Novel Higgs decays and dark matter in the exceptional supersymmetric standard model

J. P. Hall,[1,*] S. F. King,[1,†] R. Nevzorov,[2,‡] S. Pakvasa,[2,§] and M. Sher[3,‖]

[1]*School of Physics and Astronomy, University of Southampton, Southampton, UK*
[2]*Department of Physics and Astronomy, University of Hawaii, Honolulu, USA*
[3]*Physics Department, College of William and Mary, Williamsburg, USA*
(Received 29 December 2010; published 15 April 2011)

The Exceptional Supersymmetric (SUSY) Standard Model (E_6SSM) predicts three families of Higgs doublets plus three Higgs singlets, where one family develops vacuum expectation values (VEVs), while the remaining two which do not are called Inert. The model can account for the dark matter relic abundance if the two lightest Inert neutralinos, identified as the (next-to) lightest SUSY particles ((N)LSPs), have masses close to half the Z mass. In this case we find that the usual SM-like Higgs boson decays more than 95% of the time into either LSPs or NLSPs. The latter case produces a final state containing two leptons l^+l^- with an invariant mass less than or about 10 GeV. We illustrate this scenario with a set of benchmark points satisfying phenomenological constraints and the WMAP dark matter relic abundance. This scenario also predicts other light Inert chargino and neutralino states below 200 GeV, and large LSP direct detection cross sections close to current limits and observable soon at XENON100.

I. INTRODUCTION

The discovery of the Higgs boson, the last missing piece in the standard model (SM) of particle physics, is one of the main goals of upcoming accelerators. It is expected that the Higgs particle will be detected at the Large Hadron Collider (LHC) in the near future. The strategy for Higgs searches depends on the decay branching fractions of the Higgs boson to different channels. Physics beyond the standard model may affect the Higgs decay rates to SM particles and give rise to new channels of Higgs decays requiring a drastic change in the strategy for Higgs boson searches (for recent reviews of nonstandard Higgs boson decays see [1]). 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 cannot be directly detected. Invisible Higgs decay modes may occur in models with an enlarged symmetry breaking sector (Majoron models, SM with extra singlet scalar fields etc.) [2–5], in "hidden valley" models [6], in the SM with a fourth generation of fermions [7,8], in the models with compact and large extra dimensions [5,9,10], in the littlest Higgs model with T parity [11] etc.[1]

Another example wherein invisible decay modes can occur is supersymmetry (SUSY), with the lightest Higgs boson decaying into the lightest SUSY particle (LSP).

R-parity conservation ensures the stability of the LSP so that the LSP can play the role of cold dark matter (CDM) [13]. In most scenarios the LSP is the lightest neutralino, a linear combination of neutral electroweak (EW) gauginos and Higgsinos. In some regions of the parameter space the lightest Higgs boson in the Minimal Supersymmetric Standard Model (MSSM) decays into the lightest neutralino with a relatively large branching ratio, therefore giving rise to invisible final states if R parity is conserved [14].

LEP and Tevatron data still allow the neutralino LSP to be sufficiently light so that the decays of the lightest Higgs into these neutralinos is kinematically allowed. Light neutralinos can annihilate efficiently through a Z pole, resulting in a reasonable density of dark matter. Moreover, the Cryogenic Dark Matter Search (CDMS) experiment recently reported the observation of 2 events possibly due to dark matter scattering with an expected background of about 0.8 events [15]. CDMS events suggest that the mass of the dark matter particles are around 40–80 GeV while their spin-independent elastic cross section is $\sigma_{\rm SI} \approx$ few $\times 10^{-44}$ cm^2. If recent results of the CDMS experiment get confirmed then scenarios with invisible decays of the Higgs boson will become rather plausible.

Certainly 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. In the case where invisible Higgs decays dominate it is impossible to fully reconstruct a resonance and it is very challenging to identify it at the collider experiments, i.e. quantum numbers remain unknown. At e^+e^- colliders, the problems related to the observation of the invisible Higgs are less severe [3,16] since it can be tagged through the recoiling Z. Presently,

---

*jonathan.hall@soton.ac.uk
†s.f.king@soton.ac.uk
‡On leave of absence from the Theory Department, ITEP, Moscow, Russia; E-mail: nevzorov@phys.hawaii.edu
§pakvasa@phys.hawaii.edu
‖mtsher@wm.edu
[1]In the context of the so-called "nightmare scenario", in which the LHC produces nothing beyond the SM, Higgs decays into and interactions with dark matter particles were discussed in [12].

the LEP II collaborations exclude invisible Higgs masses up to 114.4 GeV [17].[2]

On the other hand, Higgs searches at hadron colliders are more difficult in the presence of such invisible decays. Previous studies have analyzed ZH and WH associated production [4,18,19] as well as $t\bar{t}H$ production [20] and $t\bar{t}VV$ ($b\bar{b}VV$) production [21] as promising channels. The possibility of observing an "invisible" Higgs boson in central exclusive diffractive production at the LHC was studied in [8]. Another proposal is to observe such an invisible Higgs in inelastic events with large missing transverse energy and two high E_T jets. In this case the Higgs boson is produced by VV fusion and has large transverse momentum resulting in a signal with two quark jets with distinctive kinematic distributions as compared to Zjj and Wjj backgrounds [10,19,22].

