Pseudo-Goldstino in Field Theory

Riccardo Argurio
Physique Théorique et Mathématique and International Solvay Institutes, Université Libre de Bruxelles, CP 231, 1050 Bruxelles, Belgium

Zohar Komargodski
School of Natural Sciences, Institute for Advanced Study, Princeton, New Jersey 08540, USA

Alberto Mariotti
Theoretische Natuurkunde and International Solvay Institutes, Vrije Universiteit Brussels, Pleinlaan 2, B-1050 Brussels, Belgium
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We consider two SUSY-breaking hidden sectors which decouple when their respective couplings to the visible particles are switched off. In such a scenario one expects to find two light fermions: the Goldstino and the pseudo-Goldstino. While the former remains massless in the rigid limit, the latter becomes massive due to radiative effects which we analyze from several different points of view. This analysis is greatly facilitated by a version of the Goldberger-Treiman relation, which allows us to write a universal nonperturbative formula for the mass. We carry out the analysis in detail in the context of gauge mediation, where we find that the pseudo-Goldstino mass is at least around the GeV scale and can be easily at the electroweak range, even in low scale models. This leads to interesting and unconventional possibilities in collider physics and it also has potential applications in cosmology.

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Introduction and summary.—In this Letter we consider models with multiple supersymmetry-breaking sectors. We assume these SUSY-breaking sectors communicate only through their respective couplings to the supersymmetric standard model (SSM). In other words, the SUSY-breaking sectors decouple when their respective couplings to the SSM are set to zero. Such models could naturally appear in string theory, where there may be several independent sources of supersymmetry breaking. They may also arise naturally in the study of quiver gauge theories. Our main objective is to study the various field-theoretic effects that are relevant in such a setup.

One may wonder whether having such SUSY-breaking sectors which interact only indirectly through the SSM is natural. Indeed, in field theory this can be perfectly natural since renormalizable contact terms may be forbidden by gauge invariance or global symmetries.

At zeroth order in the interactions with the SSM, there are obviously many massless Goldstino particles. Turning on the small couplings to the SSM, one linear combination, the true Goldstino, remains massless, while the other linear combinations get masses from tree-level and radiative effects. Denoting $f = \sqrt{(f^A)^2 + (f^B)^2}$, then the physical Goldstino and pseudo-Goldstino are given by

$$fG = f^A G^A + f^B G^B, \qquad fG' = -f^B G^A + f^A G^B. \quad (1)$$

First, consider the problem from the point of view of the universal chiral Lagrangian for spontaneously broken supersymmetry. We assume two hidden sectors, represented by the Goldstino superfields X^A and X^B, contribute in some way to the soft gaugino mass

$$\mathcal{L} = \frac{m_\lambda}{2} \int d^2\theta \left(\frac{\alpha^A}{f^A} X^A + \frac{\alpha^B}{f^B} X^B \right) W_\alpha^2. \quad (2)$$

Note that $\alpha^A + \alpha^B = 1$ by definition of m_λ. The chiral Lagrangian approach shows that the contribution from deep low momenta is quadratically sensitive to the cutoff Λ_{UV} of the chiral Lagrangian

$$m_{G'} \sim \frac{1}{16\pi^2} \frac{m_{\rm gaugino}^3}{f^2} \Lambda_{UV}^2. \quad (3)$$

Hence, the contribution is not dominated by parametrically small momenta and one has to invoke the detailed microscopic physics to determine the mass. One can nevertheless show that (3) dominates over tree-level contributions that arise due to electroweak symmetry breaking.

As an example of a microscopically well-defined setup we will analyze in detail two hidden sectors which only communicate with the SSM via gauge interactions. In this case we will find that if the two sectors have a common messenger scale and comparable SUSY-breaking scales, one can roughly estimate the mass of the pseudo-Goldstino $m_{G'}$ as $\sim$1 GeV. On the other hand, we may consider, for instance, different SUSY-breaking scales for the two sectors; then $m_{G'}$ can be easily as high as $\sim$100 GeV.

