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Non-standard higgs decays in U(1) extensions of the MSSM

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Abstract: In U(1) extensions of the Minimal Supersymmetric extension of the Standard Model there is a simple mechanism that leads to a heavy Z' boson with a mass which is substantially larger than the supersymmetry breaking scale. This mechanism may also result in a pseudoscalar state that is light enough for decays of the 125 GeV Standard Model-like Higgs boson into a pair of such pseudoscalars to be kinematically allowed. We study these decays within E_6 inspired supersymmetric models with an exact custodial symmetry that forbids tree-level flavor-changing transitions and the most dangerous baryon and lepton number violating operators. We argue that the branching ratio of the lightest Higgs boson decays into a pair of the light pseudoscalar states may not be negligibly small.

Keywords: Higgs Physics, Supersymmetric Standard Model

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Contents

1 Introduction

Searches for physics beyond the Standard Model (SM) by the LHC experiments ATLAS and CMS have set rather stringent constraints on the masses of supersymmetric (SUSY) particles and of Z' bosons. Indeed, in the case of the E_6 inspired models the LHC data exclude Z' resonances with masses $M_{Z'}$ below 2.5 TeV [1, 2]. In the simplest U(1) extensions of the Minimal Supersymmetric extension of the Standard Model (MSSM) the extra U(1) gauge symmetry is normally broken by the vacuum expectation value (VEV) of the scalar component of a superfield S, which is a singlet under the SM gauge group and carries a non-zero U(1) charge.[1] Since the VEV $\langle S \rangle$ of S and the mass $M_{Z'}$ of the Z' boson are determined by the SUSY breaking scale in these models, the multi-TeV Z' mass typically implies that other sparticles also have multi-TeV masses. Such masses typically exceed even the limits on the first and second generation squarks set by the LHC and are well above the current subTeV limits on the third generation sfermions and on additional Higgs bosons. It is therefore worthwhile considering alternative realisations of U(1) extensions, in which the Z' boson can be substantially heavier than the sparticles, and the phenomenological implications of such mechanisms.

Such scenarios may be realised when the extra U(1) gauge symmetry, U(1)$'$, is broken by two VEVs coming from the superfields S and $\overline{S}$, which are both singlets under the SM gauge group but have opposite U(1)$'$ charges. In this case the U(1)$'$ D-term contribution to the scalar potential may force the minimum of this potential to be along the D-flat direction (see, for example, [3]). As a consequence the VEVs $\langle S \rangle$ and $\langle \overline{S} \rangle$ can be much larger than the SUSY breaking scale.

[1] In the literature such states are often referred to as SM singlets and this convention will also be followed in this paper.

The simplest renormalisable superpotential of the SUSY model of the type discussed above can be written as

$$W_S = \sigma\phi S\overline{S}\,, \tag{1.1}$$

where ϕ is a scalar superfield that does not participate in the gauge interactions. When the coupling σ goes to zero the corresponding tree-level scalar potential takes the form

$$V_S = m_S^2|S|^2 + m_{\overline{S}}^2|\overline{S}|^2 + m_\phi^2|\phi|^2 + \frac{Q_S^2 g_1'^2}{2}\left(|S|^2 - |\overline{S}|^2\right)^2\,, \tag{1.2}$$

where m_S^2, $m_{\overline{S}}^2$ and m_ϕ^2 are soft SUSY breaking mass parameters squared, while g_1' is the $\mathrm{U}(1)'$ gauge coupling and Q_S is the $\mathrm{U}(1)'$ charge of the SM singlet superfields S and $\overline{S}$. In eq. (1.2) the last term is associated with the extra $\mathrm{U}(1)'$ D-term contribution. In the limit $\langle S\rangle = \langle\overline{S}\rangle$ this quartic term vanishes. If $(m_S^2 + m_{\overline{S}}^2) < 0$ then there is a run-away direction in this model, so that $\langle S\rangle = \langle\overline{S}\rangle \to \infty$. When the F-terms from the interaction in the superpotential eq. (1.1) are included this stabilizes the run-away direction and for small values of the coupling σ the SM singlet superfields tend to acquire large VEVs, i.e.

$$\langle\phi\rangle \sim \langle S\rangle \simeq \langle\overline{S}\rangle \sim \frac{1}{\sigma}\sqrt{|m_S^2 + m_{\overline{S}}^2|}\,, \tag{1.3}$$

resulting in an extremely heavy Z' boson.

Although the SUSY model mentioned above looks rather simple and elegant it also possesses an additional accidental global U(1) symmetry which can be associated with the Peccei-Quinn (PQ) symmetry [4, 5]. This symmetry is spontaneously broken by the VEVs of the SM singlet superfields resulting in a massless axion [6]. To avoid the appearance of this axion one needs to include in the superpotential of eq. (1.1) polynomial terms with respect to the superfield ϕ which explicitly break the global U(1) symmetry. If the couplings that violate the PQ symmetry are very small, then the particle spectrum of this SUSY model should contain a pseudo-Goldstone boson which can be considerably lighter than all sparticles and Higgs bosons. In fact, the corresponding pseudoscalar Higgs state may be so light that the decay of the SM-like Higgs boson into a pair of these states can be kinematically allowed.

In this article we consider such non-standard Higgs decays within well motivated E_6 inspired extensions of the MSSM, with the particular model described in section 2. We focus on scenarios with an approximate global U(1) symmetry that leads to a pseudo-Goldstone boson in the particle spectrum. The pseudo-Goldstone state in these scenarios is mainly a linear superposition of the imaginary parts of the scalar components of the SM singlet superfields ϕ, S and $\overline{S}$. The SM-like Higgs boson on the other hand is predominantly a linear superposition of the neutral components of the Higgs doublets H_u and H_d, so that the coupling of the pseudo-Goldstone state to the SM-like Higgs boson can be expected to be somewhat suppressed. However it can still lead to a non-negligible branching ratio of the lightest Higgs decays into a pair of pseudo-Goldstone bosons.

In this context it is worth noting that the decay rate of the SM-like Higgs state into a pair of pseudoscalars was intensively studied within the simplest extension of the MSSM, i.e. the Next-to-Minimal Supersymmetric extension of the SM (NMSSM). For reviews of non-standard Higgs boson decays see [7, 8] and for a more recent work see e.g. [9]. The NMSSM superpotential is given by [10–12]:

$$W_{\rm NMSSM} = \lambda S(H_d H_u) + \frac{\kappa}{3}S^3 + W_{\rm MSSM}(\mu = 0)\,, \tag{1.4}$$

where $W_{\rm MSSM}(\mu = 0)$ is the MSSM superpotential with the bilinear mass μ set to zero, and λ and κ are new NMSSM-specific couplings.[2] In the limit where the cubic coupling $\kappa = 0$ the Lagrangian of the NMSSM is invariant under the transformations of the PQ symmetry which leads to the massless axion when it is spontaneously broken by the VEV $\langle S\rangle$. If κ is rather small then the NMSSM particle spectrum involves one light scalar state and one light pseudoscalar state. In addition, if $\kappa \to 0$ and the SUSY breaking scale is of the order of a TeV, then the lightest SUSY particle (LSP) is the lightest neutralino $\tilde{\chi}^0_1$ and is predominantly singlino. In this case the LSP couplings to the SM particles are quite small resulting in a relatively small annihilation cross section for $\tilde{\chi}^0_1\tilde{\chi}^0_1 \to$ SM particles, which gives rise to a relic density that is typically much larger than its measured value. As a consequence it seems to be rather problematic to find phenomenologically viable scenarios with a light pseudoscalar in the case of the NMSSM with approximate PQ symmetry. Nevertheless a sufficiently light pseudoscalar can always be obtained by tuning the parameters of the NMSSM.

In contrast to the NMSSM the mass of the lightest neutralino in the SUSY model considered here does not become small when the PQ symmetry violating couplings vanish. Moreover, even when all PQ symmetry violating couplings are negligibly small, the LSP can be higgsino-like. This allows a reasonable value for the dark matter density to be obtained if the LSP has a mass below 1 TeV (see, for example [13, 14]). Thus the approximate PQ symmetry can lead to phenomenologically viable scenarios with a light pseudoscalar in this model.

The purpose of this paper is to study the implications of a light pseudoscalar Higgs state, which can appear in this phenomenologically interesting model, for the decays of the SM-like Higgs boson. In particular, we investigate what values of branching ratios can be expected for the SM-like Higgs decays into a pair of light pseudoscalars, taking into account the constraints arising from the model itself and the experimental restrictions due to the LHC Higgs search data.

The paper is organised as follows. In section 2 we briefly review E_6 inspired SUSY models with exact custodial $\tilde{Z}^H_2$ symmetry. In section 3 we study the breakdown of the gauge symmetry and the implications for Higgs phenomenology. In section 4 we discuss a set of benchmark scenarios that lead to the decays of the lightest Higgs boson into a pair of pseudoscalar states. Our results are summarized in section 5. In appendix A the spectrum of the neutralino states is examined.

[2] One of the motivations of the NMSSM is the dynamic generation of the supersymmetric Higgs mass parameter μ through the coupling term SH_dH_u when the singlet field S acquires a vacuum expectation value $\langle S\rangle$, i.e. $\mu = \lambda\langle S\rangle$.

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2 E_6 inspired SUSY models with exact $\tilde{Z}_2^H$ symmetry

In this section, we briefly review the E_6 inspired SUSY models with exact custodial $\tilde{Z}_2^H$ symmetry [15], which we then use to demonstrate how light pseudoscalar states can appear in SUSY models with an extra $U(1)'$ gauge symmetry. We also consider what kind of Higgs decay rates they can lead to.

The breakdown of the E_6 symmetry at high energies may lead to models based on rank-5 gauge groups with an additional $U(1)'$ factor in comparison to the SM. In this case the extra $U(1)'$ gauge symmetry is a linear combination of $U(1)_\chi$ and $U(1)_\psi$,

$$U(1)' = U(1)_\chi \cos\theta' + U(1)_\psi \sin\theta' , \tag{2.1}$$

which are defined by the breakdown of the exceptional Lie group E_6 into SO(10), $E_6 \to SO(10) \times U(1)_\psi$, and the subsequent breakdown of SO(10) into SU(5), $SO(10) \to SU(5) \times U(1)_\chi$ (for a review see refs. [16, 17]). With additional Abelian gauge symmetries it is important to ensure the cancellation of anomalies. In any model based on the subgroup of E_6 the anomalies are canceled automatically if the low-energy spectrum involves complete 27-plets. Consequently, in E_6 inspired SUSY models the particle spectrum is extended to fill out three complete 27-dimensional representations of E_6. Each 27-plet, referred to as 27_i with $i = 1, 2, 3$, contains one generation of ordinary matter, a SM singlet field S_i, that carries non-zero $U(1)'$ charge, up- and down-type Higgs doublets H_i^u and H_i^d and charged $\pm 1/3$ exotic quarks D_i and $\bar{D}_i$.

Different aspects of the phenomenology of the E_6 inspired SUSY models have been extensively studied in the past [17–28]. A few years ago the Tevatron and early LHC Z' mass limits in these models were discussed in ref. [29]. Collider signatures associated with the exotic quarks and squarks have been considered in [30]. Previously, the implications of E_6 inspired SUSY models with an additional $U(1)'$ gauge symmetry have been studied for electroweak symmetry breaking (EWSB) [31–37], neutrino physics [38, 39], leptogenesis [40, 41], electroweak (EW) baryogenesis [42, 43], the muon anomalous magnetic moment [44, 45], the electric dipole moment of the electron [46] and of the tau lepton [47], for lepton flavor violating processes like $\mu \to e\gamma$ [48] and for CP-violation in the Higgs sector [49]. The neutralino sector in E_6 inspired SUSY models was analysed in [36, 46–48, 50–57]. Such models have also been proposed as the solution to the tachyon problems of anomaly-mediated SUSY breaking, via $U(1)'$ D-term contributions [58], and have been used in combination with a generation symmetry to construct a model explaining the fermion mass hierarchy and mixing [59]. The Higgs sector and the theoretical upper bound on the lightest Higgs boson mass in the E_6 inspired SUSY models were examined in [37, 57, 60–65].

Here we focus on the E_6 inspired SUSY extension of the SM based on the low-energy SM gauge group together with an extra $U(1)_N$ gauge symmetry in which right-handed neutrinos do not participate in the gauge interactions. This corresponds to $\theta' = \arctan\sqrt{15}$ in eq. (2.1). In this Exceptional Supersymmetric Standard Model (E_6SSM) [60, 61] right-handed neutrinos may be superheavy, shedding light on the origin of the mass hierarchy in the lepton sector and providing a mechanism for the generation of the baryon asymmetry in the Universe via leptogenesis [40, 41]. E_6 inspired SUSY models with an additional

$\mathrm{U}(1)_N$ gauge symmetry have been studied in a variety of contexts. Thus they have been investigated in [39] in the context of non-standard neutrino models with extra singlets, in [50] from the point of view of $Z-Z'$ mixing, in refs. [36, 50, 51] the neutralino sector was explored, in [36, 66] the renormalisation group (RG) flow of the couplings was examined, and in [35–37] EWSB was studied. The presence of a Z' boson and of exotic quarks as predicted by the E_6SSM provides spectacular new physics signals at the LHC, which were analysed in [60–64, 67, 68]. The existence of light exotic particles also leads to the non-standard decays of the SM-like Higgs boson that were discussed in detail in [69–72]. Within the constrained version of the E_6SSM the particle spectrum and associated collider signatures were studied in [73–77] and the degree of fine tuning has recently been examined in [78]. The threshold corrections to the running gauge and Yukawa couplings in the E_6SSM and their numerical impact in the cE_6SSM were studied in detail in ref. [79]. Alternative boundary conditions that take account of D-terms from the other U(1) gauge symmetry broken at the GUT scale were considered in [80], using the first or second generation sfermion masses to constrain the GUT scale parameters. The renormalisation of the VEVs in the E_6SSM was considered in [81, 82].

