Non-linear supersymmetry and goldstino couplings to MSSM

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Abstract

We review briefly the non-linear supersymmetry formalisms in the standard realization and superfield methods. We then evaluate the goldstino couplings to the MSSM superfields and discuss their phenomenological consequences. These refer to the tree level Higgs mass and to invisible Higgs and Z boson decays. The Higgs mass is increased from its MSSM tree-level value and brought above the LEP2 mass bound for a low scale of supersymmetry breaking $\sqrt{f} \sim 2$ TeV to 7 TeV. The invisible decay rates of the Higgs and Z bosons into goldstino and neutralino are computed and shown to bring stronger constraints on f than their decays into goldstino pairs, computed previously, which are subleading in $1/f$.

1 Introduction

The interest in non-linear realizations of supersymmetry [1] goes back to the early days of supersymmetry. In the following we review the formalism of standard realization (Volkov-Akulov) of non-linear supersymmetry as well as the superfield description that was developed afterwards. As an application of the constrained superfield formalism, we compute all the couplings of the goldstino (which transforms non-linearly under Susy) to the minimal supersymmetric standard model (MSSM) fields and investigate some of their phenomenological consequences.

Spontaneous supersymmetry breaking at low energies predicts a nearly massless goldstino. This plays the role of the longitudinal component of the gravitino, which acquires a mass f/M_{Planck}, in the milli-eV range if the supersymmetry breaking scale $\sqrt{f}$ is in the multi-TeV region. By the equivalence theorem [2], it interacts with a strength $1/\sqrt{f}$ which is much stronger than the Planck suppressed couplings of the transverse gravitino, and is therefore well described by the gravity-decoupled limit of a massless Goldstone fermion. In this talk we discuss the low energy consequences of a light goldstino by assuming that $\sqrt{f}$ is an independent parameter that is a few orders of magnitude higher than the soft breaking terms in the MSSM. The picture is as follows: at high scales above $\sqrt{f}$ one has the MSSM and goldstino superfields. At TeV scales, higher than m_{soft} but below $\sqrt{f}$ one has the MSSM together with the "non-linear" goldstino. At lower energies, below superpartners masses ($\sim m_{soft}$) one is left only with the goldstino fermion coupled to the SM fields. We consider the energy region $E \sim m_{soft} < \sqrt{f}$ when goldstino effective interactions are determined by non-linear supersymmetry.

The self-interactions of the goldstino are given by the famous Volkov-Akulov action [1]. This method, geometric in nature, gives also a universal coupling of goldstino to matter through its

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energy momentum tensor, of the form $(1/f^2)T_{\mu\nu}t^{\mu\nu}$, where $T_{\mu\nu}$, $t_{\mu\nu}$ are the stress tensors of matter and of (free) goldstino, respectively [3, 4]. It was noticed however that this coupling is not the most general invariant under non-linear supersymmetry [4, 5, 6]. General invariant couplings can be derived using the so-called superfield formulations of non-linear supersymmetry. One of them promotes any ordinary field to a superfield by introducing a modified superspace that takes account of the non-linear supersymmetry transformations of the goldstino [7, 8, 9, 10]. Finally, there is the method of constrained superfields: these are usual superfields, but they are also subject to constraints that eliminate the superpartners in terms of the light degrees of freedom and of the (Weyl) goldstino [11, 12, 13, 14].

In the following we shall review briefly these methods and present the corresponding goldstino couplings induced by a given Lagrangian upon "supersymmetrizing" it in a non-linear way. We then review the constrained superfields method, and then discuss the phenomenological implications for the case of the Minimal Supersymmetric Standard Model (MSSM) which is coupled to a non-linearly transforming goldstino field. These phenomenological consequences refer to: 1) the corrections to the Higgs mass, which can be increased already at the tree level, relative to its MSSM tree-level value for a hidden sector supersymmetry scale in the region of few TeV and to 2) the invisible decays of the Higgs and Z bosons. Indeed, to the leading order in the goldstino coupling, the Higgs and Z bosons can decay into a goldstino and a (next-to-lightest NLSP) neutralino[1], with a significant decay rate, if the hidden sector supersymmetry scale is low (and the NLSP light enough). From the accurately measured Γ_Z decay rate of Z boson one can set bounds on the scale of Susy breaking $\sqrt{f}$. Regarding the Higgs decay, this partial decay rate can become comparable to that into a pair of photons in the MSSM. Moreover, it dominates over other decay channels, such as Z into a pair of goldstinos which is $\mathcal{O}(1/f^2)$ and was considered previously in the literature. These effects are presented in the last section.

2 Non-linear supersymmetry in the standard realization

To derive the non-linear supersymmetry transformation law one considers first a supersymmetry transformation $x^{\mu\,\prime} = x^\mu + i\,(\theta\sigma^\mu\overline{\xi} - \xi\sigma^\mu\overline{\theta})$; $\theta' = \theta + \xi$, and $\overline{\theta}' = \overline{\theta} + \overline{\xi}$, where the spinor indices are not shown. This transformation induces a non-linear realization on the spinors θ which can actually be generalized to arbitrary spinor fields, in particular the Weyl goldstino spinor χ, by an analogy between θ and χ; here $\theta = \kappa\chi$ where κ is introduced on dimensional grounds. The analogy gives [10]

$$\chi'(x') = \chi(x) + \frac{1}{\kappa}\,\xi, \tag{1}$$

After Taylor expanding about x one finds the transformation of the Weyl goldstino field

$$\delta\chi \;=\; \frac{1}{\kappa}\,\xi_\alpha + \kappa\,\Lambda^\mu_\xi\,\partial_\mu\,\chi, \qquad \text{where} \qquad \Lambda^\mu_\xi = i\,(\chi\sigma^\mu\bar{\xi} - \xi\sigma^\mu\bar{\chi}) \tag{2}$$

A similar relation, hermitian conjugate of that above, holds true for $\overline{\chi}$. Here Λ^μ_ξ is a goldstino field dependent translation vector and ξ_α, $\bar{\xi}^{\dot{\alpha}}$ are the (Grassmann) parameters of the transformation; κ is the goldstino decay constant; this constant has dimension of mass^{-2} and is related to the (hidden sector) supersymmetry breaking scale $\sqrt{f} = \Lambda_{susy}$:

$$\kappa = \frac{1}{\sqrt{2}\,f} = \frac{1}{\sqrt{2}\,\Lambda^2_{susy}} \tag{3}$$

We wish to construct a Lagrangian for an effective low-energy description of the goldstino and its interactions with Standard Model fields. We first consider the part of the effective action

