Phenomenology of a leptonic goldstino and invisible Higgs boson decays

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Abstract

Non-linearly realized supersymmetry, combined with the Standard Model field content and $SU(3) \times SU(2) \times U(1)$ gauge invariance, permits local dimension-six operators involving a goldstino, a lepton doublet and a Higgs doublet. These interactions preserve total lepton number if the left-handed goldstino transforms as an antilepton. We discuss the resulting phenomenology, in the simple limit where the new couplings involve only one lepton family, thus conserving also lepton flavour. Both the Z boson and the Higgs boson can decay into a neutrino and a goldstino: the present limits from the invisible Z width and from other observables leave room for the striking possibility of a Higgs boson decaying dominantly, or at least with a sizable branching ratio, via such an invisible mode. We finally comment on the perspectives at hadron and lepton colliders, and on possible extensions of our analysis.

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1.1 General theoretical framework

The first attempt to introduce supersymmetry in a particle physics model dates back to the seminal paper by Volkov and Akulov [1]: simple four-dimensional supersymmetry was non-linearly realized, the goldstino was identified with the (electron) neutrino, and a universal dimension-eight coupling of goldstino bilinears to the matter energy-momentum tensor was introduced, with strength fixed by the goldstino decay constant. Soon after, it was realized [2] that interpreting the goldstino as one of the Standard Model (SM) neutrinos was not allowed by the low-energy theorems of supersymmetry, which prescribe a much softer infrared behaviour of the goldstino amplitudes with respect to the neutrino ones. Many years later, it was also noticed that, already at the lowest order in the derivative expansion, the bilinear goldstino couplings to the SM fields are more general than the universal coupling to the energy-momentum tensor [3, 4, 5], and explicit examples were produced [6] in superstring models with D-branes.

Recently, another unexpected result was found for the single-goldstino couplings [7], again in the context of superstring models with D-branes. In the linear realization of spontaneously broken $N = 1$ supersymmetry, a single goldstino couples universally to the supersymmetry current, in a way prescribed by the supersymmetry algebra [8]. When moving to the non-linear realization [1, 9], for example by integrating out all the heavy superpartners of the goldstino and of the remaining light fields, all these couplings are expected to disappear, leaving only interactions with an even number of goldstinos. However, as recently found in [7], this is not the most general possibility: for a generic field content, there are two types of operators of dimension six, with undetermined coefficients, that contain a single goldstino coupled to a matter fermion and a gauge or scalar field.

Further restrictions are obtained by assuming the SM gauge group and field content, in addition to a gauge-singlet goldstino. At the lowest order in the goldstino decay constant $\kappa$, which sets the scale of supersymmetry breaking, the only dimension-six operator that can couple one goldstino to SM fields is:

$$O = \kappa \sum_{a=e,\mu,\tau} c_a \, O_a + \text{h.c.} , \quad O_a = \epsilon_{ij} \, l^i_a \, (\partial^\mu \tilde{G}) \, (D^\mu \phi)^j . \quad (1.1)$$

In eq. (1.1), $\kappa$ is a real coefficient of dimension (mass)$^{-2}$, whose normalization is determined, for example, by the inhomogeneous term in the goldstino transformation law, $\delta \xi \tilde{G}_a = (1/\kappa) \, \xi_a + \ldots$. The $c_a$ are generically complex dimensionless parameters, where the index $a = e, \mu, \tau$ denotes the three different lepton families. The fields appearing in the operators $O_a$ are the goldstino $\tilde{G} \sim (1,1,0)$, the lepton doublets $l^i_a = (\nu_a, e_a)^T \sim (1,2,-1/2)$ and the SM Higgs doublet $\phi^i = (\varphi^+, \varphi^0)^T \sim (1,2,+1/2)$, where the numbers in brackets denote the transformation properties with respect to $SU(3)_C \times SU(2)_L \times U(1)_Y$, and we work with two-component spinors, in the notation of ref. [10]. For convenience, it is not restrictive to work in a field basis where all kinetic terms are canonically normalized, and
the charged leptons 
\[ e^a \] 
are in a mass eigenstate basis. The symbol 
\[ \epsilon_{ij} \] 
is the \( SU(2)_L \) anti-symmetric tensor, normalized according to \( \epsilon_{12} = 1 \). Finally, \( D_\mu \) is the SM gauge-covariant derivative. In this paper, we restrict ourselves to the couplings of the goldstino with the minimal SM, including the Higgs field but neither right-handed neutrinos nor dimension-five operators inducing Majorana masses for the left-handed neutrinos: we will briefly comment on the consequences of relaxing such simplifying assumption in the final section.

### 1.2 String models with D-branes

We have already mentioned that operators of the form \((1.1)\) were recently found \([7]\) in superstring models with D-branes. Although our main considerations below will have general validity, irrespectively of the microscopic origin of the operator \((1.1)\), we describe here for illustration the main properties of the goldstino in such a D-brane string framework. Readers interested only in the phenomenological description at the effective field theory level can skip this subsection and go directly to subsection 1.3.

These higher-dimensional string constructions involve configurations of intersecting branes, combined eventually with orientifolds (for recent reviews on semirealistic D-brane models, and references to the original literature, see e.g. \([11]\)). In this context, there are (at least) two bulk supersymmetries that come into play. One (half) is spontaneously broken by the very existence of the branes, while the other (other half) is broken by the fact that the branes are at angles, or by the simultaneous presence of orientifolds. The goldstino appearing in the four-dimensional effective theory is associated with the former, not with the latter \([6, 7]\). Since the corresponding supersymmetry is broken by the presence of the branes, it can only be realized non linearly, and brane fields have no superpartners. For the simple case of two stacks of D-branes at angles, the goldstino decay constant \( \kappa \) is given by the total brane tension on their intersection. Moreover, the coefficients \( c_a \) turn out to be real, universal model-independent constants: \( c_a = 0 \) (if the lepton and the Higgs doublet come from different intersections) or \( |c_a| = 2 \). The possibility of having observable effects at the presently accessible energies is then related with the possibility of having the string scale close to the weak scale \([12]\).

