Soft Spectrum in Yukawa-Gauge Mediation

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\textbf{Abstract:} We introduce a model independent parametrization for a subclass of
gauge mediated theories, which we refer to as Yukawa-gauge mediation. Within this
formalism we study the resulting soft masses in the visible spectrum. We find general
expressions for the gaugino and scalar masses. Under generic conditions, the gaugino
mass is screened, vanishing at first order in the SUSY breaking scale.
1. Introduction

During the last few years there has been fervent activity in the study of the phenomenology of gauge mediated models of supersymmetry breaking in the SSM (see [1, 2] for reviews and references). The recent developments in finding calculable and metastable vacua with dynamical supersymmetry breaking [3] have opened up the possibility of building models for gauge mediated scenarios which are not only predictive and weakly coupled, but also have a dynamical UV completion. Typically these models reduce in the infrared to effective weakly coupled models of pure chiral fields which are generalization of O’Raifeartaigh models, whose phenomenology can be studied (see e.g. [4]).
With the incoming data of the LHC, it is relevant to study the signature of specifics mechanisms of mediation of supersymmetry breaking, irrespectively of the intricate details of the hidden sector. This has been the approach of general gauge mediation (GGM) \[5\] and subsequent works \[6\] (see for instance \[7\] for some collider studies on the GGM parameter space). In particular, it is interesting to investigate the generic structure and hierarchies in the resulting soft spectrum, in order to have smoking guns indicating particular classes of models in the complicated analysis of the soft SSM parameters.\(^1\)

The analysis of GGM \[5\] has identified the complete parameter space that is allowed by gauge mediation. A basic effective model, which already cover a relevant portion of the parameter space of general gauge mediation, is minimal gauge mediation. In minimal gauge mediation a pair of vector-like chiral fields, the messengers, are charged under the SSM gauge groups and couple in the superpotential to a spurion singlet superfield \(X\), whose scalar and \(F\)-term components take expectation values. This simple structure is realized, typically with some extra superfields and interactions, in many effective models of metastable dynamical supersymmetry breaking (see \[9\] and references therein).

In this paper we investigate a natural generalization of this scenario which we argue can often emerge as the effective theory in models with dynamical supersymmetry breaking. It is indeed natural to generalize minimal gauge mediation by replacing the spurion singlet with a dynamical field \(X\) and coupling the latter, via superpotential interaction only, to some hidden sector chiral operator. In this setup, we will assume that the field \(X\) does not get any tree level scalar or \(F\)-term vacuum expectation value. The superfield \(X\) senses the supersymmetry breaking at loop level via the superpotential interaction. In the limit in which this interaction is switched off the visible sector has a complete supersymmetric spectrum. These assumptions make the setup substantially different from minimal gauge mediation and heavily affect the consequent phenomenology.

Concrete models realizing similar scenarios have been, for instance, considered in \[10, 11\]. The purpose of this paper is to provide a general analysis of the resulting soft spectrum in the SSM for this class of models, here referred to as Yukawa-gauge mediation.

\(^1\)In this context \(R\)-symmetry plays a crucial role. If unbroken, it forbids gaugino mass generation. Moreover, it has been recently shown that \(R\)-symmetry can further control the suppression of gaugino mass in general chiral models, leading to important no go theorems \[8\].
Note that our scenario is included in GGM [5] and also in the general messenger gauge mediation of [12]. However, in our working assumption of vanishing $F$-term for $X$, the contribution to the gaugino mass computed in [12] is not present. Hence our analysis deals with the next order correction to the result of [12].

Another more technical motivation to study Yukawa-gauge mediation is that it may provide a class of models where the ratio between the gaugino and scalar masses is of one loop in a weak coupling expansion. Many models have gauginos which are lighter than the scalars because of $F$-term suppression, e.g. [13]. On the contrary, a one loop order hierarchy cannot be easily realized in general. Semi-direct gauge mediation [14] would be a natural candidate, but it suffers from gaugino screening [15, 16, 17].

Our computation shows that also in Yukawa-gauge mediation the gaugino mass is screened. This agrees with and generalizes the results of [15, 18], obtained using wave function renormalization techniques. Interestingly, we find that the gaugino mass screening is realized as a suppression in powers of the SUSY breaking scale and not in loop factors. Therefore this class of models gives a further peculiar realization of the gaugino screening phenomenon.

The organization of the paper is the following. In the next section we introduce the model and we set our parametrization for the hidden sector. In section 3 we work out the expression of the visible gaugino mass, compute the two loops integrals involved and discuss the resulting suppression. Then we add a superpotential term that breaks explicitly $R$-symmetry and we compute the resulting new contribution to the gaugino mass. In section 4 we study the contributions to the scalar mass, giving a model independent answer, and we comment on the corrections to the messenger masses. In section 5 we provide two minimal realizations of Yukawa-gauge mediation and we then conclude in section 6. We leave to the appendix most of the details of the loop integrals computations.

2. Setup of Yukawa-gauge mediation

We would like to study in a model independent way the pattern of soft masses generated by a Yukawa-gauge mediated model. The set up is the following: there is a visible sector which communicates via gauge interactions to a pair of vector-like and weakly coupled messenger chiral superfields. The messengers are charged only under the SSM gauge group (that we take to be $U(1)$). They interact through a
trilinear superpotential with a chiral field, that we call $X$. This chiral field then couples via superpotential interaction to another chiral operator $O$, of dimension two. The SUSY breaking effects are encoded in the one and two point functions of this chiral operator $O$. The resulting visible soft mass terms will be determined as functions of them.