[1] The goldstino is the LSP particle.

which contains only self-couplings of the goldstino. This must contain the standard kinetic term for a Weyl spinor with additional terms necessary to make the action invariant under the standard non-linear realization (2). An action which satisfies these criteria has been constructed by Akulov and Volkov [1], where one defines the 'vierbein'

$$E_\mu^{\ \nu} = \delta_\mu^{\ \nu} + i\,\kappa^2(\partial_\mu\chi\sigma^\nu\bar{\chi} - \chi\sigma^\nu\partial_\mu\bar{\chi}) \tag{4}$$

It can be shown using (2) that

$$\delta(\det\, E) = i\kappa\,\partial_\mu[(\chi\sigma^\mu\bar{\xi} - \xi\sigma^\mu\bar{\chi})\det E] \equiv \kappa\,\partial_\mu(\Lambda_\xi^\mu\det E)\,, \tag{5}$$

It is then natural to construct the Akulov-Volkov Lagrangian [1],

$$\mathcal{L}_{AV} = -\frac{1}{2\kappa^2}\det E \;=\; -\frac{1}{2\kappa^2} + \frac{i}{2}\,(\chi\sigma^\mu\partial_\mu\bar{\chi} - \partial_\mu\chi\sigma^\mu\bar{\chi}) + ... \tag{6}$$

so the action defined by $\mathcal{L}_{AV}$ is invariant. In the last step in (6) we expanded in powers of κ; the first term in the rhs is the cosmological constant; the second term is the usual kinetic term for the goldstino; the dots denote self-couplings proportional to the second or higher powers of κ and give the non-linear Susy extension of the kinetic energy of a Weyl spinor. The expansion actually stops because of the anticommuting properties of the goldstino spinor.

The supersymmetry algebra can be also realized non-linearly on other (non-goldstino) fields such as matter and gauge fields [3, 4]. Let ϕ_i denote some generic field, where i an index in some representation of the Lorentz group or of an internal symmetry group. One defines

$$\delta\phi_i = \kappa\;\Lambda_\xi^\mu\,\partial_\mu\phi_i\,. \tag{7}$$

This is usually referred to as the standard realization. It can be checked that this provides a representation of supersymmetry. Eq. (7) has the same form as the transformation of the goldstino (2) except for the absence of the first term. However, derivatives of ϕ_i (the field strength $F_{\mu\nu}$ and gauge covariant derivatives $D_\mu\phi_i$) do not transform covariantly according to the standard realization, even though the field ϕ_i does so. This can be avoided, by generalizing the ordinary gauge-covariant derivative and defining the non-linearly realized supersymmetry-covariant derivative

$$\mathcal{D}_\mu\phi_i \equiv (E^{-1})_\mu^{\ \nu}D_\nu\phi_i\,. \tag{8}$$

where $(E^{-1})_\mu^{\ \nu}$ is the inverse of matrix (4). If ϕ_i is a field transforming in the standard realization then $\mathcal{D}_\mu\phi_i$ also varies according to the standard realization $\delta(\mathcal{D}_\mu\phi_i) = \kappa\,\Lambda_\xi^\mu\,(\mathcal{D}_\mu\phi_i)$. A similar procedure can be applied to the field strength:

$$\mathcal{F}_{\mu\nu}^a \equiv (E^{-1})_\mu^{\ \sigma}(E^{-1})_\nu^{\ \rho}F_{\sigma\rho}^a\,, \tag{9}$$

where $F_{\mu\nu}^a$ is the ordinary field-strength. Therefore $\delta\,(\mathcal{F}_{\mu\nu}^a) = \kappa\,\Lambda_\xi^\rho\,\partial_\rho\,\mathcal{F}_{\mu\nu}^a$ and $\mathcal{F}_{\mu\nu}^a$ transforms according to the standard realization. If we expand the rhs in (9) in powers of κ, the first term will be $F_{\mu\nu}^a$, followed by appropriate couplings to the goldstino field. Using these building blocks, one can then construct an invariant effective action. If in the SM Lagrangian

$$\mathcal{L}_{SM} = \mathcal{L}_{SM}(\phi_i, D_\mu\phi_i, F_{\mu\nu}^a), \tag{10}$$

we replace all variables of $\mathcal{L}_{SM}$ by their counterparts: $F_{\mu\nu}^a \to \mathcal{F}_{\mu\nu}^a$ and $D_\mu \to \mathcal{D}_\mu$, the resulting Lagrangian itself transforms like a field in the standard realization:

$$\delta\,\left(\mathcal{L}_{SM}(\phi_i, \mathcal{D}_\mu\phi_i, \mathcal{F}_{\mu\nu}^a)\right) = \kappa\,\Lambda_\xi^\sigma\,\partial_\sigma\mathcal{L}_{SM}(\phi_i, \mathcal{D}_\mu\phi_i, \mathcal{F}_{\mu\nu}^a)\,. \tag{11}$$

We further multiply $\mathcal{L}_{SM}$ by $\det E$ to find

$$S_{eff} = \int d^4x \, \mathcal{L}_{eff} = \int d^4x \, \det E \, \mathcal{L}_{SM}(\phi_i, \mathcal{D}_\mu \phi_i, \mathcal{F}^a_{\mu\nu}) \,. \tag{12}$$

which is invariant under the supersymmetry transformation, as it can be seen by using the transformation of $\det E$ given in (5). Expanding the effective Lagrangian $\mathcal{L}_{eff}$ in powers of κ, the lowest (κ-independent) term is the Standard Model (SM) Lagrangian itself. The additional, κ-dependent terms are needed to render the action supersymmetric and they describe appropriate interactions of the SM-fields to the goldstino:

$$\mathcal{L}_{eff} = \mathcal{L}_{SM}(\phi_i, D_\mu \phi_i, F^a_{\mu\nu}) + (i\kappa^2 \chi \overleftrightarrow{\partial^\mu} \sigma^\nu \bar{\chi}) \, T_{\mu\nu} + ... \tag{13}$$

where the dots denote higher powers of κ (negligible in the low-energy limit) and $T_{\mu\nu}$ is the gauge-invariant energy-momentum tensor. The tensor coupling (13) is completely determined and it is model-independent. However, the above procedure to construct the goldstino couplings does not give the most general effective action invariant under non-linear supersymmetry. This is because additional supersymmetric terms can in principle be added to the above effective Lagrangian. Such terms, supersymmetric on their own, could come with an arbitrary relative normalization, that cannot be determined within the effective field theory.

