Nonstandard Higgs Decays and Dark Matter in the E₆SSM

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We study the decays of the lightest Higgs boson within the exceptional supersymmetric (SUSY) standard model (E₆SSM). The E₆SSM predicts three families of Higgs–like doublets plus three SM singlets that carry $U(1)_N$ charges. One family of Higgs–like doublets and one SM singlet develop vacuum expectation values. The fermionic partners of other Higgs–like fields and SM singlets form inert neutralino and chargino states. Two lightest inert neutralinos tend to be the lightest and next-to-lightest SUSY particles (LSP and NLSP). The considered model can account for the dark matter relic abundance if the lightest inert neutralino has mass close to half the $Z$ mass. In this case the usual SM-like Higgs boson decays more than 95% of the time into either LSPs or NLSPs. As a result the decays of the lightest Higgs boson into $l^±l^± + X$ might play an essential role in the Higgs searches. This scenario also predicts other light inert chargino and neutralino states below 200 GeV and large LSP direct detection cross-sections which is on the edge of observability of XENON100.

1. Introduction

Confirming the Higgs mechanism as the underlying principle of electroweak symmetry breaking is one of the main goals of upcoming accelerators. The strategy for Higgs boson searches depends on the production mechanism and on the decay branching fractions of Higgs to different channels. Physics beyond the Standard Model may lead to the modification of the Higgs signals. In particular, there exist several extensions of the Standard Model in which the Higgs boson can decay with a substantial branching fraction into particles which can not be directly detected. The presence of invisible decays modifies considerably Higgs boson searches, making Higgs discovery much more difficult. If the Higgs is mainly invisible, then the visible branching ratios will be dramatically reduced, preventing detection in the much studied channels at the LHC and the Tevatron.

At $e^+e^-$ colliders, the problems related to the observation of the invisible Higgs are less severe since it can be tagged through the recoiling $Z$. As a result the LEP II collaborations exclude invisible Higgs masses up to 114.4 GeV [1]. On the other hand, Higgs searches at hadron colliders are more difficult in the presence of such invisible decays. Previous studies have analysed $ZH$ and $WH$ associated production [2]–[5] as well as $t\bar{t}H$ production [6]–[7] and $tWV$ ($b\bar{b}VV$) production [8] as promising channels.

Here we consider the exotic decays of the lightest Higgs boson and associated novel collider signatures within the Exceptional Supersymmetric Standard Model (E₆SSM). The E₆SSM is based on the $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$ gauge group which is a subgroup of $E_6$ [9]–[10]. The additional low energy $U(1)_N$ is a linear superposition of $U(1)_X$ and $U(1)_\psi$, i.e.

$$U(1)_N = \frac{1}{4} U(1)_X + \frac{\sqrt{5}}{4} U(1)_\psi.$$  

(1)

The two anomaly–free $U(1)_\psi$ and $U(1)_X$ symmetries are defined by: $E_6 \rightarrow SO(10) \times U(1)_\psi$, $SO(10) \rightarrow SU(5) \times U(1)_X$. The extra $U(1)_N$ gauge symmetry is defined such that right–handed neutrinos do not participate in the gauge interactions. Since right–handed neutrinos have zero charges they can 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 [11].

To ensure anomaly cancellation the particle content of the E₆SSM is extended to include three complete fundamental 27 representations of $E_6$. Each 27 multiplet contains a SM family of quarks and leptons, right–handed neutrino $N^{i}_\nu$, SM-type singlet fields $S_i$ which carry non-zero $U(1)_N$ charge, a pair of $SU(2)_W$–doublets $H^d_i$ and $H^u_i$ and a pair of colour triplets of exotic quarks $D_i$ and $D_i$ which can be either diquarks (Model I) or leptoquarks (Model II) [9]–[10]. The $S_i$, $H^d_i$ and $H^u_i$ form either Higgs or inert Higgs multiplets. In addition to the complete 27 multiplets the low energy particle spectrum of the E₆SSM is supplemented by $SU(2)_W$ doublet $L_{d_4}$ and anti-doublet $\bar{L}_{d_4}$ states from extra 27' and 27" to preserve gauge coupling unification. The analysis performed in [12] shows that the unification of gauge couplings in the E₆SSM can be achieved for any
phenomenologically acceptable value of $\alpha_3(M_Z)$ consistent with the measured low energy central value. The presence of a $Z'$ boson and of exotic quarks predicted by the $E_6$SSM provides spectacular new physics signals at the LHC which were discussed in [9–10, 13–14]. Recently the particle spectrum and collider signatures associated with it were studied within the constrained version of the $E_6$SSM [13–18].