More precisely, for two stacks of \( N_1 \) coincident and \( N_2 \) coincident D-branes, intersecting in a \( 3 + 1 \) dimensional volume, we have \([7]\):

\[
\frac{1}{2 \kappa^2} = N_1 T_1 + N_2 T_2; \quad T_i = \frac{M_s^4}{4 \pi^2 g_i^2}, \quad (i = 1, 2),
\]

where \( M_s \equiv (\alpha')^{-1/2} \) is the string scale, \( T_1 \) and \( T_2 \) are the effective tensions at the intersection, and \( g_i \) \((i = 1, 2)\) is the effective four-dimensional gauge coupling on the \( i \)-th brane stack. To get an estimate of the ratio between the string scale and the supersymmetry-breaking scale, \( \sqrt{F} \equiv [1/(2\kappa^2)]^{1/4} \), we assume that the Higgs and lepton doublets come from the intersection of an abelian brane \((N_1 = 1)\) with two coincident branes \((N_2 = 2)\), describing the \( SU(2)_L \) factor of the SM gauge group. To remove the ambiguity on the
g_1, related to the hypercharge embedding, we can use the representative GUT values $g_1 \simeq g_2 \simeq 1/\sqrt{2}$: in such a case $M_s \simeq 1.6 \sqrt{F}$. For $g_1 \gg g_2$ we could reach $M_s \simeq 1.8 \sqrt{F}$. Too small values of $g_1$ would be phenomenologically incompatible with values of $\sqrt{F}$ close to the weak scale.

In our analysis, we work in the limit of global supersymmetry and neglect gravitational effects. When allowed by the above string constructions, this amounts to taking the decompactification limit in a direction transverse to the brane configuration, which suppresses the four-dimensional gravitational interactions. At finite volume, the non-linearly realized supersymmetry is in general broken when including effects coming from different locations of the bulk, thus our brane-localized goldstino is expected to acquire a mass suppressed by a volume factor. In particular, our putative goldstino should mix with the internal components of the higher-dimensional gravitino, but the resulting mass terms are expected to be small enough to be irrelevant for the purposes of the present paper.

### 1.3 Effective theory setup

In summary, we are considering an effective theory where the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge invariance is linearly realized, whilst global $N = 1$ supersymmetry is non-linearly realized. Its content amounts to the fields of the minimal, non-supersymmetric SM, plus a gauge-singlet goldstino $\tilde{G}$. Its Lagrangian is given by

$$\mathcal{L} = \mathcal{L}_{SM} - \frac{i}{2} [\tilde{G} \sigma^\mu \partial_\mu \tilde{G} - (\partial_\mu \tilde{G}) \sigma^\mu \tilde{G}] + \mathcal{O} + \ldots,$$

(1.3)

where $\mathcal{L}_{SM}$ is the renormalizable SM Lagrangian (including a possible cosmological constant term), $\mathcal{O}$ is the one-goldstino operator in (1.1), and the dots indicate additional terms, containing at least two goldstino fields and of higher order in $\kappa$.

It is important to notice that the effective theory defined by (1.1) and (1.3) conserves the total lepton number $L$, as long as we assign to the left-handed goldstino $\tilde{G}$ a total lepton number $L(\tilde{G}) = -1$. In a sense, this model partially implements the original proposal of [1]: the goldstino is not identified with a SM neutrino, but can nevertheless be regarded as a neutral gauge-singlet (anti-)lepton. The analogy between lepton number and a continuous $R$-symmetry, already noticed in [8], becomes now an identification: the conserved lepton number of our effective theory is associated with a diagonal subgroup of $U(1)_{L} \times U(1)_{R}$, where $U(1)_{L}$ corresponds to the SM lepton number, and acts on the leptons but not on the goldstino, whilst $U(1)_{R}$ is the $R$-symmetry, acting on the goldstino but not on the leptons.

The goal of the present paper is to discuss the phenomenological implications of the new operator in eq. (1.1), under the simplifying assumption that only one of the coefficients $c_\alpha$ is non-vanishing:

$$c_\alpha = c \neq 0 \quad \text{for one} \ \alpha, \quad c_a = 0 \quad \text{for} \ a \neq \alpha.$$

(1.4)
This assumption is very close in spirit to a similar one, made in most phenomenological studies of explicit $R$-parity violation in the Minimal Supersymmetric Standard Model (MSSM). The fact that in the SM the Higgs boson couples to fermions proportionally to their masses, and the geometrical setup of D-brane models, suggest that the non-vanishing coupling in (1.4) is more likely associated with the heavier lepton generations. However, in most of the following considerations we will treat all three possible choices of $\hat{a}$ in (1.4) on equal footing.

Under the assumption (1.3), the operator $O$ conserves not only the total lepton number $L$, but also the partial lepton numbers $L_a$, as long as we assign to the goldstino $\tilde{G}$ partial lepton numbers $L_a(\tilde{G}) = -1$ and $L_a(\tilde{G}) = 0$ for $a \neq \hat{a}$. This will allow us to discuss, in the central section of the paper, only processes conserving both total and partial lepton numbers, in the limit of vanishing neutrino masses. We will first review the phenomenological constraints on the new goldstino couplings coming from known physics: under our assumptions, the most stringent ones come from the LEP bounds on the invisible $Z$ width. Other constraints either are weaker or have a more ambiguous interpretation within the effective theory. We will then discuss a striking phenomenological implication of the new couplings: the possibility of having, as non-negligible or even dominant decay mode for the Higgs boson, the invisible channel consisting of a neutrino and a goldstino (or the conjugate channel). In the final section we will present our conclusions, and comment on the experimental perspectives at high-energy colliders and on possible generalizations of the present analysis.