It follows that our field theory effects surely dominate over gravity as long as $m_{3/2} \sim F/M_{PL}$ is smaller than a GeV or so. This means $\sqrt{f} \lesssim 10^9$, which covers in entirety the parameter space of models based on gauge mediation and variations thereof. On the other hand, since the field-theoretic effects can be easily as large as 100 GeV, it may be important to take them into account even in the regime of gravity mediation.

Having such heavy Goldstino-like particles in controllable low scale models potentially leads to unconventional signatures in collider physics and cosmology. Decays of SSM particles sometimes proceed predominantly into the pseudo-Goldstino and may or may not be accompanied by displaced vertices. In addition, the pseudo-Goldstino has three-body decays with observationally interesting time scales.

A recent inspiring paper [1] (see also the earlier work [2]) considers situations where the gravitational effects are significant. Consistency of supergravity Lagrangians demands the existence of universal nonrenormalizable contact terms mixing the various sectors. Assuming that this is the only source for mixing between the sectors, the authors of [1] computed the supergravity contribution to the mass of the pseudo-Goldstino. They found that the induced mass is $2m_{3/2}$. Possible corrections to this result have been studied in [3] and various interesting applications of this scenario are discussed in [4–6]. In this Letter we consider theories in the rigid limit, where these supergravity corrections are negligible.

Hidden sectors communicating with the SSM by gauge interactions.—The setup we opt to focus on is depicted in Fig. 1. We consider two SUSY-breaking theories, labeled A and B, which communicate with the SSM via gauge interactions. More precisely, when the SSM gauge couplings are set to zero, the sectors A, B decouple from each other. These decoupled theories have some global symmetry groups in which the SSM gauge group can be embedded and weakly gauged.

In essence, this is the setup of general gauge mediation [7], only that the hidden sector is assumed to consist of two decoupled field theories. When the gauge couplings are turned on, the two sectors can communicate by exchanging SSM fields. Obviously, at the zeroth order in the gauge couplings, there are two exactly massless Goldstino fermions. Our goal is to find the leading nonzero contribution in an expansion in the gauge couplings.

The mass matrix for the Goldstino system, defined by $-\frac{1}{2}G^i\mathcal{M}_{ij}G^j$ with a symmetric matrix $\mathcal{M}$, is constrained to have one zero eigenvector corresponding to the true Goldstino (1). Therefore, the matrix has to be of the form

$$\mathcal{M} = \begin{pmatrix} -\frac{f^B}{f^A}\mathcal{M}_{AB} & \mathcal{M}_{AB} \\ \mathcal{M}_{AB} & -\frac{f^A}{f^B}\mathcal{M}_{AB} \end{pmatrix}. \tag{4}$$

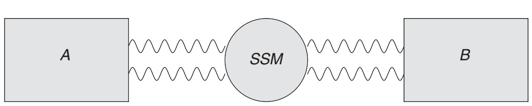

FIG. 1. Two SUSY-breaking theories communicating with the SSM via gauge interactions.

Once we have calculated $\mathcal{M}_{AB}$, the mass of the pseudo-Goldstino is determined via

$$m_{G'} = \left(\frac{f^B}{f^A} + \frac{f^A}{f^B}\right)\mathcal{M}_{AB}. \tag{5}$$

Our goal is therefore to compute the first nontrivial contribution to $\mathcal{M}_{AB}$ in an expansion in the gauge couplings. The processes contributing to $\mathcal{M}_{AB}$ consist of G^A transforming into G^B via some intermediate hidden sector and SSM fields.

It turns out that we must consider processes of order g^4. These allow for two intermediate SSM fields and are thus messenger parity invariant. The intermediate fields must be a gaugino and a gauge field or alternatively a gaugino and a D auxiliary field. This is summarized in Fig. 2.

In the absence of any particular detailed knowledge of the hidden sector we must account for the blobs formally. On the other hand, if the theory is specified and it is weakly coupled, the blobs can be computed in perturbation theory. For instance, in minimal gauge mediation the blobs are, to leading order, triangles with virtual messenger fields. Therefore, the pseudo-Goldstino obtains a mass due to three-loop corrections.

