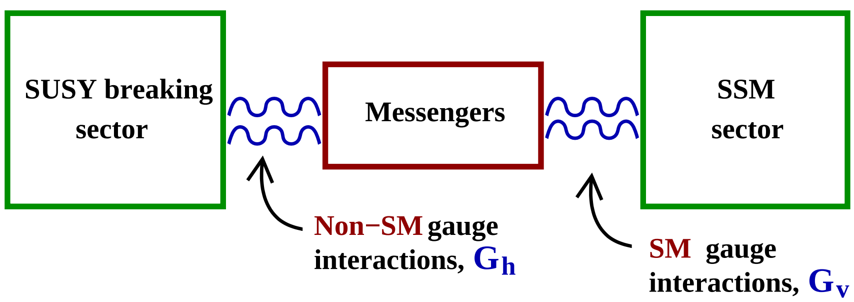

Figure 1: A cartoon of semi-direct gauge mediation. The gauge group G_h is singled-out within the hidden sector as the subgroup to which the messenger fields couple. G_v is the gauge group of the visible sector.

The idea to avoid gaugino screening in SDGM is very simple: it suffices to allow for a chiral messenger sector, in the sense that gauge symmetries prevent the presence of an explicit mass term in the superpotential for the messenger fields. Eventually, the messengers will acquire a mass (e.g. by higgsing) and disappear from the low energy spectrum. However, if there is a sufficient range for RG evolution above this scale the visible gaugino will indeed acquire a non-vanishing mass at next-to-leading order. In spirit this is very close to (and indeed inspired by) what is called Z' mediation [6], where the role of the messengers is played directly by the MSSM matter.

In what follows we first discuss the evasion of the gaugino mass screening argument both from the point of view of its original discussion in terms of wave function renormalization [2], in section 2, and by direct evaluation of Feynman diagrams [7], in section 3. In section 4 we implement our basic idea within a concrete model based on a quiver gauge theory, which can arise from D-branes at singularities, and provide additional details on the mechanism of gaugino unscreening. We end in section 5 by discussing possible phenomenological implementations of our model, also giving estimates for the gaugino and sfermion masses.

2 Gaugino (un)screening from wave function renormalization

As we mentioned in the introduction, what goes under the name of gaugino mass screening is the observation that when gauge mediation of SUSY breaking is operated by messengers, next-to-leading order corrections to the gaugino mass cancel each other. In the context in which one obtains soft masses by promoting MSSM wave function renormalization factors to (spurionic) superfields, the argument goes as follows [2].

The expression for the running physical (real) gauge coupling $R(\mu)$ is given by (for ease of comparison, we stick to the notation of [2])

$$R(\mu) = S(\mu) + S(\mu)^\dagger + \frac{T_G}{8\pi^2}\log(S(\mu) + S(\mu)^\dagger) - \sum_r \frac{T_r}{8\pi^2}\log Z_r(\mu) \ , \qquad (1)$$

where $S(\mu)$ is the holomorphic coupling and $Z_r(\mu)$ the wave function renormalization of matter fields. The sum over r is on all representations (of index T_r) of matter fields charged under the gauge group which are present below the scale μ. We recall that the presence of the logarithmic terms is in order to compensate the unphysical rescaling symmetry of the holomorphic coupling $S(\mu)$.

The running of $R(\mu)$ will experience a threshold at the scale at which the messengers stop contributing (when going towards the IR). We will call this scale μ_X. The RG running will be different above and below this scale, essentially because of the presence, above μ_X, of an extra term depending on the wave function renormalization of the messengers. Eventually, what we find for the gauge coupling for $\mu < \mu_X < \mu_0$ is

$$\begin{aligned} R(\mu) &= R(\mu_0) + \frac{b_0}{16\pi^2} \log \frac{\mu^2}{\mu_0^2} + \frac{T_G}{8\pi^2} \log \frac{\mathrm{Re} S(\mu)}{\mathrm{Re} S(\mu_0)} - \sum_r \frac{T_r}{8\pi^2} \log \frac{Z_r(\mu)}{Z_r(\mu_0)} \\ & \quad - \frac{T_M}{16\pi^2} \log \frac{\mu_X^2}{\mu_0^2} - \frac{T_M}{8\pi^2} \log \frac{Z_M(\mu_X)}{Z_M(\mu_0)} \ , \end{aligned} \tag{2}$$

