Nonlinear Supersymmetric Effective Lagrangian and Goldstino Interactions at High Energies

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Abstract

We show that the nonlinear supersymmetric effective lagrangian can be used for model-independent parameterization of the light gravitino scattering amplitude at energies up to and above the soft supersymmetry-breaking masses. This provides the most convenient framework for systematic studies of goldstino phenomenology both at low energies and in high energy colliders.

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1 Introduction

The existence of a very light gravitino is one of the characteristic features of low energy supersymmetry (SUSY) breaking [1]. If SUSY breaking occurs in the order of the weak scale, the gravitino becomes superlight with its mass several orders of magnitude smaller than an electron Volt. A superlight gravitino could have important consequences in many areas such as cosmology, astrophysics, and collider physics, and its phenomenology has been studied by many authors [2, 3] since the early works of Fayet [4]. One of the recent considerations on superlight gravitino that is of particular interest to us is gravitino production in high energy collisions. As the typical energy of the process is well above the gravitino mass, one can replace the dominant, longitudinal component of the gravitino with the goldstino [3].

In globally supersymmetric theories, two approaches exist for obtaining the goldstino coupling to the fields of the minimal supersymmetric standard model (MSSM). The model-specific approach is to integrate out the high energy modes in a given SUSY breaking model down to a certain scale above the soft SUSY breaking masses in MSSM. In general, the interactions of the goldstino with the MSSM fields are given in nonderivative form, and are model-dependent, though the vertices with a single goldstino are fixed by the Goldstino Goldberg-Treiman (GT) relation and are model-independent. As shown in Ref.[6], the correct form of goldstino-matter interaction is obtained only after careful treatment of the Feynman diagrams that involve the nondecoupling of heavy particles.

The model-independent approach of the nonlinear effective lagrangian[7, 8, 9, 10] is based on a nonlinearly realized supersymmetry in the SUSY breaking sector, and provides the most general interactions between the goldstino and a given set of fields consistent with spontaneous SUSY breaking. Here the goldstino couples derivatively, and the model-dependence on the underlying SUSY breaking mechanism appears as undetermined coefficients of the derivative expansion of goldstino fields. A clear advantage of this approach is the decoupling of heavy particles from the goldstino emission processes. The goldstino low energy theorem [11] follows automatically in the nonlinear SUSY realization.
In our previous work [10], we provided the rules for constructing nonlinear SUSY invariant operators describing goldstino couplings to matter and gauge fields in MSSM. The applications given there focus on processes involving the goldstino and the standard model (SM) particles only. Such a nonlinear goldstino lagrangian for the SM can also be obtained by explicitly integrating out the heavy SM superpartners in a given model, as has recently been done by the authors of Ref.[12] in their study of superlight gravitino production at $e^+e^-$ and hadron colliders when the SUSY particles are too heavy to be produced. It will be interesting and important to study also goldstino production at energies near or above the sparticle masses. In the nonlinear SUSY lagrangian framework, however, such a study has not been done. It is our purpose here to extend our previous analysis to the high energy regime.

After a brief review of the nonlinear SUSY realization on the MSSM fields, we study two types of scattering processes. As will be shown, nonlinear SUSY invariance plays an important role in constraining the structure of goldstino interactions.

2 The nonlinear effective lagrangian

There exist in the literature several different approaches in obtaining the nonlinear SUSY effective lagrangian [7, 9, 10]. Here we take the approach of Refs.[9, 10] which has an advantage over others in that the effective lagrangian can be written directly in terms of component fields rather than superfields as in other approaches. Because of this, it is straightforward to use this method to find the most general SUSY invariant local operators for a given set of fields. One may refer to Ref.[9, 10] for details.