Consideration of the possibility that the dominant Higgs decays will be invisible would lead to drastic changes in the strategy of Higgs boson searches. Therefore, it is a rather interesting subject of investigation as to the nature and extent of invisibility acquired by Higgs, and how it can be related to specific aspects of the models concerned, especially well motivated SUSY extensions of the SM. In this article, we consider the exotic decays of the lightest Higgs boson and associated novel collider signatures within the Exceptional Supersymmetric Standard Model (E_6SSM) [23,24]. This E_6 inspired SUSY model is based on the low-energy standard model gauge group together with an extra $U(1)_N$ gauge symmetry under which right-handed neutrinos have zero charge. In the E_6SSM the μ problem is solved in a similar way as in the Next-to-Minimal Supersymmetric Standard Model (NMSSM), but without the accompanying problems of singlet tadpoles or domain walls. Because right-handed neutrinos do not participate in the gauge interactions in this model 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 [25].

The particle spectrum of the E_6SSM contains exotic matter. In particular, it involves three SM-singlet superfields that have nonzero $U(1)_N$ charges. One of these singlets acquires a nonzero vacuum expectation value (VEV), breaking $U(1)_N$ symmetry and inducing the effective μ term and the masses of the exotic fermions. The masses of the fermion components of the two other singlet superfields (Inert neutralinos) are also related to the VEVs of the Higgs doublets. Because the Yukawa couplings that determine the strength of these interactions are constrained by the requirement of the validity of perturbation theory up to the grand unification scale the masses of the correspond-

---

[2]Similar limits could apply to the case where the Higgs decays some fraction of the time into soft lepton pairs plus missing energy, as will be the case for some of the novel Higgs decays considered in this paper.

ing Inert neutralinos are expected to be lighter than 60–65 GeV. As a result the lightest Inert neutralino tends to be the lightest SUSY particle in the spectrum. Such a neutralino can give an appropriate contribution to the dark matter density consistent with the recent observations if it has mass 35–50 GeV [26]. In this case the lightest Higgs boson decays predominantly into Inert neutralino whereas usual Higgs branching ratios are less than a few percent.

The layout of this paper is as follows. In Sec. II we briefly review the E_6SSM. In Sec. III and IV the spectrum and couplings of the Inert neutralinos, charginos and Higgs bosons are specified. The novel decays of the lightest CP-even Higgs state and dark matter constraints are discussed in Sec. V. In Sec. VI we discuss the benchmark points and the experimental constraints and predictions. Section VII summarizes and concludes the paper.

II. EXCEPTIONAL SUSY MODEL

The E_6SSM 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. The additional low energy $U(1)_N$, which is not present either in the SM or in the MSSM, is a linear superposition of $U(1)_\chi$ and $U(1)_\psi$, namely

$$U(1)_N = \frac{1}{4}U(1)_\chi + \frac{\sqrt{15}}{4}U(1)_\psi, \qquad (1)$$

where the $U(1)_\psi$ and $U(1)_\chi$ symmetries are defined by

$$E_6 \rightarrow SO(10) \times U(1)_\psi, \qquad SO(10) \rightarrow SU(5) \times U(1)_\chi.$$

Thus the E_6SSM can originate from an E_6 GUT gauge group which is broken at the GUT scale M_X. The extra $U(1)_N$ gauge symmetry is defined such that right-handed neutrinos carry zero charges.

In E_6 theories the anomalies cancel automatically; all models that are based on the E_6 subgroups and contain complete representations of E_6 should be anomaly-free. Consequently, in order to make a supersymmetric model with an extra $U(1)_N$ anomaly-free, one is forced to augment the minimal particle spectrum by a number of exotics which, together with ordinary quarks and leptons, form complete fundamental 27 representations of E_6. Thus the particle content of the E_6SSM involves at least three fundamental representations of E_6 at low energies. These multiplets decompose under the $SU(5) \times U(1)_N$ subgroup of E_6 as follows [27]:

$$27_i \rightarrow \left(10, \frac{1}{\sqrt{40}}\right)_i + \left(5^*, \frac{2}{\sqrt{40}}\right)_i + \left(5^*, -\frac{3}{\sqrt{40}}\right)_i$$
$$+ \left(5, -\frac{2}{\sqrt{40}}\right)_i + \left(1, \frac{5}{\sqrt{40}}\right)_i + (1, 0)_i. \qquad (2)$$

The first and second quantities in brackets are the $SU(5)$ representation and extra $U(1)_N$ charge, respectively, while i is a family index that runs from 1 to 3. An ordinary SM family, which contains the doublets of left-handed quarks

Q_i and leptons L_i, right-handed up-and down-quarks (u_i^c and d_i^c) as well as right-handed charged leptons, is assigned to $(10, \frac{1}{\sqrt{40}})_i + (5^*, \frac{2}{\sqrt{40}})_i$. Right-handed neutrinos N_i^c should be associated with the last term in Eq. (2), $(1, 0)_i$. Because they do not carry any charges right-handed neutrinos are expected to be superheavy allowing them to be used for both the seesaw mechanism and leptogenesis. The next-to-last term, $(1, \frac{5}{\sqrt{40}})_i$, represents SM-singlet fields S_i, which carry nonzero $U(1)_N$ charges and therefore survive down to the EW scale. The pair of $SU(2)_W$-doublets (H_i^d and H_i^u) that are contained in $(5^*, -\frac{3}{\sqrt{40}})_i$ and $(5, -\frac{2}{\sqrt{40}})_i$ have the quantum numbers of Higgs doublets. They form either Higgs or Inert Higgs $SU(2)_W$ multiplets.[3] Other components of these $SU(5)$ multiplets form color triplets of exotic quarks $\bar{D}_i$ and D_i with electric charges $-1/3$ and $+1/3$, respectively. These exotic quark states carry a $B - L$ charge $\pm 2/3$, twice that of ordinary ones. Therefore in phenomenologically viable E_6 inspired models they can be either diquarks or leptoquarks.