The superpotential in the $E_6$ inspired models involves many new Yukawa couplings that induce non–diagonal flavour transitions. To suppress these effects in the $E_6$SSM an approximate $Z_2^H$ symmetry is imposed. Under this symmetry all superfields except one pair of $H_i^d$ and $H_i^u$ (say $H_d = H_2^d$ and $H_u = H_3^d$) and one SM-type singlet field ($S = S_3$) are odd. The $Z_2^H$ symmetry reduces the structure of the Yukawa interactions to

$$W_{E_6SSM} \simeq \lambda \hat{S} (\hat{H}_u \hat{H}_d) + \lambda_{\alpha \beta} \hat{S} (\hat{H}_u^{\alpha} \hat{H}_d^{\beta}) + f_{\alpha \beta} \hat{S} (\hat{H}_d \hat{H}_u^{\beta}) + f_{\alpha \beta} \hat{S} (\hat{D}_i \hat{D}_j) + \frac{1}{2} M_{ij} \tilde{N}_i \tilde{N}_j + \mu (\hat{L}_4 \hat{L}_4) + h_{ij}^E (\hat{H}_d \hat{L}_4) \tilde{e}_j^c + h_{ij}^N (\hat{H}_u \hat{L}_4) \tilde{N}_j^c ,$$

where $\alpha, \beta = 1, 2$ and $i, j = 1, 2, 3$. The $SU(2)_W$ doublets $\hat{H}_d$ and $\hat{H}_d$ and SM-type singlet field $\hat{S}$, that are even under the $Z_2^H$ symmetry, play the role of Higgs fields. At the physical vacuum they develop vacuum expectation values (VEVs)

$$\langle H_d \rangle = \frac{1}{\sqrt{2}} \left( \begin{array}{c} v_1 \\ 0 \end{array} \right), \quad \langle H_u \rangle = \frac{1}{\sqrt{2}} \left( \begin{array}{c} 0 \\ v_2 \end{array} \right), \quad \langle S \rangle = \frac{s}{\sqrt{2}}$$

generating the masses of the quarks and leptons. Instead of $v_1$ and $v_2$ it is more convenient to use $\tan \beta = v_2/v_1$ and $v = \sqrt{v_1^2 + v_2^2} = 246$ GeV. The VEV of the SM-type singlet field, $s$, breaks the extra $U(1)_Y$ symmetry thereby providing an effective $\mu$ term as well as the necessary exotic fermion masses and also inducing that of the $Z'$ boson. In the $E_6$SSM the Higgs spectrum contains one pseudoscalar, two charged and three CP–even states. In the leading two–loop approximation the mass of the lightest CP–even Higgs boson does not exceed $150 – 155$ GeV [9].

Although $Z_2^H$ eliminates any problems related with baryon number violation and non–diagonal flavour transitions it also forbids all Yukawa interactions that would allow the exotic quarks to decay. Since models with stable charged exotic particles are ruled out by various experiments the $Z_2^H$ symmetry can only be an approximate one. From here on we assume that $Z_2^H$ symmetry violating couplings are small and can be neglected in our analysis. This assumption can be justified if we take into account that the $Z_2^H$ symmetry violating operators may give an appreciable contribution to the amplitude of $K^0 – \bar{K}^0$ oscillations and give rise to new muon decay channels like $\mu \rightarrow e^- e^+ e^-$. In order to suppress processes with non–diagonal flavour transitions the Yukawa couplings of the exotic particles to the quarks and leptons of the first two generations should be smaller than $10^{-3} – 10^{-4}$. Such small $Z_2^H$ symmetry violating couplings can be ignored in the first approximation.

2. Masses and couplings of the lightest inert neutralinos

When $Z_2^H$ symmetry violating couplings tend to zero only $H_u$, $H_d$ and $S$ acquire non-zero VEVs. In this approximation the charged components of the inert Higgsinos ($\tilde{H}_2^{+}, \tilde{H}_1^{+}, \tilde{H}_2^{-}, \tilde{H}_1^{-}$) and ordinary chargino states do not mix. The neutral components of the inert Higgsinos ($\tilde{H}_1^{0}, \tilde{H}_2^{0}, \tilde{H}_1^{0}, \tilde{H}_2^{0}$) and inert singlinos ($\tilde{S}_1, \tilde{S}_2$) also do not mix with the ordinary neutralino states. Moreover if $Z_2^H$ symmetry was exact then both the lightest state in the ordinary neutralino sector and the lightest inert neutralino would be absolutely stable. Therefore, although $Z_2^H$ symmetry violating couplings are expected to be rather small, we shall assume that they are large enough to allow either the lightest neutralino state or the lightest inert neutralino to decay within a reasonable time.