We would like to investigate how these simple assumptions constrain the allowed visible sector spectrum.

The Lagrangian of the model is
\begin{equation}
\mathcal{L} = \mathcal{L}_{\text{vis}} + \int d^4 \theta \; \Phi \mathcal{Y}_{\text{SSM}} \Phi + \tilde{\Phi} \mathcal{Y}_{\text{SSM}} \tilde{\Phi} + X^\dagger X \\
+ \int d^2 \theta \; m \Phi \tilde{\Phi} + \lambda_x X \Phi \tilde{\Phi} + \lambda_o X O + \text{h.c.} \tag{2.1}
\end{equation}

Note that it is always possible to perform a rotation of the phases of the different fields and of $O$ in such a way that the mass parameter $m$ and the couplings $\lambda_x$ and $\lambda_o$ are all real. Here we do not add an explicit mass term for $X$ in order to not introduce extra dimensionful parameters. This is the typical situation in many effective chiral models for supersymmetry breaking, where the field $X$ is a pseudo-modulus which acquires mass only through loop corrections.

In the limit in which the SUSY breaking hidden sector operator $O$ decouples from $X$, $\lambda_o \to 0$, we are left with a supersymmetric visible sector. On the other hand, supersymmetry is also restored in the limit $\lambda_x \to 0$, where the chiral field $X$ decouples from the messengers. Supersymmetry in the SSM can therefore be recovered in two different ways. This ambiguity is somehow solved by the fact that, in the expressions for the visible sector soft terms, the two couplings $\lambda_x$ and $\lambda_o$ always appear in pairs.\(^2\) One could then simply say that visible sector supersymmetry is restored in the limit $\lambda_x \lambda_o \to 0$. Finally, note that supersymmetry in the SSM is also recovered by sending to infinity the messenger masses.

The supersymmetry breaking effects can be encoded in the one and two point functions of the chiral operator $O = (O, \psi^o, F_o)$. Lorentz invariance constrains them

\(^2\)Our set up is similar to the semi-direct gauge mediation of [16]. In semi-direct gauge mediation the messengers are coupled to an hidden gauge field, which is also coupled to a non supersymmetric current. Here the messengers are coupled to a chiral field $X$ which is coupled to a non supersymmetric chiral operator $O$. The chiral field $X$ plays the role of the hidden gauge field and the chiral operator $O$ the role of the hidden sector current of semi-direct gauge mediation.
hidden sector two point functions to be of the form
\[ \langle O \rangle = O_o, \quad \langle F_o \rangle = f_o, \]
\[ \langle O(p)O(-p) \rangle = G_0(p^2), \quad \langle O(p)O^\dagger(-p) \rangle = G_1(p^2), \]
\[ \langle \psi_\alpha(p)\psi_\beta^\dagger(-p) \rangle = \epsilon_{\alpha\beta}G_2(p^2), \quad \langle \overline{\psi}_\alpha(p)\overline{\psi}_\beta^\dagger(-p) \rangle = p_\mu\sigma^\mu_{\alpha\beta}G_3(p^2), \quad (2.2) \]
\[ \langle O(p)F_o(-p) \rangle = G_4(p^2), \quad \langle O(p)F_o^\dagger(-p) \rangle = G_5(p^2), \]
\[ \langle F_o(p)F_o^\dagger(-p) \rangle = -p^2G_6(p^2), \quad \langle F_o(p)F_o(-p) \rangle = G_7(p^2), \]
where the unknown functions \( G_i(p^2) \) can depend on different scales of the hidden sector. Here \( G_2, G_4 \) and \( G_5 \) have mass dimension one, \( G_7 \) have mass dimension two, and the others are dimensionless. If supersymmetry is unbroken
\[ G_0 = G_5 = G_7 = 0, \quad G_1 = G_3 = G_6, \quad G_2 = G_4. \quad (2.3) \]

Starting from (2.1) we can derive the effective Lagrangian for \( X \)
\[ \delta \mathcal{L}_X = \lambda_o (O_o F_x + f_o x + h.c.) + \lambda_6^2 (G_6 x^\dagger \square x - iG_3 \overline{\psi}_x \sigma^\mu \partial_\mu \psi_x + G_1 F_x F_x^\dagger + G_7 x^2 + G_0 F_x F_x + G_5 F_x x^\dagger + G_4 x F_x - \frac{1}{2}G_2 \overline{\psi}_x \psi_x + h.c.) \quad (2.4) \]
One point functions for the operator \( O \) typically induce an \( F \)-term for the field \( X \), resulting in minimal gauge mediation. Hence, from now on, we assume that the one point functions of the operator \( O \) are vanishing, i.e. \( O_o = f_o = 0 \). This can be enforced by a discrete \( Z_2 \) symmetry of the hidden sector under which the operator \( O \) is charged. This symmetry is broken by the coupling of the operator \( O \) to \( X \) and to the messengers. However this \( Z_2 \) is enough to loop suppress the generation of a tadpole for the \( F \)-term of \( X \). This discrete symmetry is the analogous of the messenger parity of [19] reviewed in GGM [5].