In addition to the complete 27_i multiplets the low energy matter content of the E_6SSM can be supplemented by an $SU(2)_W$ doublet $\hat{L}_4$ and antidoublet $\hat{\bar{L}}_4$ from the extra $27'$ and $\overline{27'}$ to preserve gauge coupling unification. These components of the E_6 fundamental representation originate from $(5^*, \frac{2}{\sqrt{40}})$ of $27'$ and $(5, -\frac{2}{\sqrt{40}})$ of $\overline{27'}$ by construction.

Thus, in addition to a Z' corresponding to the $U(1)_N$ symmetry, the E_6SSM involves extra matter beyond the MSSM that forms three $5 + 5^*$ representations of $SU(5)$ plus three $SU(5)$ singlets with $U(1)_N$ charges. The analysis performed in [28] shows that the unification of gauge couplings in the E_6SSM can be achieved for any phenomenologically acceptable value of $\alpha_3(M_Z)$ consistent with the measured low energy central value, unlike in the MSSM which, ignoring the effects of high energy threshold corrections, requires values of $\alpha_3(M_Z)$ which are significantly above the experimentally measured central value. The presence of a Z' boson and of exotic quarks predicted by the E_6SSM provides spectacular new physics signals at the LHC which were discussed in [23,24,29]. Recently the particle spectrum and collider signatures associated with it were studied within the constrained version of the E_6SSM [30].

In general, the E_6 symmetry does not forbid lepton and baryon number violating operators that result in rapid proton decay. Moreover, exotic particles in E_6 inspired SUSY models give rise to new Yukawa interactions that induce unacceptably large nondiagonal flavour transitions. To suppress these effects in the E_6SSM 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 $H_d \equiv H_3^d$ and $H_u \equiv H_3^u$) and one SM-type singlet field ($S \equiv S_3$) are odd.

---

[3]We use the terminology "Inert Higgs" to denote Higgslike doublets that do not develop VEVs.

The Z_2^H symmetry reduces the structure of the Yukawa interactions to

$$W_{E_6SSM} \simeq \lambda \hat{S}(\hat{H}_u \hat{H}_d) + \lambda_{\alpha\beta} \hat{S}(\hat{H}_\alpha^d \hat{H}_\beta^u) + \tilde{f}_{\alpha\beta} \hat{S}_\alpha(\hat{H}_\beta^d \hat{H}_u)$$
$$+ f_{\alpha\beta} \hat{S}_\alpha(\hat{H}_d \hat{H}_\beta^u) + \kappa_{ij} \hat{S}(\hat{D}_i \hat{\bar{D}}_j) + h_{ij}^U(\hat{H}_u \hat{Q}_i)\hat{u}_j^c$$
$$+ h_{ij}^D(\hat{H}_d \hat{Q}_i)\hat{d}_j^c + h_{ij}^E(\hat{H}_d \hat{L}_i)\hat{e}_j^c + h_{ij}^N(\hat{H}_u \hat{L}_i)\hat{N}_j^c$$
$$+ \frac{1}{2}M_{ij}\hat{N}_i^c \hat{N}_j^c + \mu'(\hat{L}_4 \hat{\bar{L}}_4) + h_{4j}^E(\hat{H}_d \hat{L}_4)\hat{e}_j^c$$
$$+ h_{4j}^N(\hat{H}_u \hat{L}_4)\hat{N}_j^c, \tag{3}$$

where α, $\beta = 1, 2$ and i, $j = 1, 2, 3$. The $SU(2)_W$ doublets $\hat{H}_u$ 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 the Higgs fields develop VEVs:

$$\langle H_d \rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} v_1 \\ 0 \end{pmatrix}, \qquad \langle H_u \rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_2 \end{pmatrix}, \qquad \langle S \rangle = \frac{s}{\sqrt{2}}, \tag{4}$$

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 $U(1)_N$ symmetry thereby providing an effective μ term as well as the necessary exotic fermion masses and also inducing that of the Z' boson. Therefore the singlet field S must acquire a large VEV in order to avoid conflict with direct particle searches at present and past accelerators. This also requires the Yukawa couplings λ_i and κ_i to be reasonably large. If λ_i or κ_i are large enough at the GUT scale they affect the evolution of the soft scalar mass m_S^2 of the singlet field S rather strongly resulting in a negative value of m_S^2 at low energies which triggers the breakdown of the $U(1)_N$ symmetry.

Note that the surviving components from the $27'$ and $\overline{27'}$ manifest themselves in the Yukawa interactions (3) as fields with lepton number $L = \pm 1$. The corresponding mass term $\mu' L_4 \bar{L}_4$ in the superpotential (3) is not involved in the process of electroweak symmetry breaking (EWSB). Moreover this term is not suppressed by the E_6 symmetry. Therefore the parameter μ' remains arbitrary. Gauge coupling unification requires μ' to be below about 100 TeV [28]. Thus we assume that the scalar and fermion components of the superfields $\hat{L}_4$ and $\hat{\bar{L}}_4$ are very heavy so that they decouple from the rest of the particle spectrum.