One can employ two superfield formalisms to find general couplings of the goldstino to a given Lagrangian. The first of them promotes any field to a superfield, by introducing a modified superspace that takes into account the non-linear supersymmetry transformations of the goldstino. This amounts of shifting the spacetime coordinates x^μ by Λ^μ_θ given in (2), where θ are the usual superspace (Weyl) fermionic coordinates: $x^\mu \rightarrow \tilde{x}^\mu = x^\mu + \Lambda^\mu_\theta(\tilde{x})$. A similar procedure is applied to the goldstino Weyl fermion, to which one associates the goldstino superfield: $\chi(x) \rightarrow \theta/\kappa + \chi(\tilde{x})$. With these objects one can proceed to construct a supersymmetric action, using covariant derivatives as before in terms of the 'vierbein' (4) depending on the shifted coordinates $\tilde{x}$. For details see [7, 8, 9, 10]. The second formalism to find the general couplings of the goldstino to a given Lagrangian is the method of constrained superfields.

3 Non-linear supersymmetry with constrained superfields

A very convenient approach to non-linear supersymmetry is that using constrained superfields. The method preserves all the advantages of working with the standard superfield formalism. Ultimately, the role of these constraints is to eliminate the "massive" superpartners (of the SM and of the goldstino) in terms of the light degrees of freedom (SM particles and the goldstino) [11, 12, 13, 14] (see also [15, 16]). In this way a non-linear realization of supersymmetry is obtained, in which no superpartners are present anymore. As in the standard realization of Volkov-Akulov, in constructing a non-linear supersymmetric version of a given model the goldstino chiral superfield (SM gauge singlet) X plays the leading role. To see how this works and the relation to $\mathcal{L}_{AV}$ of eq. (6), consider the Lagrangian [13, 14] (f below is related to κ as in (3)):

$$\mathcal{L}_X = \int d^4\theta \, X^\dagger X - \left\{ f \int d^2\theta \, X + h.c. \right\} = |\partial_\mu \varphi|^2 + F^\dagger F + \Big(\frac{i}{2} \, \overline{\chi \sigma}^\mu \partial_\mu \chi - f \, F + h.c. \Big) \tag{14}$$

and impose a constraint on the goldstino superfield $X^2_{nl} = 0$. This constraint is solved by

$$X_{bl} \;\; = \;\; \varphi + \sqrt{2} \, \theta \chi + \theta\theta \, F, \qquad \text{where} \qquad \varphi = \frac{\chi\chi}{2\,F} \tag{15}$$

which shows explicitly how the sgoldstino (φ) is eliminated in terms of goldstino χ. The use of this X_{nl} solution in $\mathcal{L}_X$ gives, after eliminating the auxiliary fields [13, 14]:

$$\mathcal{L}_X = -\frac{1}{2\kappa^2} + i\,\overline{\chi\sigma}^\mu \partial_\mu \chi + \frac{\kappa^2}{2} \overline{\chi}^2 \partial^2 \chi^2 - \frac{\kappa^6}{2}\,\chi^2 \overline{\chi}^2\, \partial^2 \chi^2\, \partial^2 \overline{\chi}^2 \tag{16}$$

which is just the Volkov-Akulov Lagrangian $\mathcal{L}_{AV}$, that is, $\mathcal{L}_X$ is equivalent *on shell* to $\mathcal{L}_{AV}$ (this is not true off shell, due to having different degrees of freedom) [13]. Also, on shell, $F = f + \ldots.$ where f is the hidden sector SUSY breaking scale. In the infrared description of the SUSY breaking, the scalar component (sgoldstino) φ becomes a function of the goldstino χ and is therefore removed from the spectrum which contains only χ, as in $\mathcal{L}_{AV}$. This is the situation for the case of a single constrained superfield, the goldstino.

There can be additional constraints for the matter and vector superfields, which essentially integrate out some of the superpartners in terms of the light degrees of freedom. For example

$$\begin{aligned}
Q_{nl}\, X_{nl} &= 0 && \text{eliminates sfermions, leaves Weyl fermion.} \\
\overline{D}^{\dot\alpha} \left[X_{nl}\, H_{nl}^\dagger \right] &= 0 && \text{eliminates higgsino, leaves complex Higgs.} \\
X_{nl} A_{nl} - X_{nl}\, A_{nl}^\dagger &= 0 && \text{eliminates fermion, leaves real scalar} \\
W_{nl}\, X_{nl} &= 0 && \text{eliminates gauginos, leaves gauge fields.}
\end{aligned} \tag{17}$$

In this way, by adding all these constraints to $\mathcal{L}$, one obtains a supersymmetric Lagrangian in which all superpartners are "integrated out" and are absent in the infrared, being expressed in terms of the light degrees of freedom. The latter are the only degrees of freedom present in the low energy spectrum and have non-linear supersymmetry transformation rules. However, not all of the above constraints need be satisfied simultaneously. For example, in the energy region $m_{soft} \sim E \leq \sqrt{f}$, the only constraint is that of the goldstino superfield that undergoes non-linear transformation, while for the MSSM case, its superfields are unconstrained.

The next step is to find goldstino couplings to a given Lagrangian, like that of MSSM. The strategy is to write down an effective expansion of the Lagrangian (function of all superfields, including the goldstino) in powers of $1/f$ and restrict it to a given order in it. The method is general and useful to find couplings involving more than one goldstino [14]. Let us follow this idea and consider a supersymmetric theory with chiral multiplets $\Phi_i \equiv (\phi_i, \psi_i, F_i)$ and vector multiplets $V \equiv (A_\mu^a, \lambda^a, D^a)$ which couples in the most general way to X_{nl}. Its Lagrangian is

$$\begin{aligned}
\mathcal{L} &= \int d^4\theta \left[X_{nl}^\dagger X_{nl} + \Phi_i^\dagger (e^V \Phi)_i - (m_i^2/f^2)\, X_{nl}^\dagger X_{nl} \Phi_i^\dagger (e^V \Phi)_i \right] + \left\{ \int d^2\theta \left[f X_{nl} + W(\Phi_i) \right.\right. \\
&+ \left.\left. \frac{B_{ij}}{2f} X_{nl}\, \Phi_i \Phi_j + \frac{A_{ijk}}{6\,f} X_{nl} \Phi_i \Phi_j \Phi_k + \frac{1}{4}\Big(1 + \frac{2\,m_\lambda}{f} X_{nl}\Big)\, \mathrm{Tr}\, W^\alpha W_\alpha \right] + \text{h.c.} \right\},
\end{aligned} \tag{18}$$

where m_i^2, B_{ij}, A_{ijk} are soft terms, m_λ is the gaugino mass. From this, one can find the goldstino (χ) couplings to ordinary matter/gauge superfields, in the leading orders in $1/f$.