The presence of exotic matter in the E_6SSM generically leads to non-diagonal flavor transitions and rapid proton decay. A set of discrete symmetries can be imposed in order to suppress these processes [60, 61]. In this article we study the non-standard Higgs decays mentioned above within the E_6 inspired SUSY models with the extra $\mathrm{U}(1)_N$ factor in which a single discrete $\tilde{Z}_2^H$ symmetry forbids tree-level flavor-changing transitions and the most dangerous baryon and lepton number violating operators [15]. These models imply that E_6 or its subgroup is broken down to $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_\psi \times \mathrm{U}(1)_\chi$ near the Grand Unification Theory (GUT) scale, which we denote as M_X. Below this scale M_X the particle content of these SUSY models involves three copies of 27-plets and a set of M_l and $\overline{M}_l$ supermultiplets from the incomplete $27'_l$ and $\overline{27'}_l$ representations of E_6, where l runs over the different multiplets that are summarized below in table 1. All components of the complete 27_i-plets are odd under the discrete symmetry $\tilde{Z}_2^H$, while the supermultiplets $\overline{M}_l$ can be either odd or even. The supermultiplets M_l are even under the $\tilde{Z}_2^H$ symmetry, and as a consequence they can be used for the breakdown of the gauge symmetry, while preserving the discrete symmetry. To ensure that the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_\psi \times \mathrm{U}(1)_\chi$ symmetry is broken down to the $\mathrm{U}(1)_{\mathrm{em}}$ gauge group associated with electromagnetism, these supermultiplets M_l should contain the supermultiplets H_u, H_d, S and a supermultiplet called N_H^c, which has the same quantum numbers as the right-handed neutrino. Just below the GUT scale the $\mathrm{U}(1)_\psi \times \mathrm{U}(1)_\chi$ gauge symmetry is expected to be broken by the VEVs of N_H^c and $\overline{N}_H^c$ down to the $\mathrm{U}(1)_N \times Z_2^M$ in these E_6 inspired models, where $Z_2^M = (-1)^{3(B-L)}$ is a matter parity. This is possible because matter parity is a discrete subgroup of $\mathrm{U}(1)_\psi$ and $\mathrm{U}(1)_\chi$. With such a breakdown into $\mathrm{U}(1)_N \times Z_2^M$ the right-handed neutrino mass can be generated without breaking the remaining gauge symmetry and all exotic states which originate from the 27_i representations of E_6 as well as ordinary quark and lepton states survive down to low energies. In general the large VEVs $\langle N_H^c \rangle \sim \langle \overline{N}_H^c \rangle \lesssim M_X$ also induce the large Majorana masses for the right-handed neutrinos allowing them to be used for the see-saw mechanism. Since N_H^c and $\overline{N}_H^c$ acquire VEVs both of these supermultiplets must be even under the imposed $\tilde{Z}_2^H$ symmetry.

	27_i	27_i	$27'_{H_u}$ $(27'_{H_d})$	$27'_S$	$\overline{27'}_{H_u}$ $(\overline{27'}_{H_d})$	$\overline{27'}_S$	$27'_N$ $(\overline{27'}_N)$	$27'_L$ $(\overline{27'}_L)$	1
	$Q_i, u^c_i, d^c_i,$ L_i, e^c_i, N^c_i	$\overline{D}_i, D_i,$ H^d_i, H^u_i, S_i	H_u (H_d)	S	$\overline{H}_u$ $(\overline{H}_d)$	$\overline{S}$	N^c_H $(\overline{N}^c_H)$	L_4 $(\overline{L}_4)$	ϕ
$\tilde{Z}^H_2$	$-$	$-$	$+$	$+$	$-$	$\pm$	$+$	$+$	$+$
Z^M_2	$-$	$+$	$+$	$+$	$+$	$+$	$-$	$-$	$+$
Z^E_2	$+$	$-$	$+$	$+$	$-$	$\pm$	$-$	$-$	$+$

Table 1. Transformation properties of different components of E_6 multiplets under the discrete symmetries $\tilde{Z}^H_2$, Z^M_2 and Z^E_2. A '+' denotes that the field is even under the discrete symmetry, a '$-$' means that the field is odd, while '$\pm$' denotes that the field may be even or odd depending on which construction is considered.

At the TeV scale the scalar components of the superfields H_u, H_d and S play the role of Higgs fields. The VEVs of the neutral scalar components break the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_N$ gauge symmetry down to $\mathrm{U}(1)_{\mathrm{em}}$. Because of this the supermultiplets H_u, H_d and S must also be even under the $\tilde{Z}^H_2$ symmetry. In contrast, in the simplest scenario $\overline{H}_u$ and $\overline{H}_d$ are expected to be odd under this custodial symmetry so that they can combine with the superposition of the corresponding components from the 27_i, forming vectorlike states that gain masses of order M_X. The scalar component of the superfield $\overline{S}$ may also acquire a non-zero VEV breaking the $\mathrm{U}(1)_N$ symmetry. If this is the case then $\overline{S}$ has to be even under the $\tilde{Z}^H_2$ symmetry. When $\overline{S}$ is odd under the $\tilde{Z}^H_2$ symmetry then it can get combined with the superposition of the appropriate components of 27_i resulting in the formation of superheavy vectorlike states with masses $\sim M_X$. The custodial $\tilde{Z}^H_2$ symmetry allows Yukawa interactions in the superpotential that originate from $27'_l \times 27'_m \times 27'_n$ and $27'_l \times 27_i \times 27_k$ ($i, k = 1, 2, 3$ and l, m, n running over the multiplets given in table 1). It is easy to check that the corresponding set of operators does not contain any operators that lead to rapid proton decay. Since the set of supermultiplets M_l contains only one pair of doublets, H_d and H_u, the down-type quarks and charged leptons couple to just one Higgs doublet, H_d, and the up-type quarks couple to H_u only. As a result flavor-changing processes are forbidden at tree-level.

Nonetheless, if the set of $\tilde{Z}^H_2$-even supermultiplets M_l involves only H_u, H_d, S and N^c_H then the Lagrangian of the model is invariant not only with respect to the $\mathrm{U}(1)_B$ associated with baryon number conservation but also under $\mathrm{U}(1)_D$ symmetry transformations

$$D \to e^{i\alpha} D\,, \qquad \overline{D} \to e^{-i\alpha}\overline{D}\,. \tag{2.2}$$

The $\mathrm{U}(1)_D$ symmetry forbids renormalisable interactions through which the exotic quarks D can decay, thereby ensuring that the lightest exotic quark is very long-lived. Indeed, as for $\mathrm{U}(1)_B$ we expect the $\mathrm{U}(1)_D$ global symmetry to be broken by a set of non-renormalisable operators that are suppressed by inverse powers of M_X or the Planck scale M_{Pl}. While these operators allow the lightest exotic quark to decay, its lifetime tends to be considerably larger than the age of the Universe. Long-lived exotic quarks would have been produced

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during the very early epochs of the Big Bang and those that survive annihilation would subsequently have been confined in heavy hadrons forming nuclear isotopes that would be present in terrestrial matter. Various theoretical estimates [83, 84] show that if such stable relics in the mass range from 1 GeV to 10 TeV would exist in nature, today their concentration would be $\mathcal{O}(10^{-10})$ per nucleon. At the same time different experiments set strong upper limits on the relative concentrations of such nuclear isotopes from 10^{-15} to 10^{-30} per nucleon [85–87]. Therefore E_6 inspired models with very long-lived exotic quarks are ruled out.

To ensure that the lightest exotic quarks decay within a reasonable time in the simplest scenario, we supplement the set of $\tilde{Z}_2^H$ even supermultiplets M_l with L_4, where L_4 and $\overline{L}_4$ are lepton $\mathrm{SU}(2)_L$ doublet supermultiplets that originate from a pair of additional $27'_L$ and $\overline{27'}_L$. The supermultiplets L_4 and $\overline{L}_4$ should form TeV scale vectorlike states to break the $\mathrm{U}(1)_D$ symmetry and render the lightest exotic quark unstable.[3] Therefore L_4 and $\overline{L}_4$ both have to be even under the $\tilde{Z}_2^H$ symmetry. In this case the baryon and lepton number conservation implies that the exotic quarks are leptoquarks.

Here we assume that, in addition to H_u, H_d, S, L_4, $\overline{L}_4$ N_H^c and $\overline{N}_H^c$, the particle spectrum below the scale M_X involves $\tilde{Z}_2^H$-even superfields $\overline{S}$ and ϕ. The superfield ϕ does not participate in the $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_\psi \times \mathrm{U}(1)_\chi$ gauge interactions but its scalar component acquires a non-zero VEV. Taking into account that the components of the superfields $\overline{S}$ and ϕ are expected to gain TeV scale masses whereas the right-handed neutrino superfields are superheavy, the low-energy matter content in the E_6 inspired SUSY models discussed above involves

$$\begin{aligned}&(Q_i,\, u_i^c,\, d_i^c,\, L_i,\, e_i^c) + (D_i,\, \bar{D}_i) + (S_i) + (H^u_\alpha) + (H^d_\alpha)\\&+L_4 + \overline{L}_4 + S + \overline{S} + H_u + H_d + \phi\,,\end{aligned} \tag{2.3}$$

where $\alpha = 1, 2$ runs over the first two generations and $i = 1, 2, 3$ runs over all three. We have denoted here the left-handed quark and lepton doublets by Q_i and L_i, respectively, and the right-handed up- and down-type quarks and charged leptons by u_i^c, d_i^c and e_i^c. Neglecting all suppressed non-renormalisable interactions, the low-energy effective superpotential of these models can be written as

$$\begin{aligned}W =\, &\lambda S(H_u H_d) - \sigma\phi S\overline{S} + \frac{\kappa}{3}\phi^3 + \frac{\mu}{2}\phi^2 + \Lambda\phi\\&+\lambda_{\alpha\beta} S(H^d_\alpha H^u_\beta) + \kappa_{ij} S(D_i \overline{D}_j) + \tilde{f}_{i\alpha} S_i (H^d_\alpha H_u) + f_{i\alpha} S_i (H_d H^u_\alpha)\\&+g^D_{ij}(Q_i L_4)\overline{D}_j + h^E_{i\alpha} e_i^c (H^d_\alpha L_4) + \mu_L L_4 \overline{L}_4 + \tilde{\sigma}\phi L_4 \overline{L}_4 + W_{\mathrm{MSSM}}(\mu = 0)\,,\end{aligned} \tag{2.4}$$

in terms of the dimensionless couplings $\lambda, \sigma, \kappa, \lambda_{\alpha\beta}, \kappa_{ij}, \tilde{f}_{i\alpha}, f_{i\alpha}, g^D_{ij}, h^E_{i\alpha}, \tilde{\sigma}$ and the dimensionful couplings μ, μ_L and Λ, with $i, j = 1, 2, 3$ and $\alpha, \beta = 1, 2$.[4] In the limit when $\tilde{\sigma}$, κ, μ and Λ vanish the model possesses an extra global PQ symmetry. The $\mathrm{U}(1)_Y$, $\mathrm{U}(1)_N$ and PQ charges of all the matter fields are summarised in table 2.

[3] The appropriate mass term $\mu_L L_4 \overline{L}_4$ in the superpotential can be induced within SUGRA models just after the breakdown of local SUSY if the Kähler potential contains an extra term $(Z_L(L_4\overline{L}_4) + h.c)$ [88, 89].

[4] In principle, the superpotential in eq. (2.4) can also include the term $\mu_S S\overline{S}$. However this term can be eliminated by an appropriate redefinition of the superfield ϕ, i.e $\phi \to \phi - \mu_S/\sigma$.

	Q	u^c	d^c	L	e^c	S	H_u	H_d	D	$\overline{D}$	L_4	$\overline{L}_4$	$\overline{S}$	ϕ
$\sqrt{\frac{5}{3}}Q_i^Y$	$\frac{1}{6}$	$-\frac{2}{3}$	$\frac{1}{3}$	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{1}{3}$	$-\frac{1}{2}$	$\frac{1}{2}$	0	0
$\sqrt{40}Q_i^N$	1	1	2	2	1	5	-2	-3	-2	-3	2	-2	-5	0
PQ charges	-1	0	0	-1	0	-2	+1	+1	0	+2	-1	+1	0	+2

Table 2. The $U(1)_Y$, $U(1)_N$ and PQ charges of matter fields, where Q_i^N and Q_i^Y are defined with the correct E_6 normalisation factor while PQ symmetry charges are not normalized. The PQ charges are not unique, as one can add to them any contributions which are proportional to $U(1)_Y$ and $U(1)_N$ charges.

The gauge group and field content of the E_6 inspired SUSY models under consideration can originate from the 5D and 6D orbifold GUT models in which the splitting of GUT multiplets can be naturally achieved [15]. In these orbifold GUT models all GUT relations between the Yukawa couplings can get spoiled while the approximate unification of the SM gauge couplings still takes place. From eq. (2.3) it follows that extra matter beyond the MSSM fills in complete SU(5) representations in these models. As a consequence the gauge coupling unification remains almost exact in the one-loop approximation. It was also shown that in the two-loop approximation the unification of the gauge couplings in the SUSY models under consideration can be achieved for any phenomenologically acceptable value of the strong coupling $\alpha_3(M_Z)$ at the scale M_Z, consistent with the measured central low-energy value [15, 66].

For the analysis of the phenomenological implications of the SUSY models discussed above it is convenient to introduce the Z_2^E symmetry, which is defined such that $\tilde{Z}_2^H = Z_2^M \times Z_2^E$. The transformation properties of different components of the 27_i, $27'_l$ and $\overline{27'}_l$ supermultiplets under the $\tilde{Z}_2^H$, Z_2^M and Z_2^E symmetries are summarized in table 1. Since the low-energy effective Lagrangian of the E_6 inspired SUSY models studied here is invariant under the transformations of the Z_2^M and $\tilde{Z}_2^H$ symmetries, the Z_2^E symmetry associated with the exotic states is also conserved. The invariance of the Lagrangian under the Z_2^E symmetry implies that in collider experiments the exotic particles, which are odd under this symmetry, can only be created in pairs and the lightest exotic state should be absolutely stable. Using the method proposed in [90–92] it was argued that the masses of the lightest inert neutralino states,[5] which are predominantly linear superpositions of the fermion components of the superfields S_i from complete 27_i representations of E_6, do not exceed 60–65 GeV [70–72]. Because of this the corresponding states tend to be the lightest exotic particles in the spectrum.