where T_M is essentially the number of messengers, and Z_M their wave function renormalization. The constant $b_0 = 3T_G - \sum_r T_r$ is the coefficient of the one-loop beta function below the scale μ_X. The leading order contribution to the gaugino mass comes from replacing μ_X by its tree level value X and then promoting it to a spurion $X + \theta^2 F$. Next-to-leading order corrections should come from corrections to the values of μ_X and Z_M. The fact that next-to-leading order corrections vanish derives precisely from the fact that the correct expression for μ_X takes into account wave function renormalization

$$\mu_X = \frac{X}{Z_M(\mu_X)} \ . \tag{3}$$

We are then left with

$$\begin{aligned} R(\mu) &= R(\mu_0) + \frac{b_0}{16\pi^2} \log \frac{\mu^2}{\mu_0^2} + \frac{T_G}{8\pi^2} \log \frac{\mathrm{Re} S(\mu)}{\mathrm{Re} S(\mu_0)} - \sum_r \frac{T_r}{8\pi^2} \log \frac{Z_r(\mu)}{Z_r(\mu_0)} \\ & \quad - \frac{T_M}{16\pi^2} \log \frac{X^2}{\mu_0^2 Z_M(\mu_0)^2} \ , \end{aligned} \tag{4}$$

and no next-to-leading order corrections to the gaugino mass.

In a SDGM set up [3, 4] gaugino screening has dramatic consequences since the messenger mass X does *not* acquire an F-term at tree level, so there are not even leading corrections to the (visible) gaugino mass, which is then zero up until next-to-next-to-leading order [2]. This result concerns the contribution to the gaugino mass at linear order in F and at all orders in any hidden gauge or self-interaction coupling, but it does not exclude contributions of higher order in F. However, those are also all vanishing at leading order in the hidden gauge coupling [7].

Now, the question is whether it is possible to evade this argument, which seems quite general and robust. The answer is surprisingly simple: let us consider the physical gauge coupling above the messenger scale ($b'_0 = b_0 - T_M$)

$$R(\mu) = R(\mu_0) + \frac{b'_0}{16\pi^2}\log\frac{\mu^2}{\mu_0^2} + \frac{T_G}{8\pi^2}\log\frac{\mathrm{Re}S(\mu)}{\mathrm{Re}S(\mu_0)} - \sum_r \frac{T_r}{8\pi^2}\log\frac{Z_r(\mu)}{Z_r(\mu_0)} - \frac{T_M}{8\pi^2}\log\frac{Z_M(\mu)}{Z_M(\mu_0)} \; , \tag{5}$$

and now suppose that the messenger wave function $Z_M(\mu)$ experiences a SUSY breaking threshold at some scale between μ and μ_0. For instance, the messengers could couple to a hidden gauge group, whose gaugino obtains a mass. The latter can be seen as a spurionic F-term to the holomorphic hidden gauge coupling. Hence hidden gauge radiative corrections to the wave function renormalization of the messengers will propagate an F-term down to the visible gauge coupling function and gaugino screening would then not occur. This is actually exactly what happens in Z' mediation, except that the messenger's role is played by the MSSM chiral matter.

Of course, massless messengers cannot be the solution. Hence they must acquire a mass at some stage, for instance by higgsing. Well below this mass, the RG evolution is as in the case of massive messengers. However, if the messengers mass is sufficiently smaller than the scale of SUSY breaking, there is enough RG evolution between the two to produce a non-vanishing visible gaugino mass.

In essence, before higgsing the messengers will be massless because they will be in a chiral representation of the gauge groups. By consequence, they will have chiral couplings to the hidden and visible gauge groups, the wave function renormalization will be different for the two chiralities of the messengers and the scale matching crucial in obtaining the cancellation at next-to-leading order cannot be done. Below the higgsing scale, instead, one recovers a non-chiral spectrum and therefore the RG flow does not feed the visible gaugino mass anymore.

3 Direct evaluation of the gaugino mass

Here we give a different, more direct argument in favour of gaugino mass unscreening with chiral messengers.

We recall from [7] that in SDGM there are only two types of diagrams contributing to the gaugino mass, as displayed in Figure 2.