2 Phenomenology

We now discuss the phenomenology of the operator $O$ of eq. (1.1), under the assumption that, in a field basis where the charged leptons are mass eigenstates, only one of the coefficients $c_a$ is non-negligible. We can then omit the generation index $a$ and recall that

$$D_\mu \phi = \partial_\mu \phi - i \left( g \frac{\sigma^I}{2} A^I_\mu + g' \frac{1}{2} B_\mu \right) \phi,$$

(2.1)

where $\sigma^I$ ($I = 1, 2, 3$) are the Pauli matrices, $g$ and $g'$ are the $SU(2)_L$ and $U(1)_Y$ gauge coupling constants, respectively, and $A^I_\mu$ ($I = 1, 2, 3$) and $B_\mu$ the corresponding gauge bosons. We can also move to the unitary gauge, where

$$\phi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix} = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + H \end{pmatrix},$$

(2.2)

$H$ is the canonically normalized SM Higgs boson, and the vacuum expectation value $v \equiv \sqrt{2} \langle \varphi^0 \rangle$ is related with the Fermi coupling and the gauge boson masses by

$$\frac{G_F}{\sqrt{2}} = \frac{1}{2} \frac{v^2}{v} = \frac{g^2}{8 m_W^2} = \frac{g^2 + g'^2}{8 m_Z^2}.$$

(2.3)
Recalling that

\[ W_\mu^\pm = \frac{A_\mu^1 \mp i A_\mu^2}{\sqrt{2}}, \quad Z_\mu = \frac{g' B_\mu - g A_\mu^3}{\sqrt{g'^2 + g^2}}, \]

we can then write

\[ \mathcal{O} = \left[ \frac{\kappa c}{\sqrt{2}} \nu (\partial^\mu \tilde{G})(\partial_\mu H) + h.c. \right] \]

\[ - \left[ i \frac{\kappa c}{\sqrt{2}} m_Z \nu (\partial^\mu \tilde{G}) Z_\mu + h.c. \right] \]

\[ + \left[ i \frac{\kappa c}{\sqrt{2}} m_W e (\partial^\mu \tilde{G}) W_\mu^+ + h.c. \right] \]

\[ - \left[ i \frac{\kappa c \sqrt{g'^2 + g^2}}{2 \sqrt{2}} \nu (\partial^\mu \tilde{G}) Z_\mu H + h.c. \right] \]

\[ + \left[ i \frac{\kappa c \sqrt{g'^2 + g^2}}{2} e (\partial^\mu \tilde{G}) W_\mu^+ H + h.c. \right]. \]

Before examining the phenomenological consequences of the interactions (2.5)-(2.9), it is useful to recall the general experimental limits on the supersymmetry breaking scale. Under the assumption that all superpartners of the goldstino and of the SM particles are sufficiently heavy to have escaped detection, which corresponds to the non-linear realization, the most stringent bounds come from the processes \( e^+ e^- \rightarrow \tilde{G} \tilde{G} \gamma \) at LEP2 [13, 14, 15], and give \( \sqrt{F} > 238 \text{ GeV} \) (95\% c.l.). The study of the processes \( pp \rightarrow \tilde{G} \tilde{G} \gamma, \tilde{G} \tilde{G} \text{ jet} \) at the Tevatron collider [13, 16] has led so far to a published bound [17] \( \sqrt{F} > 221 \text{ GeV} \) (95\% c.l.), lower than the LEP2 bound. Finally, indirect bounds from the muon anomalous magnetic moment [18], from flavour physics [19] and from astrophysics and cosmology [20] are less stringent than the present collider bounds \(^1\).

Inspection of eqs. (2.5)-(2.9) shows that the new local interactions in (2.6) and (2.7) can describe non-standard contributions to measured processes, where an antineutrino is replaced by a goldstino (a neutrino by an antigoldstino) in the final state. These contributions occur already at order \( \kappa \) in the amplitude, thus \( \kappa^2 \) in the cross-section or in the decay rate. In addition, the new local interactions in (2.5), (2.8) and (2.9) can describe, at the same order in \( \kappa \), decays of the SM Higgs boson into final states containing leptons and goldstinos.

2.1 Constraints from known physics

Leaving the study of Higgs decays for the next subsection, we now discuss the phenomenological constraints coming from known physics. Whenever possible, it will be convenient to express these constraints in terms of the auxiliary mass parameter

\[ M \equiv \frac{1}{\sqrt{\left| \kappa c \right|}} = 2^{1/4} \sqrt{\frac{F}{|c|}}. \]

\(^1\)At least this is true in the absence of (light) superpartners of the SM particles and of the goldstino, as is the case in the present theoretical framework.
The interaction in (2.6) describes the decays \( Z \to \nu \tilde{G} \) and \( Z \to \bar{\nu} \tilde{G} \). Their signal would be an additional contribution to the \( Z \) partial width into invisible products, besides the one associated with the SM neutrinos:

\[
\Delta \Gamma_{\text{inv}}(Z) = \Gamma(Z \to \nu \tilde{G}) + \Gamma(Z \to \bar{\nu} \tilde{G}) = \frac{m_Z^5}{192 \pi M^4}.
\] (2.11)

The present upper bound on exotic contributions to the invisible \( Z \) width is [21]:

\[
\Delta \Gamma_{\text{inv}}(Z) < 2.0 \text{ MeV}, \quad (95\% \text{ c.l.}).
\] (2.12)

Plugging this into eq. (2.11), we obtain:

\[
M > 270 \text{ GeV}.
\] (2.13)

For specific values of \( c \), we can extract a bound on the supersymmetry breaking scale and compare it with the collider bounds. For example, taking \( |c| = 2 \) as suggested by the results of [7], we get \( \sqrt{F} > 320 \) GeV, slightly stronger than the present collider bounds \(^2\).

The interaction in (2.7) describes the decays \( W^+ \to \ell^+ \tilde{G} \) and \( W^- \to \ell^- \tilde{G} \), where \( \ell = e, \mu, \tau \) according to the choice of \( \tilde{a} \) in (1.4), with partial widths (neglecting the charged lepton mass):

\[
\Gamma(W^+ \to \ell^+ \tilde{G}) = \Gamma(W^- \to \ell^- \tilde{G}) = \frac{m_W^5}{192 \pi M^4}.
\] (2.14)

The constraints from \( W \) decays are similar to those from \( Z \) decays, but weaker. Under the assumption (1.4), exotic \( W \) decays of the type (2.14) could produce violations of lepton universality, via their additional contributions to one of the leptonic widths. Taking \( M \) at its lower bound (2.13), we would find \( \Delta \Gamma_\ell(W) \simeq 1 \text{ MeV} \), corresponding to \( \Delta BR_\ell(W) \simeq 4.5 \times 10^{-4} \), still below the present precision of the LEP2 and Tevatron experiments [21].