In the field basis ($\tilde{H}_2^{0}, \tilde{H}_1^{0}, \tilde{S}_2, \tilde{H}_1^{+}, \tilde{H}_1^{0}, \tilde{S}_1$) the mass matrix of the inert neutralino sector takes a form

$$M_{IN} = \begin{pmatrix} A_{22} & A_{21} \\ A_{12} & A_{11} \end{pmatrix}, \quad A_{\alpha \beta} = -\frac{1}{\sqrt{2}} \begin{pmatrix} 0 & \lambda_{\alpha \beta} \sin \beta \\ \lambda_{\beta \alpha} \sin \beta & 0 \end{pmatrix} \begin{pmatrix} \tilde{f}_{\alpha \beta} v \sin \beta \\ f_{\beta \alpha} v \cos \beta \end{pmatrix} ,$$

(4)
where $A_{\alpha\beta}$ are $3 \times 3$ sub-matrices and $A_{12} = A_{21}^T$. In the basis of inert chargino interaction states ($H_2^+, H_1^+, \tilde{H}_2^-, \tilde{H}_1^-$) the corresponding mass matrix can be written as

$$M_{IC} = \begin{pmatrix} 0 & C^T \\ C & 0 \end{pmatrix}, \quad C_{\alpha\beta} = \frac{1}{\sqrt{2}} \lambda_{\alpha\beta} s,$$

(5)

where $C_{\alpha\beta}$ are $2 \times 2$ sub-matrices.

In our analysis we require the validity of perturbation theory up to the GUT scale that constrains the allowed range of all Yukawa couplings. We also choose $s$ and $\lambda_{\alpha\beta}$ so that the masses of all inert chargino states are larger than 100 GeV and $Z'$ boson is relatively heavy ($M_{Z'} \gtrsim 865$ GeV). The restrictions specified above set very strong limits on the masses of the lightest inert neutralinos. In particular, our numerical analysis indicates that the lightest and second lightest inert neutralinos ($\chi_1^0$ and $\chi_2^0$) are typically lighter than $60 - 65$ GeV [19, 20]. Therefore the lightest inert neutralino tends to be the lightest SUSY particle in the spectrum and can play the role of dark matter. The neutralinos $\chi_0^0$ and $\chi_2^0$ are predominantly inert singlinos. Their couplings to the $Z$-boson can be rather small so that such inert neutralinos would remain undetected at LEP.

In order to clarify the results of our numerical analysis, it is useful to consider a simple scenario when

$$\lambda_{\alpha\beta} = \lambda_0 \delta_{\alpha\beta}, \quad f_{\alpha\beta} = f_0 \delta_{\alpha\beta}, \quad \tilde{f}_{\alpha\beta} = \tilde{f}_0 \delta_{\alpha\beta}.$$

In this case the mass matrix of inert neutralinos reduces to the block diagonal form. In the limit where $f_0 = \tilde{f}_0$ one can easily prove using the method proposed in [21] that there are theoretical upper bounds on the masses of the lightest and second lightest inert neutralino states. The corresponding theoretical restrictions are

$$|m_{\chi_1^0}|^2 \lesssim \mu_0^2 \left[ m_{\chi_+^2}^2 + \frac{f_0^2 v^2}{2} \left( 1 + \sin^2 2\beta \right) \right] - \sqrt{\left( m_{\chi_+^2}^2 + \frac{f_0^2 v^2}{2} \left( 1 + \sin^2 2\beta \right) \right)^2 - \frac{f_0^4 v^4 \sin^2 2\beta}{2}},$$

(6)

where $m_{\chi_+^2} = \lambda_0 s/\sqrt{2}$ are the masses of the inert charginos. The value of $\mu_0$ decreases with increasing $|m_{\chi_+^2}|$ and $\tan \beta$. At large values of $|m_{\chi_+^2}|$ and $\tan \beta$, Eq. (6) simplifies resulting in

$$|m_{\chi_1^0}|^2 \lesssim \frac{f_0^2 v^4 \sin^2 2\beta}{4 \left( m_{\chi_+^2}^2 + \frac{f_0^2 v^2}{2} \left( 1 + \sin^2 2\beta \right) \right)}.$$

(7)

The theoretical restriction on $|m_{\chi_1^0}|$ achieves its maximal value around $\tan \beta \approx 1.5$. For this value of $\tan \beta$ the requirement of the validity of perturbation theory up to the GUT scale implies that $f_1 = \tilde{f}_1 = f_2 = \tilde{f}_2$ are less than 0.6. As a consequence the lightest inert neutralinos are lighter than $60 - 65$ GeV for $|m_{\chi_+^2}| > 100$ GeV.