It can be checked that these couplings are equivalent to those obtained "in the standard way" by the equivalence theorem [2], from a theory with the corresponding explicit soft breaking (see below) in which the goldstino couples as:

$$(1/f)\ \partial^\mu \chi\ J_\mu = \ -(1/f)\, \chi\ \partial^\mu J_\mu + (\text{total space-time derivative}), \tag{19}$$

Here J_μ is the supercurrent of the theory corresponding to that in (18) in which the goldstino is essentially replaced by the spurion S, according to $S \equiv m_{soft}\theta^2 \leftrightarrow (m_{soft}/f) X_{nl}$ [14], with the corresponding explicit soft breaking terms:

$$\mathcal{L}' = \int d^4\theta \left[1 - m_i^2\,\theta^2\overline{\theta}^2\right] \Phi_i^\dagger (e^V\Phi)_i + \int d^2\theta \left[W(\Phi_i) - (1/2)\,B_{ij}\,\theta^2\Phi_i\Phi_j - (1/6)\,A_{ijk}\,\theta^2\,\Phi_i\Phi_j\Phi_k + \frac{1}{4}\,(1-2m_\lambda\theta^2)\,\mathrm{Tr}\,W^\alpha W_\alpha\right] + \text{h.c.}\,, \tag{20}$$

With this, eq.(19) shows that, on-shell, all goldstino couplings are proportional to soft terms. Indeed, the supercurrent of (20) and its divergence are given by (with $\mathcal{D}_{\mu,ij} = \delta_{ij}\,\partial_\mu + i\,g\,A_\mu^a\,T_{ij}^a$)

$$\begin{aligned} J_\alpha^\mu &= -[\sigma^\nu\overline{\sigma}^\mu\psi_i]_\alpha\,[\mathcal{D}_{\nu,ij}\phi_j]^\dagger + i\,[\sigma^\mu\overline{\psi}_i]_\alpha F_i - \frac{1}{2\sqrt{2}}[\sigma^\nu\overline{\sigma}^\rho\,\sigma^\mu\overline{\lambda}^a]_\alpha\,F_{\nu\rho}^a + \frac{i}{\sqrt{2}}D^a\,[\sigma^\mu\overline{\lambda}^a]_\alpha \\ \partial_\mu J_\alpha^\mu &= \psi_{i,\alpha}\,(m_i^2{\phi^\dagger}_j + B_{ij}\phi_j + (1/2)A_{ijk}\phi_j\phi_k\,) + \frac{m_\lambda}{\sqrt{2}}\left[(\sigma^{\mu\nu})_\alpha^{\ \beta}\,\lambda_\beta^a F_{\mu\nu}^a + D^a\,\lambda_\alpha^a\right]. \end{aligned} \tag{21}$$

so all goldstino couplings are proportional to the soft terms. From (19), (21) one then recovers the couplings with one goldstino. However, the constrained superfields formalism in (18) has an advantage over this "standard procedure" in that it can be applied even when evaluating couplings with more than one goldstino. As mentioned, this is done by writing all effective operators (involving X_{nl}) to a fixed order in $1/f$. It is more difficult to find these from (20). It should be mentioned however, that this method does not take into account possible derivative couplings of the goldstino to the matter, which can arise, through the Ferrara-Zumino current.

4 Goldstino couplings to the MSSM and their implications.

In the following we explore some phenomenological applications of the above formalism. We only impose the constraint on the goldstino superfield, ie this is the only field with a non-linear supersymmetry transformation in the energy region $E \sim m_{soft} < \sqrt{f}$. At energy scales below m_{soft}, constraints similar to that of goldstino superfield must be applied to the MSSM superfields themselves, corresponding to integrating out the corresponding superpartners in terms of the light degrees of freedom. Nevertheless, the problems that we address are not affected at tree level by the additional constraints on the MSSM superfields such as quarks and leptons superfields. These can actually be imposed at a later stage on the results found. For related applications of goldstino interactions and phenomenology see also references [17] to [22].

4.1 Effective Lagrangian of goldstino couplings to MSSM

Using the method of constrained superfields we determine the general couplings of the goldstino to MSSM superfields that involve one and two goldstino spinors. As mentioned, we do not take into account the derivative couplings of the goldstino to matter, which can in principle arise. Here, the only difference from the ordinary MSSM is in the supersymmetry breaking sector. Supersymmetry is broken spontaneously via a vacuum expectation value of F, fixed by its equation of motion (see later). The Lagrangian of the MSSM coupled to the goldstino is then

$$\begin{aligned} \mathcal{L}_H &= \sum_{i=1,2} c_i \int d^4\theta\; X_{nl}^\dagger X_{nl}\; H_i^\dagger\, e^{V_i}\, H_i + \sum_\Phi c_\Phi \int d^4\theta\; X_{nl}^\dagger X_{nl}\;\Phi^\dagger e^V\,\Phi + \frac{B}{f}\int d^2\theta\, X_{nl}\, H_1\, H_2 \\ &+ \frac{A_u}{f}\int d^2\theta\, X_{nl}\, H_2\, Q\, U^c + \frac{A_d}{f}\int d^2\theta\, X_{nl}\, Q\, D^c\, H_1 + \frac{A_e}{f}\int d^2\theta\, X_{nl}\, L\, E^c\, H_1 + h.c. \\ &+ \sum_{i=1}^{3}\frac{1}{16\,g_i^2\,\zeta}\frac{2\,m_{\lambda_i}}{f}\int d^2\theta\, X_{nl}\,\mathrm{Tr}\,[\,W^\alpha\,W_\alpha]_i + h.c. \end{aligned} \tag{22}$$

where an effective expansion in $1/f$ is assumed. Here $c_\Phi = -m_\Phi^2/f^2$, $\Phi : Q, U^c, D^c, L, E^c$; m_Φ are soft masses which we can choose to be all equal to m_0; ζ cancels the Trace factor.