The non-renormalizable nature of our new \( d = 6 \) operator suggests that its effects have a strong, power-like suppression when the typical energy scale of the processes under consideration is much smaller than \( M \). However, we should check that precisely measured low-energy processes, such as \( \mu \) and \( \tau \) decays, cannot give constraints stronger than (2.13). To give an idea of the sensitivity in \( \mu \) and \( \tau \) decays, we recall the present experimental precision [22] on the respective rates (or, equivalently, on the lifetimes):

\[
\left| \frac{\Delta \Gamma_\mu}{\Gamma_\mu} \right| \sim 2 \times 10^{-5}, \quad \left| \frac{\Delta \Gamma_\tau}{\Gamma_\tau} \right| \sim \frac{1}{300}.
\] (2.15)

Exotic contributions to \( \mu \) and \( \tau \) decays can be originated by the new charged current interactions of eq. (2.7). The new Feynman diagrams involve a goldstino at the place of an antineutrino (or an antigoldstino at the place of a neutrino) on an external line, thus

\(^2\)Collider bounds on \( \sqrt{F} \) in the non-linear realization have a very mild dependence on the free parameters of the \( \mathcal{O}(\kappa^2) \) four-fermion couplings between two SM fermions and two goldstinos [3, 4, 13, 7]. This dependence is negligible for the present analysis.
The contributions to the decay rates then at most as
\[ \left| \frac{\Delta \Gamma_\ell}{\Gamma_\ell} \right| \lesssim \frac{m_\ell^2}{g^2 M^4}, \quad (\ell = \mu, \tau). \] (2.16)

Taking \( M \) at its lower bound (2.13), we would find \( \Delta \Gamma_\mu/\Gamma_\mu \lesssim 3 \times 10^{-8} \) and \( \Delta \Gamma_\tau/\Gamma_\tau \lesssim 10^{-5} \), orders of magnitude below the present experimental sensitivity.

Another process sensitive to the new couplings in (2.6) and (2.7) is \( e^+ e^- \to \gamma + \text{nothing} \) at LEP2. With two goldstinos in the final state, this process has an amplitude \( \mathcal{O}(\kappa^2) \), and is used \([13],[14],[15]\) to establish the model-independent lower bound on the supersymmetry-breaking scale. With one neutrino and one goldstino (one antineutrino and one antigoldstino) in the final state, the amplitude for this process occurs at \( \mathcal{O}(\kappa) \), and may receive two contributions: the one from (2.6), corresponding to Z exchange in the s-channel, is always present under our assumptions; the one from (2.7), corresponding to W exchange in the t-channel, is present if and only if the new coupling involves the first generation (\( \hat{a} = e \)). To be conservative, we can estimate the bound on \( M \), defined in (2.10), by assuming that (2.6) does contribute, but (2.7) does not. According to the formalism of \([23]\), we can approximate the differential cross-section for the processes \( e^+ e^- \to \gamma \nu \tilde{G} \) and \( e^+ e^- \to \gamma \overline{\nu G} \), dominated by the soft and collinear part of the photon spectrum \( (x_\gamma \ll 1 \text{ and/or } \sin^2 \theta_\gamma \ll 1) \), by:
\[ \frac{d\sigma}{dx_\gamma d\cos \theta_\gamma} \approx \sigma_0(\tilde{s}) \frac{\alpha}{\pi} \frac{1 + (1 - x_\gamma)^2}{x_\gamma \sin^2 \theta_\gamma}, \] (2.17)

where \( x_\gamma \) is the fraction of the beam energy carried by the photon, \( \theta_\gamma \) is the scattering angle of the photon with respect to the direction of the incoming electron in the centre-of-mass frame, \( \tilde{s} = (1 - x_\gamma) s \), \( \sqrt{s} \) is the energy in the centre-of-mass frame, and
\[ \sigma_0 \equiv \sigma(e^+ e^- \to \nu \tilde{G}) + \sigma(e^+ e^- \to \overline{\nu G}). \] (2.18)

In the conservative case where only the coupling (2.6) is present, we can write
\[ \sigma_0(\tilde{s}) = \frac{12 \pi s^2}{m_Z^4} \frac{\Gamma_e \Delta \Gamma_{\text{inv}}(Z)}{(\tilde{s} - m_Z^2)^2 + \Gamma_Z^2 m_Z^2}, \] (2.19)

where \( \Gamma_e \equiv \Gamma(Z \to e^+ e^-) \simeq 84 \text{ MeV} \). We can approximate the present LEP2 sensitivity \([15]\) by requiring that, for \( \sqrt{s} = 207 \text{ GeV} \), \( x_\gamma > 0.05 \) and \( |\cos \theta_\gamma| < 0.95 \), it is \( \sigma(e^+ e^- \to \gamma \nu \tilde{G}) + \sigma(e^+ e^- \to \gamma \overline{\nu G}) < 0.1 \text{ pb} \). Plugging \( M > 270 \text{ GeV} \) from eq. (2.13) into the expression (2.11) for \( \Delta \Gamma_{\text{inv}}(Z) \), and making use of eqs. (2.17)–(2.19), we find
\[ \sigma(e^+ e^- \to \gamma \nu \tilde{G}) + \sigma(e^+ e^- \to \gamma \overline{\nu G}) < 0.011 \text{ pb}, \] (2.20)
We can try to extract additional constraints on $M$ by considering high-energy SM processes where virtual goldstinos are exchanged on the internal lines. To extract reliable constraints, however, we should be sure that there are no additional local operators contributing to the total amplitude for the same process at $O(\kappa^2)$. In the absence of a microscopic theory, such bounds should be regarded only as order of magnitude estimates.

An example is $e^+e^- \rightarrow W^+W^-$ at LEP2. At the classical level, and in the limit where the electron mass is neglected, the SM amplitude receives contributions from three Feynman diagrams: one with the $t$-channel exchange of the electron neutrino, and two with the $s$-channel exchange of the photon and the $Z$. If the new interaction (2.7) does not involve the first lepton family ($\hat{a} = \mu, \tau$), there is of course no exotic contribution, thus no constraint. Otherwise ($\hat{a} = e$), the new interaction generates an $O(\kappa^2)$ amplitude, corresponding to a Feynman diagram with goldstino exchange in the $t$-channel, that interferes with the SM diagrams. To estimate the order of magnitude bound from $W$-pair production at LEP2, we decompose the total cross-section $\sigma_{WW}$ for the (on-shell) process as:

$$\sigma_{WW} = \sigma_{SM} + \sigma_{\kappa^2} + \sigma_{\kappa^4}. \quad (2.21)$$