Although Z_2^H eliminates any problems related with baryon number violation and nondiagonal flavor 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 [31] the Z_2^H symmetry must be broken. At the same time, the breakdown of Z_2^H should not give rise to operators that would lead to rapid proton decay. There are two ways to overcome this problem: the Lagrangian must be

invariant with respect to either a Z_2^L symmetry, under which all superfields except leptons are even (Model I), or a Z_2^B discrete symmetry, which implies that exotic quark and lepton superfields are odd whereas the others remain even (Model II). If the Lagrangian is invariant under the Z_2^L symmetry, then the terms in the superpotential which permit exotic quarks to decay and are allowed by the E_6 symmetry can be written in the form

$$W_1 = g_{ijk}^Q \hat{D}_i (\hat{Q}_j \hat{Q}_k) + g_{ijk}^q \hat{\bar{D}}_i \hat{d}_j^c \hat{u}_k^c, \tag{5}$$

that implies that exotic quarks are diquarks. If Z_2^B is imposed then the following couplings are allowed:

$$W_2 = g_{ijk}^E \hat{e}_i^c \hat{D}_j \hat{u}_k^c + g_{ijk}^D (\hat{Q}_i \hat{L}_j) \hat{\bar{D}}_k. \tag{6}$$

In this case baryon number conservation requires the exotic quarks to be leptoquarks.

III. INERT CHARGINOS AND NEUTRALINOS

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 $\mu \to e^- e^+ e^-$. In order to suppress processes with non-diagonal flavor 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.

In this approximation and given the previous assumption that only H_u, H_d and S acquire nonzero VEVs the charged components of the Inert Higgsinos ($\tilde{H}_2^{u+}$, $\tilde{H}_1^{u+}$, $\tilde{H}_2^{d-}$, $\tilde{H}_1^{d-}$) and ordinary chargino states do not mix. The neutral components of the Inert Higgsinos ($\tilde{H}_1^{d0}$, $\tilde{H}_2^{d0}$, $\tilde{H}_1^{u0}$, $\tilde{H}_2^{u0}$) 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, the lighter of the two being the stable LSP.

In the field basis ($\tilde{H}_2^{d0}$, $\tilde{H}_2^{u0}$, $\tilde{S}_2$, $\tilde{H}_1^{d0}$, $\tilde{H}_1^{u0}$, $\tilde{S}_1$) the mass matrix of the Inert neutralino sector takes a form

$$M_{\text{IN}} = \begin{pmatrix} A_{22} & A_{21} \\ A_{12} & A_{11} \end{pmatrix}, \tag{7}$$

where $A_{\alpha\beta}$ are 3×3 submatrices given by [26]

$$A_{\alpha\beta} = -\frac{1}{\sqrt{2}} \begin{pmatrix} 0 & \lambda_{\alpha\beta}s & \tilde{f}_{\beta\alpha}v\sin\beta \\ \lambda_{\beta\alpha}s & 0 & f_{\beta\alpha}v\cos\beta \\ \tilde{f}_{\alpha\beta}v\sin\beta & f_{\alpha\beta}v\cos\beta & 0 \end{pmatrix}, \tag{8}$$

so that $A_{\alpha\beta} = A_{\beta\alpha}^T$. In the basis of Inert chargino interaction states ($\tilde{H}_2^{u+}$, $\tilde{H}_1^{u+}$, $\tilde{H}_2^{d-}$, $\tilde{H}_1^{d-}$) the corresponding mass matrix can be written as

$$M_{\text{IC}} = \begin{pmatrix} 0 & C^T \\ C & 0 \end{pmatrix}, \qquad C_{\alpha\beta} = \frac{1}{\sqrt{2}} \lambda_{\alpha\beta}s. \tag{9}$$

where $C_{\alpha\beta}$ are 2×2 submatrices. From Eqs. (7)–(9) one can see that in the exact Z_2^H symmetry limit the spectrum of the Inert neutralinos and charginos in the E_6SSM can be parametrized in terms of

$$\lambda_{\alpha\beta}, \qquad f_{\alpha\beta}, \qquad \tilde{f}_{\alpha\beta}, \qquad \tan\beta, \qquad s. \tag{10}$$

In other words the masses and couplings of the Inert neutralinos are determined by 12 Yukawa couplings, which can be complex, $\tan\beta$ and s. Four of the Yukawa couplings mentioned above, i.e. $\lambda_{\alpha\beta}$, as well as the VEV of the SM-singlet field s set the masses and couplings of the Inert chargino states. Six off-diagonal Yukawa couplings define the mixing between the two families of the Inert Higgsinos and singlinos.

In the following analysis we shall choose the VEV of the SM-singlet field to be large enough ($s \gtrsim 2400$ GeV) so that the experimental constraints on Z' boson mass ($M_{Z'} \gtrsim 865$ GeV) and $Z - Z'$ mixing are satisfied. In order to avoid the LEP lower limit on the masses of Inert charginos the Yukawa couplings $\lambda_{\alpha\beta}$ are chosen so that all Inert chargino states are heavier than 100 GeV. In addition, we also require the validity of perturbation theory up to the GUT scale and that constrains the allowed range of all Yukawa couplings.