An equivalent formulation consists in encoding the SUSY breaking effect directly in the two point functions of the chiral field \( X \). At first order in the insertion of the hidden sector two point functions \( G_i \), the propagators for the field \( X \) are
\[ \langle F_x(p)F_x(-p) \rangle = \lambda_o^2 G_0(p^2), \quad \langle F_x(p)F_x^\dagger(-p) \rangle = \lambda_o^2 G_1(p^2), \]
\[ \langle \psi_x(p)\psi_x^\dagger(-p) \rangle = \frac{\lambda_6^2 G_2(p^2)\delta_{\alpha\beta}}{p^2}, \quad \langle \overline{\psi}_x(p)\overline{\psi}_x^\dagger(-p) \rangle = \frac{\lambda_o^2 G_3(p^2)p_\mu\sigma^\mu_{\alpha\beta}}{p^2}, \]
\[ \langle x(p)F_x(-p) \rangle = -\frac{\lambda_5^2 G_4(p^2)}{p^2}, \quad \langle x(p)F_x^\dagger(-p) \rangle = -\frac{\lambda_o^2 G_5(p^2)}{p^2}, \]
\[ \langle x(p)x^\dagger(-p) \rangle = -\frac{\lambda_7^2 G_6(p^2)}{p^2}, \quad \langle x(p)x(-p) \rangle = \frac{\lambda_7^2 G_7(p^2)}{p^4}, \quad (2.5) \]
where we have assumed no mass terms for $X$.

Finally, note that in the ansatz (2.1) there is a $U(1)_R$ symmetry under which the operator $O$ has $R$-charge two and the field $X$ has $R$-charge zero. This symmetry should be broken down to $Z_2$ in order to generate visible majorana gaugino mass. Here the only sources for $R$-symmetry breaking are the two point functions $G_i$. Non-vanishing $G_2, G_4, G_5$ break this $R$-symmetry to $Z_2$, $G_0$ breaks it to $Z_4$, while the other $G_i$ preserve it.

3. Gaugino mass

In this section we compute the two loop contributions to gaugino mass, showing the presence of gaugino mass screening in Yukawa-gauge mediation. Details about computation of the loop integrals are reported in the appendix A. Here we only discuss the different contributions, the non-trivial cancellations that occur and the resulting soft term.

There are six non-equivalent graphs contributing to the gaugino mass. They are depicted in figure 1. Note that these involve only the functions $G_2, G_4$ and $G_5$. As already observed, these are indeed the only two point functions among (2.2) which break $R$-symmetry down to $Z_2$.

![Figure 1: The six graphs contributing to visible gaugino mass at leading loop order. The internal line involving a bubble is a propagator of components of the chiral superfield $X$. All the remaining internal lines correspond to propagators of the messengers.](image)

There are two graphs with $G_5$

\begin{align}
    m_{[a,5]} &= 8g^2\lambda^2_2\lambda_0^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2 G_5(k^2)}{k^2[l^2 + m^2]^2[(l - k)^2 + m^2]} , \\
    m_{[b,5]} &= -4g^2\lambda^2_2\lambda_0^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{l \cdot (l - k) G_5(k^2)}{k^2[l^2 + m^2]^2[(l - k)^2 + m^2]^2} .
\end{align}
Performing the analytic integration over the loop momentum $l$ it is easy to check that the two contributions cancel each other precisely.\footnote{The structure of the integrals is exactly as those in \cite{16}.} We conclude that the two point function $G_3$ does not contribute to gaugino mass at the leading loop order.

Of the remaining four graphs in figure 1, two involve $G_2$ and two $G_4$. Their explicit expression can be found in appendix A. They combine in pairs to give the following integrals in terms of the difference $G_2 - G_4$

$$m_{[a]} = -8g^2 \lambda_2^2 \lambda_0 \frac{1}{2(2\pi)^4} \int \frac{d^4k}{k^2} \frac{d^4l}{l^2} \frac{m^2}{k^2} (G_2(k^2) - G_4(k^2)),$$

$$m_{[b]} = -4g^2 \lambda_2^2 \lambda_0 \frac{1}{2(2\pi)^4} \int \frac{d^4k}{k^2} \frac{d^4l}{l^2} \frac{m^2}{k^2} (G_2(k^2) - G_4(k^2)),$$

The integration over the loop momentum $l$ can be done analytically and we can write

$$m_{[a,b]} = -4g^2 \lambda_2^2 \lambda_0 \frac{1}{16\pi^2 m^2} \int \frac{d^4k}{k^2} \frac{1}{4(2\pi)^4} L_{[a,b]}(k^2/m^2) (G_2(k^2) - G_4(k^2)),$$  \hspace{1cm} (3.3)

where the kernels $L_{[a,b]}$ are defined and computed in appendix A. The gaugino mass is then given by

$$m_\lambda = m_{[a]} + m_{[b]} = -4g^2 \lambda_2^2 \lambda_0 \frac{1}{16\pi^2} \frac{1}{m^2} \int \frac{d^4k}{k^2} \frac{1}{4(2\pi)^4} L(k^2/m^2) (G_2(k^2) - G_4(k^2)),$$  \hspace{1cm} (3.6)

with

$$L(x) = L_{[a]}(x) + L_{[b]}(x) = \frac{x(4 + x) + 4 \sqrt{x(4 + x)} \ Arcth \left( \frac{2}{\sqrt{4 + x}} \right)}{x(4 + x)^2}. \hspace{1cm} (3.7)$$

This is our general result for the two loop visible gaugino mass in Yukawa-gauge mediation.