Here $\sigma_{WW}$ is the SM cross-section, $\sigma_{\kappa^2}$ is the $O(\kappa^2)$ interference contribution, and $\sigma_{\kappa^4}$ is the $O(\kappa^4)$ contribution from the square of the exotic amplitude, which should be strongly suppressed with respect to the previous one. An approximate formula for $|\sigma_{\kappa^2}|$, valid only at the first non-trivial order in the expansion parameter $(2 m_W/\sqrt{s})$, but sufficient for an order of magnitude estimate, is:

$$|\sigma_{\kappa^2}| \simeq \frac{\alpha}{768 \sin^2 \theta_W \cos^2 \theta_W} \frac{s}{M^4}. \quad (2.22)$$

Taking $M$ at its lower bound (2.13), and $\sqrt{s} = 206.6$ GeV, where the LEP2 average [21] is $\sigma^{(LEP)}_{WW} = 17.28 \pm 0.27$ pb, from (2.22) we find $\sigma_{\kappa^2} \simeq 0.18$ pb, well below one standard deviation. We have checked that the complete theoretical expression for $|\sigma_{\kappa^2} + \sigma_{\kappa^4}|$ gives a similar constraint, and that data at lower values of $\sqrt{s}$ are less restrictive.

It is conceivable that the new physics originating the operator (1.1) gives rise to additional operators in the effective theory at the weak scale, for example local operators involving only four SM fermions. In particular, four-fermion operators involving at least two left-handed leptons of type $\hat{a}$ can be generated by quantum corrections (e.g. box diagrams) in the presence of the interactions (2.6) and (2.7). In the generic framework of models with a low supersymmetry-breaking scale, the question of four-SM-fermion operators was addressed in [24]. It was found that a supersymmetry-breaking scale as low as the direct experimental bound can naturally coexist with the level of suppression of the dangerous four-fermion operators required by the LEP and Tevatron bounds [22].
fermions are localized at the same brane intersection, giving rise to dimension six effective operators. However, deciding whether those bounds could be applicable in the present context would require an explicit realistic string construction. To be conservative, we should consider only four-fermion operators involving two left-handed leptons of type $\tilde{a}$. The strength of the corresponding dimension-six effective operator depends on the intersection angle, and becomes maximal for orthogonal branes. In the case of coincident branes, the dimension-six effective operator vanishes, and the leading contribution to four-fermion contact interactions comes at dimension eight. The limit on the string scale can then be extracted from Bhabha scattering, leading to $M_s > 490$ GeV \[26\]. This is roughly comparable with the direct bound \[24\], taking into account that, as argued in section 1.2, we expect $M_s \sim \left(\frac{1}{6} \sqrt{\frac{1}{6}} \right) M$.

To conclude, we notice that the new interactions in (2.6) and (2.7) do not contribute to Veltman’s $\rho$ parameter at the one-loop level. There are, however, quadratically divergent one-loop contribution to $[\Pi_{VV}^{\text{new}}(m_V^2) - \Pi_{VV}^{\text{new}}(0)]/m_V^2$, $(V = Z, W)$, originated by gauge boson self-energy diagrams with a goldstino and a lepton of type $\tilde{a}$ on the internal lines. Choosing an ultraviolet cutoff $\Lambda \sim M$ suggests that the natural value of these contributions is $O(m_V^2)/(64 \pi^2 M^2)$. Even for $M$ at its lower bound \[24\] and $V = Z$, we find an $O(1.8 \times 10^{-4})$ contribution, corresponding to $\delta S = O(1.6 \times 10^{-2})$ or $\delta \tilde{\epsilon}_3 = O(1.4 \times 10^{-4})$, in agreement with the present bounds \[22\].

### 2.2 Higgs boson decays

The constraints discussed in the previous section leave room for very interesting signatures of the new goldstino interactions in Higgs decays.

The interaction in (2.5) describes the invisible decays $H \to \nu \tilde{G}$ and $H \to \bar{\nu} \tilde{G}$, at a rate:

$$\Gamma_{\text{inv}}(H) = \Gamma(H \to \nu \tilde{G}) + \Gamma(H \to \bar{\nu} \tilde{G}) = \frac{m_H^5}{64 \pi M^4}.$$ \hspace{1cm} (2.23)

Notice the similarity between eq. (2.23) and the universal formula \[27\] expressing the decay rate for a massive superparticle into its massless superpartner and a goldstino. Indeed, using the spinor algebra, integration by parts and the equations of motion, the three-point coupling (2.5) can be rewritten as $[1/(2 \sqrt{2})] \kappa c m_H^3 \tilde{G} \nu H + \text{h.c.}$, as if the spin-0 Higgs doublet and the spin-1/2 lepton doublet were members of a single chiral supermultiplet in the linear realization. It is a curious coincidence that, in the D-brane models of \[7\], Higgs and lepton doublets sit in the same multiplet, but of a supersymmetry different from the one associated with the goldstino $\tilde{G}$.

The interactions in (2.8) and (2.9) could describe the decays $H \to Z \nu \tilde{G}$, $Z \bar{\nu} \tilde{G}$, $W^- l^+ \tilde{G}$, $W^+ l^- \tilde{G}$, which however are strongly suppressed by phase space with respect to the two-body decays in (2.23), and will be neglected here.

The present LEP2 bound on a Higgs boson with SM production cross-section but dominant invisible decays \[23\] is very close to the bound on the SM Higgs \[20\], $m_H >
In the following we will then consider values of the Higgs boson mass above the SM bound, even if the precise bound on $m_H$ for a Higgs boson with both SM and invisible decay modes could be slightly weaker. We also remind the reader that, in models with a low supersymmetry breaking scale, and in contrast with the MSSM, a SM-like Higgs boson is not bound to be light (for a recent discussion, see e.g. [30]).

Taking into account the constraint (2.13) and the experimental lower bound on $m_H$, we can now study the possible phenomenological relevance of the new invisible Higgs decay modes of eq. (2.23). We computed the Higgs branching ratios, as functions of $m_H$ and $M$, by including the new invisible channels in the program HDECAY [31]. The results are displayed in figs. 1–4. Fig. 1 shows the most important Higgs branching ratios as functions of $m_H$, for two representative values of $M$: 300 and 600 GeV. Fig. 2 shows the same branching ratios, but as functions of $M$, for four representative values of $m_H$: 115, 140, 200 and 400 GeV. Fig. 3a shows contours of $BR_{inv}(H) = \Gamma_{inv}(H)/\Gamma_{tot}(H)$ in the $(M, m_H)$ plane. Fig. 3b shows the ratio $\Gamma_H/m_H$, as a function of $m_H$, in the large mass region, in the SM case and for two representative values of $M$: 300 and 500 GeV (notice that the 500 GeV line is already very close to the SM one). Figs. 4a and 4b show the total Higgs width $\Gamma_H$ and the branching ratio $BR(H \rightarrow \gamma \gamma)$ as functions of $m_H$, in the SM case and for the same representative values of $M$ as in Fig. 3b.