The nonlinear effective lagrangian is given in the form:

$$I_{\text{eff}} = \int d^4x \, \det A \, L_{\text{eff}}(\mathcal{D}_\mu \chi, \mathcal{D}_\mu \bar{\chi}, \phi^i, \mathcal{D}_\mu \phi^i, \mathcal{F}_{\mu
u})$$

where $L_{\text{eff}}$ is a gauge invariant function of the standard realization basic building blocks.
defined by

\[ D_\mu \chi = (A^{-1})_\mu ^\nu \partial_\nu \chi, \]
\[ D_\mu \phi^i = (A^{-1})_\mu ^\nu (\partial_\nu \delta_{ij} + T^a_{ij} A_a^\nu) \phi^j, \]
\[ F^a_{\mu\nu} = (A^{-1})_\mu ^\alpha (A^{-1})_\nu ^\beta F^a_{\alpha\beta}, \quad (2) \]

where \( \chi \) is the goldstino, \( \phi^i \) denotes generic scalar or fermion fields, \( A^a_\nu \) and \( F^a_{\alpha\beta} \) are respectively the gauge field and field strength tensor, and \( T^a \) the gauge group generators. The goldstino self-interaction is simply given by the Volkov-Akulov action \( \mathcal{L}_{AV} \) [3],

\[ \mathcal{L}_{AV} = -F^2 \det A \quad (3) \]

where \( F \) is the goldstino decay constant and \( A \) is the Volkov-Akulov vierbein defined by

\[ A^\nu_\mu = \delta^\nu_\mu + \frac{i}{2F^2 \chi} \partial_\mu \sigma^\nu \bar{\chi}. \quad (4) \]

It is convenient to catalog the terms in the effective Lagrangian, \( \mathcal{L}_{\text{eff}} \), by an expansion in the number of goldstino fields which appear when the Volkov-Akulov vierbein is set to unity. Then we have

\[ \mathcal{L}_{\text{eff}} = \left[ \mathcal{L}_\text{(0)} + \mathcal{L}_\text{(1)} + \mathcal{L}_\text{(2)} + \cdots \right], \quad (5) \]

where the subscript \( n \) on \( \mathcal{L}_\text{(n)} \) denotes that each independent SUSY invariant operator in that set begins with \( n \) Goldstino fields.

For the MSSM fields, \( \mathcal{L}_\text{(0)} \) is obtained from the MSSM lagrangian by replacing the ordinary gauge covariant derivatives with the SUSY-gauge covariant derivatives and the ordinary field strength tensor with the SUSY covariant field strength tensor:

\[ \mathcal{L}_\text{(0)} = \mathcal{L}_{\text{MSSM}}(\phi, D_\mu \phi, F_{\mu\nu}). \quad (6) \]

Note that \( \mathcal{L}_\text{(0)} \) is independent of the underlying SUSY breaking dynamics, and the goldstino dependence arises only from higher dimension terms in the matter covariant derivatives and the SUSY covariant field strength tensor. In particular, the vertices having two on-shell
goldstinos arise from the SUSY and gauge invariant kinetic terms for the matter and gauge fields,

\[ \mathcal{L}_{(0)}^{\chi} = -i \bar{\psi}^i \sigma^\mu D_\mu \psi^i - (\mathcal{D}^\mu \phi^i)^\dagger (\mathcal{D}_\mu \phi^i) - \frac{1}{4} F_{\mu \nu}^a F^{a \mu \nu} \]  \hspace{1cm} (7)

The terms in \( \mathcal{L}_{(1)} \) describe, at the one goldstino level, the standard single goldstino derivative coupling to the supercurrent,

\[ \mathcal{L}_{(1)} = \frac{1}{F} \mathcal{D}_\mu \chi^\mu J_\alpha^\mu + \text{h.c.} \]  \hspace{1cm} (8)

where \( J_\alpha^\mu \) is the supercurrent. Like \( \mathcal{L}_{(0)} \), \( \mathcal{L}_{(1)} \) is model-independent. Explicitly, the goldstino coupling to the fermion-scalar pair \((\psi, \phi)\) and the gauge-gaugino pair \((A_\mu, \lambda)\) are given by

\[ \mathcal{L}_{(1)} = \frac{1}{F} \partial_\mu \chi \sigma^\nu \sigma^\rho \phi \sigma^\mu \phi D_\nu \psi^\dagger M_{\mu \nu} \]  \hspace{1cm} (9)

where \( D_\mu \phi \) is the gauge covariant derivative. Throughout this paper, the goldstinos are taken to be on-shell external particles.