The presence of lightest exotic particles with masses below 60 GeV gives rise to new decay channels of the SM-like Higgs boson. Moreover, if these states are heavier than 5–10 GeV (i.e. approximately above the bottom quark pair threshold) then the SM-like Higgs state decays predominantly into the lightest inert neutralinos while the total branch-

[5] We use the terminology "inert Higgs" to denote the H_α^u, H_α^d and S, whose scalar components do not develop VEVs. The fermionic components of these supermultiplets form inert neutralino and chargino states.

ing ratio into SM particles gets strongly suppressed. Nowadays such scenarios are basically ruled out. On the other hand if the lightest exotic particles have masses below 5 GeV their couplings to the SM-like Higgs state are small so that problems with non-standard Higgs decays can be avoided. However, as their couplings to the gauge bosons, quarks and leptons are also very small this results in a cold dark matter density that is much larger than its measured value because the corresponding annihilation cross section tends to be small. The simplest phenomenologically viable scenarios imply that the lightest inert neutralinos are substantially lighter than 1 eV.[6] This can be achieved if $f_{\alpha\beta} \sim \tilde{f}_{\alpha\beta} \lesssim 10^{-6}$. In this case the lightest exotic particles form hot dark matter (dark radiation) in the Universe but give only a very minor contribution to the dark matter density.

The Z_2^M symmetry conservation ensures that R-parity is also conserved in the SUSY models discussed above. As also the Z_2^E symmetry is conserved there are two states possible that can be stable. This is either the lightest R-parity even exotic state or the lightest R-parity odd state with $Z_2^E = +1$. In the E_6 inspired models studied here the stable state tends to be the lightest ordinary neutralino state (i.e. the lightest neutralino state with $Z_2^E = +1$). Like in the MSSM, this state may account for all or for some of the observed cold dark matter density.

As mentioned before, in the simplest case the sector responsible for the breakdown of the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_N$ gauge symmetry involves H_u, H_d and S. For this case the Higgs sector of the E_6 inspired SUSY models with the extra $\mathrm{U}(1)_N$ factor was explored in [60]. If CP-invariance is preserved then the Higgs spectrum in these models contains three CP-even, one CP-odd and two charged states. The singlet dominated CP-even state is always almost degenerate with the Z' gauge boson. In contrast to the MSSM, the lightest Higgs boson in these models can be heavier than $110 - 120$ GeV even at tree-level. In the two-loop approximation the lightest Higgs boson mass does not exceed $150 - 155$ GeV [60]. Recently, the RG flow of the Yukawa couplings and the theoretical upper bound on the lightest Higgs boson mass in these models were analysed in the vicinity of the quasi-fixed point [94] that appears as a result of the intersection of the invariant and quasi-fixed lines [95]. It was argued that near the quasi-fixed point the upper bound on the mass of the SM-like Higgs boson is rather close to 125 GeV [94].

The qualitative pattern of the Higgs spectrum in the E_6 inspired SUSY models with the extra $\mathrm{U}(1)_N$ gauge symmetry and minimal Higgs sector is determined by the Yukawa coupling λ. When $\lambda < g_1'$, where g_1' is the gauge coupling associated with the $\mathrm{U}(1)_N$ gauge symmetry, the singlet dominated CP-even state is very heavy and decouples from the rest of the spectrum, which makes the Higgs spectrum indistinguishable from the one in the MSSM. If $\lambda \gtrsim g_1'$ the Higgs spectrum has an extremely hierarchical structure, which is rather similar to the one that arises in the NMSSM with the approximate PQ symmetry [96–98]. As a consequence the mass matrix of the CP-even Higgs sector can be diagonalised using a perturbative expansion [97–101]. In this case the mass of the second lightest CP-even Higgs state is set by the Z' boson mass, while the heaviest CP-even, CP-odd and charged states are almost degenerate and lie beyond the multi-TeV range.

[6] The presence of very light neutral fermions in the particle spectrum might have interesting implications for neutrino physics (see, for example [93]).

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3 The Higgs sector

3.1 The Higgs potential and gauge symmetry breaking

As was mentioned in the previous section, the sector responsible for breaking the gauge symmetry in the SUSY model under consideration includes two Higgs doublets, H_u and H_d, as well as the SM singlet fields S, $\overline{S}$ and ϕ.

The interactions between these fields are determined by the structure of the gauge interactions and by the superpotential in eq. (2.4). The resulting Higgs potential reads

$$\begin{aligned}
V &= V_F + V_D + V_{\rm soft} + \Delta V\,,\\
V_F &= \lambda^2|S|^2(|H_d|^2+|H_u|^2) + \left|\lambda(H_dH_u) - \sigma\phi\overline{S}\right|^2 + \sigma^2|\phi|^2|S|^2 +\\
&\quad + \left|-\sigma(S\overline{S}) + \kappa\phi^2 + \mu\phi + \Lambda\right|^2\,,\\
V_D &= \sum_{a=1}^{3}\frac{g_2^2}{8}\left(H_d^\dagger\sigma_a H_d + H_u^\dagger\sigma_a H_u\right)^2 + \frac{g'^2}{8}\left(|H_d|^2-|H_u|^2\right)^2 +\\
&\quad + \frac{g_1'^2}{2}\left(\tilde{Q}_{H_d}|H_d|^2 + \tilde{Q}_{H_u}|H_u|^2 + \tilde{Q}_S|S|^2 - \tilde{Q}_S|\overline{S}|^2\right)^2\,,\\
V_{\rm soft} &= m_S^2|S|^2 + m_{\overline{S}}^2|\overline{S}|^2 + m_{H_d}^2|H_d|^2 + m_{H_u}^2|H_u|^2 + m_\phi^2|\phi|^2\\
&\quad + \left[\lambda A_\lambda S(H_uH_d) - \sigma A_\sigma\phi(S\overline{S}) + \frac{\kappa}{3}A_\kappa\phi^3 + B\frac{\mu}{2}\phi^2 + \xi\Lambda\phi + h.c.\right],
\end{aligned} \tag{3.1}$$

where $H_d^T = (H_d^0,\, H_d^-)$, $H_u^T = (H_u^+,\, H_u^0)$ and $(H_dH_u) = H_u^+H_d^- - H_u^0H_d^0$, and $\tilde{Q}_{H_d}$, $\tilde{Q}_{H_u}$ and $\tilde{Q}_S$ are the effective $\mathrm{U}(1)_N$ charges of H_d, H_u and S. Furthermore σ_a $(a = 1, 2, 3)$ denote the three Pauli matrices. At tree-level the Higgs potential in eq. (3.1) is described by the sum of the first three terms. V_F and V_D contain the F-and D-term contributions that do not violate SUSY. The terms in the expression for V_D are proportional to the $\mathrm{SU}(2)_L$, $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_N$ gauge couplings, i.e. g_2, g' and g_1', respectively. The values of the gauge couplings g_2 and g' at the EW scale are well known, whereas the low-energy value of the extra $\mathrm{U}(1)_N$ coupling g_1' and the effective $\mathrm{U}(1)_N$ charges of H_d, H_u and S can be calculated assuming gauge coupling unification [60]. The soft SUSY breaking terms are collected in $V_{\rm soft}$ and include the soft masses $m_{H_d}^2$, $m_{H_u}^2$, m_S^2, $m_{\overline{S}}^2$ and m_ϕ^2, the trilinear couplings A_λ, A_σ and A_κ, the bilinear coupling B and a linear coupling ξ. The term ΔV in eq. (3.1) contains the loop corrections to the Higgs effective potential. In SUSY models the most significant contribution to ΔV comes from the loops involving the top-quark and its superpartners, the stops.

At the physical minimum of the scalar potential, eq. (3.1), the Higgs fields develop VEVs

$$\begin{aligned}
&< H_d > = \frac{1}{\sqrt{2}}\begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \qquad < H_u > = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_2 \end{pmatrix},\\
&< S > = \frac{s_1}{\sqrt{2}}\,, \qquad < \overline{S} > = \frac{s_2}{\sqrt{2}}\,, \qquad < \phi > = \frac{\varphi}{\sqrt{2}}\,.
\end{aligned} \tag{3.2}$$

Using the short-hand notation $\partial V/\partial\Phi_{(\Phi=\langle\Phi\rangle)} \equiv \partial V/\partial\langle\Phi\rangle$, the minimum conditions for the Higgs potential of eq. (3.1) read,

$$\begin{aligned}
\frac{\partial V}{\partial s_1} &= m_S^2\, s_1 - \frac{\lambda A_\lambda}{\sqrt{2}} v_1 v_2 - \frac{\sigma A_\sigma}{\sqrt{2}}\varphi s_2 + \left(\frac{\sigma}{2}s_1 s_2 - \frac{\kappa}{2}\varphi^2 - \frac{\mu}{\sqrt{2}}\varphi - \Lambda\right)\sigma s_2 \\
&\quad + \frac{\sigma^2}{2}\varphi^2 s_1 + \frac{g_1'^2}{2}\left(\tilde{Q}_{H_d} v_1^2 + \tilde{Q}_{H_u} v_2^2 + \tilde{Q}_S s_1^2 - \tilde{Q}_S s_2^2\right)\tilde{Q}_S s_1 \\
&\quad + \frac{\lambda^2}{2}(v_1^2+v_2^2)s_1 + \frac{\partial\Delta V}{\partial s_1} = 0\,, \\
\frac{\partial V}{\partial s_2} &= m_{\bar{S}}^2\, s_2 - \frac{\sigma A_\sigma}{\sqrt{2}}\varphi s_1 + \left(\frac{\sigma}{2}s_1 s_2 - \frac{\kappa}{2}\varphi^2 - \frac{\mu}{\sqrt{2}}\varphi - \Lambda\right)\sigma s_1 \\
&\quad + \frac{\sigma^2}{2}\varphi^2 s_2 - \frac{g_1'^2}{2}\left(\tilde{Q}_{H_d} v_1^2 + \tilde{Q}_{H_u} v_2^2 + \tilde{Q}_S s_1^2 - \tilde{Q}_S s_2^2\right)\tilde{Q}_S s_2 \\
&\quad + \frac{\lambda\sigma}{2} v_1 v_2 \varphi + \frac{\partial\Delta V}{\partial s_2} = 0\,, \\
\frac{\partial V}{\partial \varphi} &= m_\varphi^2\, \varphi - \frac{\sigma A_\sigma}{\sqrt{2}} s_1 s_2 + B\mu\varphi + \sqrt{2}\xi\Lambda + \frac{\kappa A_\kappa}{\sqrt{2}}\varphi^2 + \frac{\sigma^2}{2}(s_1^2+s_2^2)\varphi \\
&\quad - 2\left(\frac{\sigma}{2}s_1 s_2 - \frac{\kappa}{2}\varphi^2 - \frac{\mu}{\sqrt{2}}\varphi - \Lambda\right)\left(\kappa\varphi + \frac{\mu}{\sqrt{2}}\right) + \frac{\lambda\sigma}{2}v_1 v_2 s_2 + \frac{\partial\Delta V}{\partial\varphi} = 0\,, \\
\frac{\partial V}{\partial v_1} &= m_{H_d}^2\, v_1 - \frac{\lambda A_\lambda}{\sqrt{2}} s_1 v_2 + \frac{\lambda\sigma}{2} v_2 s_2 \varphi + \frac{\lambda^2}{2}(v_2^2+s_1^2)v_1 + \frac{\bar{g}^2}{8}\left(v_1^2 - v_2^2\right)v_1 \\
&\quad + \frac{g_1'^2}{2}\left(\tilde{Q}_{H_d} v_1^2 + \tilde{Q}_{H_u} v_2^2 + \tilde{Q}_S s_1^2 - \tilde{Q}_S s_2^2\right)\tilde{Q}_{H_d} v_1 + \frac{\partial\Delta V}{\partial v_1} = 0\,, \\
\frac{\partial V}{\partial v_2} &= m_{H_u}^2 v_2 - \frac{\lambda A_\lambda}{\sqrt{2}} s_1 v_1 + \frac{\lambda\sigma}{2} v_1 s_2 \varphi + \frac{\lambda^2}{2}(v_1^2+s_1^2)v_2 + \frac{\bar{g}^2}{8}\left(v_2^2 - v_1^2\right)v_2 \\
&\quad + \frac{g_1'^2}{2}\left(\tilde{Q}_{H_d} v_1^2 + \tilde{Q}_{H_u} v_2^2 + \tilde{Q}_S s_1^2 - \tilde{Q}_S s_2^2\right)\tilde{Q}_{H_u} v_2 + \frac{\partial\Delta V}{\partial v_2} = 0\,,
\end{aligned} \tag{3.3}$$

where $\bar{g} = \sqrt{g_2^2 + g'^2}$. Instead of specifying v_1, v_2, s_1 and s_2, it is more convenient to use

$$\tan\beta = v_2/v_1 \qquad \text{and} \qquad \tan\theta = s_2/s_1\,. \tag{3.4}$$

The VEV v is given by the electroweak scale, $v = \sqrt{v_1^2 + v_2^2} \approx 246\,\text{GeV}$, and $s = \sqrt{s_1^2 + s_2^2}$ sets the Z' mass, as discussed below.

Initially the Higgs sector involves fourteen degrees of freedom. However four of them are massless Goldstone modes. They are swallowed by the $W^\pm$, Z and Z' gauge bosons. The charged $W^\pm$ bosons gain masses via the interaction with the neutral components of the Higgs doublets H_u and H_d just in the same way as in the MSSM, resulting in $M_W = \frac{g_2}{2}v$. On the other hand the mechanism of the neutral gauge boson mass generation differs substantially. Let the Z' and Z states be the gauge bosons associated with the group $U(1)_N$ and with the SM-like Z boson, respectively. Then the $Z - Z'$ mass-squared matrix is given by

$$M_{ZZ'}^2 = \begin{pmatrix} \frac{\bar{g}^2}{4}v^2 & \Delta^2 \\ \Delta^2 & g_1'^2 v^2\left(\tilde{Q}_{H_d}^2\cos^2\beta + \tilde{Q}_{H_u}^2\sin^2\beta\right) + g_1'^2\tilde{Q}_S^2 s^2 \end{pmatrix}, \tag{3.5}$$

where

$$\Delta^2 = \frac{\bar{g}g_1'}{2}v^2\left(\tilde{Q}_{H_d}\cos^2\beta - \tilde{Q}_{H_u}\sin^2\beta\right).$$

The fields S and $\overline{S}$ must acquire large VEVs, *i.e.* $s_1 \simeq s_2 \gg 1\,\text{TeV}$, to ensure that the extra $U(1)_N$ gauge boson is sufficiently heavy. Then the mass of the lightest neutral gauge boson Z_1 is very close to $M_Z = \bar{g}v/2$, whereas the mass of Z' is set by $M_{Z'} \approx g_1'\tilde{Q}_S\, s$.