We can see from the various figures that, for values of $M$ close to the lower bound of eq. (2.13), the invisible Higgs decay modes can be the dominant ones. This behaviour persists for moderate values of $M$, especially for $m_H \sim 130$ GeV, where we can still have an $\mathcal{O}(10\%)$ invisible branching ratio for $M \sim 750$ GeV. When the invisible decay modes dominate, we could see an effect both in a dedicated search for invisible Higgs decays, and, indirectly, by measuring a deficit in the Higgs branching ratios into SM channels, or a total Higgs width larger than the SM prediction.

We should recall at this point that in models where $\sqrt{F}$ is close to the weak scale but supersymmetry is linearly realized, the goldstino can pick up small components along the neutralinos (gauginos and higgsinos). These can in turn induce $Z\bar{G}G$ couplings to goldstino pairs, and lead to the decays $Z \rightarrow \bar{G}G$ or $H \rightarrow G\bar{G}$ at a non-negligible rate. Here, however, we work in the non-linear realization, and we assumed that the goldstino is a pure SM singlet, thus we consistently neglected such a possibility. The non-linear realization also permits, in principle, a dimension-seven gauge-invariant operator [7] proportional to $[\phi^\dagger \phi \left( \partial^\mu \bar{G} \right) \sigma^{\mu\nu} \left( \partial_\nu \bar{G} \right) + \text{h.c.}]$, which would violate total lepton number and lead to $H \rightarrow G\bar{G}$ decays. Here, according to the results of explicit calculations in D-brane models [7], we assumed that the above dimension-seven operator is absent.

The possibility of invisible Higgs decays was also considered in other theoretical frameworks. One is the existence of a fourth generation, with negligible mixing with the first three, and a specific mass spectrum: this may allow for $H \rightarrow N \bar{N}$ decays [32], where $N$ is a 50-80 GeV neutrino, but is only marginally allowed by electroweak precision data [33] [22]. Another one is the MSSM with non-universal gaugino masses, which may allow for a large $H \rightarrow \tilde{\chi}^0 \tilde{\chi}^0$ branching ratio, where $\tilde{\chi}^0$ is the lightest neutralino. A third one
case, however, the invisible channel can never be the dominant decay mode. A final one is the possibility of Higgs decays into Majorons in non-minimal supersymmetric models with spontaneously broken $R$-parity [35]: in this case, at the price of rather complicated constructions, the invisible mode could be dominant.

3 Conclusions and outlook

In this paper we examined the phenomenological implications of the dimension-six operator (1.1), allowed by non-linear supersymmetry coupled to the minimal non-supersymmetric SM. We worked in the limit where the new couplings involve only one lepton family and neutrino masses are vanishing, so that total and partial lepton numbers are conserved in perturbation theory. We showed that the most stringent phenomenological constraints on the mass scale $M$ of the new interactions come from invisible $Z$ decays, and give $M > 270$ GeV. We also examined other constraints from measured processes, and found that either they are less stringent or they have a more ambiguous physical interpretation. We finally discussed the most striking phenomenological signature originated by the new interactions: the possibility for the Higgs boson to decay into the invisible channel neutrino + goldstino, or the conjugate one, with non-negligible or even dominant branching ratio.

The phenomenological scenario discussed in this paper can be tested at the Tevatron, at the LHC, and at a possible international $e^+e^-$ linear collider (ILC) with $\sqrt{s} \gtrsim 500$ GeV [36]. These colliders can improve the existing bounds on anomalous interactions among SM fermions and on anomalous contributions to $W$-pair production, albeit with the interpretation ambiguities already mentioned in subsection 2.1. More importantly, Run II of the Tevatron and the LHC can probe higher values of the supersymmetry breaking scale $\sqrt{F}$, by looking for single photon (single jet) plus missing transverse energy, as originated by the production of goldstino-antigoldstino pairs in association with a photon (jet). However, extracting reliable experimental bounds on $\sqrt{F}$ at hadron colliders seems to be more difficult than expected by preliminary theoretical studies [16]. The latter estimated the Tevatron Run I sensitivity at $\sqrt{F} \sim 310$ GeV, whereas the actual experimental bound [17] is only $\sqrt{F} > 221$ GeV. Preliminary theoretical estimates [16] for Run II of the Tevatron and for the LHC were foreseeing a sensitivity up to $\sqrt{F} \sim 410$ GeV and $\sqrt{F} \sim 1.6$ TeV, respectively, but taking into account the actual detector environment and the very similar shapes of signal and background may lead to a significant downgrade of these estimates. The two main processes of interest for the ILC are: goldstino pair production in association with a photon, which should be sensitive to values of $\sqrt{F}$ of the order of $\sqrt{s}$; the production of neutrino-goldstino pairs in association with a photon, for which one could repeat the LEP2 analysis of eqs. (2.17)–(2.19), eventually dropping the simplifying approximations.

As for invisible Higgs decays, looking for direct or indirect signals at the LHC looks
quite challenging but not impossible. For example, having dominant decays into invisible channels does not seem to preclude the possibility of Higgs detection, even if it makes it more difficult: existing proposals try to exploit associated $ZH$ production \cite{37}, associated $t\bar{t}H$ production \cite{38}, production via $WW$ fusion with tagged forward jets \cite{39}, and central exclusive diffractive production \cite{40}. Indirect signals would correspond to measuring deviations from the SM predictions for the Higgs branching ratios into SM channels and for the total Higgs width. Direct and indirect detection of an invisible Higgs decay channel should be much simpler at an ILC with sufficient energy and luminosity.

There are many generalizations of the present analysis that would be interesting to consider. A more systematic study of the processes considered in section 2.1 could be performed. The possibility of generating the operator $\mathcal{O}$ starting from an underlying linear realization could be also explored: one may be led to a two-Higgs model with some operators breaking lepton number and $R$-symmetry to some diagonal subgroup; alternatively, taking into account the hints coming from brane-world models, one may consider a model with fully broken extended supersymmetry and additional massive states besides those of the MSSM.