The theoretical and experimental restrictions specified above set very strong limits on the masses and couplings of the lightest Inert neutralinos. In particular, our numerical analysis indicates that the lightest and second lightest Inert neutralinos are always light. They typically have masses below 60–65 GeV. These neutralinos are predominantly Inert singlinos. From our numerical analysis it follows that the lightest and second lightest Inert neutralinos might have rather small couplings to the Z boson so that any possible signal which these neutralinos could give rise to at LEP would be extremely suppressed. As a consequence such Inert neutralinos would remain undetected. At the same time four other Inert neutralinos, which are approximately linear superpositions of neutral components of Inert Higgsinos, are normally heavier than 100 GeV.

A. The diagonal Inert Yukawa approximation

In order to clarify the results of our numerical analysis, it is useful to consider a few simple cases that give some

analytical understanding of our calculations. In the simplest case when all off-diagonal Yukawa couplings vanish, considered in [26],

$$\lambda_{\alpha\beta} = \lambda_\alpha \delta_{\alpha\beta}, \qquad f_{\alpha\beta} = f_\alpha \delta_{\alpha\beta}, \qquad \tilde{f}_{\alpha\beta} = \tilde{f}_\alpha \delta_{\alpha\beta},$$

the mass matrix of Inert neutralinos reduces to the block diagonal form while the masses of the Inert charginos are given by

$$m_{\chi_\alpha^\pm} = \frac{\lambda_\alpha}{\sqrt{2}} s. \tag{11}$$

In the limit where $f_\alpha = \tilde{f}_\alpha$ one can easily prove using the method proposed in [32] that there are theoretical upper bounds on the masses of the lightest and second lightest Inert neutralino states. The corresponding theoretical restrictions are

$$|m_{\chi_\alpha^0}|^2 \lesssim \mu_\alpha^2 = \frac{1}{2}\left[|m_{\chi_\alpha^\pm}|^2 + \frac{f_\alpha^2 v^2}{2}(1 + \sin^2 2\beta) - \sqrt{\left(|m_{\chi_\alpha^\pm}|^2 + \frac{f_\alpha^2 v^2}{2}(1 + \sin^2 2\beta)\right)^2 - f_\alpha^4 v^4 \sin^2 2\beta} \right]. \tag{12}$$

The value of μ_α decreases with increasing $|m_{\chi_\alpha^\pm}|$ and $\tan\beta$, hence reaching its maximum value of $\frac{f_\alpha}{\sqrt{2}} v$ for $m_{\chi_\alpha^\pm} \to 0$ and $\tan\beta \to 1$. At large values of $|m_{\chi_\alpha^\pm}|$ and $\tan\beta$, Eq. (12) simplifies resulting in

$$|m_{\chi_\alpha^0}|^2 \lesssim \frac{f_\alpha^4 v^4 \sin^2 2\beta}{4(|m_{\chi_\alpha^\pm}|^2 + \frac{f_\alpha^2 v^2}{2}(1 + \sin^2 2\beta))}. \tag{13}$$

Equations (12) and (13) demonstrate that the upper bound on the mass of the lightest Inert neutralino also depends on the values of Yukawa couplings f_α and $\tilde{f}_\alpha$. For relatively small values of $\tan\beta$, the theoretical restrictions on f_α and $\tilde{f}_\alpha$, due to the requirement that the perturbation theory is valid up to the GUT scale, become weaker with increasing $\tan\beta$. However, at large values of $\tan\beta$ the upper bounds on $|m_{\chi_\alpha^0}|$ become rather small according to Eqs. (12) and (13). When $\tan\beta$ tends to unity, μ_α^2 also decreases because the constraints on f_α and $\tilde{f}_\alpha$ become more and more stringent. The theoretical restrictions on $|m_{\chi_\alpha^0}|$ achieve their maximal value around $\tan\beta \simeq 1.5$. 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 $|m_{\chi_\alpha^\pm}| > 100$ GeV.

The Inert neutralino mass matrix (7) and (8) can be diagonalized using the neutralino mixing matrix defined by

$$N_i^a M^{ab} N_j^b = m_i \delta_{ij}, \qquad \text{no sum on } i. \tag{14}$$

In the limit where off-diagonal Yukawa couplings vanish and $\lambda_\alpha s \gg f_\alpha v, \tilde{f}_\alpha v$ the eigenvalues of the Inert neutralino mass matrix can be easily calculated (see [26]). The masses of the four heaviest Inert neutralinos are set by the masses of Inert chargino states

$$m_{\chi_{3,4,5,6}^0} \simeq \pm m_{\chi_\alpha^\pm} - \frac{\tilde{f}_\alpha f_\alpha v^2 \sin 2\beta}{4 m_{\chi_\alpha^\pm}}. \tag{15}$$

The masses of the two lightest Inert neutralinos are determined by the values of the Yukawa couplings $\tilde{f}_\alpha$ and f_α:

$$m_{\chi_\alpha^0} \simeq \frac{\tilde{f}_\alpha f_\alpha v^2 \sin 2\beta}{2 m_{\chi_\alpha^\pm}}. \tag{16}$$

These are naturally small and hence good candidates for being the LSP and NLSP since $m_{\chi_\alpha^\pm} \sim s$ from Eq. (11) and hence $m_{\chi_\alpha^0} \sim v^2/s$ as observed in [26].