Using a formulation similar to the one in \cite{20}, we can rewrite this non-vanishing contribution as

$$m_\lambda = -g^2 \lambda_2^2 \lambda_0 \frac{1}{16\pi^2} \frac{1}{m^2} \int \frac{d^4k}{k^2} \frac{d^4l}{l^2} \frac{m^2}{k^2} L(k^2/m^2)(\{Q_\alpha, \bar{Q}_\alpha, F_o O\}). \hspace{1cm} (3.8)$$

Note that arguments in \cite{20} show that the combination $G_2 - G_4 \simeq \{Q_\alpha, \bar{Q}_\alpha, F_o O\}$ is subleading in the SUSY breaking scale. On the other hand, $G_5$ is at leading order \cite{20}, but its contribution to the gaugino mass cancels out as we have just discussed. Hence, whatever the integral over $k^2$, we expect the gaugino mass to be suppressed in this class of models.
We conclude that, generically, in Yukawa-gauge mediation the gaugino mass receives contributions that are at the leading loop order (at least two loops), but subleading in the SUSY breaking scale. This is in agreement with the gaugino mass screening result derived via analytic continuation in superspace techniques \[15\], through which only the leading effects in the SUSY breaking scale are captured. In our computation we obtain an expression encoding the next order contributions in the SUSY breaking scale, for the generic set up explained in section 2.

3.1 Breaking $R$-symmetry explicitly

The Lagrangian in section 2 has a $U(1)_R$ symmetry under which $R[\mathcal{O}] = 2$ and $R[X] = 0$. As a consequence, the two point functions that can enter into the expression for the visible gaugino mass are only the ones that break the $R$-symmetry down to $Z_2$. Note however that the cancellation of the $G_5$ contribution and the resulting suppression in the SUSY breaking scale of the gaugino mass is not a direct consequence of $R$-symmetry.

Nevertheless, we could modify the previous scenario by adding an explicit $R$-symmetry breaking term in the superpotential for the chiral superfield $X$. The most generic renormalizable superpotential is

$$\Delta W = \frac{m_x}{2} X^2 + \frac{\lambda_{\text{def}}}{3} X^3. \quad (3.9)$$

This breaks completely the $R$-symmetry and, as a consequence, other $G_i$ can now enter into the gaugino mass computation.

Figure 2: The additional graph contributing to visible gaugino mass with the superpotential deformation $\Delta W$.

After a careful analysis, one can show that the only extra diagram that is now generated at two loop is the one in figure 2. This gives an extra contribution to the
gaugino mass

\[ \Delta m_\lambda = 4g^2 \lambda_x \lambda_{def} \lambda_o \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{mG_7(k^2)}{[k^2 + m_x^2][l^2 + m^2]^3} \]

which is typically not subleading in the SUSY breaking scale.

This is not in disagreement with the result of [15]. In fact, in this particular setting, we are effectively generating a tadpole for the \( F \)-term of \( X \) at one loop and the contribution to the gaugino mass is expected to be proportional to this loop generated \( F \)-term.

4. Scalar masses

In this section we compute the induced soft masses for the scalars of the SSM.

The model at hand is a subclass of general gauge mediation [5]. We can express the scalar masses as a loop integral over a combination of the functions \( C_i \)

\[ m^2_{s_f} = -g^4 \int \frac{d^4p}{(2\pi)^4} \frac{1}{p^2} (3C_1(p^2) - 4C_{1/2}(p^2) + C_0(p^2)) , \]  

where the \( C_i \) can be computed by considering the quantum corrections to the propagators of each component of the vector superfield. Following the same strategy of [16], we can compute the \( C_i \) as functions of the unknown \( G_i \) and, exchanging the order of integration and performing some of the loop integrals, re-express the scalar masses as one loop integral of the functions \( G_i \) convoluted with a non-trivial kernel.

We argue, following the reasoning in [12], that the final answer is of the form

\[ m^2_{s_f} = \frac{g^4 \lambda_x^2 \lambda_o}{(16\pi^2)^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^4} S(k^2/m^2)(Q^4(\langle O^1O \rangle) = \]

\[ = \frac{g^4 \lambda_x^2 \lambda_o^2}{(16\pi^2)^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} S(k^2/m^2)(G_1(k^2) + 2G_3(k^2) + G_6(k^2)) , \]  

thus depending only on the combination \( G_1 - 2G_3 + G_6 \).

In order to compute the kernel \( S(k^2) \) it is sufficient to focus only on the contribution to (4.1) given by \( G_1 \). To every \( C_i \) corresponds a two loop integral, involving \( G_1 \), that we denote as

\[ C_i^1(p^2) = \lambda_x^2 \lambda_o^2 \int \frac{d^4l}{(2\pi)^4} \frac{d^4k}{(2\pi)^4} S_{i,1}(l, k, p, m) G_1(k^2) . \]
These expressions and in particular the functions $S_{i,1}$ are computed in appendix B. Matching the formula (4.1) and (4.2)

$$\frac{\lambda_x^2 \lambda_o^2}{(16\pi^2)^2} \int \frac{d^4k}{(2\pi)^4k^2} S(k^2) G_1(k^2) = -\int \frac{d^4p}{(2\pi)^4p^2} \left( C_0^1(p^2) - 4C_{1/2}^1(p^2) + 3C_1^1(p^2) \right),$$

we can identify the kernel

$$\frac{1}{(16\pi^2)^2} \frac{S(k^2)}{k^2} = -\int \frac{d^4p}{(2\pi)^4p^2} \frac{d^4l}{(2\pi)^4 l^2} (S_{0,1}(l,k,p,m) - 4S_{1/2,1}(l,k,p,m) + 3S_{1,1}(l,k,p,m)).$$

(4.4)

This integration can be performed only numerically, order by order in an expansion in $k^2$.