After eliminating the auxiliary fields one obtains new couplings $\mathcal{L}^{new}$ beyond those of the usual onshell, supersymmetric part of MSSM. The onshell Lagrangian is

$$\mathcal{L}^{new} \equiv \mathcal{L}^{aux}_{F\,(1)} + \mathcal{L}^{aux}_{F\,(2)} + \mathcal{L}^{aux}_{D} + \mathcal{L}^{extra}_{m} + \mathcal{L}^{extra}_{g} \tag{23}$$

The first two terms are obtained from eliminating the auxiliary F fields; $\mathcal{L}^{aux}_{F,(1)}$ recovers all MSSM soft terms (not shown) plus a cosmological term; for $\mathcal{L}^{aux}_{F,(2)}$ one has, up to $\mathcal{O}(1/f^3)$

$$\begin{aligned}
\mathcal{L}^{aux}_{F\,(2)} &= \Big\{ \frac{\overline{\chi}\overline{\chi}}{2\,f^2}\Big[\mu\big(m_1^2+m_2^2\big)\,h_1.h_2-\big(m_1^2+m_Q^2+m_D^2\big)h_1.\phi_Q\phi_D-\big(m_1^2+m_L^2+m_E^2\big)h_1.\phi_L\phi_E \\
&- \;\big(m_2^2+m_Q^2+m_U^2\big)\phi_Q\phi_U.h_2+\big(B\,h_2-A_d\,\phi_Q\phi_D-A_e\,\phi_L\,\phi_E\big)^\dagger\big(\mu h_2-\phi_Q\phi_D-\phi_L\,\phi_E\big) \\
&+ \;\big(B\,h_1-A_u\,\phi_Q\,\phi_U\big)^\dagger\big(\mu\,h_1-\phi_Q\,\phi_U\big)+\big(A_d\,\phi_D\,h_1-A_u\,h_2\,\phi_U\big)^\dagger\big(\phi_D\,h_1-h_2\,\phi_U\big) \\
&+ \;A_d\left(|\phi_Q.h_1|^2+|\phi_E\,h_1|^2\right)+A_u\,|h_2.\phi_Q|^2+A_e\,|\phi_L.h_1|^2\Big]+h.c.\Big\}-\frac{1}{f^2}\Big|B\,h_1.h_2 \\
&+ A_u h_2.\phi_Q\,\phi_U+A_d\phi_Q\,\phi_D.h_1+A_e\phi_L\,\phi_E.h_1+\frac{m_{\lambda_i}}{2}\lambda_i\lambda_i+\big(m_1^2|h_1|^2+m_2^2|h_2|^2+m_\Phi^2|\phi_\Phi|^2\big)\Big|^2 \\
&- \;\frac{1}{f}\,\Big[m_1^2\,\overline{\chi}\overline{\psi}_{h_1}\,h_1+m_2^2\,\overline{\chi}\overline{\psi}_{h_2}\,h_2+m_\Phi^2\,\overline{\chi}\overline{\psi}_\Phi\,\phi_\Phi+h.c.\Big]+\mathcal{O}(1/f^3)
\end{aligned} \tag{24}$$

which contains Weyl goldstino couplings to MSSM fields, but also couplings independent of χ such as new quartic Higgs couplings! A summation is understood over the SM group indices: $i=1,2,3$ in the gaugino term and over $\Phi=Q,U^c,D^c,L,E^c$ in the mass terms; appropriate contractions among $SU(2)_L$ doublets are understood for holomorphic products, when the order displayed is relevant. There are leading interactions $\mathcal{O}(1/f)$ in the last line which are actually dimension-four in fields. Similar couplings exist at $\mathcal{O}(1/f^2)$ and involve scalar and gaugino fields. Finally, the Yukawa matrices are restored in (24) by replacing $\phi_Q\phi_D\to\phi_Q\gamma_d\phi_D$, $\phi_Q\phi_U\to\phi_Q\gamma_u\phi_U$, $\phi_L\phi_E\to\phi_L\gamma_e\phi_E$.

There are also new couplings from terms involving the auxiliary components of the vector superfields of the SM. Eliminating them one finds

$$\mathcal{L}^{aux}_D = \frac{-1}{2}\sum_{i=1}^{3}\Big[\tilde{D}^a_i+\frac{1}{4\,f^2}\left(m_{\lambda_i}\,\chi\chi+h.c.\right)\tilde{D}^a_i+\frac{1}{\sqrt{2}\,f}\big(m_{\lambda_i}\,\chi\,\lambda^a_i+h.c.\big)\Big]^2+\mathcal{O}(1/f^3) \tag{25}$$

which brings gaugino-goldstino and Higgs-goldstino couplings; $\tilde{D}^a_i$ are MSSM auxiliary fields

$$\begin{aligned}
\tilde{D}_1 &= -1/2\,g_1\big(-h_1^\dagger h_1+h_2^\dagger h_2+1/3\;\phi_Q^\dagger\phi_Q-4/3\;\phi_U^\dagger\phi_U+2/3\;\phi_D^\dagger\phi_D-\phi_L^\dagger\phi_L+2\,\phi_E^\dagger\phi_E\big) \\
\tilde{D}_2^a &= -1/2\,g_2\left(h_1^\dagger\sigma^a h_1+h_2^\dagger\sigma^a h_2+\phi_Q^\dagger\sigma^a\phi_Q+\phi_L^\dagger\sigma^a\phi_L\right) \\
\tilde{D}_3^a &= -1/2\,g_3\left(\phi_Q^\dagger\,t^a\phi_Q-\phi_U^\dagger\,t^a\phi_U-\phi_D^\dagger\,t^a\phi_D\right)
\end{aligned} \tag{26}$$

Further, the total Lagrangian also contains new couplings, not induced by the auxiliary fields and absent in the MSSM. In the matter sector these are:

$$
\begin{aligned}
\mathcal{L}_m^{extra} = & \frac{1}{4f^2}|\partial_\mu(\chi\chi)|^2 + \Big(\frac{i}{2}\overline{\chi}\overline{\sigma}^\mu\,\partial_\mu\chi + h.c.\Big)\\
- & \sum_{i=1}^{2}\frac{m_i^2}{f^2}\Big\{\overline{\chi}\overline{\psi}_{h_i}\,\chi\psi_{h_i} + \Big[\frac{i}{2}\,(\overline{\chi}\,\overline{\sigma}^\mu\,\chi)\,(h_i^\dagger\,\mathcal{D}_\mu\,h_i) + \frac{i}{2}|h_i|^2\,\overline{\chi}\,\overline{\sigma}^\mu\partial_\mu\chi + h.c.\Big]\Big\}\\
- & \Big[m_i^2 \to m_\Phi^2, H_i \to \Phi\Big] + \Big\{\ \frac{B}{f}\ \Big[\frac{1}{2\,f}\ \chi\chi\,\psi_{h_1}.\psi_{h_2} - h_1.(\chi\psi_{h_2}) - (\chi\psi_{h_1}).h_2\Big]\\
+ & \frac{A_u}{f}\Big[\frac{1}{2\,f}\chi\chi\,\big(\,h_2.\psi_Q\,\psi_U + \psi_{h_2}.\phi_Q\,\psi_U + \psi_{h_2}.\psi_Q\,\phi_U\big) - \chi\,(h_2.\phi_Q\,\psi_U + h_2.\psi_Q\,\phi_U\\
+ & \psi_{h_2}.\phi_Q\,\phi_U)\Big] + \Big[\frac{A_d}{f}\,\Big(\frac{1}{2\,f}\ \chi\chi\,(\psi_Q\,\psi_D.h_1 + \,\phi_Q\,\psi_D.\psi_{h_1} + \psi_Q\,\phi_D.\psi_{h_1})\\
- & \chi\,(\phi_Q\,\psi_D.h_1 + \psi_Q\,\phi_D.h_1 + \phi_Q\,\phi_D.\psi_{h_1})\Big) + (D\to E, L\to Q)\Big] + h.c.\Big\} + \mathcal{O}(1/f^3). \quad (27)
\end{aligned}
$$