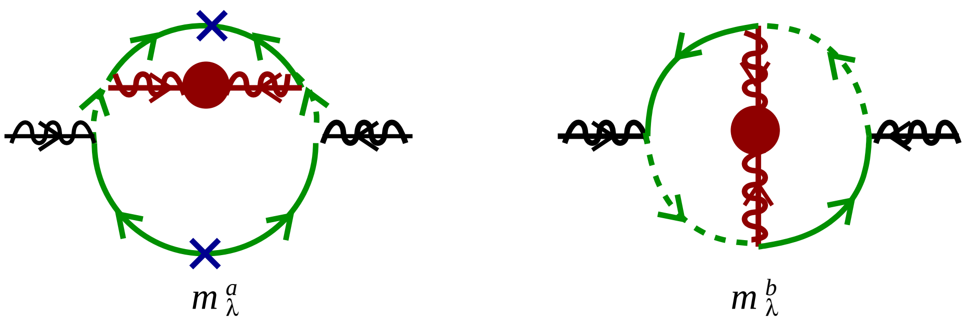

Figure 2: The diagrams contributing to the gaugino mass. The external line corresponds to the visible gaugino λ, the internal line with the blob attached corresponds to the propagator of the hidden gaugino (the blob encodes the exact hidden sector non-supersymmetric correction to the propagator), while all other internal lines correspond to messenger fields. The left diagram has two (supersymmetric) mass insertions, each one represented by a cross on the corresponding messenger fermionic line.

The result of gaugino mass screening (at leading order in the hidden gauge coupling, but to all orders in F) comes about by noticing that the two diagrams cancel each other exactly at zero momentum, and independently of the SUSY breaking current insertion on the hidden gaugino (chiral) propagator [7].

When messengers are chiral, and hence massless, the cancellation no longer holds for a trivial reason: one cannot write the first diagram, m^a_λ, since it would involve mass insertions on the fermionic messenger lines. Then, the visible gaugino mass is non zero and given by the massless limit of the second diagram, with messengers of only a single chirality running in the loop. A very similar diagram appears indeed in the context of Z' mediation [6].

It is quite straightforward to evaluate explicitly this diagram. From each chiral messenger, in the limit in which we can consider it massless, we obtain a contribution given by (suppressing group theory factors)

$$m_\lambda = 4g_v^2 g_h^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{l\cdot(l-k)}{l^4(l-k)^4} \frac{g_h^2 M B(k^2/M^2)}{k^2 + (g_h^2 M B(k^2/M^2))^2} \ . \qquad (6)$$

In the above, g_v and g_h are the couplings of the visible and the hidden gauge groups, respectively, in which the messengers are bifundamentals, M is a scale related to the SUSY breaking dynamics in the hidden sector, while $B(k^2/M^2)$ is the chiral correlator of the hidden sector fermionic current that determines the mass of the hidden gaugino λ_h, see [7, 8]. Note that we have resummed the hidden gaugino chiral propagator in order to avoid IR divergences.

First of all we evaluate by standard techniques the kernel

$$\int \frac{d^4l}{(2\pi)^4} \frac{l\cdot(l-k)}{l^4(l-k)^4} = \frac{1}{(4\pi)^2}\frac{1}{k^2} \ . \qquad (7)$$

Note that this is the correct $m \to 0$ limit of the kernel that one writes for messengers of mass m (after factoring out a power of m^2), and that was computed in [7].

We can now use this result to compute the visible gaugino mass. For definiteness, we approximate $B(k^2/M^2)$ by a step function, and we take the hidden gaugino mass to be $m_{\lambda_h} = g_h^2 M B(0)$ as we set ourselves in the regime $m_{\lambda_h} \ll M$. We get from eq. (6)

$$\begin{aligned} m_\lambda &= 4\frac{g_v^2 g_h^2}{(4\pi)^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} \frac{g_h^2 M B(k^2/M^2)}{k^2 + (g_h^2 M B(k^2/M^2))^2} \\ &= 4\frac{\alpha_v}{4\pi}\frac{\alpha_h}{4\pi} \int_0^{M^2} dk^2 \frac{m_{\lambda_h}}{k^2 + m_{\lambda_h}^2} \sim \frac{\alpha_v}{4\pi}\frac{\alpha_h}{4\pi} m_{\lambda_h} \log\frac{M^2}{m_{\lambda_h}^2} \ . \end{aligned} \qquad (8)$$

As already noticed, the messengers will eventually get a mass by some model-dependent dynamical mechanism (e.g. higgsing, confinement). However, assuming that the dynamical mass scale scale is much smaller than the hidden gaugino mass, the above expression will only have negligible corrections. Even in the case where the messengers' acquired mass is of the same order of m_{λ_h}, but still much smaller than M, it can be shown that the expression above will be corrected at most by an $\mathcal{O}(1)$ factor.