The final answer is that the kernel $S(k^2/m^2)$ is a positive function of $k^2/m^2$, with the following expansion for small momenta

$$S(x \to 0) = 4x - \frac{4}{9}x^2 + \ldots$$

(4.5)

Moreover, $S(k^2/m^2)$ behaves logarithmically at large momenta, ensuring the convergence of the integral.

The expression at small momenta provides a consistency check of our computation. Indeed, we can recover the case of minimal gauge mediation by setting $G_6 = G_3 = 0$ and $G_1 \simeq |F|^2(2\pi)^4\delta^4(k)$. In this particular limit the diagrams involved in the computation are exactly the diagrams of minimal gauge mediation, at first order in the supersymmetry breaking parameter $F$. Plugging this ansatz in (4.2) the integral over $k^2$ can be done trivially and our result for $S(k^2/m^2)$ gives

$$m_{sf}^2 \simeq \frac{g^4\lambda_x^2 \lambda_o^2 |F|^2}{(16\pi^2)^2 |m|^2},$$

which coincides with minimal gauge mediation if we consider that the effective $F$-term is $F_{eff} = F\lambda_x \lambda_o$.

We conclude that (4.2) is the model independent expression for the scalar masses in Yukawa-gauge mediation. The sign of the scalar masses is then determined by the sign of the combination $G_1 - 2G_3 + G_6$.

4.1 Messenger mass corrections

For completeness we can compute the corrections to the messenger mass matrix induced by the coupling with the hidden sector.

At one loop the diagonal entries of the messenger mass matrix are corrected by the graphs in figure 3, which read

$$\delta m_{\phi\phi}^2 = -\lambda_x^2 \lambda_o^2 \int \frac{dk^4}{(2\pi)^4} \frac{G_1 - 2G_3 + G_6}{[k^2 + m^2]}.$$
Figure 3: Graphs contributing to the diagonal element of the mass matrix for the messengers.

The supertrace over the messengers sector receives correction proportional to this contribution. Notice that it has the opposite sign with respect to the visible sector scalar squared masses. This is a common feature of gauge mediated models with messengers [21].

The off diagonal one loop correction is

$$\delta m^2_{\Phi\Phi} = 2\lambda^2 \lambda_o \int \frac{dk^4}{(2\pi)^4} \frac{m(G_2 - G_4 - G_5)}{k^2[k^2 + m^2]}.$$  (4.7)

These expressions should be taken into account in explicit models of Yukawa-gauge mediation to check that the messengers do not turn tachionic due to quantum corrections.

5. Examples of Yukawa-gauge mediation

5.1 Example 1: a toy model

A very basic model of Yukawa-gauge mediation can be obtained introducing one single extra chiral field $Y$ coupled to a spurion $M_0 = m_0 + \theta^2 f$ and setting $O = MY$.

In such a way

$$W_o = \lambda_o MXY$$  (5.1)

has simply the form of a mass term for $X$ and $Y$. The Lagrangian for the chiral field $Y$ is

$$\mathcal{L} = \int d^4\theta \; Y^\dagger Y + \int d^2\theta \; \frac{1}{2} M_0 Y^2 + h.c.$$  (5.2)

A parity symmetry for $Y$, that forbids $F$-term generation for $X$, is present. Here we have two supersymmetric scales $M$ and $m_0$, together with a SUSY breaking scale $f$. 

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which we take all real for simplicity. We will consider the SUSY breaking scale $\sqrt{f}$ to be always the smallest.

Note that this is clearly only a toy model since the interaction of $X$ with the hidden sector is effectively a mass term and, furthermore, the hidden sector consists of free chiral fields only. Nevertheless, this toy model is sufficient to highlight the basic features of Yukawa-gauge mediation that we have found with the general analysis, i.e. a suppression of the two loop gaugino mass in powers of the SUSY breaking scale.

All the functions $G_i$ in this model, at the lowest non-vanishing order in the SUSY breaking scale, are generated at tree level. The quantities we need to give an estimate of the gaugino mass are:

$$G_2(p^2) - G_4(p^2) = M^2 \frac{m_0 f^2}{[p^2 + m_0^2]^3},$$

$$G_5(p^2) = -M^2 \frac{m_0 f}{[p^2 + m_0^2]^2},$$

$$G_7(p^2) = M^2 \frac{m_0 f}{[p^2 + m_0^2]^2}. \quad (5.3)$$

We have also written the explicit expression for $G_5$ to illustrate that it is first order in the SUSY breaking scale. However, as proven earlier, this does not contribute to the visible gaugino mass.