3.2 The Higgs boson spectrum

For the analysis of the Higgs boson spectrum we use eq. (3.3) for the extrema to express the soft masses $m_{H_d}^2$, $m_{H_u}^2$, m_S^2, $m_{\overline{S}}^2$ and m_ϕ^2 in terms of s, v, φ, β, θ and other parameters. Because of the conversation of the electric charge, the charged components of the Higgs doublets are not mixed with the neutral Higgs fields. They form a separate sector, the spectrum of which is described by a 2×2 mass matrix. The determinant of this matrix is zero and results in the appearance of two Goldstone states, i.e.

$$G^- = H_d^-\cos\beta - H_u^{+*}\sin\beta \tag{3.6}$$

and its charge conjugate (that are absorbed into the longitudinal degrees of freedom of the $W^\pm$ gauge bosons) and of two charged Higgs states,

$$H^+ = H_d^{-*}\sin\beta + H_u^+\cos\beta \tag{3.7}$$

with mass

$$m_{H^\pm}^2 = \frac{\sqrt{2}\lambda s}{\sin 2\beta}\left(A_\lambda\cos\theta - \frac{\sigma\varphi}{\sqrt{2}}\sin\theta\right) - \frac{\lambda^2}{2}v^2 + \frac{g_2^2}{4}v^2 + \Delta_\pm\,, \tag{3.8}$$

where $\Delta_\pm$ denotes the loop corrections to $m_{H^\pm}^2$.

The imaginary parts of the neutral components of the Higgs doublets and the imaginary parts of the two SM singlet fields S and $\overline{S}$ compose two neutral Goldstone states

$$\begin{aligned} G &= \sqrt{2}(\mathrm{Im}\,H_d^0\cos\beta - \mathrm{Im}\,H_u^0\sin\beta)\,,\\ G' &= \sqrt{2}(\mathrm{Im}\,S\cos\theta - \mathrm{Im}\,\overline{S}\sin\theta)\cos\gamma - \sqrt{2}(\mathrm{Im}\,H_u^0\cos\beta + \mathrm{Im}\,H_d^0\sin\beta)\sin\gamma\,, \end{aligned} \tag{3.9}$$

which are swallowed by the Z and Z' bosons, as well as three physical states. In eq. (3.9) we have introduced

$$\tan\gamma = \frac{v}{2s}\sin 2\beta\,. \tag{3.10}$$

In the field basis $(P_1,\, P_2,\, P_3)$, where

$$\begin{aligned} P_1 &= \sqrt{2}(\mathrm{Im}\,H_u^0\cos\beta + \mathrm{Im}\,H_d^0\sin\beta)\cos\gamma + \sqrt{2}(\mathrm{Im}\,S\cos\theta - \mathrm{Im}\,\overline{S}\sin\theta)\sin\gamma\,,\\ P_2 &= \sqrt{2}\left(\mathrm{Im}\,S\sin\theta + \mathrm{Im}\,\overline{S}\cos\theta\right)\,,\\ P_3 &= \sqrt{2}\mathrm{Im}\,\phi\,, \end{aligned} \tag{3.11}$$

the mass matrix of the CP-odd Higgs sector takes the form

$$\tilde{M}^2 = (\tilde{M}_{ij}^2)\,,\quad i,j=1,2,3\,, \tag{3.12}$$

with

$$
\begin{aligned}
\tilde{M}_{11}^2 &= \frac{\sqrt{2}\lambda s}{\sin 2\beta \cos^2\gamma}\left(A_\lambda \cos\theta - \frac{\sigma\varphi}{\sqrt{2}}\sin\theta\right) + \tilde{\Delta}_{11}\,, \\
\tilde{M}_{12}^2 &= \tilde{M}_{21}^2 = \frac{\lambda v}{\sqrt{2}\cos\gamma}\left(A_\lambda \sin\theta + \frac{\sigma\varphi}{\sqrt{2}}\cos\theta\right) + \tilde{\Delta}_{12}\,, \\
\tilde{M}_{13}^2 &= \tilde{M}_{31}^2 = \frac{\lambda\sigma v s}{2\cos\gamma}\sin\theta + \tilde{\Delta}_{13}\,, \\
\tilde{M}_{22}^2 &= \frac{2\sigma\varphi}{\sin 2\theta}\left(\frac{A_\sigma}{\sqrt{2}} + \frac{\kappa}{2}\varphi + \frac{\mu}{\sqrt{2}} + \frac{\Lambda}{\varphi}\right) \\
&\quad + \frac{\lambda v^2 \sin 2\beta}{\sqrt{2}s\sin 2\theta}\left(A_\lambda \sin^3\theta - \frac{\sigma\varphi}{\sqrt{2}}\cos^3\theta\right) + \tilde{\Delta}_{22}\,, \\
\tilde{M}_{23}^2 &= \tilde{M}_{32}^2 = \sigma s\left(\frac{A_\sigma}{\sqrt{2}} - \kappa\varphi - \frac{\mu}{\sqrt{2}}\right) - \frac{\lambda\sigma}{4}v^2\sin 2\beta\cos\theta + \tilde{\Delta}_{23}\,, \\
\\
\tilde{M}_{33}^2 &= \frac{\sigma s^2}{2\sqrt{2}\varphi}A_\sigma \sin 2\theta - 2B\mu - 3\frac{\kappa A_\kappa}{\sqrt{2}}\varphi - \sqrt{2}(\xi+\mu)\frac{\Lambda}{\varphi} + \sigma\kappa s^2\sin 2\theta - \frac{\kappa\mu}{\sqrt{2}}\varphi \\
&\quad -4\kappa\Lambda + \frac{\sigma\mu s^2}{2\sqrt{2}\varphi}\sin 2\theta - \frac{\lambda\sigma s}{4\varphi}v^2\sin\theta\sin 2\beta + \tilde{\Delta}_{33}\,.
\end{aligned}
\tag{3.13}
$$

In eqs. (3.13) the $\tilde{\Delta}_{ij}$ $(i,j=1,2,3)$ denote loop corrections. Since in the models under consideration s must be much larger than v, it follows that γ goes to zero. Moreover, since in phenomenologically acceptable SUSY models the supersymmetry breaking scale also tends to be considerably larger than v, the mixing between P_1 and the other two pseudoscalar states P_2 and P_3 is somewhat suppressed. So one CP-odd mass eigenstate is predominantly P_1. The other two CP-odd mass eigenstates are mainly made up of linear superpositions of the imaginary parts of the SM singlet fields S, $\overline{S}$ and ϕ, i.e. of P_2 and P_3. In other words, as the off-diagonal entries $\tilde{M}_{12}^2$, $\tilde{M}_{13}^2 \ll \tilde{M}_{11}^2$, the mass matrix, eqs. (3.12) and (3.13), can be diagonalised analytically. In particular, the mass of the CP-odd Higgs eigenstate, that is predominantly P_1, is set by $\tilde{M}_{11}^2$. As a consequence this CP-odd state and the charged physical Higgs states are expected to be approximately degenerate.

The mass matrix, eqs. (3.12) and (3.13), is diagonalised by means of a unitary transformation U that relates the components of the CP-odd Higgs basis eq. (3.11) to the corresponding Higgs mass eigenstates A_i $(i=1,2,3)$,

$$
\begin{pmatrix} P_1 \\ P_2 \\ P_3 \end{pmatrix} = U \begin{pmatrix} A_1 \\ A_2 \\ A_3 \end{pmatrix}. \tag{3.14}
$$

The pseudoscalar mass eigenstates are labeled according to increasing absolute value of mass, where A_1 is the lightest CP-odd Higgs state and A_3 the heaviest. At tree-level and neglecting all terms proportional to λv one obtains

$$
\begin{aligned}
m_{A_3}^2 &\simeq \max\left\{\frac{\sqrt{2}\sigma A_\sigma \varphi}{\sin 2\theta \cos^2\delta}\,, \quad \frac{\sqrt{2}\lambda s}{\sin 2\beta}\left(A_\lambda \cos\theta - \frac{\sigma\varphi}{\sqrt{2}}\sin\theta\right)\right\},\\
m_{A_2}^2 &\simeq \min\left\{\frac{\sqrt{2}\sigma A_\sigma \varphi}{\sin 2\theta \cos^2\delta}\,, \quad \frac{\sqrt{2}\lambda s}{\sin 2\beta}\left(A_\lambda \cos\theta - \frac{\sigma\varphi}{\sqrt{2}}\sin\theta\right)\right\},\\
m_{A_1}^2 &\simeq \cos^2\delta\left[-2B\mu - 3\frac{\kappa A_\kappa}{\sqrt{2}}\varphi - \sqrt{2}\xi\frac{\Lambda}{\varphi} + \frac{9}{4}\sigma\kappa s^2\sin 2\theta\right.\\
&\qquad\left. +\sqrt{2}\frac{\sigma\mu s^2}{\varphi}\sin 2\theta + \frac{\sigma s^2\Lambda}{2\varphi^2}\sin 2\theta\right],
\end{aligned}
\tag{3.15}
$$

where we have defined

$$\tan\delta \simeq \frac{s}{2\varphi}\sin 2\theta\,. \tag{3.16}$$

In this case the lightest CP-odd mass eigenstate is a linear combination of P_2 and P_3,

$$A_1 \simeq -P_2\sin\delta + P_3\cos\delta\,. \tag{3.17}$$

In the limit where the previously discussed global U(1) symmetry violating couplings κ, μ and Λ vanish, the mass of the lightest CP-odd Higgs boson goes to zero.

The CP-even Higgs sector involves the $\mathrm{Re}\, H_d^0$, $\mathrm{Re}\, H_u^0$, $\mathrm{Re}\, S$, $\mathrm{Re}\, \overline{S}$ and $\mathrm{Re}\,\phi$. In the field space basis (S_1, S_2, S_3, S_4, S_5), where

$$
\begin{aligned}
\mathrm{Re}\, S &= (S_1\cos\theta + S_2\sin\theta + s_1)/\sqrt{2}\,,\\
\mathrm{Re}\, \overline{S} &= (-S_1\sin\theta + S_2\cos\theta + s_2)/\sqrt{2}\,,\\
\mathrm{Re}\,\phi &= (S_3 + \varphi)/\sqrt{2}\,,\\
\mathrm{Re}\, H_d^0 &= (S_5\cos\beta - S_4\sin\beta + v_1)/\sqrt{2}\,,\\
\mathrm{Re}\, H_u^0 &= (S_5\sin\beta + S_4\cos\beta + v_2)/\sqrt{2}\,,
\end{aligned}
\tag{3.18}
$$

the mass matrix of the CP-even Higgs sector takes the form

$$M^2 = (M_{ij}^2)\,,\quad i,j = 1,\ldots,5\,, \tag{3.19}$$

where

$$
\begin{aligned}
M_{11}^2 &= g_1'^2\tilde{Q}_S^2 s^2 - \frac{\sigma^2 s^2}{2}\sin^2 2\theta + \sqrt{2}\sigma A_\sigma\varphi\sin 2\theta\\
&\quad+\left(\kappa\sigma\varphi^2 + \sqrt{2}\sigma\mu\varphi + 2\sigma\Lambda\right)\sin 2\theta + \frac{\lambda A_\lambda}{2\sqrt{2}s}v^2\cos\theta\sin 2\beta\\
&\quad-\frac{\lambda\sigma\varphi}{4s}v^2\sin\theta\sin 2\beta + \Delta_{11}\,,\\
M_{12}^2 &= M_{21}^2 = \frac{\sigma^2 s^2}{4}\sin 4\theta - \sqrt{2}\sigma A_\sigma\varphi\cos 2\theta\\
&\quad-\left(\kappa\sigma\varphi^2 + \sqrt{2}\sigma\mu\varphi + 2\sigma\Lambda\right)\cos 2\theta + \frac{\lambda A_\lambda}{2\sqrt{2}s}v^2\sin\theta\sin 2\beta\\
&\quad+\frac{\lambda\sigma\varphi}{4s}v^2\cos\theta\sin 2\beta + \Delta_{12}\,,
\end{aligned}
$$