Equivalently, the single Goldstino interactions can be written in nonderivative form. For SUSY QED with a massless fermion carrying unit charge, this is simply given by

\[ \mathcal{L}_{(1)\text{ND}} = \frac{m_\phi^2}{F} \chi \psi \phi^\dagger \frac{im_\chi}{\sqrt{2} F} \chi \sigma^{\mu \nu} F_{\mu \nu} - \frac{em_\lambda}{\sqrt{2} F} \phi \chi \lambda \]  \hspace{1cm} (10)

where the first two are the standard GT trilinear couplings, whereas the presence of the quartic term is required for the two forms of the single Goldstino interactions to be equivalent [14].

Unlike \( \mathcal{L}_{(0)} \) and \( \mathcal{L}_{(1)} \), the rest of the terms in the effective lagrangian are model-dependent. The operators in \( \mathcal{L}_{(2)} \), for example, can be written in the form

\[ \mathcal{L}_{(2)} = \frac{1}{F^2} \mathcal{D}_\mu \chi^\alpha \mathcal{D}_\nu \chi^\dagger M_{\alpha \beta}^{\mu \nu} + \cdots \]  \hspace{1cm} (11)

where \( M_{\alpha \beta}^{\mu \nu} \) denotes operators composed of the MSSM fields and containing arbitrary coefficients.
3 Fermion-antifermion annihilation

In this section, we examine the validity of the nonlinear effective lagrangian description of the fermion antifermion annihilation process $\psi \bar{\psi} \rightarrow \chi \bar{\chi}$, both in the zero energy limit and at energies above the sfermion mass. For simplicity, the matter wyle fermion is assigned a conserved global $U(1)$ charge and is massless. Nonlinear SUSY invariance will be seen to place very useful constraints on the low energy operators.

We start with the model-independent sfermion exchange contribution shown in Fig. 1, using the nonderivative GT coupling of Eq.(10). In the low energy limit, an expansion in $1/m^2_\phi$ gives rise to the effective local operators,

$$m_\phi^2 \frac{F^2}{\phi^2} (\chi \psi) (\bar{\chi} \bar{\psi}) + \frac{1}{F^2} (\bar{\chi} \bar{\psi}) \partial^2 (\chi \psi) + \cdots,$$

(12)

where terms suppressed by powers of $\partial^2/m^2_\phi$ have been dropped. Note that nonlinear SUSY invariance forbids the presence of the first term which leads to incorrect energy dependence of the amplitude in the zero momentum limit of the goldstino. The second term, on the other hand, is allowed by the goldstino low energy theorem [11] and consistent with nonlinear SUSY invariance as will be shown below. The low energy goldstino decoupling as dictated by nonlinear SUSY realization thus requires the presence of new interactions to cancel out the first term above.

As an example, we consider the toy model of Ref.[15]. The interaction terms of the model relevant for our process are given by

$$L_{int} = m_\phi^2 \frac{F^2}{\phi^2} (\chi \psi \phi^* + \bar{\chi} \bar{\psi} \phi) - m_\phi^2 \frac{F^2}{\phi^2} (\chi \psi)(\bar{\chi} \bar{\psi}).$$

(13)

Note that the first two terms are the standard GT trilinear interactions which, after integrating out the sfermion, give rise to the effective local operators of Eq.(12). The quartic operator in the model cancels the first term in Eq.(12), leading to the effective low energy interaction well below the sfermion mass [15],

$$O_{eff} = \frac{1}{F^2} (\bar{\chi} \bar{\psi}) \partial^2 (\chi \psi).$$

(14)
This operator is consistent with nonlinear SUSY realization as we now show.