We start from (3.6)

$$m_\lambda = -4g^2 \lambda_x^2 \lambda_o^2 \frac{1}{16 \pi^2 m^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} L(k^2/m^2) (G_2(k^2) - G_4(k^2)) \quad (5.4)$$

with $L(k^2/m^2)$ given by (3.7) and we plug in the explicit expression of $G_2(p^2) - G_4(p^2)$. In the limit $m \gg m_0 \gg \sqrt{f}$, the gaugino mass at leading order evaluates to

$$m_\lambda \approx -\frac{g^2 \lambda_x^2 \lambda_o^2 M^2 f^2}{(16\pi^2)^2 m^2 m_0^3}. \quad (5.5)$$

This is at two loop order, but it is at second order in the SUSY breaking scale $f$. Thus we see the suppression we have explained in the previous sections.

In the case $R$-symmetry is broken explicitly by the deformation (3.9), the gaugino mass has the extra correction (3.10) involving $G_7$. In the same limit as above, with also $m_x \sim m$, this reads

$$\Delta m_\lambda \approx -\frac{g^2 \lambda_x \lambda_{def} \lambda_o^2 M^2 m_0^2 f}{(16\pi^2)^2 m m_0^3}. \quad (5.6)$$
Observe that as expected such a correction is still at two loop order, but at leading order in the SUSY breaking scale. This dominates over (5.5) in the range of parameters \( \frac{f}{m_0} < \frac{m^3}{m_0^2} \).

For completeness we can also give an estimate of the sfermion masses generated in this toy model. The combination of \( G_i \) functions we need is

\[
G_1(p^2) - 2G_3(p^2) + G_6(p^2) = \frac{M^2 (p^2 - m_0^2) f^2}{p^2 (m_0^2 + p^2)^3} \tag{5.7}
\]

This expression dies off for momenta larger than \( m_0 \), that we can use as an effective cutoff in the convolution with the kernel \( S(k^2/m^2) \) in (4.2). In the limit \( m \gg m_0 \) we can approximate the kernel as in (4.5) and we find

\[
m_{s^2} \sim -\frac{g^4 \lambda^2 \lambda^2_0}{(16\pi^2)^3} \frac{M^2 f^2}{m^2 m_0^2} \tag{5.8}
\]

The sfermion masses are negative in this model, therefore it should be considered only as a toy model. Note however that these masses are generated at second order in \( f \). Hence, besides the loop factor, we find a suppression in powers of the SUSY breaking scale for the gaugino mass compared with the sfermion masses, as expected.

### 5.2 Example 2

A less basic model of Yukawa-gauge mediation is characterized by two extra chiral fields \( Y \) and \( Z \) coupled to a spurion \( M_0 = m_0 + \theta^2 f \). Precisely we set \( O = YZ \). The superpotential term in (2.1) is therefore of the form

\[
W_0 = \lambda_0 XYZ, \tag{5.9}
\]

without any dimensionful parameter. In this model the coupling of \( X \) with the hidden chiral operator is indeed a Yukawa interaction.

The Lagrangian for the two extra superfields \( Y \) and \( Z \) is

\[
\mathcal{L} = \int d^4\theta \ (Y^\dagger Y + Z^\dagger Z) + \int d^2\theta \ \frac{1}{2} M_0 (Y^2 + Z^2) + h.c. \tag{5.10}
\]

For simplicity we have introduced only one supersymmetric scale \( m_0 \) and the SUSY breaking scale \( f \), which we take real. Also this model possesses the discrete parity symmetry discussed in the introduction that forbids the radiative generation of a tadpole for the \( F \)-term of \( X \).
We now show by computing the gaugino mass at order $\lambda^2$ that this simple model reproduces the qualitative features we have highlighted with the general analysis. At leading order in $f$, $G_2 - G_4$ is given by

$$G_2(p^2) - G_4(p^2) = 2m_0 \int \frac{dk^4}{(2\pi)^4} \frac{f^3}{[k^2 + m_0^2]^3[(k-p)^2 + m_0^2]^2}$$

$$= \frac{f^3}{16\pi^2 m_0^3} g_{2-4}(p^2/m_0^2)$$

(5.11)

where

$$g_{2-4}(x) = -\frac{x(2 - x)(4 + x) - 8(1 + x)^2 \sqrt{x(4 + x)} \arctanh(\sqrt{x/4+x})}{x^2(4 + x)^3}. \quad (5.12)$$

Plugging this expression into (3.6) with (3.7) one can read off the gaugino mass. In the specific case where the messengers mass is larger than all the other mass scales in the problem, i.e. $m \gg m_0 \gg \sqrt{f}$,

$$m_\lambda \simeq -4 \frac{g^2 \lambda_x \lambda_0^2}{(16\pi^2)^3} \frac{f^3}{8m^2 m_0^3}. \quad (5.13)$$

This contribution is at the leading order in loop factors (three loops), but suppressed in the SUSY breaking scale up to the third order. This is even a larger suppression than we would expect on the basis of the general study in section 3. It would be nice to analyze if this is a generic feature of models with only chiral fields and renormalizable interactions in the hidden sector.