The inert neutralino mass matrix (1) can be diagonalized using the neutralino mixing matrix defined by

$$N_i^a M^{ab} N^b_j = m_{ij}, \quad \text{no sum on } i.$$  

(8)

In the limit where off-diagonal Yukawa couplings vanish and $\lambda_0 s \gg f_0 v$, $\tilde{f}_0 v$ the eigenvalues of the inert neutralino mass matrix can be easily calculated (see [22]). In particular, the masses of two lightest inert neutralino states ($\chi_1^0$ and $\chi_2^0$) are given by

$$m_{\chi_1^0} \simeq \frac{\tilde{f}_0 f_0 v^2 \sin 2\beta}{2 m_{\chi_+^2}}.$$  

(9)

From Eq. (9) one can see that the masses of $\chi_1^0$ and $\chi_2^0$ are determined by the values of the Yukawa couplings $f_0$ and $\tilde{f}_0$.

The lightest inert neutralino states are made up of the following superposition of interaction states

$$\chi_1^0 = N_1^a \tilde{H}_2^0 + N_2^a \tilde{H}_2^0 + N_3^a \tilde{S}_2 + N_4^a \tilde{H}_1^0 + N_5^a \tilde{H}_1^0 + N_6^a \tilde{S}_1,$$

(10)

Using the above lightest and second lightest inert neutralino compositions it is straightforward to derive the couplings of these states to the $Z$-boson. In general the part of the Lagrangian that describes the interactions of $Z$ with $\chi_1^0$ and $\chi_2^0$ can be presented in the following form:

$$\mathcal{L}_{Z\chi} = \sum_{\alpha, \beta} \frac{M_Z}{2v} Z_{\mu} \left( \frac{\bar{\nu}_\alpha \gamma_\mu \gamma_5 \chi_1^0}{\lambda_{\alpha\beta} s} \right) R_{Z\alpha\beta}, \quad R_{Z\alpha\beta} = N_{\alpha}^1 N_{\beta}^1 + N_{\alpha}^2 N_{\beta}^2 + N_{\alpha}^4 N_{\beta}^4 + N_{\alpha}^5 N_{\beta}^5.$$  

(11)
In the case where off–diagonal Yukawa couplings go to zero while \(\lambda_\alpha s \gg f_\alpha v\), \(\tilde{f}_\alpha v\) the relative couplings of the lightest and second lightest inert neutralino states to the Z-boson are given by

\[
R_{Z\alpha\beta} = R_{Z\alpha\alpha} \delta_{\alpha\beta}, \quad R_{Z\alpha\alpha} = \frac{v^2}{2m_{\chi_1}^2} \left( f_{\alpha}^2 \cos^2 \beta - \tilde{f}_{\alpha}^2 \sin^2 \beta \right). \tag{12}
\]

One can see that the couplings of \(\chi_1^0\) and \(\chi_2^0\) to the Z-boson can be very strongly suppressed or even tend to zero. This happens when \(|f_\alpha| \cos \beta = |\tilde{f}_\alpha| \sin \beta\), which is when \(\chi_0^\alpha\) contains a completely symmetric combination of \(\tilde{H}_0^\alpha\) and \(\tilde{H}_0^\alpha\). Eq. (12) also indicates that the couplings of \(\chi_1^0\) and \(\chi_2^0\) to Z are always small when inert charginos are rather heavy or \(\tilde{f}_\alpha\) and \(f_\alpha\) are small (i.e. \(m_{\chi_0^\alpha} \rightarrow 0\)).

Although \(\chi_1^0\) and \(\chi_2^0\) might have extremely small couplings to Z, their couplings to the lightest CP–even Higgs boson \(h_1\) cannot be negligibly small if the corresponding states have appreciable masses. If all Higgs states except the lightest one are considerably heavier than the EW scale the mass matrix of the CP–even Higgs sector can be diagonalised using the perturbation theory \([23]–[27]\). Then the effective Lagrangian that describes the interactions of the inert neutralinos with the lightest CP-even Higgs eigenstate takes the form

\[
L_{h_1\chi} \simeq \sum_{i,j} (-1)^{\theta_i + \theta_j} X_{ij}^{h_1} \left( \psi_i^{0T} (-i\gamma_5)^{\theta_i + \theta_j} \psi_j^{0} \right) h_1, \quad X_{ij}^{h_1} = -\frac{1}{\sqrt{2}} \left( F_{ij} \cos \beta + \tilde{F}_{ij} \sin \beta \right), \tag{13}
\]

where \(i,j = 1, 2, \ldots 6\) and

\[
F_{ij} = f_{11} N_i^6 N_j^5 + f_{12} N_i^6 N_j^1 + f_{21} N_i^3 N_j^5 + f_{22} N_i^3 N_j^1,
\]

\[
\tilde{F}_{ij} = f_{11} N_i^6 N_j^4 + f_{12} N_i^6 N_j^2 + f_{21} N_i^3 N_j^4 + f_{22} N_i^3 N_j^2.
\]