$$
\begin{aligned}
M_{13}^2 &= M_{31}^2 = \sigma^2\varphi s\cos 2\theta - \frac{\lambda\sigma}{4}v^2\sin\theta\sin 2\beta + \Delta_{13}\,,\\
M_{14}^2 &= M_{41}^2 = \frac{g_1'^2}{2}\tilde{Q}_S(\tilde{Q}_{H_u} - \tilde{Q}_{H_d})sv\sin 2\beta - \frac{\lambda A_\lambda}{\sqrt{2}}v\cos\theta\cos 2\beta\\
&\quad -\frac{\lambda\sigma}{2}\varphi v\sin\theta\cos 2\beta + \Delta_{14}\,,\\
M_{15}^2 &= M_{51}^2 = g_1'^2\tilde{Q}_S(\tilde{Q}_{H_d}\cos^2\beta + \tilde{Q}_{H_u}\sin^2\beta)sv - \frac{\lambda A_\lambda}{\sqrt{2}}v\cos\theta\sin 2\beta\,,\\
&\quad +\lambda^2 vs\cos^2\theta - \frac{\lambda\sigma}{2}\varphi v\sin\theta\sin 2\beta + \Delta_{15}\,,\\
M_{22}^2 &= \frac{\sigma^2 s^2}{2}\sin^2 2\theta + \frac{\sqrt{2}\sigma A_\sigma\varphi}{\sin 2\theta}\cos^2 2\theta + \left(\kappa\sigma\varphi^2 + \sqrt{2}\sigma\mu\varphi + 2\sigma\Lambda\right)\frac{\cos^2 2\theta}{\sin 2\theta}\\
&\quad +\frac{\lambda A_\lambda v^2}{2\sqrt{2}s\cos\theta}\sin^2\theta\sin 2\beta - \frac{\lambda\sigma\varphi v^2}{4s\sin\theta}\cos^2\theta\sin 2\beta + \Delta_{22}\,,\\
M_{23}^2 &= M_{32}^2 = -\frac{\sigma A_\sigma}{\sqrt{2}}s + \sigma^2\varphi s\sin 2\theta - \sigma s\left(\kappa\varphi + \frac{\mu}{\sqrt{2}}\right)\\
&\quad +\frac{\lambda\sigma}{4}v^2\cos\theta\sin 2\beta + \Delta_{23}\,,\\
M_{24}^2 &= M_{42}^2 = \left(-\frac{\lambda A_\lambda}{\sqrt{2}}v\sin\theta + \frac{\lambda\sigma}{2}\varphi v\cos\theta\right)\cos 2\beta + \Delta_{24}\,,\\
M_{25}^2 &= M_{52}^2 = \frac{\lambda^2}{2}sv\sin 2\theta + \left(-\frac{\lambda A_\lambda}{\sqrt{2}}v\sin\theta + \frac{\lambda\sigma}{2}\varphi v\cos\theta\right)\sin 2\beta + \Delta_{25}\,,\\
M_{33}^2 &= \frac{\sigma A_\sigma s^2}{2\sqrt{2}\varphi}\sin 2\theta - \sqrt{2}\xi\frac{\Lambda}{\varphi} + \frac{\kappa A_\kappa}{\sqrt{2}}\varphi + \mu\left(\frac{\sigma s^2}{2\sqrt{2}\varphi}\sin 2\theta + 3\frac{\kappa\varphi}{\sqrt{2}} - \frac{\sqrt{2}\Lambda}{\varphi}\right)\\
&\quad +2\kappa^2\varphi^2 - \frac{\lambda\sigma s}{4\varphi}v^2\sin\theta\sin 2\beta + \Delta_{33}\,,\\
M_{34}^2 &= M_{43}^2 = \frac{\lambda\sigma}{2}sv\sin\theta\cos 2\beta + \Delta_{34}\,,\\
M_{35}^2 &= M_{53}^2 = \frac{\lambda\sigma}{2}sv\sin\theta\sin 2\beta + \Delta_{35}\,,\\
M_{44}^2 &= \frac{\sqrt{2}\lambda s}{\sin 2\beta}\left(A_\lambda\cos\theta - \frac{\sigma\varphi}{\sqrt{2}}\sin\theta\right) + \left(\frac{\bar{g}^2}{4} - \frac{\lambda^2}{2}\right)v^2\sin^2 2\beta\\
&\quad +\frac{g_1'^2}{4}(\tilde{Q}_{H_u} - \tilde{Q}_{H_d})^2 v^2\sin^2 2\beta + \Delta_{44}\,,\\
M_{45}^2 &= M_{54}^2 = \left(\frac{\lambda^2}{4} - \frac{\bar{g}^2}{8}\right)v^2\sin 4\beta + \frac{g_1'^2}{2}v^2(\tilde{Q}_{H_u} - \tilde{Q}_{H_d})\times\\
&\quad \times(\tilde{Q}_{H_d}\cos^2\beta + \tilde{Q}_{H_u}\sin^2\beta)\sin 2\beta + \Delta_{45}\,,\\
M_{55}^2 &= \frac{\lambda^2}{2}v^2\sin^2 2\beta + \frac{\bar{g}^2}{4}v^2\cos^2 2\beta + g_1'^2 v^2(\tilde{Q}_{H_d}\cos^2\beta + \tilde{Q}_{H_u}\sin^2\beta)^2 + \Delta_{55}\,.
\end{aligned}
\tag{3.20}
$$

In eq. (3.20) the Δ_{ij} denote the loop corrections. The components of the CP-even Higgs basis, eq. (3.18), are related to the CP-even Higgs mass eigenstates h_i $(i = 1, \ldots, 5)$ by virtue of a unitary transformation,

$$\begin{pmatrix} S_1 \\ S_2 \\ S_3 \\ S_4 \\ S_5 \end{pmatrix} = \tilde{U} \begin{pmatrix} h_1 \\ h_2 \\ h_3 \\ h_4 \\ h_5 \end{pmatrix}, \tag{3.21}$$

where again the CP-even Higgs eigenstates are labeled according to increasing absolute value of mass, with h_1 being the lightest CP-even Higgs state and h_5 the heaviest.

If all SUSY breaking parameters as well as $\lambda s \sim \sigma s \sim \sigma\varphi \sim M_S$ are considerably larger than the EW scale, all masses of the CP-even Higgs states except for the lightest Higgs boson mass are determined by the SUSY breaking scale M_S. Because the minimal eigenvalue of the mass matrix, eqs. (3.19)–(3.20), is always less than its smallest diagonal element the lightest Higgs state in the CP-even sector (approximately S_5) remains always light irrespective of the SUSY breaking scale, i.e. $m^2_{h_1} \lesssim M^2_{55}$ like in the MSSM and NMSSM. In the interactions with other SM particles this state manifests itself as a SM-like Higgs boson if $M_S \gg M_Z$.

In the limit where $\lambda \sim \sigma \to 0$, the off-diagonal tree-level entries M^2_{24}, M^2_{25}, M^2_{34} and M^2_{35} of the mass matrix, eqs. (3.19)–(3.20), become negligibly small. At the same time, according to eq. (1.3), the diagonal entry M^2_{11} that is set by the mass of the Z' boson tends to be substantially larger than M^2_S, i.e. $M^2_{11} \simeq M^2_{Z'} \sim M^2_S/\sigma^2$, whereas $\cos 2\theta$ almost vanishes in this case. Indeed, combining the first and the second equations for the extrema eq. (3.3) one obtains the following tree-level expression for $\cos 2\theta$

$$\cos 2\theta \simeq \frac{m^2_{\overline{S}} - m^2_S}{m^2_{\overline{S}} + m^2_S + \sigma^2\varphi^2 + g'^2_1\tilde{Q}^2_S s^2} \sim \frac{M^2_S}{M^2_{Z'}} \sim \sigma^2\,, \tag{3.22}$$

which becomes vanishingly small when $M_{Z'} \gg M_S$ and/or when $\sigma \sim \lambda \to 0$. In this case the hierarchical structure of the mass matrix, eqs. (3.19)–(3.20), implies that the mass of the Z' boson and the mass of the heaviest CP-even Higgs particle associated with S_1 are almost degenerate. Thus the heaviest CP-even Higgs state can be integrated out. The mass of another CP-even state that is predominantly S_4, the mass of the CP-odd state that corresponds to P_1, and the masses of the charged Higgs states are also almost degenerate in this limit. Assuming that the Higgs state that is mainly S_4 is the second heaviest CP-even Higgs state, neglecting all terms that are proportional to the global U(1)-violating couplings (κ, μ and Λ) and setting $\cos 2\theta = 0$, one obtains the following approximate analytic expressions for the tree-level masses of the three lightest CP-even Higgs bosons,

$$\begin{aligned} m^2_{h_{3,2}} &\simeq \frac{\sigma^2 s^2}{4}\left[1 + \frac{A_\sigma}{\sqrt{2}\sigma\varphi} \pm \left|1 - \frac{A_\sigma}{\sqrt{2}\sigma\varphi}\right| \sqrt{1 + 16\,\frac{\varphi^2}{s^2}}\right], \\ m^2_{h_1} &\simeq \frac{\bar{g}^2}{4} v^2 \cos^2 2\beta \simeq M^2_Z \cos^2 2\beta\,. \end{aligned} \tag{3.23}$$

Note that in the scenario under consideration the tree-level expression for the SM-like Higgs mass $m^2_{h_1}$ is essentially the same as in the MSSM.

4 Non-standard Higgs decays

We now focus on that region of the parameter space that corresponds to the approximate global U(1) symmetry mentioned above, where we have a light pseudoscalar. Since our primary concern in our numerical investigation is to study non-standard Higgs decays and we do not assume any breaking pattern among the soft masses, most of the sfermion masses do not play a significant role. We therefore choose the SUSY breaking parameters that control the masses of the sfermions to be well above the TeV scale thus comfortably evading limits set by the LHC and decoupling them from the spectrum. We do, however, adjust the stop mass parameters to get a Higgs mass of $125 - 126\,\mathrm{GeV}$. The gaugino masses are also chosen to be heavy enough to give a Higgsino dark matter candidate and to evade the LHC limit on the gluino mass.

Additionally we assume that the SM singlet superfields $\overline{S}$, S and ϕ acquire very large VEVs inducing multi-TeV masses of the Z' boson. Our analysis of the Higgs sector in the previous section indicates that the Higgs spectrum in general has a very hierarchical structure when all SUSY breaking parameters are sufficiently large, i.e. above about 1 TeV. In this limit only the SM-like Higgs boson and the lightest CP-odd Higgs state associated with the spontaneously broken approximate global U(1) symmetry can be considerably lighter than 1 TeV.

The presence of a light pseudoscalar Higgs state in the particle spectrum can result in non-standard decays of the lightest CP-even Higgs boson in the U(1) extensions of the MSSM under consideration. Expanding the Higgs potential, eq. (3.1), about its physical minimum one obtains the trilinear coupling that describes the interaction of the lightest CP-even Higgs scalar with the lightest pseudoscalar Higgs states. At tree-level the corresponding part of the Lagrangian can be written as

$$\mathcal{L}_{h_1A_1A_1} = -G_{h_1A_1A_1} h_1 A_1 A_1\,, \tag{4.1}$$

with the trilinear Higgs couplings $G_{h_1A_1A_1}$ which is given by the rather lengthy expression

$$\begin{aligned}
G_{h_1A_1A_1} = {} & \tilde{U}_{51}\Bigg\{U_{11}^2\Bigg[\frac{\lambda^2}{4}v\cos^2\gamma(1+\cos^2 2\beta)+\frac{\lambda^2}{2}v\sin^2\gamma\cos^2\theta-\frac{\bar{g}^2}{8}v\cos^2\gamma\cos^2 2\beta \\
& +\frac{1}{2}\left(\frac{\lambda A_\lambda}{\sqrt{2}}\cos\theta-\frac{\lambda\sigma}{2}\varphi\sin\theta\right)\sin 2\gamma+\frac{g_1'^2}{2}v\left(\tilde{Q}_{H_d}\cos^2\beta+\tilde{Q}_{H_u}\sin^2\beta\right)\times \\
& \times\left(\tilde{Q}_{H_d}\sin^2\beta\cos^2\gamma+\tilde{Q}_{H_u}\cos^2\beta\cos^2\gamma+\tilde{Q}_S\sin^2\gamma\cos 2\theta\right)\Bigg] \\
& +U_{11}U_{21}\Bigg[\frac{\lambda^2}{2}v\sin 2\theta\sin\gamma+g_1'^2\tilde{Q}_S v\left(\tilde{Q}_{H_d}\cos^2\beta+\tilde{Q}_{H_u}\sin^2\beta\right)\sin\gamma\sin 2\theta \\
& +\left(\frac{\lambda A_\lambda}{\sqrt{2}}\sin\theta+\frac{\lambda\sigma}{2}\varphi\cos\theta\right)\cos\gamma\Bigg]+\frac{\lambda\sigma}{2}\sin\theta\,U_{11}U_{31}(s\cos\gamma+v\sin 2\beta\sin\gamma) \\
& +U_{21}^2\left[\frac{\lambda^2}{2}v\sin^2\theta-\frac{g_1'^2}{2}\tilde{Q}_S v\cos 2\theta\left(\tilde{Q}_{H_d}\cos^2\beta+\tilde{Q}_{H_u}\sin^2\beta\right)\right] \\
& -\frac{\lambda\sigma}{2}v\sin 2\beta\cos\theta\,U_{21}U_{31}\Bigg\}+\tilde{U}_{41}\Bigg\{U_{11}^2\Bigg[\left(-\frac{\lambda^2}{8}+\frac{\bar{g}^2}{16}\right)v\cos^2\gamma\sin 4\beta \\
& +\frac{g_1'^2}{4}v\sin 2\beta(\tilde{Q}_{H_u}-\tilde{Q}_{H_d})\Big(\tilde{Q}_{H_d}\sin^2\beta\cos^2\gamma+\tilde{Q}_{H_u}\cos^2\beta\cos^2\gamma
\end{aligned}$$