If we introduce also the deformation (3.9), the additional graph involving $G_7$ contributes to the gaugino mass. At the lowest order in $f$

$$G_7(p^2) = 2m_0 \int \frac{dk^4}{(2\pi)^4} \left( \frac{f^2}{[k^2 + m_0^2]^2[(k-p)^2 + m_0^2]^2} + \frac{f^2}{[k^2 + m_0^2]^3[(k-p)^2 + m_0^2]^2} \right)$$

$$= \frac{2f^2}{16\pi^2 m_0^3} g_7(p^2/m_0^2), \quad (5.14)$$

with

$$g_7(x) = \frac{x^2(4 + x) + 4x \sqrt{x(4 + x)} \arctanh(\sqrt{x/4+x})}{x^2(4 + x)^2}. \quad (5.15)$$

We can therefore evaluate the mass of the gaugino in this setup with $m_x \sim m$

$$\Delta m_\lambda \simeq -4 \frac{g^2 \lambda_x \lambda_0^2 \lambda_0^2}{(16\pi^2)^3} \frac{f^2}{mm_x^2}. \quad (5.16)$$

Note that this contribution is at three loops and at second order in the SUSY breaking scale, so it results suppressed one order less with respect to (5.13). As in the previous toy model, this contribution dominates over the latter if $\frac{f}{m_0^2} < \frac{m_0^3}{mm_x}$.
6. Conclusions

In this paper we have studied in a model independent formalism a specific subclass of gauge mediation, which we referred to as Yukawa-gauge mediation. This is characterized by messenger fields which couple to a singlet chiral field $X$ which, in turn, couples only via superpotential interaction to a chiral operator $\mathcal{O}$ that parametrizes the hidden sector. Assuming a parity symmetry that protects $\mathcal{O}$ from taking vevs, the phenomenology is encoded in its two point functions. This class of models is interesting since they can emerge naturally as the low energy effective theory of dynamical supersymmetry breaking models, and their phenomenology can be very different from the minimal gauge mediation case.

In this set up we studied the resulting soft masses of the gaugino and of the scalars of the SSM. We found that generically the gaugino mass is suppressed in powers of the SUSY breaking scale. The scalar masses are instead typically generated at the leading order. We have explicitly illustrated these features in two explicit examples.

It is interesting to make a parallel between our results on gaugino screening and what happens in semi-direct gauge mediation (SDGM) [14]. The two loop diagrammatic cancellation we find for $G_5$ is the same that occurs in SDGM [16]. In SDGM the two loop gaugino mass is vanishing at all orders in the SUSY breaking scale and the next non-vanishing contribution is at higher loop orders. In Yukawa-gauge mediation, instead, there survives a two loop contribution to the gaugino mass (proportional to $G_2 - G_4$), but this is subleading in an expansion in the SUSY breaking scale. The gaugino mass is hence screened in both classes of models, but in a substantially different manner.

In our analysis we also studied the case of a superpotential deformation that can lead to a gaugino mass of the same order in the SUSY breaking scale than the scalar masses. This realizes a scenario with a hierarchy of one loop factor between gaugino and sfermion masses. It would be relevant to investigate if this particular feature, which results very difficult to realize in general, can be obtained in other and more generic setups with weakly coupled chiral fields.

A strategy that has been proposed in order to avoid gaugino screening in SDGM is to consider chiral messengers [22]. It would be interesting to reconsider our analysis under this different assumption. The results of [22] suggest that the gaugino mass would be unscreened also in Yukawa-gauge mediation, but a careful investigation is necessary.
Finally we would like to comment on a different ansatz for Yukawa-gauge mediation. Instead of parametrizing the coupling of $X$ with the hidden sector as in (2.1), one could take a superpotential of the form $W = \lambda_o X X \tilde{O}$. As long as we consider only one point functions of $\tilde{O}$, the general analysis of the previous sections is valid and one can still parametrize the SUSY breaking effects as in (2.5).\footnote{The only caveat is in the definition of the functions $G_i$, which must now be computed according to (2.5).} Going to the next order in $\lambda_o$ the two point functions of $\tilde{O}$ play a role, but we expect that our results similarly extend with minimal adjustments also to this case.

More generally, it would be interesting to study the extended setup where $X$ interacts with hidden sector operators both via linear and quadratic superpotential terms.

Acknowledgments

We are grateful to Riccardo Argurio, Gabriele Ferretti, Dan Thompson and Brian Wecht for useful comments on the draft. F.G. is an Aspirant of FWO-Vlaanderen and A.M. is a Postdoctoral Researcher of FWO-Vlaanderen. This work is supported in part by the FWO-Vlaanderen through the project G.0114.10N, and in part by the Belgian Federal Science Policy Office through the Interuniversity Attraction Pole IAP VI/11.

A. Gaugino mass loop integrals

We have six graphs contributing to the gaugino mass. They are reported in figure 1. Two involve $G_2$. The first of those gives:

$$m_{[a,2]} = -8g^2 \lambda_x^2 \lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2 G_2(k^2)}{k^2 [l^2 + m^2]^3 (l - k)^2 + m^2}.$$  

The second contribution with $G_2$ is:

$$m_{[b,2]} = -4g^2 \lambda_x^2 \lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2 G_2(k^2)}{k^2 [l^2 + m^2]^2 (l - k)^2 + m^2}.$$  

Other two graphs involve $G_4$, they evaluate respectively to

$$m_{[a,4]} = 8g^2 \lambda_x^2 \lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2 G_4(k^2)}{k^2 [l^2 + m^2]^3 (l - k)^2 + m^2}.$$
and
\[ m_{[a,4]} = 4g^2\lambda_x^2\lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2G_4(k^2)}{k^2[l^2 + m^2][l - k]^2 + m^2]}. \]

The different factor of 2 in graphs with the a and b topology is due to the internal loops. In fact, in \( m_{[a,2]} \) there is a fermionic loop involving \( G_2 \), which brings a factor of 2. In \( m_{[a,4]} \) there are two possible equivalent ways to insert the internal loop involving \( G_4 \) in the scalar propagator.