In the processes of Fig. 2 the external Goldstinos are at zero momentum. One can interpret the blobs as three-point functions of the supercurrent and two insertions of operators of the linear current multiplet. In other words, the pertinent correlation functions are of the form $\langle S_{\nu\alpha}(x)j_\mu(y)\bar{j}_{\dot\alpha}(z)\rangle$, $\langle S_{\nu\alpha}(x)J(y)\bar{j}_{\dot\alpha}(z)\rangle$, $\langle S_{\nu\alpha}(x)j_\mu(y)j_\beta(z)\rangle$, $\langle S_{\nu\alpha}(x)J(y)j_\beta(z)\rangle$. For our purposes we need the external state to be a zero-momentum Goldstino; hence, the correlation functions above should be studied only in the limit of large x (much larger than any other scale in the problem).

In this large x limit the three-point functions above simplify dramatically. The reason is that at very low energies the supercurrent flows to the Goldstino particle $S^{A,B}_{\mu\alpha} \sim f^{A,B}\sigma_{\mu\alpha\dot\alpha}\bar{G}^{A,B\dot\alpha}$ and therefore the large x limit corresponds to inserting a zero-momentum Goldstino in the correlation function. This is the same as acting with the supercharge on the vacuum and thus these three-point functions are related to two-point functions of the form $\langle[\bar{Q}_{\dot\gamma}, j_\mu(y)\bar{j}_{\dot\alpha}(z)]\rangle$, $\langle[\bar{Q}_{\dot\gamma}, J(y)\bar{j}_{\dot\alpha}(z)]\rangle$, $\langle[\bar{Q}_{\dot\gamma}, j_\mu(y)j_\beta(z)]\rangle$, $\langle[\bar{Q}_{\dot\gamma}, J(y)j_\beta(z)]\rangle$. These two-point functions, in turn, appear in the calculations of soft masses in gauge mediation. We adopt notation similar to that in general gauge mediation [7]:

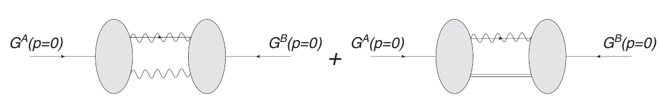

FIG. 2. At order g^4 the sectors A and B can communicate via two intermediate SSM fields. We must also add the diagrams with the gauginos flowing in the opposite direction.

$$
\begin{aligned}
\langle J^{A,B}(p)J(-p)^{A,B}\rangle &= C_0^{A,B}(p^2),\\
\langle j_\alpha^{A,B}(p)\bar{j}_{\dot\alpha}^{A,B}(-p)\rangle &= -\sigma^\mu_{\alpha\dot\alpha}p_\mu C_{1/2}^{A,B}(p^2),\\
\langle j_\mu^{A,B}(p)j_\nu^{A,B}(-p)\rangle &= -(p^2\eta_{\mu\nu}-p_\mu p_\nu)C_1^{A,B}(p^2),\\
\langle j_\alpha^{A,B}(p)j_\beta^{A,B}(-p)\rangle &= \epsilon_{\alpha\beta}B_{1/2}^{A,B}(p^2).
\end{aligned}
\tag{6}
$$

The discussion above shows that the leading order contribution to the pseudo-Goldstino mass should be captured by the functions in (6). A quick way to derive the precise relations between these two-point functions and three-point functions is to start by recalling the effective quadratic action for the vector multiplet:

$$
\begin{aligned}
\frac{1}{g^2}\mathcal{L}_{\rm eff} &= \frac{1}{2}C_0^A D^2 - iC_{1/2}^A\lambda\sigma^\mu\partial_\mu\bar\lambda - \frac{1}{4}C_1^A F_{\mu\nu}F^{\mu\nu}\\
&\quad - \frac{1}{2}(B_{1/2}^A\lambda\lambda + \text{c.c.}) + A\leftrightarrow B.
\end{aligned}
\tag{7}
$$

This breaks supersymmetry if $B_{1/2}\neq 0$ and if the C's are not all equal. However, it can be supersymmetrized by adding terms linear in the Goldstino as follows:

$$
\begin{aligned}
\frac{1}{g^2}\mathcal{L}_{\rm eff}^{\text{one } G} &= \frac{1}{\sqrt{2}f^A}(C_0^A - C_{1/2}^A)G^A\sigma^\mu\partial_\mu\bar\lambda D\\
&\quad + \frac{i}{\sqrt{2}f^A}(C_1^A - C_{1/2}^A)G^A\sigma_\nu\partial_\mu\bar\lambda F^{\mu\nu}\\
&\quad + \frac{iB_{1/2}^A}{\sqrt{2}f^A}\left(G^A\lambda D - \frac{i}{2}\lambda\sigma^\mu\bar\sigma^\nu G^A F_{\mu\nu}\right)\\
&\quad + A\leftrightarrow B.
\end{aligned}
\tag{8}
$$