Again one can see that the masses of the lightest Inert neutralino states decrease with increasing $\tan\beta$ and chargino masses. In this approximation the lightest Inert neutralinos are made up of the following superposition of interaction states:

$$\tilde{\chi}_\alpha^0 = N_\alpha^1 \tilde{H}_2^{d0} + N_\alpha^2 \tilde{H}_2^{u0} + N_\alpha^3 \tilde{S}_2 + N_\alpha^4 \tilde{H}_1^{d0}$$
$$+ N_\alpha^5 \tilde{H}_1^{u0} + N_\alpha^6 \tilde{S}_1, \tag{17}$$

where

$$N_1^1 = N_1^2 = N_1^3 = 0,$$

$$N_1^4 \simeq -\frac{f_1 v \cos\beta}{\lambda_1 s},$$

$$N_1^5 \simeq -\frac{\tilde{f}_1 v \sin\beta}{\lambda_1 s},$$

$$N_1^6 \simeq 1 - \frac{1}{2}\left(\frac{v}{\lambda_1 s}\right)^2 [f_1^2 \cos^2\beta + \tilde{f}_1^2 \sin^2\beta]; \tag{18}$$

$$N_2^1 \simeq -\frac{f_2 v \cos\beta}{\lambda_2 s},$$

$$N_2^2 \simeq -\frac{\tilde{f}_2 v \sin\beta}{\lambda_2 s},$$

$$N_2^4 = N_2^5 = N_2^6 = 0,$$

$$N_2^3 \simeq 1 - \frac{1}{2}\left(\frac{v}{\lambda_2 s}\right)^2 [f_2^2 \cos^2\beta + \tilde{f}_2^2 \sin^2\beta].$$

From Eq. (18) it becomes clear that the lightest and second lightest Inert neutralinos are mostly Inert singlinos.

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 χ_1^0 and χ_2^0, can be presented in the following form:

$$\mathcal{L}_{Z\chi\chi} = \sum_{\alpha,\beta} \frac{M_Z}{2\upsilon} Z_\mu (\chi_\alpha^{0T} \gamma_\mu \gamma_5 \chi_\beta^0) R_{Z\alpha\beta},$$

$$R_{Z\alpha\beta} = N_\alpha^1 N_\beta^1 - N_\alpha^2 N_\beta^2 + N_\alpha^4 N_\beta^4 - N_\alpha^5 N_\beta^5. \tag{19}$$

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

$$R_{Z\alpha\beta} = R_{Z\alpha\alpha} \delta_{\alpha\beta},$$

$$R_{Z\alpha\alpha} = \frac{\upsilon^2}{2m_{\chi_\alpha^\pm}^2} (f_\alpha^2 \cos^2\beta - \tilde{f}_\alpha^2 \sin^2\beta). \tag{20}$$

Equation (20) demonstrates that the couplings of χ_1^0 and χ_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 χ_α^0 contains a completely symmetric combination of $\tilde{H}_\alpha^{d0}$ and $\tilde{H}_\alpha^{u0}$. Equation (20) also indicates that the couplings of χ_1^0 and χ_2^0 to Z are always small when Inert charginos are rather heavy or $\tilde{f}_\alpha$ and f_α are small (i.e. $m_{\chi_\alpha^0} \to 0$).

B. Δ_{27} and pseudo-Dirac lightest neutralino states

In order to provide an explanation of the origin of the approximate Z_2^H symmetry that singles out the third family of Higgs doublets and singlets, and to account for tribimaximal mixing and other features of the quark and lepton spectrum, a Δ_{27} family symmetry has been applied to the E$_6$SSM [33].[4] The addition of a Δ_{27} family symmetry implies an Inert neutralino mass matrix with $A_{11} \approx A_{22} \approx 0$, where $A_{\alpha\beta}$ are defined in Eq. (7), leading to approximately degenerate lightest neutralino states with a pseudo-Dirac structure.

When all flavor diagonal Yukawa couplings $\lambda_{\alpha\alpha}, f_{\alpha\alpha}$ and $\tilde{f}_{\alpha\alpha}$ exactly vanish, i.e. $A_{11} = A_{22} = 0$, all Inert Higgsinos and singlinos form Dirac states. In this limit the Lagrangian of the E$_6$SSM is invariant under a $U(1)$ global symmetry. The fermion components of the Inert Higgs superfields transform under this symmetry as follows:

$$\tilde{S}_1 \to e^{i\alpha} \tilde{S}_1, \qquad \tilde{H}_1^u \to e^{i\alpha} \tilde{H}_1^u, \qquad \tilde{H}_1^d \to e^{i\alpha} \tilde{H}_1^d,$$

$$\tilde{S}_2 \to e^{-i\alpha} \tilde{S}_2, \qquad \tilde{H}_2^u \to e^{-i\alpha} \tilde{H}_2^u, \qquad \tilde{H}_2^d \to e^{-i\alpha} \tilde{H}_2^d. \tag{21}$$

In the above limiting case the lightest Inert neutralino is a Dirac state formed predominantly by $\tilde{S}_1$ and $\tilde{S}_2$. In this case

---

[4]The corresponding mass terms come from the product $(3 \times 3 \times 3_f)(3 \times \bar{3}'_f)$, where 3, 3_f and $\bar{3}'_f$ are triplet representations of Δ_{27}. The 27_i multiplets that contain quarks and leptons form 3 representations of the Δ_{27} group. 3_f and $\bar{3}'_f$ contain flavon fields that break Δ_{27}. In the considered model the nonzero mass of the lightest Inert neutralino state is induced by the symmetric invariant that appears in the $(3 \times 3 \times 3)$ decomposition of the Δ_{27} triplet representation (i.e. $123 + 231 + 312 + 213 + 321 + 132$) [34].