In Eq. (13) \(\psi_0^i = (-i\gamma_5)^0 \alpha_0^i\) is the set of inert neutralino eigenstates with positive eigenvalues, while \(\theta_i\) equals 0 (1) if the eigenvalue corresponding to \(\chi_0^i\) is positive (negative). The inert neutralinos are labeled according to increasing absolute value of mass, with \(\psi_0^1\) being the lightest inert neutralino and \(\psi_0^6\) the heaviest.

In the limit when off–diagonal Yukawa couplings that determine the interactions of the inert Higgs fields with \(H_u, H_d\) and \(S\) vanish and \(\lambda_\alpha s \gg f_\alpha v\), \(\tilde{f}_\alpha v\) one obtains

\[
X_{\alpha\beta}^{h_1} \simeq \frac{|m_{\chi_0^\alpha}|}{v} \delta_{\alpha\beta}, \tag{14}
\]

where \(\alpha, \beta = 1, 2\), labeling the two light, mostly inert singlino states. These simple analytical expressions for the couplings of the SM–like Higgs boson to the lightest and second lightest inert neutralinos are not as surprising as they may first appear. When the Higgs spectrum is hierarchical, the VEV of the lightest CP–even state is responsible for all light fermion masses in the E6SSM. As a result we expect that their couplings to SM–like Higgs can be written as usual as being proportional to the mass divided by the VEV.

### 3. Exotic Higgs decays and Dark Matter

In our analysis we require that the lightest inert neutralino account for all or some of the observed dark matter relic density, which is measured to be \(\Omega_{\text{CDM}} h^2 = 0.1099 \pm 0.0062 \) \([28]\). If a theory predicts a greater relic density of dark matter than this then it is ruled out, assuming standard pre-BBN cosmology. A theory that predicts less dark matter cannot be ruled out in the same way but then there would have to be other contributions to the dark matter relic density.

In the limit where all non-SM fields other than the two lightest inert neutralinos are heavy (\(\gtrsim\) TeV) the lightest inert neutralino state in the E6SSM results in too large density of dark matter. Indeed, because the mass of \(\chi_0^1\) is inversely proportional to the masses of inert charginos the lightest inert neutralinos tend to be very light \(|m_{\chi_0^1}| \ll M_Z\). As a result the couplings of \(\chi_0^1\) to gauge bosons, Higgs states, quarks (squarks) and leptons (sleptons) are quite small leading to a relatively small annihilation cross section for \(\chi_0^1 \chi_0^1 \rightarrow \text{SM particles}\). Since the dark matter number density is inversely proportional to the annihilation cross section at the freeze-out temperature the lightest inert neutralino state gives rise to a relic density that is typically much larger than
The lightest Higgs boson exchange. Thus in the leading approximation the spin–independent part of scattering, which is associated with the spin-independent cross section, is mediated mainly by the lightest inert neutralino to quarks (leptons) and squarks (sleptons) are suppressed, the cross section in the E\text{–}boson annihilation proceeds through the Z–boson resonance, i.e. \( 2|m_{\chi_0^0}| \approx M_Z \).

Because the scenarios that result in the reasonable density of dark matter imply that \( \chi_0^0 \) and \( \chi_0^0 \) have large couplings to the lightest Higgs boson which are much larger than the \( b \)–quark Yukawa coupling and the decays of the lightest Higgs boson into these inert neutralinos are kinematically allowed, the SM–like Higgs boson decays predominantly into \( \chi_1^0 \) and \( \chi_2^0 \). The corresponding partial decay widths are given by

\[
\Gamma(h_1 \rightarrow \chi_0^0 \chi_0^0) = \frac{\Delta_{\alpha\beta}}{8\pi m_{h_1}} \left( X_{h_1}^{\alpha\beta} + X_{h_1}^{\alpha\beta} \right)^2 \left[ m_{h_1}^2 - (|m_{\chi_0^0}| + (-1)^{\theta_\alpha + \theta_\beta} |m_{\chi_0^0}|)^2 \right] \times \sqrt{1 - \frac{|m_{\chi_0^0}|^2}{m_{h_1}^2}} - \frac{|m_{\chi_0^0}|^2}{2 m_{h_1}^2} - \frac{|m_{\chi_0^0}|^2}{2 m_{h_1}^2} - \frac{4 m_{\chi_0^0}^2 |m_{\chi_0^0}|^2}{m_{h_1}^4},
\]