$$
\begin{aligned}
&+\tilde{Q}_S \sin^2\gamma\cos 2\theta\Big)\Big] + \frac{g_1'^2}{2}\tilde{Q}_S(\tilde{Q}_{H_u}-\tilde{Q}_{H_d})v\sin 2\beta\sin\gamma\sin 2\theta\, U_{11}U_{21}\\
&+\frac{\lambda\sigma}{2}v\cos 2\beta\sin\gamma\sin\theta\, U_{11}U_{31} - \frac{g_1'^2}{4}\tilde{Q}_S(\tilde{Q}_{H_u}-\tilde{Q}_{H_d})v\sin 2\beta\cos 2\theta\, U_{21}^2\\
&-\frac{\lambda\sigma}{2}v\cos 2\beta\cos\theta\, U_{21}U_{31}\Big\} + \tilde{U}_{31}\Big\{U_{11}^2\Big[-\frac{\lambda\sigma}{4}s\sin\theta\sin 2\beta\cos^2\gamma + \frac{\sigma^2}{2}\varphi\sin^2\gamma\\
&-\frac{\lambda\sigma}{4}v\sin 2\gamma\sin\theta - \frac{\sigma}{2}\sin 2\theta\sin^2\gamma\left(\frac{A_\sigma}{\sqrt{2}}+\kappa\varphi+\frac{\mu}{\sqrt{2}}\right)\Big]\\
&+U_{11}U_{21}\left[\frac{\lambda\sigma}{2}v\cos\theta\cos\gamma + \sigma\left(\frac{A_\sigma}{\sqrt{2}}+\kappa\varphi+\frac{\mu}{\sqrt{2}}\right)\sin\gamma\cos 2\theta\right]\\
&+U_{21}^2\left[\frac{\sigma^2}{2}\varphi + \frac{\sigma}{2}\left(\frac{A_\sigma}{\sqrt{2}}+\kappa\varphi+\frac{\mu}{\sqrt{2}}\right)\sin 2\theta\right] - \sigma\kappa s\, U_{21}U_{31}\\
&+\kappa U_{31}^2\left(\kappa\varphi+\frac{\mu}{\sqrt{2}}-\frac{A_\kappa}{\sqrt{2}}\right)\Big\} + \tilde{U}_{21}\Big\{U_{11}^2\Big[-\frac{\lambda\sigma}{4}\varphi\sin 2\beta\cos^2\gamma\cos\theta\\
&+\frac{\lambda^2}{4}s\cos^2\gamma\sin 2\theta + \frac{\lambda A_\lambda}{2\sqrt{2}}\sin 2\beta\cos^2\gamma\sin\theta + \frac{\sigma^2}{4}s\sin^2\gamma\sin 2\theta\Big] \qquad\qquad (4.2)\\
&+U_{11}U_{31}\left[\frac{\lambda\sigma}{2}v\cos\gamma\cos\theta + \sigma\left(\frac{A_\sigma}{\sqrt{2}}-\kappa\varphi-\frac{\mu}{\sqrt{2}}\right)\sin\gamma\cos 2\theta\right]\\
&+\frac{\sigma^2}{4}s\sin 2\theta\, U_{21}^2 + \sigma\left(\frac{A_\sigma}{\sqrt{2}}-\kappa\varphi-\frac{\mu}{\sqrt{2}}\right)\sin 2\theta\, U_{21}U_{31} + \frac{\sigma}{2}\left(\sigma s\sin 2\theta+\kappa s\right)U_{31}^2\Big\}\\
&+\tilde{U}_{11}\Big\{U_{11}^2\Big[\frac{\lambda\sigma}{4}\varphi\sin 2\beta\cos^2\gamma\sin\theta + \frac{\lambda^2}{2}s\cos^2\gamma\cos^2\theta + \frac{\lambda A_\lambda}{2\sqrt{2}}\sin 2\beta\cos^2\gamma\cos\theta\\
&+\frac{g_1'^2}{2}\tilde{Q}_S s\left(\tilde{Q}_{H_d}\sin^2\beta\cos^2\gamma + \tilde{Q}_{H_u}\cos^2\beta\cos^2\gamma + \tilde{Q}_S\sin^2\gamma\cos 2\theta\right)\Big]\\
&+\left[-\frac{\lambda\sigma}{2}v\cos\gamma\sin\theta + \sigma\left(\kappa\varphi+\frac{\mu}{\sqrt{2}}-\frac{A_\sigma}{\sqrt{2}}\right)\sin\gamma\sin 2\theta\right]U_{11}U_{31}\\
&+\left(g_1'^2\tilde{Q}_S^2-\frac{\sigma^2}{2}\right)s\sin\gamma\sin 2\theta\, U_{11}U_{21} + \left[\frac{\sigma^2}{2}-\frac{g_1'^2}{2}\tilde{Q}_S^2\right]s\cos 2\theta\, U_{21}^2\\
&+\sigma\left(\frac{A_\sigma}{\sqrt{2}}-\kappa\varphi-\frac{\mu}{\sqrt{2}}\right)\cos 2\theta\, U_{21}U_{31} + \frac{\sigma^2}{2}s\cos 2\theta\, U_{31}^2\Big\}\,.
\end{aligned}
$$

If $m_{A_1} \lesssim 60\,\mathrm{GeV}$ then the CP-even Higgs boson with mass around 125 GeV can decay into a pair of two lightest pseudoscalar Higgs bosons A_1, through the interaction given in eq. (4.1). The corresponding partial decay width is given by

$$
\Gamma(h_1 \to A_1A_1) = \frac{G^2_{h_1A_1A_1}}{8\pi m_{h_1}}\sqrt{1-\frac{4m^2_{A_1}}{m^2_{h_1}}}\,. \tag{4.3}
$$

To compare the partial width of the non-standard Higgs decay of the SM-like Higgs state, eq. (4.3), with the Higgs decay rates into the SM particles, we specify a set of benchmark points (see table 4). For each benchmark scenario a code that is automatically generated by FlexibleSUSY [102] is used[7] (based on SARAH [103–107] and SOFTSUSY [108, 109]) to determine the spectrum of the masses. The complete one-loop self energies are included in the determination of all masses in the model, and leading two-loop contributions to the CP-even and CP-odd Higgs bosons from the NMSSM ($O(\alpha_t \alpha_s)$ and $O(\alpha_b \alpha_s)$) [110] and the MSSM ($O(\alpha_t^2)$, $O(\alpha_b \alpha_\tau)$, $O(\alpha_b^2)$, $O(\alpha_\tau^2)$ and $O(\alpha_t \alpha_b)$) [111–115] are included by using files provided by Pietro Slavich. The additional corrections that may arise due to our model are expected to be small, as either the new particles do not couple directly to the involved particles or their contributions are small due to suppressed couplings and/or large masses.

The couplings and branching ratios of the lightest CP-even Higgs state were also obtained by calling routines generated by FlexibleSUSY in a small extension of the automatically generated code from FlexibleSUSY. FlexibleSUSY uses SARAH-4.2.1 to derive analytical expressions, which is independent of the derivation used to obtain the expressions presented here. Therefore we were able to do an independent check of the mass matrices presented in the previous section and of the coupling $G_{h_1 A_1 A_1}$ given in eq. (4.2) by comparing our code from FlexibleSUSY numerically against an alternative Mathematica code based on those expressions.

Additionally, we cross-checked the thus obtained total width and branching ratios for the lightest CP-even Higgs boson against the ones obtained from the code HDECAY [116–119], respectively its extension eHDECAY [120, 121]. The Fortran program HDECAY, originally written for the calculation of the decay widths and branching ratios of the SM and the MSSM Higgs bosons, has been extended to allow for the possiblity to change the couplings of the SM Higgs boson by global modification factors. Using the modification factors, i.e. the ratios of the couplings of the lightest CP-even Higgs boson in our model to the SM particles with respect to the corresponding couplings of the SM Higgs boson of same mass, as inputs in HDECAY, we are able to compute the decay rates of h_1 into SM particle final states. Implementing in addition the partial decay width $h_1 \to A_1 A_1$ generated by FlexibleSUSY, HDECAY can be used to also calculate all branching ratios.[8] This procedure allows us to profit from the state-of-the-art QCD corrections implemented in HDECAY, which can be taken over to our model.[9] It should be noted, that in the loop-induced couplings to gluons and photons, respectively, the SUSY-related loops are not taken into account, however, as the option of applying coupling modifications in HDECAY only applies to the SM Higgs boson. As in our scenarios the sfermion and charged Higgs boson masses are heavy, the change should be marginal only. The branching ratios for the decay $h_1 \to A_1 A_1$ given for our benchmark scenarios in table 2, are the ones obtained by this procedure.

To simplify our analysis we set $B = \mu = \xi = 0$ and $\Lambda = 0$. This does not change the physics we are investigating here and still leaves us with one crucial Peccei-Quinn violating coupling κ. We also fix $\sigma = 0.1$ and $\tan\beta \simeq 10$. A large value of $\tan\beta$ allows us to maxime

[7] We used an adapted version of FlexibleSUSY-1.0.2 which contains updates that will appear in the new version FlexibleSUSY-1.0.3. The generated (and modified) code can be supplied on request.

[8] Note, that also the decay $h_1 \to Z A_1$ is in principle possible. In all benchmarks scenarios, however, this decay is kinematically closed or strongly suppressed.

[9] The electroweak corrections cannot be taken over. They are consistently included only for the SM part of the decay widths, cf. [120, 121].

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the tree-level mass value of the lightest CP-even Higgs boson, so that the experimental value of $\sim 125\,\text{GeV}$ can be obtained more easily. The small value for σ leads to VEVs that can be much heavier than the SUSY breaking scale, cf. eq. (1.3). Then, in order to find an appropriate set of benchmark points, we vary λ, κ, A_κ, A_λ, A_σ, A_t, $m^2_{Q_3}$, $m^2_{u^c_3}$, φ, s and θ. In all our benchmark scenarios m_S^2 and $m_{H_u}^2$ are negative which ensures that the Higgs fields S and H_u acquire VEVs that result in non-zero VEVs for the other Higgs fields H_d, $\overline{S}$ and φ. This should trigger the breaking of the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y \times \mathrm{U}(1)_N$ symmetry down to the $\mathrm{U}(1)_{\mathrm{em}}$. The soft scalar masses associated with the superpartners of the left-handed and right-handed components of the top quark and the mixing in the stop sector are chosen such that the SM-like Higgs state has a mass of approximately $125 - 126\,\text{GeV}$.

To construct benchmark scenarios which are consistent with cosmological observations, it is important to guarantee that they lead to relic densities $\Omega_{\mathrm{CDM}}h^2$ that are not larger than the result given by PLANCK [122]:

$$\Omega_{\mathrm{CDM}}h^2 = 0.1187 \pm 0.0017\,. \tag{4.4}$$

A theory predicting a greater relic density than the PLANCK result is basically ruled out, assuming standard pre-BBN cosmology. A theory that predicts less dark matter cannot be ruled out in the same way, but would require to have other contributions to the dark matter relic density. Since the dark matter density is inversely proportional to the annihilation cross section at the freeze-out temperature, this cross section has to be sufficiently large. In the E_6 inspired SUSY models considered here, the cold dark matter density is formed by the lightest neutralino states. At first glance the neutralino sector in these models is more complicated than the one in the MSSM. Indeed, it contains eight states which is twice as large as the one for the MSSM. However an analysis of the corresponding neutralino mass matrix, which is specified in the appendix A, indicates that this matrix given in the basis $(\tilde{H}^0_d,\ \tilde{H}^0_u,\ \tilde{W}_3,\ \tilde{B},\ \tilde{B}',\ \tilde{S}\cos\theta - \tilde{\overline{S}}\sin\theta,\ \tilde{S}\sin\theta + \tilde{\overline{S}}\cos\theta,\ \tilde{\phi})$ has a rather simple structure. The $\tilde{H}^0_d$, $\tilde{H}^0_u$, $\tilde{W}_3$, $\tilde{B}$ denote the fermion components of the corresponding Higgs doublet and gauge fields, $\tilde{B}'$ the gaugino related to the Z' vector superfield, and finally, $\tilde{S}, \tilde{\overline{S}}, \tilde{\phi}$ the fermion components of the SM singlet Higgs superfields $S, \overline{S}$ and ϕ. From eq. (A.5) it follows that the masses of two neutralino states, that are linear superpositions of $\tilde{B}'$ and $\tilde{S}\cos\theta - \tilde{\overline{S}}\sin\theta$, are controlled by the Z' boson mass. In the limit where λ is small and $M_{Z'} \gg M_S$ these states decouple from the rest of the sparticle spectrum. As can be seen from eq. (A.5), two other neutralino eigenstates are formed by the linear superpositions of $\tilde{S}\sin\theta + \tilde{\overline{S}}\cos\theta$ and $\tilde{\phi}$. In the situation where κ is much smaller than λ and σ, the masses of these eigenstates are approximately

$$m_{\chi^0_{5,6}} \simeq \frac{\sigma\varphi}{2\sqrt{2}}\left(1 \pm \sqrt{1 + 4\,\frac{s^2}{\varphi^2}}\right) \sim M_S\,. \tag{4.5}$$

When λ is small and $M_S \gg M_Z$ the mixing of these states with the MSSM-like neutralino (the superposition of $\tilde{H}^0_d$, $\tilde{H}^0_u, \tilde{W}_3$ and $\tilde{B}$) is also strongly suppressed. Thus, if the corresponding neutralino states are not the lightest ones, then they can be ignored in first approximation. This permits us to reduce the 8×8 matrix, eqs. (A.3)–(A.6), to a 4×4 mass matrix which is rather similar to the one in the MSSM. In the limit $M_S \gg M_Z$ the masses of the MSSM-like neutralino states, which are predominantly bino, wino and higgsino, are set by M_1, M_2 and $\lambda s \cos\theta/\sqrt{2}$, respectively.

The qualitative pattern of the neutralino spectrum discussed above reveals that for sufficiently small values of λ the lightest neutralino tends to be a higgsino dominated state. If this higgsino-like state is lighter than 1 TeV, then it leads to a cold dark matter density that is less than the observed value [13, 14]. Therefore, in all benchmark scenarios specified in table 4, the value of the coupling λ is chosen such that the lightest neutralino is predominantly higgsino with mass below 1 TeV.