In the effective nonlinear goldstino lagrangian, two dimension-eight operators can be written down to describe this low energy scattering. They are given by [10, 15]:

\[ L_{\psi\bar{\psi}\chi\bar{\chi}} = -\frac{1}{4F^2} \left( \chi \overset{\leftrightarrow}{\partial^2} \sigma^\mu \bar{\sigma}^\nu \bar{\chi} \right) \left( \psi \overset{\leftrightarrow}{\partial^2} \sigma^\nu \sigma^\mu \bar{\psi} \right) + \frac{C_{ff}}{F^2} \left( \psi \partial^\mu \chi \right) \left( \bar{\psi} \partial^\mu \bar{\chi} \right), \tag{15} \]

where the first term comes from the SUSY invariant fermion kinetic term in \( L_{(0)} \) and the second is a model-dependent operator in \( L_{(2)} \) with an arbitrary coefficient. The effective interaction of Eq.(14) is recovered from Eq.(15) when \( C_{ff} = -2 \) and is thus SUSY invariant.

Now at energies of the order of the sfermion mass, the scattering amplitude develops a dependence on the mass of the superpartner. This is seen in the model of Ref.[15] as arising from the same two diagrams due to sfermion exchange and the contact term (Fig. 2). Explicitly, the effective interaction is now given by,

\[ \mathcal{M}_{ff} = \frac{2}{F^2} (\partial_\mu \chi \partial^\mu \psi) (\bar{\chi} \bar{\psi}) - \frac{2}{F^2} (\bar{\chi} \bar{\psi}) \frac{\partial^2}{\partial^2 - m_\phi^2} (\partial_\mu \chi \partial^\mu \psi), \tag{16} \]

which reduces to Eq.(14) in the zero energy limit.

It is clear that the validity of the nonlinear effective lagrangian description of Eq.(15) breaks down at the sfermion mass scale, where the amplitude is expected to develop a dependence on the sfermion mass. This seems to suggest that the nonlinear effective lagrangian is a double expansion in \( 1/m_\phi \) and \( 1/F \), and that its usefulness is limited below the relevant sparticle mass.
Interestingly, however, this seemingly disappointing feature of the nonlinear effective lagrangian can be remedied by including the induced contributions from $\mathcal{L}_{(1)}$. The contact term of Eq. (13) plus the sfermion exchange diagram arising from $\mathcal{L}_{(1)}$ can be shown to give the same amplitude as in Eq. (16) (see Fig. 2). Note that in the zero energy limit, the sfermion decouples from the nonlinear lagrangian and one is left with the contact interaction of the first diagram on the R.H.S. of Fig. 2. The complete nonlinear effective lagrangian description of $\mathcal{L}_{(0)} + \mathcal{L}_{(1)} + \mathcal{L}_{(2)} + \cdots$ for goldstino emission processes remains valid well above the soft SUSY breaking masses, and it is an expansion only in $1/F$.

4 Photon-photon annihilation

As another example, we consider photon photon annihilation into a pair of goldstinos. Within the MSSM, this process receives contribution only from photino exchange (Fig. 3). With the GT coupling of Eq. (10), one can easily find the induced effective interaction at energies well below the photino mass $m_\lambda$,

$$-\frac{i}{2F^2}F^{\mu\nu}F_{a\mu}\chi^\alpha\sigma^\beta\partial_\nu\chi + \frac{m_\lambda}{8F^2}\chi(F_{\mu\nu}F^{\mu\nu} + \frac{i}{2}\epsilon^{\mu\nu\alpha\beta}F_{\mu\nu}F_{\alpha\beta}) + h.c. \quad (17)$$

where terms suppressed by powers of $E/m_\lambda$ are dropped. The first term above corresponds to the dimension-eight operator from the expansion of the photon kinetic term in the nonlinear
effective lagrangian. The second term can not be made invariant under the nonlinear SUSY transformation, and has to be canceled by other contributions from outside of the MSSM.