where \( \Delta_{\alpha\beta} = \frac{1}{2} (1) \) for \( \alpha = \beta \) (\( \alpha \neq \beta \)). On the other hand the large coupling of \( \chi_0^0 \) to the lightest Higgs state give rise to the relatively large \( \chi_0^0 \)–nucleon elastic–scattering cross section. Since in the E\text{–}SSM the couplings of the lightest inert neutralino to quarks (leptons) and squarks (sleptons) are suppressed, the \( \chi_0^0 \)–nucleon elastic scattering, which is associated with the spin-independent cross section, is mediated mainly by the \( t \)–channel lightest Higgs boson exchange. Thus in the leading approximation the spin–independent part of \( \chi_0^0 \)–nucleon cross section in the E\text{–}SSM takes the form

\[
\sigma_{SI} = \frac{4m_r^2 m_N^2 |X_{h_1}^{s}|^2 F^N|^2}{\pi v^2 m_{h_1}^4},
\]

\[
m_r = \frac{m_{\chi_0^0} m_N}{m_{\chi_0^0} + m_N}, \quad F^N = \sum_{q=u,d,s} f_{Tq}^N + \frac{2}{27} \sum_{Q=c,b,t} f_{TQ}^N,
\]

where

\[
m_N f_{Tq}^N = \langle N | m_{\bar{q} q} | N \rangle, \quad f_{TQ}^N = 1 - \sum_{q=u,d,s} f_{Tq}^N.
\]

As one can see from Eq. \( \text{[15]} \) the value of \( \sigma_{SI} \) depends rather strongly on the hadronic matrix elements, i.e. the coefficients \( f^N_{Tq} \), that are related to the \( \pi \)–nucleon \( \sigma \) term and the spin content of the nucleon. As a consequence \( \sigma_{SI} \) varies over a wide range (see Table 1). Recently the CDMSII and XENON100 collaborations set upper limits on \( \sigma_{SI} \) \cite{31, 32}.

In order to illustrate the features of the E\text{–}SSM mentioned above we specify a set of benchmark points (see Table 1). For each benchmark scenario we calculate the spectrum of the inert neutralinos, inert charginos and Higgs bosons as well as the branching ratios of the decays of the lightest CP-even Higgs state and the dark matter relic density. In Table 1 the masses of the heavy Higgs states are computed in the leading one–loop approximation. In the case of the lightest Higgs boson mass the leading two–loop corrections are taken into account. In order to construct benchmark scenarios that are consistent with cosmological observations we restrict our considerations to low values of \( \tan \beta \lesssim 2 \) that allows to obtain \( |m_{\chi_1^0}| \sim |m_{\chi_2^0}| \sim M_Z/2 \). However, even for \( \tan \beta \lesssim 2 \) the lightest inert neutralino states can get appreciable masses only if inert chargino mass

\[1\text{When \( f_{\alpha\beta} \), \( f_{\alpha\beta} \) \rightarrow 0 \ the masses of \( \chi_0^0 \) and \( \chi_0^0 \) tend to zero and inert singlino states essentially decouple from the rest of the spectrum. In this limit the lightest non-decoupled neutralino may be rather stable and can play the role of dark matter \cite{28}. The presence of very light neutral fermions in the particle spectrum might have interesting implications for the neutrino physics (see, for example \cite{29}).}
Table I: Benchmark scenarios. The branching ratios and decay widths of the lightest Higgs boson, the masses of the Higgs states, inert neutralinos and charginos as well as the couplings of $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$ are calculated for $s = 2400$ GeV, $m_Q = m_U = M_S = 700$ GeV, $X_t = \sqrt{3}M_S$ that correspond to $m_{h^0} \simeq M_{\tilde{g}'} \simeq 890$ GeV.