To make the theory fully supersymmetric, in addition to (8), we need to add terms bilinear in the Goldstinos, and terms with derivatives acting on the Goldstinos. In order to compute $\mathcal{M}_{AB}$, (8) suffices. The procedure we have invoked here is a supersymmetric reincarnation of the Goldberger-Treiman relation.

From here to derive the mass of the pseudo-Goldstino we only need to carry out the contractions using the vertices in (8). After the dust settles, we find that the leading order contribution to the mass of the pseudo-Goldstino is

$$
\begin{aligned}
m_{G'} &= \frac{g^4}{2}\left(\frac{1}{(f^A)^2}+\frac{1}{(f^B)^2}\right)\int\frac{d^4p}{(2\pi)^4}B_{1/2}^A(C_0^B - 4C_{1/2}^B\\
&\quad + 3C_1^B) + A\leftrightarrow B.
\end{aligned}
\tag{9}
$$

Note that the combination of the C functions in the integrand is precisely the one appearing in the formula for the soft scalar mass in gauge mediation. The discussion in [8] shows that $C_0 - 4C_{1/2} + 3C_1$ behaves at most like $1/p^4$ at large momentum and it is also possible to prove that $B_{1/2}$ scales at most like $1/p$ at large momentum. Consequently, the integral is UV convergent.

The computation above has been greatly simplified by the structure of the matrix (4), which allowed us to compute $m_{G'}$ only in terms of $\mathcal{M}_{AB}$. As a consistency check, we can also compute the diagonal elements of the mass matrix $\mathcal{M}_{AA}$ and $\mathcal{M}_{BB}$. In order to do this one must take into account also the corrections to (8) quadratic in each of the Goldstinos.

We can now estimate (9) crudely. Assume both hidden sectors have some typical supersymmetric scale M and the SUSY-breaking scales are f^A, f^B. To leading order in the SUSY-breaking scales we would get

$$
m_{G'} \sim \frac{g^4}{(16\pi^2)^3}\left(\frac{f^A}{f^B}+\frac{f^B}{f^A}\right)\left(\frac{f^A}{M}+\frac{f^B}{M}\right).
\tag{10}
$$

If the two SUSY-breaking scales are comparable this leads to the estimate

$$
m_{G'} \sim \frac{g^4}{(16\pi^2)^3}\frac{f}{M} \sim \frac{g^2}{(16\pi^2)^2}m_{\rm soft} \sim 1\ \text{GeV}.
\tag{11}
$$

We have included a factor of $\mathcal{O}(10)$ due to the sum over the gauge sector of the SSM.

However, we can also entertain other possibilities. For instance, consider a situation where the fundamental supersymmetric scales in the two sectors are comparable but the SUSY-breaking scales are different. To be concrete we assume that $f^A \gg f^B$ (the soft parameters thus mostly originate in sector A). In this case, the formula (10) predicts an enhancement of $m_{G'}$ by f^A/f^B. This ratio, however, cannot be arbitrarily large because at some point the backreaction of the SSM on the hidden sector B becomes too large and our formalism breaks down. By computing the sGoldstino vacuum expectation value in sector B, we can estimate that the backreaction is surely tame for $f^A/f^B \ll 10^3$. (For this estimate we have assumed the mass of the sGoldstino is around $f^{A,B}/M$.) Thus, we can easily imagine the pseudo-Goldstino picking a mass at the electroweak range. Note that such a (perhaps surprisingly) large mass for the pseudo-Goldstino is achieved effortlessly and ubiquitously in low scale models, where corrections from supergravity are completely negligible.

One can also evaluate (9) explicitly in a variety of simple realizations of gauge mediation, for instance, if the two hidden sectors are copies of minimal gauge mediation.