4 A model of chiral messengers

In this section we present a model that implements the ideas developed above. Our goal is not to present a complete phenomenologically viable model, but to show that the idea discussed in the previous sections can find a concrete realization.

The several gauge groups and chiral superfields needed in a model of SDGM can be easily encoded in a quiver gauge theory that can actually be found among those arising from D-branes at Calabi-Yau singularities [9]. The specific model we consider here can be obtained, for instance, by considering fractional D3-branes at a del Pezzo 3 singularity, and is depicted in figure 3.

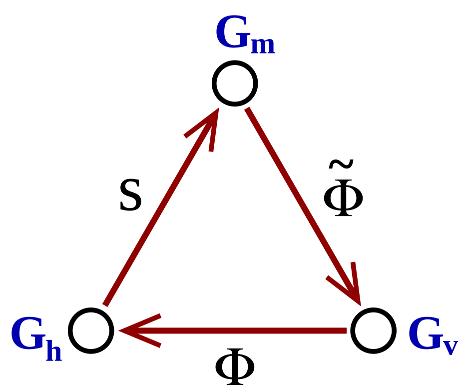

Figure 3: The quiver gauge theory arising at a dP_3 Calabi-Yau singularity describing the messenger sector and its interactions. Visible matter fields are attached to the group G_v, and can be engineered in terms of flavor D7-branes at the singularity. The SUSY breaking dynamics couples instead only to G_h.

There are three gauge groups, whose ranks should be one and the same to avoid gauge anomalies. In the following we assume $G_v = G_h = G_m = SU(5)$, having in mind applications to GUT theories. There are three different bifundamental superfields: the chiral superfield Φ charged under G_v and G_h, the chiral superfield $\tilde{\Phi}$ charged under G_m and G_v, and an extra superfield S charged under G_h and G_m. The superfields Φ and $\tilde{\Phi}$ transform in the 5 resp. the $\bar{5}$ of the GUT group G_v. In addition, there is a (unique) superpotential term

$$W = yS\tilde{\Phi}\Phi \ . \tag{9}$$

Of course, besides these fields there will be other chiral superfields charged only under the visible gauge group G_v, forming the chiral matter of the visible sector. Similarly, there will be extra dynamical fields affecting only the hidden sector G_h, that will give rise to supersymmetry breaking but whose detailed dynamics will not be addressed here. The presence of these fields will always be understood in the following but we will concentrate only on those fields charged under at least two of the groups. The third gauge group, G_m, is needed to make the whole theory free of gauge anomalies.

Notice in passing that one could also reduce the number of messengers as seen by the visible sector by replacing G_h and G_m with some lower rank group and attaching to them enough extra matter to cancel their cubic anomalies (or, in the case of $SU(2)$ the global anomaly arising from an odd number of fields).

The transition from the chiral messenger model at higher energies to a model where the messengers are eventually massive is done by giving a diagonal VEV v to S. Once S has a VEV, the messengers obtain a supersymmetric mass equal to yv, and G_h and G_m are higgsed to a diagonal $SU(5)$.

In the absence of SUSY breaking the off-diagonal combination of the two gauginos λ_h and λ_m would get a Dirac mass by mixing with the fermion in S, while the diagonal gaugino would stay massless.

The story changes if SUSY breaking is present and affects G_h. Then, before higgsing, the G_h gaugino λ_h already has a (Majorana) mass. After higgsing it will mix (negligibly if we assume $g_h v, g_m v \ll m_{\lambda_h}$) with the G_m gaugino λ_m, however the messengers Φ and $\tilde{\Phi}$, even if massive, will still couple to a different gaugino. Hence, to leading order, the contribution to the visible gaugino mass will be the one discussed in the previous section.