Note the presence of interactions that are dimension-four in fields like $(B/f)\,h_1\chi\psi_{h_2}$, relevant for phenomenology at low f. Finally, there are also new couplings in the gauge sector

$$
\begin{aligned}
\mathcal{L}_g^{extra} \ = & \sum_{i=1}^{3}\frac{m_{\lambda_i}}{2\,f}\ \Big[\frac{\chi\chi}{-2\,f}\,\Big(2\,i\,\lambda^a\sigma^\mu\,\Delta_\mu\,\overline{\lambda}^a - \frac{1}{2}\,F_{\mu\nu}^a\,F^{a\,\mu\nu} - \frac{i}{4}\epsilon^{\mu\nu\rho\sigma}\,F_{\mu\nu}^a\,F_{\rho\sigma}^a\Big)\\
- & \sqrt{2}\,\chi\sigma^{\mu\nu}\lambda^a\,F_{\mu\nu}^a\Big]_i + h.c. + \mathcal{O}(1/f^3), \qquad (28)
\end{aligned}
$$

with $i = 1, 2, 3$ the gauge group index and $\sigma^{\mu\nu} = i/4\,(\sigma^\mu\overline{\sigma}^\nu - \sigma^\nu\overline{\sigma}^\mu)$.

$\mathcal{L}^{new}$ together with the *onshell* part of the purely supersymmetric part of the MSSM Lagrangian give the final action which contains all the goldstino couplings to the MSSM with one and two goldstinos. Some of their phenomenological implications are discussed below.

4.2 MSSM Higgs potential and mass corrections

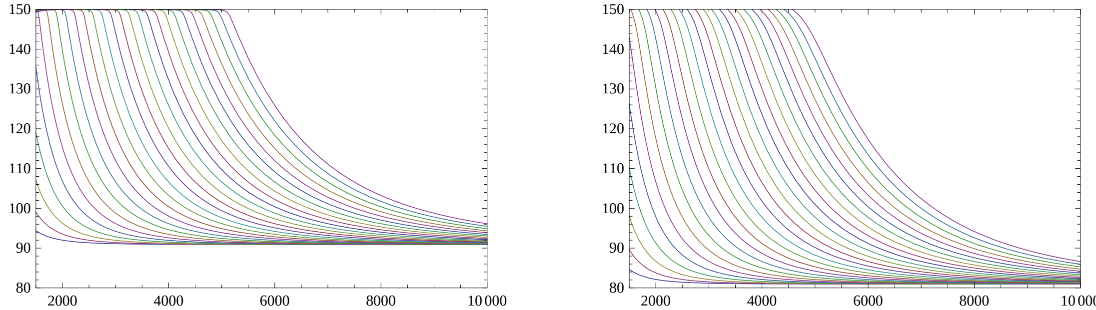
From $\mathcal{L}^{new}$ the full scalar potential is also identified. For the Higgs sector this is

$$
\begin{aligned}
V \ = \ & f^2 + \big(|\mu|^2 + m_1^2\big)\ |h_1|^2 + \big(|\mu|^2 + m_2^2\big)|h_2|^2 + \big(B\,h_1.h_2 + h.c.\big) \qquad (29)\\
+ & \frac{1}{f^2}\Big|m_1^2\,|h_1|^2 + m_2^2\,|h_2|^2 + B\,h_1.h_2\Big|^2 + \frac{g_1^2 + g_2^2}{8}\,\Big[|h_1|^2 - |h_2|^2\Big]^2 + \frac{g_2^2}{2}\,|h_1^\dagger\,h_2|^2 + \mathcal{O}(1/f^3)
\end{aligned}
$$

Notice the first term in the last line, generated onshell, which is dimension-four in fields and not present in the MSSM. Its origin is ultimately due to the constraint on the goldstino superfield and elimination of the sgoldstino in terms of light degrees of freedom (Weyl goldstino). When f is large, with soft masses fixed, one recovers the usual MSSM limit. This scalar potential brings corrections to the value of the Higgs mass of the MSSM. Their exact form can be found in [21]; at large $u \equiv \tan\beta$ with m_A fixed one finds

$$
\begin{aligned}
m_h^2 \ &= \ \Big[m_Z^2 + \mathcal{O}(1/u^2)\Big] + \frac{v^2}{2\,f^2}\,\Big[(2\,\mu^2 + m_Z^2)^2 + \mathcal{O}(1/u^2)\Big] + \mathcal{O}(1/f^3)\\
m_H^2 \ &= \ \Big[m_A^2 + \mathcal{O}(1/u^2)\Big] + \frac{1}{f^2}\,\mathcal{O}(1/u^2) + \mathcal{O}(1/f^3) \qquad (30)
\end{aligned}
$$

Therefore, m_h is increased from its MSSM tree-level value. This increase is driven by a large μ and apparently is of supersymmetric origin, but the quartic Higgs couplings giving this effect involved combinations of soft masses (see (29)). These soft masses combined to give, at the EW minimum, the μ-dependent increase in (30). As seen from Figs. 1 a), b), there is a reasonable range in the parameter space for which the Higgs mass can reach the LEP2 bound of 114.4 GeV even at the tree level, due to effective corrections associated with integrating out the scalar and



(a) m_h in function of $\sqrt{f}$, μ parameter

(b) m_h in function of $\sqrt{f}$, μ parameter

Figure 1: The tree-level Higgs masses (GeV) as functions of $\sqrt{f}$ (in GeV); here $m_A = 150$ GeV and m_h increases as μ varies from 400 to 3000 GeV in steps of 100 GeV. In (a) $\tan\beta = 50$ and in (b) $\tan\beta = 5$, showing a milder dependence on $\tan\beta$ than in MSSM.

auxiliary field of the goldstino superfield. This result is important since it does not require new physics in the visible sector to solve the discrepancy between MSSM tree level bound and the LEP2 bound. This is done for the case of a low enough supersymmetry breaking scale, $\sqrt{f} \sim 2$ to 7 TeV. This effect could also be relevant for a scenario in which the Higgs mass is found at a mass well above the 114.4 GeV bound[2], in which case the MSSM requires large quantum corrections (and a significant amount of fine tuning of the EW scale).