Phenomenology of Goldstinos.—In the scenario presented in this note, the pseudo-Goldstino is generically the next-to-lightest supersymmetric particle, with the lightest supersymmetric particle being of course the very light gravitino. The pseudo-Goldstino is not stable and its decay can be analyzed via the chiral Lagrangian. For instance, the terms responsible for the gaugino mass (2) give rise to vertices of the form $\sim G\sigma_\mu\bar\sigma_\nu\lambda F^{\mu\nu}$ which induce three-body decays of the pseudo-Goldstino into two photons and the true Goldstino. There are also some very important vertices with two Goldstinos. In fact, the naive estimate based on dimensional analysis fails due to an exact cancellation between the different vertices. An analogous story takes place in the couplings to the SM fermions.

One is left with the following estimate of the decay width into two standard model fermions and the true Goldstino [5]:

$$\Gamma_{G'\to Gf\bar{f}} \sim \frac{m_{G'}^9}{10^5 f_{\rm eff}^4}\left(\frac{(m_{\tilde{f}}^A)^2 \tan\theta - (m_{\tilde{f}}^B)^2 \cot\theta}{m_{\tilde{f}}^2}\right)^2. \quad (12)$$

We denote $\tan\theta = f^B/f^A$ and $(m_{\tilde{f}}^{A,B})^2$ are the contributions to the mass of the slepton from the two hidden sectors, such that $(m_{\tilde{f}}^A)^2 + (m_{\tilde{f}}^B)^2 = m_{\tilde{f}}^2$. There is a similar width to decay into two photons and the true Goldstino.

Consider theories with two general SUSY-breaking scales $f^A \geq f^B$. Assuming again, for simplicity, that the messenger scales in the two sectors are comparable and taking $m_{\tilde{f}}^{A,B} \propto f^{A,B}/M$, we find

$$\tau \sim 10^{21} \ {\rm sec}\left(\frac{f_{\rm eff}}{10^{10}\ {\rm GeV}^2}\right)^4\left(\frac{f^B}{f^A}\right)^7. \quad (13)$$

To derive the estimate above we have taken the mass of the pseudo-Goldstino to be $m_{G'} \sim f^A/f^B$ GeV. This gives rise to many different possibilities. For instance, when the pseudo-Goldstino is around the weak or TeV scale (i.e., $f^A/f^B \sim 10^{2-3}$), models of low scale mediation $\sqrt{f} \sim 10^{4-5}$ GeV give a lifetime of the order of a few seconds. Still keeping the pseudo-Goldstino at the weak-TeV scale, we can also choose $\sqrt{f_{\rm eff}} \sim 10^8$ GeV which leads to lifetimes of the order of 10^{23-24} secs. Both of these time scales have potentially interesting observable consequences [9]. One can of course imagine many other scenarios stemming from (13), including scenarios with lighter pseudo-Goldstino.

One can also easily imagine many unconventional collider manifestations of the setup here. One obvious consequence of having two different hidden sectors is that the relation between the decay time of the lightest observable-sector supersymmetric particle (LOSP) and the scale of SUSY-breaking is no longer universally determined. This can have several different consequences.

For instance, consider two hidden sectors with comparable messenger scales but with a possible hierarchy in the SUSY-breaking scales. From the couplings (2) we see that the gaugino is equally likely to decay to either of the Goldstinos (since the dependence on f cancels and only the supersymmetric scale remains). Therefore, if the LOSP is bino- or winolike, and it is heavier than the pseudo-Goldstino, many of the processes of the SSM will terminate in a heavy, long-lived, pseudo-Goldstino (the decay can be prompt or there can be displaced vertices). This also comes accompanied by an isolated photon from the last step of the decay. Having such an invisible heavy particle as missing energy is clearly different from conventional scenarios of gauge mediation where the missing energy is carried away by practically massless objects. It is also distinguishable from gravity mediation, where the LOSP is stable on collider time scales and therefore, if it is a gaugino, no isolated photons are expected.

The very brief remarks above are just to demonstrate that unusual collider and cosmological signatures are definitely possible. Clearly, it will be interesting to investigate the various possibilities further. It is also important to study more general hidden sector paradigms, beyond gauge mediation.

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