The low energy limit of this process has recently been studied in a toy model concerning the energy dependence of the cross section \[ \text{[6]} \]. In this model, beside the standard GT coupling of Eq.\( (10) \), there also exist a scalar \( S \) and a pseudoscalar \( P \) coupled to two goldstinos. The \( S \) and \( P \) are the scalar partners of the goldstino. The relevant terms of the lagrangian in the model are given by

\[
\mathcal{L}_{\text{int}} = \frac{i}{\sqrt{2}} \frac{m_\lambda}{F} \chi \sigma^{\mu \nu} \lambda F_{\mu \nu} - \frac{1}{2\sqrt{2}} \frac{m_S^2}{F} S \chi \chi + \frac{i}{2\sqrt{2}} \frac{m_P^2}{F} P \chi \chi

- \frac{1}{2\sqrt{2}} \frac{m_\lambda}{F} [SF_{\mu \nu}F^{\mu \nu} - \frac{1}{2} \epsilon^{\mu \nu \alpha \beta} PF_{\mu \nu}F_{\alpha \beta}] + \text{h.c.} \quad (18)
\]

where \( m_S \) and \( m_P \) are the masses of the \( S \) and \( P \). In the zero energy limit, the contribution due to \( S \) and \( P \) exchange cancels the SUSY noninvariant term from the photino exchange in Eq. \( (17) \). The total low energy interaction is now given by,

\[
\mathcal{O}_{\gamma \gamma} = -\frac{i}{2F^2} F^{\mu \nu} F_{\alpha \mu} \chi \sigma^\alpha \partial_\nu \chi, \quad (19)
\]

which is independent of the masses of the photino and the \( S, P \).

From the nonlinear goldstino lagrangian standpoint, as a consequence of the derivative coupling nature of the goldstino, the heavy sparticles always decouple from the low energy processes and one only needs to consider the contributions of the contact interactions. At the dimension-eight level, there exists only one such operator responsible for the low energy process \( \gamma \gamma \to \chi \overline{\chi} \). This operator arises from the photon kinetic term in \( \mathcal{L}_{(0)} \) and is exactly given by Eq. \( (19) \), in accord with the explicit model calculation.
We now turn to the high energy limit of the process and study the validity of the nonlinear lagrangian. In the model of Ref. [6], the contributing diagrams are due to the exchange of photino, $S$, and $P$ respectively (see Fig. 4). Using the interactions of Eq. (18), the high energy amplitude is found to be

$$M^{ND}_{\gamma\gamma} = -\frac{i}{2F^2} (F^{\mu\nu}\chi) \sigma^\alpha \frac{m^2_{\lambda}}{m^2_{\lambda} - \partial^2} \left[(\partial_\nu \chi) F_{\alpha\mu}\right]$$

$$- \frac{m_{\lambda}}{8F^2} (F_{\mu\nu}\chi) \frac{m^2_{\lambda}}{m^2_{\lambda} - \partial^2} \left[\chi \left(F^{\mu\nu} + \frac{i}{2} \epsilon^{\mu\nu\alpha\beta} F_{\alpha\beta}\right)\right] + \text{h.c.}$$

$$+ \frac{m_{\lambda}}{8F^2} \left[F^{\mu\nu} F_{\mu\nu} \frac{m^2_{S}}{m^2_{S} - \partial^2} \left(\chi\chi\right) + \frac{i}{2} \epsilon^{\mu\nu\alpha\beta} F_{\mu\nu} F_{\alpha\beta} \frac{m^2_{P}}{m^2_{P} - \partial^2} \left(\chi\chi\right)\right] + \text{h.c.}, \quad (20)$$

which reduces to Eq. (19) in the zero momentum limit.

In the framework of the nonlinear goldstino lagrangian, the contact contribution from the model-independent interaction in $\mathcal{L}_{(0)}$ (see Eq. (19)) does not give the complete scattering amplitude at high energies. Including the photino exchange contribution arising from $\mathcal{L}_{(1)}$ now becomes necessary but is not sufficient by itself. The toy model lagrangian of Eq. (18) suggests that the goldstino couplings to its scalar partners be included in the model-dependent piece of the nonlinear lagrangian $\mathcal{L}_{(2)}$, which will give new contributions to the high energy photon-photon annihilation process. The coefficients of the derivative couplings of the goldstino to $S$ and $P$ can be determined in the model from the decay amplitudes of $S \rightarrow \chi\chi$ and $P \rightarrow \chi\chi$. They are given by

$$\mathcal{L}^{S,P}_{(2)} = -\frac{1}{\sqrt{2}F} (S - iP) \partial_\mu \chi \partial^\mu \chi + \text{h.c.} \quad (21)$$