Another important generalization, in analogy with the studies of $R$-parity breaking in the MSSM, would be the discussion of the flavour-dependent phenomenology in the presence of more than one non-negligible $c_a$ ($a = e, \mu, \tau$). This could be also related with attempts at addressing the flavour problem within semirealistic brane models \cite{11}. In such a case, the operator $\mathcal{O}$ still conserves the total lepton number but violates lepton flavour. However, lepton flavour violation (LFV) appears in the very special form dictated by the operator (1.1), which would deserve a dedicated analysis of the processes with LFV. The discussion of neutrino masses and neutrino oscillations would presumably be non-trivial. Introducing Majorana masses for the left-handed neutrinos via dimension-five operators, we would have no new dimension-six operators besides (1.1), but we would need to consider both the latter and neutrino masses as sources of LFV. Introducing right-handed neutrinos into the model, to generate Dirac neutrino masses via Yukawa couplings, would force us to consider an additional dimension-six operator besides (1.1), of the form

$$\tilde{\mathcal{O}} = i \kappa \sum_{a=e,\mu,\tau} \tilde{c}_a \tilde{\mathcal{O}}_a + h.c., \quad \tilde{\mathcal{O}}_a = \nu^c_a \sigma^\mu (\bar{\nu}^c \tilde{G}) B_\mu \nu,$$

where $B_{\mu \nu}$ is the $U(1)_Y$ field strength and string considerations suggest this time $\tilde{c}_a = 0$ or $|\tilde{c}_a| = \sqrt{2}$. All this, however, goes beyond the aim of the present paper and is left for future work.
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References

[1] D. V. Volkov and V. P. Akulov, JETP Lett. 16 (1972) 438 and Phys. Lett. B 46 (1973) 109.

[2] W.A. Bardeen, unpublished; B. de Wit and D. Z. Freedman, Phys. Rev. D 12 (1975) 2286 and Phys. Rev. Lett. 35 (1975) 827.

[3] A. Brignole, F. Feruglio and F. Zwirner, JHEP 9711 (1997) 001 [arXiv:hep-th/9709111].

[4] T. E. Clark, T. Lee, S. T. Love and G. Wu, Phys. Rev. D 57 (1998) 5912 [arXiv:hep-ph/9712353].

[5] M. A. Luty and E. Ponton, Phys. Rev. D 57 (1998) 4167 [arXiv:hep-ph/9706268v3].

[6] I. Antoniadis, K. Benakli and A. Laugier, Nucl. Phys. B 631 (2002) 3 [arXiv:hep-th/0111209].

[7] I. Antoniadis and M. Tuckmantel, [arXiv:hep-th/0406010]

[8] P. Fayet, Phys. Lett. B 70 (1977) 461.

[9] E. A. Ivanov and A. A. Kapustnikov, J. Phys. A 11 (1978) 2375; T. Uematsu and C. K. Zachos, Nucl. Phys. B 201 (1982) 250; S. Ferrara, L. Maiani and P. C. West, Z. Phys. C 19 (1983) 267; S. Samuel and J. Wess, Nucl. Phys. B 221 (1983) 153; J. Wess, Nonlinear realization of the N=1 supersymmetry, in Quantum Theory of Particles and Fields, B. Jancewicz and J. Lukierski eds., World Scientific, 1983, p. 223, Karlsruhe preprint 83-0101; T. E. Clark and S. T. Love, Phys. Rev. D 54 (1996) 5723 [arXiv:hep-ph/9608243].

[10] J. Wess and J. Bagger, Supersymmetry and supergravity, 2nd edition, Princeton University Press, Princeton, NJ, 1992.

[11] A. M. Uranga, Fortsch. Phys. 51, 879 (2003); D. Cremades, L. E. Ibanez and F. Marchesano, [arXiv:hep-ph/0212048] E. Kiritsis, Fortsch. Phys. 52, 200 (2004) [arXiv:hep-th/0310001]; D. Lüst, Class. Quant. Grav. 21, S1399 (2004) [arXiv:hep-th/0401156].

[12] J. D. Lykken, Phys. Rev. D 54 (1996) 3693 [arXiv:hep-th/9603133]; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, Phys. Lett. B 436 (1998) 257 [arXiv:hep-ph/9804398]; G. Shiu and S. H. H. Tye, Phys. Rev. D 58 (1998) 106007 [arXiv:hep-th/9805157].

[13] O. Nachtmann, A. Reiter and M. Wirbel, Z. Phys. C 27 (1985) 577.
[15] A. Heister et al. [ALEPH Collaboration], Eur. Phys. J. C 28 (2003) 1; P. Achard et al. [L3 Collaboration], Phys. Lett. B 587 (2004) 16 arXiv:hep-ex/0402002; J. Abdallah et al. [DELPHI Collaboration], arXiv:hep-ex/0406019.

[16] A. Brignole, F. Feruglio, M. L. Mangano and F. Zwirner, Nucl. Phys. B 526 (1998) 136 [Erratum-ibid. B 582 (2000) 759] arXiv:hep-ph/9801329.

[17] T. Affolder et al. [CDF Collaboration], Phys. Rev. Lett. 85 (2000) 1378 arXiv:hep-ex/0003026; D. Acosta et al. [CDF Collaboration], Phys. Rev. Lett. 89 (2002) 281801 arXiv:hep-ex/0205057.

[18] S. Ferrara and E. Remiddi, Phys. Lett. B 53 (1974) 347; R. Barbieri and L. Maiani, Phys. Lett. B 117 (1982) 203; A. Georges and P. Le Doussal, Nucl. Phys. B 255 (1985) 532; F. del Aguila, Phys. Lett. B 160 (1985) 87; A. Mendez and F. X. Orteu, Nucl. Phys. B 256 (1985) 181 and Phys. Lett. B 163 (1985) 167; A. Brignole, E. Perazzi and F. Zwirner, JHEP 9909 (1999) 002 arXiv:hep-ph/9904367.

[19] A. Brignole and A. Rossi, Nucl. Phys. B 587 (2000) 3 arXiv:hep-ph/0006030.