Let us see this is in a bit more detail. Our superfields are $S = v + \theta\sigma + \ldots$, $\Phi = \phi + \theta\psi + \ldots$ and $\tilde{\Phi} = \tilde{\phi} + \theta\tilde{\psi} + \ldots$. After higgsing, the terms in the SUSY Lagrangian bilinear in the relevant fermions are

$$\mathcal{L} \supset \mathcal{L}_{ferm} = i\sqrt{2}g\phi^*\psi\lambda_h - i\sqrt{2}gv\sigma\lambda_h + i\sqrt{2}gv\sigma\lambda_m - i\sqrt{2}g\tilde{\phi}^*\tilde{\psi}\lambda_m + yv\tilde{\psi}\psi \ . \quad (10)$$

For convenience, we have set the two couplings g_h and g_m to the same value g and have once again dropped the group theory factors. To the above Lagrangian, we have to add the SUSY breaking mass for the hidden gaugino

$$\mathcal{L}' = \tfrac{1}{2} m_{\lambda_h} \lambda_h \lambda_h \ . \tag{11}$$

The lagrangian $\mathcal{L}_{ferm}+\mathcal{L}'$ characterizes the fermionic sector of the theory. Note that the higgsing scale gv can be different from the messenger mass scale yv. Moreover, we have not yet assumed any specific relation between the different scales gv, yv and m_{λ_h}.

We are interested in computing the contributions to the visible gaugino mass. The diagrams are the ones in figure 2. Observe that the visible gaugino couples to the messengers, whose fermions have a Dirac mass in $\mathcal{L}_{ferm}$. The messengers then couple to the hidden and messenger gauginos λ_h and λ_m. The shorter way to perform the computation is to invert the quadratic part of the Lagrangian for the fermions (λ_h, λ_m, σ) and extract the two point functions

$$\begin{aligned}
\mathcal{B}_{hh} &\equiv \langle \lambda_h \lambda_h \rangle = \frac{m_{\lambda_h}(k^2+2g^2v^2)^2}{k^2(k^2+4g^2v^2)^2 + m_{\lambda_h}^2 (k^2+2g^2v^2)^2} \\
\mathcal{B}_{mm} &\equiv \langle \lambda_m \lambda_m \rangle = \frac{4g^4v^4 m_{\lambda_h}}{k^2(k^2+4g^2v^2)^2 + m_{\lambda_h}^2 (k^2+2g^2v^2)^2} \\
\mathcal{B}_{hm} &\equiv \langle \lambda_h \lambda_m \rangle = \frac{2g^2v^2 m_{\lambda_h}(k^2+2g^2v^2)}{k^2(k^2+4g^2v^2)^2 + m_{\lambda_h}^2 (k^2+2g^2v^2)^2} \ .
\end{aligned} \tag{12}$$

Note that all of them vanish in the supersymmetric limit $m_{\lambda_h} = 0$.

The two point functions computed above enter into the computation of the diagrams of figure 2 as the blobs in the internal lines. It is easy to see that in the diagram to the left of Figure 1, the internal gaugino line is $\langle \lambda_h \lambda_m \rangle$, while there are two diagrams corresponding to the one on the right, one with a $\langle \lambda_h \lambda_h \rangle$ line and the other with a $\langle \lambda_m \lambda_m \rangle$ line.

With a computation similar to the one in [7] we can obtain the resulting contribution to the visible gaugino mass

$$m_\lambda = 8 g_v^2 g^2 \int \frac{d^4k}{(2\pi)^4} \left(L_a(k^2,(yv)^2)\mathcal{B}_{hm} + L_b(k^2,(yv)^2) \frac{\mathcal{B}_{hh}+\mathcal{B}_{mm}}{2} \right) \ , \tag{13}$$

where L_a and L_b have been computed in [7], and we recall that $L_a = -L_b$. The expression (rescaled by $1/m^2$ with respect to [7]) is

$$L_b(k^2, m^2) = \frac{1}{2(4\pi)^2}\left(\frac{1}{k^2} + \frac{1}{k^2+4m^2} - \frac{16m^4}{[k^2(k^2+4m^2)]^{3/2}}\text{arctanh}\sqrt{\frac{k^2}{k^2+4m^2}}\right) . \tag{14}$$

Plugging the above kernel back into eq. (13), we finally obtain for the gaugino mass

$$m_\lambda = 4g_v^2 g^2 \int \frac{d^4k}{(2\pi)^4} L_b(k^2, (yv)^2)\left(\mathcal{B}_{hh} + \mathcal{B}_{mm} - 2\mathcal{B}_{hm}\right) . \tag{15}$$

Recalling the explicit form of the two point functions (12), one sees that the gaugino mass is logarithmically UV divergent. This is expected since we introduced an explicit soft supersymmetry breaking term (11). The natural ultraviolet cut off is the supersymmetry breaking scale M.