4.3 Invisible Higgs and Z bosons decays

The couplings of the goldstino field to the MSSM fields have some interesting consequences. One is that for a light neutralino NLSP (goldstino is the LSP) there is the possibility of a decay of the neutral higgses into a goldstino and the NLSP χ_1^0. The coupling Higgs-goldstino-neutralino is only suppressed by $1/f$. It arises from the following terms in $\mathcal{L}^{new}$ and from the terms in the supersymmetric part of the usual MSSM Lagrangian, hereafter denoted $\mathcal{L}_0^{onshell}$:

$$\begin{aligned}\mathcal{L}^{new} + \mathcal{L}_0^{onshell} \supset & -\frac{1}{f}\left[m_1^2\,\chi\psi_{h_1^0}\,h_1^{0\,*} + m_2^2\,\chi\psi_{h_2^0}\,h_2^{0\,*}\right] - \frac{B}{f}\left[\chi\psi_{h_2^0}\,h_1^0 + \chi\psi_{h_1^0}\,h_2^0\right] \\ & - \frac{1}{f}\sum_{i=1,2}\frac{m_{\lambda_i}}{\sqrt{2}}\,\tilde{D}_i^a\,\chi\lambda_i^a - \frac{1}{\sqrt{2}}\Big[g_2\lambda_2^3 - g_1\lambda_1\Big]\Big[h_1^{0\,*}\psi_{h_1^0} - h_2^{0\,*}\psi_{h_2^0}\Big] + h.c. \end{aligned} \tag{31}$$

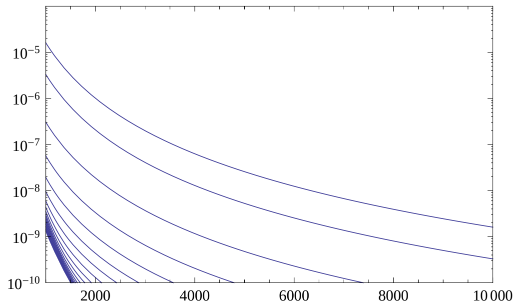
If the NLSP is light enough, $m_{\chi_1^0} < m_h$, then h^0, H^0 can decay into it plus a goldstino which has a mass of order $f/M_{Planck} \sim 10^{-3}$ eV[3]. The decay rate is

$$\Gamma_{h^0 \to \chi_1^0\,\chi} = \frac{m_h}{16\,\pi\,f^2}\,\Big|\sum_{k=1}^{4}\delta'_k\,\mathcal{Z}_{1k}\Big|^2\left(1 - \frac{m_{\chi_1^0}^2}{m_{h^0}^2}\right)^2 \tag{32}$$

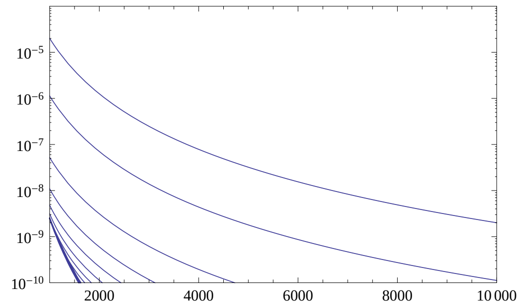
where $\mathcal{Z}$ is the matrix diagonalising the neutralino mass matrix. The partial decay rate has corrections coming from both higgsino ($\mathcal{Z}_{13}$, $\mathcal{Z}_{14}$) and gaugino fields ($\mathcal{Z}_{11}$, $\mathcal{Z}_{12}$), since they both acquire a goldstino component. The gaugino correction arises after gaugino-goldstino mixing, SUSY and EW symmetry breaking, (as shown by m_{λ_i}, m_Z dependence in δ'_k with the latter provided in [21]) and was not included in previous similar studies [22, 24, 25].

[2] In the MSSM the EW scale fine tuning is minimized at $m_h \approx 114 \pm 2$ GeV, beyond this value it increases exponentially, to a level of one part in 1000 for $m_h \approx 122$ GeV [27]. See also [28].

[3] If this is not true, the decay of neutralino into h^0 and goldstino takes place, examined in [25].



(a) $\Gamma_{h^0 \to \chi \psi_x}$ in function of $\sqrt{f}$.



(b) $\Gamma_{h^0 \to \chi \psi_x}$ in function of $\sqrt{f}$.

Figure 2: The partial decay rate of $h^0 \to \chi \chi_1^0$ for (a): $\tan\beta = 50$, $m_{\lambda_1} = 70$ GeV, $m_{\lambda_2} = 150$ GeV, μ increases from 50 GeV (top curve) by a step 50 GeV, $m_A = 150$ GeV. Compare against Fig. 1 (a) corresponding to a similar parameters range. At larger μ, m_h increases, but the partial decay rate decreases. Similar picture is obtained at low $\tan\beta \sim 5$. (b): As for (a) but with $\tan\beta = 5$. Compare against Fig. 1 (b). The total SM decay rate, for $m_h \sim 114$ GeV, is of order 10^{-3}, thus the branching ratio in the above cases becomes comparable to that of SM Higgs going into $\gamma\gamma$ (see Fig. 2 in [23]).

In Fig. 2 the partial decay rate is presented for various values of μ, m_A and $m_{\lambda_{1,2}}$. A larger decay rate requires a light $\mu \sim \mathcal{O}(100)$ GeV, when the neutralino χ_1^0 has a larger higgsino component. At the same time an increase of m_h above the LEP bound requires a larger value for μ, close to $\mu \approx 700$ GeV if $\sqrt{f} \approx 1.5$ TeV, and $\mu \approx 850$ GeV if $\sqrt{f} \approx 2$ TeV, see Figure 1 (a). The results in Figure 2 show that the partial decay rate can be significant ($\sim 3 \times 10^{-6}$ GeV), if we recall that the total SM Higgs decay rate (for $m_h \approx 114$ GeV) is about 3×10^{-3} GeV, with a branching ratio of $h^0 \to \gamma\gamma$ of 2×10^{-3}, (Figure 2 in [23]). Thus the branching ratio of the process can be close to that of SM $h^0 \to \gamma\gamma$. The decay is not very sensitive to $\tan\beta$ (Figure 2 (b)), due to the extra contribution (beyond MSSM) from the quartic Higgs coupling. It would be interesting to analyze the above decay rate at the one-loop level, for a more careful comparison to SM Higgs decays rates.