[20] P. Fayet, Phys. Lett. B 69 (1977) 489; M. Fukugita and N. Sakai, Phys. Lett. B 149 (1982) 23; A. Bouquet and C. E. Vayonakis, Phys. Lett. B 116 (1982) 219; M. Nowakowski and S. D. Rindani, Phys. Lett. B 348 (1995) 115 arXiv:hep-ph/9410262; T. Moroi, H. Murayama and M. Yamaguchi, Phys. Lett. B 303 (1993) 289; T. Gherghetta, Nucl. Phys. B 485 (1997) 25 arXiv:hep-ph/9607448; J. A. Grifols, R. Mohapatra and A. Riotto, Phys. Lett. B 401 (1997) 283 arXiv:hep-ph/9610458; A. Brignole, F. Feruglio and F. Zwirner, Nucl. Phys. B 501 (1997) 332 arXiv:hep-ph/9703286.

[21] The LEP Collaborations ALEPH, DELPHI, L3, OPAL, the LEP Electroweak Working Group, the SLD Electroweak and Heavy Flavour Groups, arXiv:hep-ex/0312023.

[22] S. Eidelman et al. [Particle Data Group Collaboration], Phys. Lett. B 592 (2004) 1.

[23] E. A. Kuraev and V. S. Fadin, Sov. J. Nucl. Phys. 41 (1985) 466 [Yad. Fiz. 41 (1985) 733]; G. Altarelli and G. Martinelli, Physics at LEP, CERN Yellow Report 86-02 (J. Ellis and R. Peccei eds.), Vol.1, p.47; O. Nicosini and L. Trentadue, Phys. Lett. B 196 (1987) 551, Z. Phys. C 39 (1988) 479 and Nucl. Phys. B 318 (1989) 1.

[24] A. Brignole, F. Feruglio and F. Zwirner, Phys. Lett. B 438 (1998) 89 arXiv:hep-ph/9805282.
[26] D. Bourilkov, Phys. Rev. D 62, 076005 (2000) [arXiv:hep-ph/0002172]; M. Acciarri et al. [L3 Collaboration], Phys. Lett. B 489, 81 (2000) [arXiv:hep-ex/0005028].

[27] N. Cabibbo, G. R. Farrar and L. Maiani, Phys. Lett. B 105, 155 (1981).

[28] ALEPH, DELPHI, L3 and OPAL Collaborations, The LEP Working Group for Higgs Boson Searches, arXiv:hep-ex/0107032 J. Abdallah et al. [DELPHI Collaboration], Eur. Phys. J. C 32 (2004) 475 [arXiv:hep-ex/0401022]; A. Sopczak, arXiv:hep-ph/0408047.

[29] ALEPH, DELPHI, L3 and OPAL Collaborations, The LEP Working Group for Higgs Boson Searches, Phys. Lett. B 565 (2003) 61 [arXiv:hep-ex/0306033].

[30] A. Brignole, J. A. Casas, J. R. Espinosa and I. Navarro, Nucl. Phys. B 666 (2003) 105 [arXiv:hep-ph/0301121].

[31] A. Djouadi, J. Kalinowski and M. Spira, Comput. Phys. Commun. 108 (1998) 56 [arXiv:hep-ph/9704448].

[32] K. Belotsky, D. Fargion, M. Khlopov, R. Konoplich and K. Shibaev, Phys. Rev. D 68 (2003) 054027 [arXiv:hep-ph/0210153].

[33] V. A. Novikov, L. B. Okun, A. N. Rozanov and M. I. Vysotsky, JETP Lett. 76 (2002) 127 [Pisma Zh. Eksp. Teor. Fiz. 76 (2002) 158] [arXiv:hep-ph/0203132].

[34] G. F. Giudice, R. Rattazzi and J. D. Wells, Nucl. Phys. B 595 (2001) 250 [arXiv:hep-ph/0002178]; I. Antoniadis and R. Sturani, Nucl. Phys. B 631 (2002) 66 [arXiv:hep-th/0201166]; J. L. Hewett and T. G. Rizzo, JHEP 0308 (2003) 028 [arXiv:hep-ph/0202155]; M. Battaglia, D. Dominici, J. F. Gunion and J. D. Wells, arXiv:hep-ph/0402062.

[35] F. de Campos, M. A. Garcia-Jareno, A. S. Joshipura, J. Rosiek and J. W. F. Valle, Nucl. Phys. B 451 (1995) 3 [arXiv:hep-ph/9502237]; M. Hirsch, J. C. Romao, J. W. F. Valle and A. Villanova del Moral, arXiv:hep-ph/0407269.

[36] http://www.interactions.org/linearcollider/~.

[37] R. M. Godbole, M. Guchait, K. Mazumdar, S. Moretti and D. P. Roy, Phys. Lett. B 571 (2003) 184 [arXiv:hep-ph/0304137].
[38] B. P. Kersevan, M. Malawski and E. Richter-Was, Eur. Phys. J. C 29 (2003) 541; M. Malawski, arXiv:hep-ph/0407160.

[39] O. J. P. Eboli and D. Zeppenfeld, Phys. Lett. B 495 (2000) 147 arXiv:hep-ph/0009158; B. Di Girolamo, A. Nikitenko, L. Neukermans, K. Mazumdar and D. Zeppenfeld, in ‘Physics at TeV Colliders’ (Les Houches, France, 21 May - 1 Jun 2001), p. 42; S. Abdullin et al., CMS-NOTE-2003-033.

[40] K. Belotsky, V. A. Khoze, A. D. Martin and M. G. Ryskin, Eur. Phys. J. C 36 (2004) 503 arXiv:hep-ph/0406037.
Figure 1: The most important Higgs branching ratios, as functions of $m_H$, for two representative values of $M$: 300 and 600 GeV.
Figure 2: The most important Higgs branching ratios, as functions of $M$, for four representative values of $m_H$: 115, 140, 200 and 400 GeV.
Figure 3: (a) Contours of the invisible Higgs branching ratio $BR_{\text{inv}}(H)$ in the $(M, m_H)$ plane. (b) The ratio $\Gamma_H/m_H$, as a function of $m_H$ in the large mass region, in the SM and for two representative values of $M$: 300 and 500 GeV.

Figure 4: The total Higgs width $\Gamma_H$ and the branching ratio $BR(H \rightarrow \gamma \gamma)$, as functions of $m_H$ in the intermediate mass region, in the SM and for two representative values of $M$: 300 and 500 GeV.