| i | ii | iii | iv | v |
|---|----|-----|----|---|
| $\lambda$ | 0.6 | 0.6 | 0.468 | 0.468 | 0.468 |
| $\tan(\beta)$ | 1.7 | 1.564 | 1.5 | 1.5 | 1.5 |
| $A_\lambda$ | 1600 | 1600 | 600 | 600 | 600 |
| $m_{h^\pm} \simeq m_A \simeq m_{h^0}$/GeV | 1977 | 1990 | 1145 | 1145 | 1145 |
| $m_{h^0}$/GeV | 133.1 | 134.8 | 115.9 | 115.9 | 115.9 |
| $\lambda_{22}$ | 0.094 | 0.0001 | 0.094 | 0.001 | 0.468 |
| $\lambda_{21}$ | 0 | 0.06 | 0 | 0.079 | 0.05 |
| $\lambda_{12}$ | 0 | 0.06 | 0 | 0.080 | 0.05 |
| $\lambda_{11}$ | 0.059 | 0.0001 | 0.059 | 0.001 | 0.08 |
| $f_{22}$ | 0.53 | 0.001 | 0.53 | 0.04 | 0.05 |
| $f_{21}$ | 0.05 | 0.476 | 0.053 | 0.68 | 0.9 |
| $f_{12}$ | 0.05 | 0.466 | 0.053 | 0.68 | 0.002 |
| $f_{11}$ | 0.53 | 0.001 | 0.53 | 0.04 | 0.002 |
| $f_{12}$ | 0.53 | 0.001 | 0.53 | 0.04 | 0.002 |
| $f_{11}$ | 0.53 | 0.001 | 0.53 | 0.04 | 0.05 |
| $m_{\tilde{\chi}_1^0}$/GeV | 33.62 | -36.69 | 35.42 | -45.08 | -46.24 |
| $m_{\tilde{\chi}_2^0}$/GeV | 47.78 | 36.88 | 51.77 | 55.34 | 46.60 |
| $m_{\tilde{\chi}_3^0}$/GeV | 108.0 | -103.11 | 105.3 | -133.3 | 171.1 |
| $m_{\tilde{\chi}_4^0}$/GeV | -152.1 | 103.47 | -152.7 | 136.9 | -171.4 |
| $m_{\tilde{\chi}_2^\pm}$/GeV | 163.5 | 139.80 | 162.0 | 178.4 | 805.4 |
| $m_{\tilde{\chi}_2^0}$/GeV | -200.8 | -140.35 | -201.7 | -192.2 | -805.4 |
| $m_{\tilde{\chi}_2^\mp}$/GeV | 100.1 | 101.65 | 100.1 | 133.0 | 125.0 |
| $m_{\tilde{\chi}_2^\pm}$/GeV | 159.5 | 101.99 | 159.5 | 136.8 | 805.0 |
| $\Omega h^2$ | 0.0109 | 0.107 | 0.107 | 0.0324 | 0.000065 |
| $R_{Z_11}$ | -0.144 | -0.132 | -0.115 | -0.0217 | -0.0224 |
| $R_{Z_12}$ | 0.051 | 0.0043 | -0.045 | -0.0020 | -0.213 |
| $R_{Z_22}$ | -0.331 | -0.133 | -0.288 | -0.0524 | -0.0226 |
| $\sigma_{Zf}/10^{-44}$ cm$^2$ | 1.7-7.1 | 2.0-8.2 | 3.5-14.2 | 6.0-24.4 | 6.1-25.0 |
| Br($h \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0$) | 57.8% | 49.1% | 76.3% | 83.4% | 49.3% |
| Br($h \rightarrow \tilde{\chi}_1^0\tilde{\chi}_1^0$) | 0.34% | 3.5 x 10^{-11} | 0.26% | 7.6 x 10^{-9} | 3.0 x 10^{-8} |
| Br($h \rightarrow \tilde{\chi}_2^0\tilde{\chi}_2^0$) | 39.8% | 49.2% | 20.3% | 12.3% | 47.9% |
| Br($h \rightarrow \tilde{b}\tilde{b}$) | 1.87% | 1.5% | 2.83% | 3.95% | 2.58% |
| Br($h \rightarrow \tau\tau$) | 0.196% | 0.166% | 0.30% | 0.41% | 0.27% |
| $\Gamma_{tot}$/MeV | 141.2 | 169.0 | 82.0 | 58.8 | 90.1 |

Eigenstates are light, i.e. $m_{\tilde{\chi}_1^\pm} \simeq 100 - 200$ GeV. We demonstrate (see benchmark point (v) in Table 1) that the scenarios with only one light inert chargino mass eigenstate may lead to the dark matter density consistent with cosmological observations.

When $\tan(\beta) \lesssim 2$ the mass of the lightest CP–even Higgs boson is very sensitive to the choice of the coupling $\lambda(M_t)$. In particular, to satisfy LEP constraints $\lambda(M_t)$ must be larger than $g_1' \simeq 0.47$, where $g_1'$ is the low energy $U(1)_X$ gauge coupling. If $\lambda \gtrsim g_1'$ the vacuum stability requires all Higgs states except the lightest one to be considerably heavier than the EW scale so that the qualitative pattern of the Higgs spectrum is rather
similar to the one which arises in the PQ symmetric NMSSM\cite{26,27,33}. In this case the lightest Higgs state manifests itself in the interactions with gauge bosons and fermions as a SM–like Higgs boson.