Finally, two graphs involving \( G_5 \) contribute to the gaugino mass as
\[ m_{[a,5]} = 8g^2\lambda_x^2\lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2G_5(k^2)}{(l - k)^2 + m^2}] \]
and
\[ m_{[b,5]} = -4g^2\lambda_x^2\lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{l \cdot (l - k)G_5(k^2)}{(l - k)^2 + m^2]}^2. \]

The mismatching factor of 2 is as for \( G_4 \). The difference in sign for \( m_{[b,5]} \) is because of the fermion propagators, that involve sigma-matrices and not masses, as for \( m_{[b,4]} \).

As explained in the main text, these two contributions cancel each other as in [16], giving no correction to the visible gaugino mass.

There are therefore four graphs that contribute to the gaugino mass. They combine in pairs to give the following integrals in terms of the difference \( G_2 - G_4 \)
\[ m_{[a]} = -8g^2\lambda_x^2\lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2(G_2(k^2) - G_4(k^2))}{k^2[l^2 + m^2][l - k]^2 + m^2]}, \]
\[ m_{[b]} = -4g^2\lambda_x^2\lambda_o^2 \int \frac{d^4k}{(2\pi)^4} \int \frac{d^4l}{(2\pi)^4} \frac{m^2(G_2(k^2) - G_4(k^2))}{k^2[l^2 + m^2][l - k]^2 + m^2]}^2. \]

The integration over the \( l \) momentum can be done analytically and we can write
\[ m_{[a,b]} = -4g^2\lambda_x^2\lambda_o^2 \frac{1}{16\pi^2m^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} L_{[a,b]}(k^2/m^2) (G_2(k^2) - G_4(k^2)), \]
where
\[ L_{[a]}(k^2/m^2) = 2 \int \frac{d^4l}{l^2 + m^2]_3[(l - k)^2 + m^2]}, \]
\[ L_{[b]}(k^2/m^2) = - \int \frac{d^4l}{l^2 + m^2]}_2[(l - k)^2 + m^2]. \]

Using the Feynman parametrization, exchanging the order of integration between the loop integral and the Feynman parameter integral one finally obtains
\[ L_{[a]}(x) = \frac{x(2 + x)(4 + x) - 8\sqrt{x(4 + x)} \text{Arct}h\left(\sqrt{\frac{x}{4 + x}}\right)}{x^2(4 + x)^2}, \]
\[ L_{[b]}(x) = \frac{2x(4 + x) - 4(2 + x)\sqrt{x(4 + x)} \text{Arct}h\left(\sqrt{\frac{x}{4 + x}}\right)}{x^2(4 + x)^2}. \]
The gaugino mass is then given by
\[
m_\lambda = -4g^2\lambda^2\lambda_0^2 \frac{M}{16\pi^2 m^2} \int \frac{d^4k}{(2\pi)^4} \frac{1}{k^2} L(k^2/m^2) (G_2(k^2/M^2) - G_4(k^2/M^2)) ,
\]
with
\[
L(x) = L_{[\alpha]} + L_{[\beta]} = \frac{x(4 + x) - 4\sqrt{x(4 + x)} \text{Arctanh} \left( \frac{\sqrt{x}}{1 + x} \right)}{x^2(4 + x)}.
\]

B. Computation of the kernel \( S(k^2/m^2) \)

We consider the contribution to the \( C_i \) functions, i.e. the propagator of the component of the vector superfields, given by the \( G_1 \) function. The Feynman diagrams are

![Feynman Diagrams](image)

Figure 4: Graphs in which \( G_1 \) is contributing to the visible sfermion masses.

the one in figure 4, and they evaluate to

\[
\begin{align*}
DG^1_a &= 4 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2] [(l - p)^2 + m^2]}, \\
DG^1_b &= -2 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2] [(l - k - p)^2 + m^2] [(l - p)^2 + m^2]}, \\
\lambda G^1_a &= -4 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{(l - p)^\mu \sigma_{\alpha\dot{\alpha}} G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2] [(l - k - p)^2 + m^2]}, \\
AG^1_a &= 4 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{(2l - p)_\mu (2l - p)_\nu G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2] [(l - k - p)^2 + m^2] [(l - p)^2 + m^2]}, \\
AG^1_b &= 2 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{(2l - p)_\mu (2l - 2k - p)_\nu G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2] [(l - k - p)^2 + m^2] [(l - p)^2 + m^2]}, \\
AG^1_d &= -4 \int \frac{dk^4}{(2\pi)^4} \int \frac{dl^4}{(2\pi)^4} \frac{\eta_{\mu\nu} G_1(k^2)}{(l^2 + m^2)^2 [(l - k)^2 + m^2]}. \\
\end{align*}
\]

(B.1)
The functions $C_i^1$ are simple combinations of these expressions, for instance $C_0^1 = \lambda_2^2 \lambda_2^2 (DG_{a_1}^1 + DG_{b_1}^1)$. From this we can obtain the functions $S_i, 1$ and the kernel $S(k^2/m^2)$. The terms in (B.1) are exactly as in [16], but with a sign difference in $DG_{b_1}^1$, and the computation of the kernel can be done along the same lines.

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