In summary, we choose the following values for our benchmark scenarios:

The soft SUSY breaking left- and right-handed mass parameters:

$$\begin{aligned} m^2_{Q_{1,2}} = m^2_{u^c_{1,2}} = m^2_{L_{1,2,3}} = m^2_{e^c_{1,2,3}} = m^2_{L_4} = m^2_{\overline{L}_4} = 100\ \text{TeV}^2\,, \\ m^2_D = m^2_{\overline{D}} = m^2_{H^d_{1,2}} = m^2_{H^u_{1,2}} = 4\ \text{TeV}^2 \end{aligned} \tag{4.6}$$

The soft SUSY breaking gaugino mass parameters:

$$M_1 = 600\ \text{GeV}\,, \quad M_2 = 1.2\ \text{TeV}\,, \quad M_3 = 3.6\ \text{TeV} \tag{4.7}$$

The coupling values:

$$\mu = B = \xi = 0, \quad \Lambda = 0, \tag{4.8}$$

$$\kappa_{ij} = 0.5\,\delta_{ij} \tag{4.9}$$

$$f_{i\alpha} = \tilde{f}_{i\alpha} = g^D_{ij} = h^E_{ij} = 0 \tag{4.10}$$

$$\tilde{\sigma} = 0, \ \mu_L = 10\ \text{TeV} \tag{4.11}$$

The mixing angle: $\tan\beta^{\overline{\text{DR}}}(M_Z) = 10\,,$ (4.12)

where the masses in eq. (37) are associated with the soft scalar masses of the scalar components of superfields listed in eq. (7) and δ_{ij} in eq. (40) denotes the Kronecker δ. All other values are specified in table 4, with the exception of those trilinear couplings that are zero. The parameters are chosen such that $m_h \approx 125$ GeV. The SM parameters are chosen as

$$\begin{aligned} \alpha_{\text{em}}(M_Z) &= 1/127.916\,, & \alpha_s(M_Z) &= 0.1184\,, \\ M_Z &= 91.1876\ \text{GeV}\,, & M_W &= 80.404\ \text{GeV}\,, \\ m_t &= 173.18\ \text{GeV}\,, & m_b(m_b)^{\overline{\text{MS}}} &= 4.2\ \text{GeV}\,, \quad m_\tau = 1.777\,\text{GeV}\,. \end{aligned} \tag{4.13}$$

As in all benchmark scenarios we took care to choose m_{A_1} small enough that the decays of the SM-like Higgs state into a pair of the lightest pseudoscalar Higgs bosons are kinematically allowed, it is important to ensure that the lightest CP-odd Higgs boson could not have been detected in the collider experiments to date. In this context it is worth noting that in the case of a hierarchical structure of the Higgs spectrum, which is caused by a large SUSY breaking scale, the couplings of this pseudoscalar state to the SM particles are naturally suppressed. Indeed, as was pointed out in subsection 3.2, the lightest CP-odd Higgs boson is predominantly a superposition of the imaginary parts of the SM singlet fields $\overline{S}$, S and ϕ. These do not couple directly to the SM particles and, furthermore, the

channel	best fit value	$2 \times 1\sigma$ error
$VH \to Vbb$	0.97	± 1.06
$H \to \tau\tau$	1.02	± 0.7
$H \to \gamma\gamma$	1.14	± 0.4
$H \to WW$	0.78	± 0.34
$H \to ZZ$	1.11	± 0.46

Table 3. The combined ATLAS and CMS signal rates with errors for the $bb, \tau\tau, \gamma\gamma, WW$ and ZZ final states. Apart from the bb final state, where Higgs-strahlung VH ($V = W, Z$) is the production channel, they are based on the inclusive production cross section. Details can be found in refs. [123] and [124].

mixing between these singlet fields and the neutral components of the Higgs doublets is rather small in this case. As a consequence in all benchmark scenarios presented in tables 4 and 5 the absolute value of the relative coupling of the lightest pseudoscalar Higgs to the Z-boson and the SM-like Higgs state, $R_{ZA_1h_1}$, is always smaller than $10^{-3} - 10^{-4}$. All other couplings of the lightest CP-odd Higgs boson to the SM particles are also extremely small. This state therefore has escaped detection in past and present collider experiments.

Due to the hierarchical structure of the Higgs spectrum the lightest SM-like Higgs state has couplings close to the SM values. Its coupling values R_{XXh_1} to the SM particles $X = V, f_u, f_d$ ($V \equiv W, Z,\ f_u \equiv u, c, t,\ f_d \equiv d, s, b, e, \mu, \tau$) relative to the ones of the SM Higgs boson, are given in table 5 for the various benchmarks scenarios. The h_1 couplings deviate by at most 4% from the SM Higgs couplings. These coupling values feed into the production processes and the branching ratios of the SM-like state h_1. It has to be made sure that the μ-values, i.e. the ratios of h_1 production cross section times branching ratio normalized to the corresponding SM values, agree within the respective errors with the μ-values reported by the LHC experiments for the various final states. For definiteness we take here the values given in refs. [123] and [124]. We follow the procedure of ref. [125] and combine the signal rates and errors of the two experiments according to eq. (5) in [126]. We require the h_1 μ-values to be within 2 times the 1σ interval around the respective best fit value. The combined signal rates and errors are given in table 3. For the calculation of the μ-values we have assumed the dominant production cross section to be given by gluon fusion. Subsequently, we have approximated the ratios of the h_1 and the SM gluon fusion production cross sections by the ratio of their decay widths Γ_{gg} into a gluon pair. The gluon decay widths have been calculated with `HDECAY` in both cases, as outlined above. The program includes the higher order QCD corrections to this decay.[10] The approximation of the production cross section ratios by the decay width ratios is valid within about 10–20% depending on the scenario [9]. The μ-value for h_1 into the final state XX is hence given by

$$\mu_{XX}(h_1) \approx \frac{\Gamma_{gg}(h_1) BR(h_1 \to XX)}{\Gamma_{gg}(H^{\rm SM}) BR(H^{\rm SM} \to XX)} , \tag{4.14}$$

where $H^{\rm SM}$ denotes the SM Higgs boson with the same mass as h_1. The branching ratios

[10] For details and a recent discussion, see e.g. [120, 121, 127].

BR again have been obtained with HDECAY. For the $b\bar{b}$ final state, however, Higgs-strahlung Vh_1 ($V = Z, W$) is the production channel to be used, in accordance with the experiments. This cross section is given by the SM value multiplied with the squared coupling ratio R_{VVh_1} to gauge bosons. For the $b\bar{b}$ final state we hence have

$$\mu_{b\bar{b}}(h_1) \approx \frac{\sigma_{Vh_1} BR(h_1 \to b\bar{b})}{\sigma_{VH^{\rm SM}} BR(H^{\rm SM} \to b\bar{b})} = R^2_{VVh_1} \frac{BR(h_1 \to b\bar{b})}{BR(H^{\rm SM} \to b\bar{b})} \,. \tag{4.15}$$

Let us now turn to the discussion of the five benchmark scenarios BMA-BME, summarized in tables 4 and 5. Table 4 lists the parameter values defining the scenarios (in addition to those values common to all benchmarks, given in eqs. (4.6)–(4.12)) and the corresponding scalar Higgs masses m_{h_1}–m_{h_5} and the pseudoscalar ones m_{A_1}–m_{A_3}, with the SM-like Higgs boson given by the lightest scalar h_1. In table 5 we give for each benchmark scenario the coupling $G_{h_1A_1A_1}$, relevant for the non-standard decay $h_1 \to A_1A_1$, and the ratios R_{XXh_1} of the couplings of the SM-like h_1 with respect to the ones of the SM Higgs boson for the couplings to a pair of massive gauge bosons V, of up-type quarks f_u and of down-type quarks f_d. The table contains the μ-values in the LHC Higgs search final states and finally the partial decay width and the branching ratio for the non-standard decay of h_1 into a pair of lightest pseudoscalars, along with the h_1 total decay width. For all scenarios, the μ-values are within 2 times the 1σ error interval around the experimentally measured values.[11] As can be inferred from table 5, the branching ratios of the non-standard Higgs decays h_1 into A_1A_1 vary from $10^{-3}\,\%$ to 21% for the various benchmarks. The table shows, that their size depends rather strongly on the absolute value of the coupling κ. Reasonably large branching ratios ($\gtrsim 1\%$) of these Higgs decays can be obtained for κ values of $\mathcal{O}(0.01)$. Thus we obtain in benchmark scenario BMA with $\kappa = 0.03$ a branching ratio of $\sim 8\%$. In scenario BMC, where also $\kappa = 0.03$, the parameters are further optimized to get a large branching ratio, resulting in $BR(h_1 \to A_1A_1) \approx 0.21$. The scenario BMD has $\kappa = 0.0273$ and a non-standard branching ratio of 0.017. The branching ratio is smaller here, as by increasing A_λ, resulting in a larger charged Higgs mass, we are getting closer to the SM-limit, as can be seen from the coupling ratios to the SM particles, which are almost one. Therefore, the non-standard coupling $G_{h_1A_1A_1}$ is smaller compared to BMA and BMC and hence also the corresponding branching ratio. Both in scenario BMB and BME the κ value is chosen to be small, $\kappa = 0.001$. This results in $BR(h_1 \to A_1A_1) = 4.4 \cdot 10^{-5}$ in scenario BMB. In BME the VEV s has been increased and hence the mass of the Z' which is $M_{Z'} \approx 6\,$TeV here, resulting in an even smaller branching ratio. With h_1 non-standard branching ratios of $\mathcal{O}(10^{-5})$ and μ rates that are compatible with the SM, this scenario will not be quickly ruled out by the LHC.[12]

For the non-standard decays to take place, the pseudoscalar mass must be small enough. As follows from eqs. (3.15) such a small m_{A_1} can always be obtained by tun-

[11]The difference in the μ-values for the WW and ZZ final states arises from different branching ratios. Although the coupling modifications are the same for these final states, the branching ratios differ, as in eHDECAY electroweak corrections are included in the decay width in the term linear in the SM amplitude. For details, see [120, 121].

[12]The lightest CP-odd Higgs state that originates from decays of the SM-like Higgs boson, predominantly decays into either a pair of b-quarks or τ-leptons giving rise to four fermion final states.

	BMA	BMB	BMC	BMD	BME
λ	0.100	0.090	0.100	0.090	0.090
σ	0.1	0.1	0.1	0.1	0.1
κ	0.03	0.001	0.03	0.0273	0.001
A_λ [GeV]	600	2222	800	2222	4444
A_σ [GeV]	1200	1200	1400	1200	2400
A_κ [GeV]	1013	1000	1023	1026	2200
A_t [GeV]	-1186	-1171	-5944	-1170	-1163
$m^2_{Q_3}$ [GeV]^2	$3.0 \cdot 10^6$	$1.0 \cdot 10^6$	$3.0 \cdot 10^6$	$2.0 \cdot 10^6$	$1.0 \cdot 10^6$
$m^2_{u^c_3}$ [GeV]^2	$2.0 \cdot 10^8$	$1.0 \cdot 10^8$	$2.0 \cdot 10^8$	$8.0 \cdot 10^7$	$8.0 \cdot 10^7$
φ [TeV]	6	6	6	6	12
s [TeV]	8	8	8	8	16
$\tan\theta$	0.91	0.91	0.91	0.91	0.91
$M_1 = M_1'$ [GeV]	600	600	600	600	1200
$m_{\chi^0_1}$ [GeV]	420	376	419	377	761
m^2_S [GeV]^2	$-7.023 \cdot 10^5$	$-1.51 \cdot 10^6$	$-6.287 \cdot 10^5$	$-1.135 \cdot 10^6$	$-2.282 \cdot 10^6$
$m^2_{\bar{S}}$ [GeV]^2	$1.303 \cdot 10^6$	$1.918 \cdot 10^6$	$1.400 \cdot 10^6$	$1.669 \cdot 10^6$	$3.655 \cdot 10^6$
m^2_ϕ [GeV]^2	$6.145 \cdot 10^4$	$1.292 \cdot 10^5$	$1.353 \cdot 10^5$	$7.045 \cdot 10^4$	$5.219 \cdot 10^5$
$m^2_{H_d}$ [GeV]^2	$8.267 \cdot 10^5$	$7.132 \cdot 10^6$	$1.306 \cdot 10^6$	$7.023 \cdot 10^6$	$2.684 \cdot 10^7$
$m^2_{H_u}$ [GeV]^2	$-2.419 \cdot 10^6$	$-1.063 \cdot 10^6$	$-2.977 \cdot 10^6$	$-2.448 \cdot 10^5$	$-1.597 \cdot 10^5$
$m_{Z'}$ [TeV]	2.956	2.964	2.956	2.961	5.939
$m_{H^\pm}$ [GeV]	799	2550	1057	2550	5123
m_{A_1} [GeV]	35.37	33.01	51.53	18.97	28.41
m_{A_2} [GeV]	791	1206	1051	1168	2430
m_{A_3} [GeV]	1159	2547	1257	2548	5122
m_{h_1} [GeV]	126.156	125.76	125.881	125.699	126.225
m_{h_2} [GeV]	387	460	267	393	852
m_{h_3} [GeV]	791	795	1016	925	1569
m_{h_4} [GeV]	936	2547	1051	2548	5122
m_{h_5} [GeV]	3099	3093	3125	3104	6188

Table 4. Parameters defining the benchmark points BMA-BME, with the associated Higgs masses.

ing the parameter A_κ. The degree of tuning is illustrated in figure 1, where the region of the parameter space that leads to a sufficiently small m_{A_1} for non-standard Higgs decays, is shown in the κ–A_κ plane. One can see that with decreasing κ, the range of A_κ that results in a sufficiently small m_{A_1} becomes considerably wider. This corresponds to a smaller degree of fine-tuning required to obtain a sufficiently light pseudoscalar.[13] Therefore, in

[13]The branching ratios for this figure have been estimated using generated routines from `FlexibleSUSY`, which allows for a fast scan over the parameter range. We have checked for a few points that the branching ratios from this simple estimate agree with the ones from `HDECAY` well enough for our purposes. The

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	BMA	BMB	BMC	BMD	BME
$R_{f_u f_u h_1}$	-0.9974	-0.9997	0.9888	0.9994	-0.9999
$R_{f_d f_d h_1}$	-1.0412	-1.0036	1.0142	1.0032	-1.0013
R_{VVh_1}	-0.99789	-0.9997	0.9891	0.9995	-0.9999
$R_{ZA_1h_1}$	-1.608 $\cdot 10^{-4}$	1.626 $\cdot 10^{-4}$	6.634 $\cdot 10^{-7}$	-1.082 $\cdot 10^{-4}$	8.142 $\cdot 10^{-5}$
$G_{h_1 A_1 A_1}$ [GeV]	-1.2704	-0.0270	2.5782	0.5136	-0.0149
μ_{bb}	0.9684	1.0317	0.7883	0.9841	1.0314
$\mu_{\tau\tau}$	0.8731	0.9350	0.7865	0.9838	0.9347
μ_{ZZ}	0.7558	0.8741	0.7480	0.9765	0.8786
μ_{WW}	0.7247	0.8383	0.7483	0.9766	0.8422
$\mu_{\gamma\gamma}$	0.8180	0.9461	0.7482	0.9766	0.9499
$BR(h_1 \to A_1 A_1)$	0.0818	4.415 $\cdot 10^{-5}$	0.2078	0.0172	1.391 $\cdot 10^{-5}$
$\Gamma(h_1 \to A_1 A_1)$ [GeV]	4.215 $\cdot 10^{-4}$	1.967 $\cdot 10^{-7}$	1.206 $\cdot 10^{-3}$	7.960 $\cdot 10^{-5}$	6.270$\cdot 10^{-8}$
Γ_{tot} [GeV]	5.154 $\cdot 10^{-3}$	4.456 $\cdot 10^{-3}$	5.805$\cdot 10^{-3}$	4.618 $\cdot 10^{-3}$	4.451 $\cdot 10^{-3}$

Table 5. For the benchmarks BMA-BME, the couplings and coupling ratios of the SM-like scalar Higgs h_1 as well as the μ-values for the LHC Higgs search channels. The ratio $R_{ZA_1h_1}$ denotes the ratio of the coupling in the SUSY model under consideration with respect to the corresponding coupling in the MSSM. The last three lines show the branching ratio and partial width for the decay $h_1 \to A_1 A_1$ and the total width.

order to get a pseudoscalar with mass around $40 - 60\,$GeV for the κ values of $\mathcal{O}(0.03)$ of BMA, BMC and BMD, a fine-tuning of order 1% is needed. It turns out that BME is strongly fine-tuned ($\sim$ 1%) as well. This is because the SUSY breaking scale is doubled compared to the other scenarios and the value of A_κ is twice the value of A_κ in benchmark BMB. For scenario BMB, on the other hand, the fine-tuning is $\lesssim$ 10%.