Our benchmark scenarios indicate that in the case when $m_{\chi^0_2} \sim |m_{\chi_0^0}| \sim M_Z/2$ the SM–like Higgs boson decays more than 95% of the time into $\chi^0_1$ and $\chi^0_2$ while the total branching ratio into SM particles varies from 2% to 4%. When the masses of the lightest and second lightest inert neutralinos are close or they form a Dirac (pseudo–Dirac) state (see benchmark scenarios (ii) and (v) in Table 1) then the decays of the lightest Higgs boson into $\chi^0_1$ and $\chi^0_2$ lead to the missing $E_T$ in the final state. Thus these decay channels give rise to a large invisible branching ratio of the SM–like Higgs boson. If the mass difference between the second lightest and the lightest inert neutralino is 10 GeV or more, then some of the decay products of a $\chi^0_1$ that originates from a SM-like Higgs boson decay might be observed at the LHC. In our analysis we assume that all scalar particles, except for the lightest Higgs boson, are heavy and that the couplings of the inert neutralino states to quarks, leptons and their superpartners are relatively small. As a result the second lightest inert neutralino decays into the lightest one and a fermion–antifermion pair mainly via a virtual $Z$. In our numerical analysis we did not manage to find any benchmark scenario with $m_{\chi_1^0} - |m_{\chi_2^0}| \gtrsim 20$ GeV leading to reasonable values of $\Omega_{CDM} h^2$. Hence we do not expect any observable jets at the LHC associated with the decay of a $\chi^0_1$ produced through a Higgs decay. However, it might be possible to detect $\mu^+\mu^-$ pairs that come from the exotic decays of the lightest CP–even Higgs state mentioned above.

In Table 1 benchmark scenarios (i), (iii), (iv) can lead to these relatively energetic muon pairs in the final state of the SM-like Higgs decays. Since the Higgs branching ratios into SM particles are rather suppressed, the decays of the lightest CP–even Higgs state into $t\bar{t}l^+l^- + X$ might play an essential role in Higgs searches.

In Table 1 we also specify the interval of variations of $\sigma_{SI}$ for each benchmark scenario. The lower limit on $\sigma_{SI}$ corresponds to $f_{TaN}^N = 0$ while the upper limit implies that $f_{TaN}^N = 0.36$. From Table 1 and Eq. (10) it also becomes clear that $\sigma_{SI}$ decreases when $m_{h_1}$ grows. Since in all of the benchmark scenarios presented in Tables 1 the lightest inert neutralino is relatively heavy ($|m_{\chi_1^0}| \sim M_Z/2$), allowing for a small enough dark matter relic density, the coupling of $\chi^0_1$ to the lightest CP-even Higgs state is always large giving rise to a $\chi^0_1$–nucleon spin-independent cross-section which is on the edge of observability of XENON100.

In addition to the exotic Higgs decays and large LSP direct detection cross-sections, the scenarios considered here imply that at least two of the inert neutralino states that are predominantly the fermion components of the inert Higgs doublet superfields and one of the inert chargino states should have masses below 200 GeV. Because these states are almost inert Higgsinos they couple rather strongly to $W$ and $Z$–bosons. Thus at hadron colliders the corresponding inert neutralino and chargino states can be produced in pairs via off-shell $W$ and $Z$–bosons. Since they are light their production cross sections at the LHC are not negligibly small. After being produced inert neutralino and chargino states sequentially decay into the LSP and pairs of leptons and quarks resulting in distinct signatures that can be discovered at the LHC in the near future.

Acknowledgments

R.N. would like to thank FTPI, University of Minnesota for its hospitality and M. A. Shifman, K. A. Olive, A. I. Vainshtein, A. Mustafayev, W. Vinci, P. A. Bolokhov, P. Koroteev for fruitful discussions. Authors are grateful to X. R. Tata, M. Muehlleitner, L. Clavelli, D. Stockinger, D. J. Miller, D. G. Sutherland, J. P. Kumar, D. Marfatia, K. R. Dienes, B. D. Thomas for valuable comments and remarks. The work of R.N. and S.P. was supported by the National Science Foundation PHY-0755262. S.F.K. acknowledges partial support from the STFC Rolling Grant ST/G000557/1. J.P.H. is thankful to the STFC for providing studentship funding.

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