Unfortunately we cannot compute the integral (15) exactly. However, we can study some interesting limits. The combination multiplying the kernel L_b in eq. (15) is

$$\mathcal{B}_{hh} + \mathcal{B}_{mm} - 2\mathcal{B}_{hm} = \frac{m_{\lambda_h} k^4}{k^2(k^2+4g^2v^2)^2 + m_{\lambda_h}^2 (k^2+2g^2v^2)^2} . \tag{16}$$

First, the supersymmetric case is recovered for $m_{\lambda_h} \to 0$. In this limit the combination (16) and in particular each of the two point functions in (12) vanishes.

Gaugino screening can be recovered letting $gv \gg M$ with arbitrary yv. This corresponds to a higgsing at very high scale. The resulting effective theory is like the one studied in [7]. In this limit the integral (15) is vanishing as $\sim M^4/g^4v^4$.

In any other limit, and in particular as long as M is the highest scale in the model, it is obvious that (16) does not vanish and that there will be a contribution to the visible gaugino mass at this order.

We can consider the regime where $gv \ll m_{\lambda_h}$ and also $yv \ll m_{\lambda_h}$. In this limit the higgsing can be considered as a subdominant effect with respect to SUSY breaking, at least as far as the fermionic sector is concerned. This can be seen by analyzing the leading contribution for the two point functions

$$\mathcal{B}_{hh} = \frac{m_{\lambda_h}}{k^2 + m_{\lambda_h}^2} \quad , \quad \mathcal{B}_{mm} = \frac{m_{\lambda_h} 4g^4v^4}{k^4(k^2+m_{\lambda_h}^2)} \quad , \quad \mathcal{B}_{hm} = \frac{m_{\lambda_h} 2g^2v^2}{k^2(k^2+m_{\lambda_h}^2)} . \tag{17}$$

At leading order only $\mathcal{B}_{hh}$ contributes to the integral (15), and we can compute it as

$$m_\lambda = \frac{4g_v^2 g^2 m_{\lambda_h}}{(4\pi)^4}\int_0^{M^2} dk^2 \frac{1}{k^2+m_{\lambda_h}^2} \sim \frac{4g_v^2 g^2 m_{\lambda_h}}{(4\pi)^4}\log\frac{M^2}{m_{\lambda_h}^2} , \tag{18}$$

which is the same result as in section 3. The first corrections to this expression can be easily computed expanding the kernel (14) and the combination (16) and performing the integral. They scale like $(yv)^2/m_{\lambda_h}$ and $(gv)^2/m_{\lambda_h}$. The analytic result (18) is then robust for $gv \ll m_{\lambda_h}$ and $yv \ll m_{\lambda_h}$.

In the most general case, $gv, yv, m_{\lambda_h} \ll M$, we can perform the integral (15) numerically. One eventually obtains a result which is essentially of the form (18) with the log factor replaced by a smaller $\mathcal{O}(1)$ factor.

This calculable model makes it clear that unscreening of the visible gaugino mass is possible, and actually still holds even when taking into account that the messengers eventually do get a mass. What is important of course is that there is ultimately a sizable hierarchy between the scale of higgsing, v and the SUSY breaking scale M.

One might be turned off by the fact that the scale of the VEV v, which has to be small with respect to the SUSY breaking scale M, is essentially introduced by hand.[1] In fact, our model presents itself an alternative possibility. In the absence of a higgsing, the gauge group G_m will confine at a scale Λ_m which we can naturally take to be smaller than or of the order of m_{λ_h}. At energies well above Λ_m the theory is chiral and the arguments of the previous section, with massless chiral messengers, apply. At energies below Λ_m the theory confines and we turn to an effective description. The group G_m has 5 colors and 5 flavors, hence we can set ourselves on the baryonic branch of the moduli space, where the meson superfields have zero VEVs. The mesons will act as composite messengers

$$\tilde{\Phi}_{\text{comp}} = \frac{1}{\Lambda_m} S \tilde{\Phi} \ . \tag{19}$$

and both messengers, Φ and $\tilde{\Phi}_{\text{comp}}$, will get a mass of order $y\Lambda_m$. We are now in a situation of semi-direct gauge mediation in all similar to the one described in [7], so we expect to have no contribution to the visible gaugino mass below this scale.

Clearly in this strongly coupled model the transition around Λ_m is less under

[1] The VEV v is essentially a (Goldstone) modulus. One can try to fix it in several ways, the most straightforward being promoting the $SU(5)$ groups to $U(5)$ (compensating the mixed anomalies by a Green-Schwarz-like mechanism) and then turning on a FI parameter for the off-diagonal $U(1)_{h-m}$.

control, but the gross features of the visible soft spectrum should be quite similar to the previous case.