the LSP and its antiparticle have opposite charges with respect to the extra global $U(1)$ and this might lead to so-called asymmetric dark matter (ADM) [35–37]. In the framework of the ADM scenario there can be an asymmetry between the density of dark matter particles and their anti-particles in the early Universe similar to that for ordinary baryons. This may have a considerable effect on the relic density calculations [36]. In particular, if an asymmetry exists between the number density of dark matter particles and their antiparticles in the early Universe, then one can get an appreciable dark matter density even if the dark matter particle-antiparticle annihilation cross section is very large like in the case of baryons. Moreover, if most of the dark matter antiparticles are eliminated by annihilation with their particles then such an ADM scenario does not have the usual indirect signatures associated with the presence of dark matter (e.g. there is no high energy neutrino signal from annihilations in the Sun etc.). At the same time, a relatively high concentration of dark matter particles can build up in the Sun altering heat transport in the solar interior and affecting the low energy neutrino fluxes [37].

In practice the Δ_{27} scenario tells us that we are somewhat away from the above limiting case, with a broken global $U(1)$ symmetry leading to almost degenerate pseudo-Dirac lightest neutralinos, where the relic density of the LSP can be calculated by standard methods. It will turn out that the LSP cannot be too light (must be of order $M_Z/2$) in order not to have too high a cosmological relic density. At the same time we will see that the two lightest neutralinos cannot be too heavy in order for perturbation theory to be valid up to the GUT scale. In practice this means that in realistic scenarios the two lightest Inert neutralino states are rather close in mass. The Δ_{27} scenario provides a natural explanation of this successful neutralino mass pattern. It is worth noting that the results from the previous section can be reinterpreted in terms of this scenario. Specifically in the case where $A_{11} = A_{22} = 0$ and $A_{21} = A_{12}$ a block diagonalization of the Inert neutralino mass matrix (7) yields $A_{22} \to A'_{22} = -A_{21}$ and $A_{11} \to A'_{11} = A_{21}$ (with $A_{21} = A_{12} \to A'_{21} = A'_{12} = 0$). This only corresponds to a redefinition of the generations 1 and 2 and does not mix fields of different hypercharge. This provides the following dictionary between these two scenarios: $\lambda'_{11} = -\lambda'_{22} = \lambda_{21}; f'_{11} = -f'_{22} = f_{21}; \tilde{f}'_{11} = -\tilde{f}'_{22} = \tilde{f}_{21}$. Rewriting the Inert neutralino mass matrix in this block diagonal form also makes it clear that the R_{Z12} coupling vanishes in this limit, as it did in the Sec. III A.

C. Scenario with one light family of Inert Higgsinos

Another limit that it is worth considering corresponds to the case where one pair of Inert Higgs doublets decouples from the rest of the spectrum. This occurs when either the corresponding states are extremely heavy ($\gtrsim 1$ TeV) or they have rather small couplings to other Inert Higgs fields. When $\tilde{H}_2^{d0}$ and $\tilde{H}_2^{u0}$ decouple, the Inert neutralino mass

matrix (7) reduces to a 4×4 matrix. If $\lambda_{11}s \gg f_{\alpha 1}v$, $\tilde{f}_{\alpha 1}v$, the Inert Higgs states associated with $\hat{H}_1^{d0}$ and $\hat{H}_1^{u0}$ can be integrated out. Then the resulting 2×2 mass matrix can be written as follows:

$$M_{\mathrm{IS}} = \frac{v^2 \sin 2\beta}{4m_{\chi_1^\pm}} \begin{pmatrix} 2\tilde{f}_{11}f_{11} & \tilde{f}_{11}f_{21} + f_{11}\tilde{f}_{21} \\ \tilde{f}_{11}f_{21} + f_{11}\tilde{f}_{21} & 2\tilde{f}_{21}f_{21} \end{pmatrix}. \quad (22)$$

The masses of the lightest and second lightest Inert neutralinos, which are predominantly superpositions of the Inert singlinos $\tilde{S}_1$ and $\tilde{S}_2$, are given by

$$m_{\chi_1^0, \chi_2^0} = \frac{v^2 \sin 2\beta}{4m_{\chi_1^\pm}} [\tilde{f}_{11}f_{11} + \tilde{f}_{21}f_{21}$$

$$\pm \sqrt{(f_{11}^2 + f_{21}^2)(\tilde{f}_{11}^2 + \tilde{f}_{21}^2)}]. \quad (23)$$

From Eq. (23) it is easy to see that the substantial masses of the lightest and second lightest Inert neutralinos can be induced even if only one family of the Inert Higgsinos couples to S_1 and S_2.