The high energy scattering amplitude can now be calculated in the nonlinear lagrangian by including the contact contribution (Eq. (19)), the photino exchange, and the $S$ and $P$ exchanges. It is given by

$$M^{D}_{\gamma\gamma} = -\frac{i}{2F^2} F^{\mu\nu} F_{\alpha\mu} \chi \sigma^\alpha \partial_\nu \chi$$

$$- \frac{i}{2F^2} (F^{\mu\nu}\chi) \sigma^\alpha \frac{\partial^2}{m^2_{\lambda} - \partial^2} \left[(\partial_\nu \chi) F_{\alpha\mu}\right]$$
Figure 4: Equivalence of amplitudes in photon-photon annihilation. The dotted vertices are from the nonderivative lagrangian, while those on the R.H.S. denoted by circled cross are from the nonlinear effective lagrangian.

\[
\frac{m_\lambda}{8F^2} (F_{\mu\nu}\chi) \frac{\partial^2}{m_\lambda^2 - \partial^2} \left[ \chi(F^{\mu\nu} + \frac{i}{2} \epsilon^{\mu\nu\alpha\beta} F_{\alpha\beta}) \right] + h.c.
\]

\[
+ \frac{m_\lambda}{8F^2} \left[ F^{\mu\nu} F_{\mu\nu} \frac{\partial^2}{m_\lambda^2 - \partial^2} (\chi\chi) + \frac{i}{2} \epsilon^{\mu\nu\alpha\beta} F_{\mu\nu} F_{\alpha\beta} \frac{\partial^2}{m_\lambda^2 - \partial^2} (\chi\chi) \right] + h.c. . \tag{22}
\]

Though the two amplitudes in Eq. (20) and Eq. (22) come from completely different schemes, it is straightforward to see that they are indeed identical (see Fig. 4). Note that when some of the superpartners (say $S$ and $P$) become very heavy compared to the typical energy transfer in the process, these particles automatically decouple in the nonlinear lagrangian and one only needs to consider the contact interaction and the diagram involving the light superpartner (say $\lambda$) exchange. For the nonderivative goldstino coupling in the model of Ref.[6], the heavy particles do not decouple and need to be included in the diagrams. We have thus explicitly shown that the nonlinear goldstino lagrangian provides a valid description of the annihilation process $\gamma\gamma \rightarrow \chi\chi$ at energies up to and above the soft breaking sparticle masses. Contributions from $\mathcal{L}_{(0)}$, $\mathcal{L}_{(1)}$, and $\mathcal{L}_{(2)}$ all need to be taken into account at high energies.

5 Conclusion

If SUSY is broken around the weak scale, the accompanying gravitino will be superlight, and direct gravitino production in high energy collisions becomes feasible. We have demonstrated
that the nonlinear goldstino lagrangian provides a valid description of the goldstino interactions both below and above the soft SUSY breaking masses, though the latter requires the superpartners to be properly taken into account. At energies well below the soft masses, the superpartners decouple from the low energy theory and one recovers the effective nonlinear goldstino lagrangian for the SM. Generally speaking, particles much heavier than the typical energy of the goldstino always decouple in the nonlinear lagrangian approach, and one only needs to consider light particle effects. In contrast, with nonderivative goldstino coupling as is often the case in explicit models of SUSY breaking, heavy particles do not decouple and their exchange effects need to be explicitly included in the diagrams. The nonlinear goldstino lagrangian thus provides the most economical description of the supersymmetry breaking sector.

Before we understand the origin of supersymmetry breaking, the effective nonlinear goldstino lagrangian provides the most convenient framework for model-independent and systematic studies of goldstino phenomenology at both low energies and in high energy collisions. As all the operators in the nonlinear lagrangian are organized in expansions of a single parameter, the inverse of the goldstino decay constant $1/F$, phenomenological studies of the goldstino can yield valuable information on the supersymmetry breaking scale, and may also provide helpful guidelines in seeking the realistic model of supersymmetry breaking.

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