Both figure 1 and tables 4 and 5 demonstrate that the branching ratios of the decays of the SM-like Higgs boson into a pair of pseudoscalar Higgs states become smaller when κ decreases. This is not a surprising result. Indeed, in the PQ symmetric limit, when the global PQ symmetry is only broken spontaneously, the coupling $G_{h_1 A_1 A_1}$ is set by (see for example [128])

$$G_{h_1 A_1 A_1} \simeq \frac{m_{h_1}^2}{2M_{PQ}}\,\varepsilon\,, \tag{4.16}$$

where M_{PQ} is the PQ symmetry breaking scale and ε represents a suppression associated with the mixing between the SM-like Higgs state and heavy CP-even Higgs states that induce the breakdown of the PQ symmetry. Equation (4.16) also determines the size of $G_{h_1 A_1 A_1}$ when the PQ violating couplings are rather small and m_{A_1} is naturally very light. A simple estimate using the values of the VEVs of the SM singlet fields given in table 4 indicates that the absolute value of $G_{h_1 A_1 A_1}$ is expected to be considerably smaller than 1 GeV when $\kappa \to 0$. Thus one can expect that small values of κ lead to small branching

deviations are due to approximation made for the h_1 decays into SM particles, which in our simple estimate were set equal to the ones of a SM Higgs boson with same mass. Further differences arise due to the inclusion of higher order corrections in `HDECAY`.

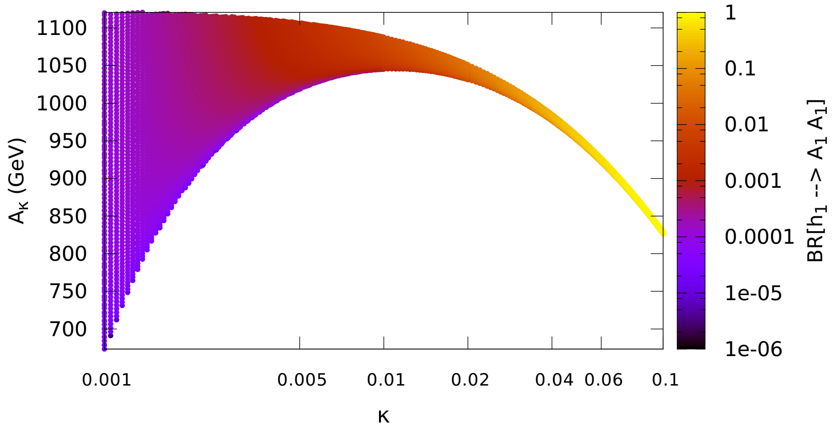

Figure 1. Colour contours of the branching ratio for $h_1 \to A_1 A_1$ in the κ–A_κ plane. An adapted version of `FlexibleSUSY` was used to calculate the mass spectra and the branching ratios. All other parameters are fixed to the values of BMA. For each value of κ there is a lower limit on A_κ where the the BR is zero because the pseudoscalar is too heavy and an upper limit above which there is a pseudoscalar tachyon.

ratios of the non-standard Higgs decays. When $\kappa \sim 0.001$ a pseudoscalar state with mass around $40 - 60\,\mathrm{GeV}$ can be obtained with very little fine-tuning. However in this case, the branching ratios of the non-standard Higgs decays are of the order of 10^{-4} and below. At the same time when the PQ violating parameters are sufficiently large so that one should fine-tune m_{A_1} for $h_1 \to A_1 A_1$ to be kinematically allowed then there can be additional contributions to $G_{h_1 A_1 A_1}$ from the explicit PQ violating terms which can make its absolute value even larger than $m_{h_1}^2/(2M_{PQ})$. This is why we can obtain a large partial width for this decay when $\kappa \gtrsim 0.01$.

We can hence summarize that in the model, that we introduced in order to reduce the fine-tuning, from which simple U(1)-extended SUSY models suffer, the non-standard decay of the SM-like Higgs state h_1 into a pair of pseudoscalars is possible. Small κ values lead naturally to light pseudoscalar masses without fine-tuning. In this case, however, the non-standard branching ratios are tiny, and it will be difficult to test these exotic decays at the LHC. For larger absolute values of κ, the corresponding branching ratio can reach the level of a few per cent up to $\mathcal{O}(20\%)$. This can be achieved, however, only at the price of fine-tuning, in order to get a low enough pseudoscalar mass in this case for the decay to be kinematically allowed.

5 Conclusions

In this paper we considered the non-standard decays of the SM-like Higgs state within U(1) extensions of the MSSM in which the extra U(1) gauge symmetry is broken by two VEVs of the SM singlet fields S and $\overline{S}$ with opposite U(1) charges. Because in these models S and $\overline{S}$ can acquire very large VEVs along the D-flat direction, the Z'-boson

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can naturally be substantially heavier than the SUSY particles. This allows us to satisfy experimental constraints and to alleviate the fine-tuning associated with the Z'-boson. Such U(1) extensions of the MSSM can possess an approximate global U(1) symmetry that gets spontaneously broken by the VEVs of S and $\overline{S}$ leading to a pseudo-Goldstone boson in the particle spectrum. If this pseudo-Goldstone state is considerably lighter than the lightest CP-even Higgs boson then it may give rise to the decays of the SM-like Higgs particle into a pair of pseudoscalars.

Here we studied such decays within well motivated SUSY extensions of the SM based on the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_N \times Z_2^M$ group which is a subgroup of E_6. The low-energy matter content of these E_6 inspired models includes three 27 representations of E_6, a pair of $SU(2)_L$ doublets L_4 and $\overline{L}_4$, a pair of the SM singlets S and $\overline{S}$ with opposite $U(1)_N$ charges, as well as an SM singlet ϕ that does not participate in the gauge interactions. To suppress flavor changing processes at tree-level and to forbid the most dangerous baryon and lepton number violating operators an extra $\tilde{Z}_2^H$ discrete symmetry is imposed. We analysed the spectrum of the CP-even and CP-odd Higgs bosons in these E_6 inspired models assuming that all SUSY breaking parameters are of the order of the TeV scale. Expanding the Higgs potential we obtained an analytical expression for the coupling of the lightest CP-even Higgs state h_1 to a pair of the lightest Higgs pseudoscalars A_1. The dependence of the branching ratio of the non-standard Higgs decay, $h_1 \to A_1 A_1$, on the parameters in these SUSY models was examined. For simplicity, we assumed that there is only one dimensionless coupling κ in the superpotential that explicitly violates the global U(1) symmetry. When κ vanishes the global U(1) symmetry is restored.

In order to illustrate the results of our analysis we specified a set of benchmark scenarios with the SM-like Higgs mass around $125-126$ GeV. To ensure that the obtained benchmark points are consistent with the measured value of the cold dark matter density we chose the parameters such that the lightest neutralino is mainly higgsino with a mass below 1 TeV. In this case the dark matter density tends to be smaller than its observed value. The results of our analysis indicate that the couplings of the lightest pseudoscalar Higgs to the SM particles are always quite small. As a consequence, although this pseudoscalar state can be rather light, it could escape detection at former and present collider experiments.

We argued that the branching ratio of the decays of the lightest CP-even Higgs boson into a pair of the lightest pseudoscalars depends rather strongly on the absolute value of the coupling κ. For absolute values of κ which are substantially smaller than 0.01 the branching ratio of the non-standard Higgs decay in a light pseudoscalar pair decreases considerably. Indeed, in the limit when κ goes to zero the global U(1) symmetry is spontaneously broken and the coupling of the lightest Higgs pseudoscalar to the SM-like Higgs becomes extremely suppressed. As a result the branching ratio of the non-standard Higgs decay tends to be negligibly small. Therefore although for $\kappa \sim 0.001$ the lightest Higgs pseudoscalar with a mass of $40-60$ GeV can be obtained without fine-tuning the branching ratio of the SM-like Higgs decays into a pair of the lightest CP-odd states is smaller than 10^{-4}. Decays with such small branching ratios will be difficult to be tested at the LHC.

When $\kappa \gtrsim 0.01$ the branching ratio of the non-standard Higgs decays can be larger than 1%. Nonetheless a fine tuning of at least 1% is required in this case to obtain a lightest pseudoscalar state with mass of $40-60\,\mathrm{GeV}$. After being produced from the decay of the lightest CP-even Higgs boson, the lightest CP-odd Higgs states sequentially decay into a pair of either b-quarks or τ-leptons. Thus these decays of the lightest CP-even Higgs boson result in four fermion final states.

We have found that with a modest fine tuning of A_κ one can obtain scenarios with $h_1 \to A_1 A_1$ branching ratios of $\simeq 1\%$ and acceptable values of dark matter relic density.

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A The neutralino mass matrix

After the breaking of the gauge symmetry in the E_6 inspired SUSY models under consideration all superpartners of the gauge and Higgs bosons get non-zero masses. Because the extra vector superfield associated with the Z' boson and the extra SM singlet Higgs superfields S, $\overline{S}$ and ϕ are electromagnetically neutral they do not contribute any extra particles to the chargino spectrum. As a consequence the chargino mass matrix and its eigenvalues remain almost the same as in the MSSM, i.e.

$$m^2_{\chi^\pm_{1,2}} = \frac{1}{2}\Big[M_2^2 + \mu_{\rm eff}^2 + 2M_W^2 \pm \sqrt{(M_2^2 + \mu_{\rm eff}^2 + 2M_W^2)^2 - 4(M_2\mu_{\rm eff} - M_W^2 \sin 2\beta)^2}\Big], \tag{A.1}$$

where M_2 is the $\mathrm{SU}(2)_L$ gaugino mass and

$$\mu_{\rm eff} = \frac{\lambda s \cos\theta}{\sqrt{2}}\,. \tag{A.2}$$

The non-observation of the lightest chargino at the collider experiments implies that $|M_2|$, $|\mu_{\rm eff}| \gtrsim 100\,\mathrm{GeV}$.

The neutralino sector of the SUSY models under consideration involves four extra neutralinos besides the four MSSM ones. One of them, $\tilde{B}'$, is an extra gaugino coming from the Z' vector superfield. Three other states are the fermion components $\tilde{S}$, $\tilde{\overline{S}}$ and $\tilde{\phi}$ of the SM singlet Higgs superfields S, $\overline{S}$ and ϕ. In the basis ($\tilde{H}^0_d$, $\tilde{H}^0_u$, $\tilde{W}_3$, $\tilde{B}$, $\tilde{B}'$, $\tilde{S}\cos\theta - \tilde{\overline{S}}\sin\theta$, $\tilde{S}\sin\theta + \tilde{\overline{S}}\cos\theta$, $\tilde{\phi}$) the neutralino mass matrix can be written as

$$M_{\tilde{\chi}^0} = \begin{pmatrix} A & C^T \\ C & B \end{pmatrix}, \tag{A.3}$$

with

$$A = \begin{pmatrix} 0 & -\frac{\lambda s\cos\theta}{\sqrt{2}} & \frac{gv}{2}\cos\beta & -\frac{g'v}{2}\cos\beta \\ -\frac{\lambda s\cos\theta}{\sqrt{2}} & 0 & -\frac{gv}{2}\sin\beta & \frac{g'v}{2}\sin\beta \\ \frac{gv}{2}\cos\beta & -\frac{gv}{2}\sin\beta & M_2 & 0 \\ -\frac{g'v}{2}\cos\beta & \frac{g'v}{2}\sin\beta & 0 & M_1 \end{pmatrix}, \tag{A.4}$$

$$B = \begin{pmatrix} M_1' & g_1'\tilde{Q}_S s & 0 & 0 \\ g_1'\tilde{Q}_S s & \frac{\sigma\varphi}{\sqrt{2}}\sin 2\theta & -\frac{\sigma\varphi}{\sqrt{2}}\cos 2\theta & 0 \\ 0 & -\frac{\sigma\varphi}{\sqrt{2}}\cos 2\theta & -\frac{\sigma\varphi}{\sqrt{2}}\sin 2\theta & -\frac{\sigma s}{\sqrt{2}} \\ 0 & 0 & -\frac{\sigma s}{\sqrt{2}} & \sqrt{2}\kappa\varphi + \mu \end{pmatrix}, \tag{A.5}$$

$$C = \begin{pmatrix} \tilde{Q}_{H_d} g_1' v\cos\beta & \tilde{Q}_{H_u} g_1' v\sin\beta & 0 & 0 \\ -\frac{\lambda v}{\sqrt{2}}\sin\beta\cos\theta & -\frac{\lambda v}{\sqrt{2}}\cos\beta\cos\theta & 0 & 0 \\ -\frac{\lambda v}{\sqrt{2}}\sin\beta\sin\theta & -\frac{\lambda v}{\sqrt{2}}\cos\beta\sin\theta & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \tag{A.6}$$

where M_1, M_2 and M_1' are the soft SUSY breaking gaugino masses for $\tilde{B}$, $\tilde{W}_3$ and $\tilde{B}'$, respectively. Here we neglect the Abelian gaugino mass mixing M_{11} between $\tilde{B}$ and $\tilde{B}'$. The top-left 4×4 block of the mass matrix in eq. (A.3) contains the neutralino mass matrix of the MSSM where the parameter μ is replaced by $\mu_{\rm eff}$. The lower right 4×4 submatrix represents extra neutralino states in this SUSY model.

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