Another coupling that is also present in the leading order $(1/f)$ is that of goldstino to Z_μ boson and to a neutralino. Depending on the relative mass relations, it can bring about a decay of Z_μ (χ_j^0) into χ_j^0 (Z_μ) and a goldstino, respectively. The relevant terms are

$$\begin{aligned} \mathcal{L}^{new} + \mathcal{L}_0^{onshell} \supset & -\frac{1}{4}\, \overline{\psi}_{h_1^0} \overline{\sigma}^\mu \psi_{h_1^0}\, (g_2 V_2^3 - g_1\, V_1)_\mu + \frac{1}{4}\, \overline{\psi}_{h_2^0} \overline{\sigma}^\mu \psi_{h_2^0}\, (g_2 V_2^3 - g_1\, V_1)_\mu \Big\} \\ & - \sum_{i=1}^{2} \frac{m_{\lambda_i}}{\sqrt{2}\, f}\, \chi\, \sigma^{\mu\nu}\, \lambda_i^a\, F^a_{\mu\nu,\, i} + h.c. \end{aligned} \tag{33}$$

where the last term was generated in (28) (i labels the gauge group). Since the higgsinos acquired a goldstino component ($\propto \chi/f$) via mass mixing, the first line above induces additional $\mathcal{O}(1/f)$ couplings of the higgsino to goldstino and to $Z_\mu = (1/g)\,(g_2 V_2^3 - g_1\, V_1)_\mu$ with $g^2 = g_1^2 + g_2^2$. Then, if $m_{\chi_1^0}$ is lighter than Z_μ then a decay of the latter into $\chi_1^0 + \chi$ is possible. The decay rate of this process is (with $j = 1$):

$$\Gamma_{Z \to \chi \chi_j^0} = \frac{m_Z^5}{32\pi f^2} \Big[\zeta_1 |w_j|^2 + \zeta_2\, |v_j|^2 + \zeta_3\, (w_j\, v_j^* + w_j^*\, v_j)\Big] \Big(1 - \frac{m_{\chi_j}^2}{m_Z^2}\Big)^2 \tag{34}$$

with $\zeta_1 = 2(2 + r^2)\,\mu^2/m_Z^2$, $\zeta_2 = 2(8 + r^2)(1 + 2r^2)$, $\zeta_3 = -2(4 + 5r^2)\mu/m_Z$ where $r = m_{\chi_j}/m_Z$. This decay rate should be within the LEP error for Γ_Z, which is 2.3 MeV [26] (ignoring theoretical uncertainties which are small). From this, one finds a lower bound for $\sqrt{f}$, which can be as high as $\sqrt{f} \approx 700$ GeV for the parameter space considered previously in Figure 1, while generic values are $\sqrt{f} \sim \mathcal{O}(400)$ GeV. Therefore the results for the increase of m_h, that needed a value for $\sqrt{f}$ in the TeV region, escape this constraint. This constraint does not apply if the lightest neutralino has a mass larger than m_Z (this can be arranged for example by a larger m_{λ_1}), when the opposite decay ($\chi_j \to Z\,\chi$) takes place.

There also exists the interesting possibility of an invisible decay of Z_μ gauge boson into a pair of goldstino fields, that we review here [6, 14, 18]. This is induced by the following terms in the Lagrangian, after the Higgs field acquires a VEV:

$$\begin{aligned}\mathcal{L}^{new} + \mathcal{L}_0^{onshell} \supset\ & \Big\{ \frac{1}{4\,f^2}\,\overline{\chi}\overline{\sigma}^\mu \chi\ (g_2 V_2^3 - g_1\,V_1)_\mu\,(m_1^2\,v_1^2/2 - m_2^2\,v_2^2/2) \\ & - \frac{1}{4}\,\overline{\psi}_{h_1^0}\overline{\sigma}^\mu \psi_{h_1^0}\ (g_2 V_2^3 - g_1\,V_1)_\mu + \frac{1}{4}\,\overline{\psi}_{h_2^0}\overline{\sigma}^\mu \psi_{h_2^0}\ (g_2 V_2^3 - g_1\,V_1)_\mu \Big\} + h.c. \quad (35)\end{aligned}$$

The decay rate is then

$$\Gamma_{Z\to\chi\chi} = \frac{m_Z}{24\,\pi\,g^2}\left[\frac{m_Z^4}{2\,f^2}\right]^2 \cos^2 2\beta \tag{36}$$

in agreement with previous results obtained for $B = 0$ [6, 14, 18]. The decay rate is independent of m_A and should be within the LEP error for Γ_Z (2.3 MeV [26]). One can then easily see that the increase of the Higgs mass above the LEP bound (114.4 GeV) seen earlier in Figure 1 is consistent with the current bounds for this decay rate, which thus places only mild constraints on f, below the TeV scale (≈ 200 GeV) [6, 18]. This is due to the extra suppression powers of $1/f$ relative to the decay rates evaluated earlier in this section. A similar analysis can also be done for the decay of the lightest higss into a pair of goldstinos, but it does not bring any significant constraints for f, being strongly suppressed ($\sim 1/f^4$).

5 Conclusions

We reviewed briefly the non-linear supersymmetry in the standard realization of Volkov-Akulov and in the constrained superfield formalism. The latter method can be easily applied to the MSSM model, to derive the couplings of the MSSM superfields to the goldstino fermion. This is possible for an energy region above m_{soft} and below the scale of supersymmetry breaking $\sqrt{f}$ which is assumed to be a few times higher than the soft masses. In this energy range the only field that undergoes non-linear transformations is the goldstino field. Using this idea, all goldstino couplings to the MSSM were computed up to and including the order $\mathcal{O}(1/f^2)$. Regarding phenomenology, one interesting consequence is the increase of the lightest Higgs mass from its MSSM tree level value, which becomes close to the LEP2 bound for $\sqrt{f} \sim 2-7$ TeV. This value is consistent with the current constraints from invisible Z and Higgs decays. Finally, it is important to emphasize the main differences of this model to the MSSM: (1) the presence of an additional fermionic light mode, namely the goldstino (or equivalently the longitudinal gravitino) which is SM singlet and (2) the existence of an extra symmetry, namely non-linear supersymmetry.

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