5 Sfermion masses and phenomenology

We have not yet discussed sfermion masses. This is a model-dependent issue and chiral messenger models as the ones we have discussed here have a potential problem, in this respect.

In SDGM one gets, generically, a non-supersymmetric contribution to the messenger mass squared which provides a non-vanishing supertrace. If this contribution is negative it can overwhelm the supersymmetric messenger mass and make the messengers tachyonic. If on the other hand the contribution is positive, there can instead be problems with the sfermions of the visible sector that will generically acquire negative squared masses, since the latter is proportional to minus the supertrace of the messenger mass matrix squared [7] (see also [10]).

Notice that the contribution to the supertrace depends on the hidden sector current correlators C_s^h [7, 8], which are thus not directly related to the correlator B^h entering the expressions for the gaugino mass.

Let us discuss the two above possibilities in turn. If the messenger supertrace is positive, the sfermions are all tachyonic. In this scenario, all we can do is to find a mechanism to suppress the sfermion masses. There are several such mechanisms, for instance sequestering [11] (see also [12]) by a large extra dimension [13]–[16] or by coupling to a conformal sector [17, 18], deconstruction [19], or holographic gauge mediation [20, 21]. All such models eventually lead to gaugino mediation, where the sfermion squared masses are positive and are generated by RG flow below the scale of the visible gaugino. Hence, in the framework of gaugino mediation, we could use our model of chiral SDGM to generate the gaugino mass in the first place. Otherwise, the phenomenology is the same as the one of a generic gaugino mediated scenario.

In the other scenario, where the messenger supertrace is negative, the sfermion squared masses are positive but the messengers must have a sufficiently positive SUSY mass to compensate for the negative supertrace. In this case, some more tuning is needed. Indeed, we need to ensure that the messengers are not tachyonic by enforcing the bound

$$yv > \frac{\alpha_h}{4\pi} M \ . \tag{20}$$

We can now compare gaugino versus generic sfermion masses m_{sf} (again not paying attention to the group theory factors that can be easily reinstated). The computation of the sfermion masses is unaffected by the fact that the messengers are chiral and we can borrow the result from [7] where we set to yv the messengers mass

$$m_{sf}^2 \sim \left(\frac{\alpha_v}{4\pi}\right)^2 \left(\frac{\alpha_h}{4\pi}\right)^2 M^2 \log \frac{M^2}{(yv)^2} \ , \tag{21}$$

where the limit $yv \ll M$ is understood. This is to be compared with the visible gaugino mass where we have

$$m_\lambda \sim \left(\frac{\alpha_v}{4\pi}\right) \left(\frac{\alpha_h}{4\pi}\right) m_{\lambda_h} \log \frac{M^2}{m_{\lambda_h}^2} \ . \tag{22}$$

Assuming that the log factors and other corrections are of order unity, we see that the ratio between gaugino and sfermion masses is given by[2]

$$\frac{m_\lambda}{m_{sf}} \sim \frac{m_{\lambda_h}}{M} \sim \frac{\alpha_h}{4\pi} \ . \tag{23}$$

It is then possible to achieve a not too split visible spectrum if $\frac{\alpha_h}{4\pi}$ is not too small. A reasonable ordering of scales that one could aim for is the following

$$m_\lambda < m_{sf} < m_{\lambda_h}, gv < yv < M \ , \tag{24}$$

with one or two orders of magnitude between each scale. For instance, we could take $\alpha_h/4\pi \sim 10^{-2}$. If we start then with $m_\lambda \sim 10^2$ GeV, we get $m_{sf} \sim 10^4$ GeV, $m_{\lambda_h} \sim g_h v \sim 10^5$ GeV, $yv \sim 10^6$ GeV and finally $M \sim 10^7$ GeV. (The values given for the visible sector particles are to be considered as boundary conditions for the MSSM RG flow, as usual.) Note that we have to require that the coupling y is at the edge of perturbativity, $y^2/(4\pi)^2 \lesssim 1$. We conclude that even in the case of negative supertrace it is possible to produce a soft spectrum which is not too hierarchical, although in a small region of the parameter space.

[2]Here and above we have assumed that the hidden sector correlators are still perturbative, hence the factors of $1/4\pi$. If they arise directly from strongly coupled dynamics one should omit this extra factor.

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