Using Eq. (19) one can also calculate the couplings of χ_1^0 and χ_2^0 to the Z boson:

$$R_{Z11} = \frac{v^2}{2m_{\chi_1^\pm}^2} [(f_{11}\cos\theta + f_{21}\sin\theta)^2\cos^2\beta - (\tilde{f}_{11}\cos\theta + \tilde{f}_{21}\sin\theta)^2\sin^2\beta],$$

$$R_{Z22} = \frac{v^2}{2m_{\chi_1^\pm}^2} [(f_{21}\cos\theta - f_{11}\sin\theta)^2\cos^2\beta - (\tilde{f}_{21}\cos\theta - \tilde{f}_{11}\sin\theta)^2\sin^2\beta],$$

$$R_{Z12} = R_{Z21} = \frac{v^2}{2m_{\chi_1^\pm}^2}\left[\left(\frac{1}{2}(f_{21}^2 - f_{11}^2)\sin 2\theta + f_{11}f_{21}\cos 2\theta\right)\cos^2\beta\right.$$

$$\left. - \left(\frac{1}{2}(\tilde{f}_{21}^2 - \tilde{f}_{11}^2)\sin 2\theta + \tilde{f}_{11}\tilde{f}_{21}\cos 2\theta\right)\sin^2\beta\right], \quad (24)$$

where $\tan 2\theta = (\tilde{f}_{11}f_{21} + \tilde{f}_{21}f_{11})/(\tilde{f}_{21}f_{21} - \tilde{f}_{11}f_{11})$. Again from Eqs. (24) it follows that R_{Z11}, R_{Z22} and R_{Z12} are typically small since $m_{\chi_\alpha^\pm} \sim s$ from Eq. (11) and hence they are proportional to v^2/s^2. However this assumes the lightest Inert chargino is rather heavy. Alternatively the couplings may be small due to a cancellation between different contributions in Eqs. (24), and/or the f couplings being small (i.e. $m_{\chi_1^0, \chi_2^0} \to 0$).

The simple hierarchical structure of the spectrum of the Inert neutralinos considered above allows us to highlight an interesting scenario which does not normally appear in the simplest SUSY extensions of the SM such as the MSSM and NMSSM. When $\tilde{f}_{11} = f_{21} = 0$ the diagonal entries of the mass matrix (22) vanish leading to the formation of a Dirac lightest Inert neutralino state. In this case the Lagrangian of the model is invariant under extra $U(1)$ global symmetry transformations[5] $\tilde{S}_1 \to e^{i\alpha}\tilde{S}_1$, $\hat{H}_1^u \to e^{-i\alpha}\hat{H}_1^u$, $\hat{H}_1^d \to e^{i\alpha}\hat{H}_1^d$, $\tilde{S}_2 \to e^{-i\alpha}\tilde{S}_2$. In fact if the E_6SSM possess such an exact $U(1)$ global symmetry, then the spectrum of the Inert neutralinos contains a set of Dirac states only.

IV. HIGGS MASSES AND COUPLINGS

The presence of light Inert neutralinos in the particle spectrum of the E_6SSM makes possible the decays of the Higgs bosons into these exotic final states. In this and the next section we argue that such decays may result in

the modification of the SM-like Higgs signal at current and future colliders. Since our main concern in this paper is the decays of the SM-like lightest Higgs boson, we shall ignore the effects of the Inert Higgs scalars and pseudoscalars which do not mix appreciably with the scalar sector responsible for EWSB. We also assume that all the Inert bosons are heavier than the SM-like Higgs boson.

The sector responsible for the EWSB in the E_6SSM includes two Higgs doublets H_u and H_d as well as the SM-singlet field S. The Higgs effective potential can be written in the following form:

$$V = V_F + V_D + V_{\mathrm{soft}} + \Delta V,$$

$$V_F = \lambda^2|S|^2(|H_d|^2 + |H_u|^2) + \lambda^2|(H_dH_u)|^2,$$

$$V_D = \frac{g_2^2}{8}(H_d^\dagger\sigma_a H_d + H_u^\dagger\sigma_a H_u)^2 + \frac{g'^2}{8}(|H_d|^2 - |H_u|^2)^2$$

$$+ \frac{g_1'^2}{2}(\tilde{Q}_1|H_d|^2 + \tilde{Q}_2|H_u|^2 + \tilde{Q}_S|S|^2)^2,$$

$$V_{\mathrm{soft}} = m_S^2|S|^2 + m_1^2|H_d|^2 + m_2^2|H_u|^2$$

$$+ [\lambda A_\lambda S(H_uH_d) + \mathrm{H.c.}], \quad (25)$$

where g_2, $g' = \sqrt{3/5}g_1$ and g_1' are the low energy $SU(2)_W$, $U(1)_Y$ and $U(1)_N$ gauge couplings while $\tilde{Q}_1$, $\tilde{Q}_2$ and $\tilde{Q}_S$ are the effective $U(1)_N$ charges of H_d, H_u and S. The term ΔV represents the contribution from loop corrections to the Higgs effective potential. Here $H_d^T = (H_d^0, H_d^-)$, $H_u^T = (H_u^+, H_u^0)$ and $(H_dH_u) = H_u^+H_d^- - H_u^0H_d^0$.

---

[5]Similar results can be obtained for $f_{11} = \tilde{f}_{21} = 0$.

Initially the EWSB sector involves 10 degrees of freedom. However four of them are massless Goldstone modes which are swallowed by the $W^\pm$, Z and Z' gauge bosons that gain nonzero masses. In the limit where $s \gg v$ the masses of the $W^\pm$, Z and Z' gauge bosons are given by

$$M_W = \frac{g_2}{2} v, \qquad M_Z \simeq \frac{\bar{g}}{2} v, \qquad M_{Z'} \simeq g_1' \tilde{Q}_S s,$$

where $\bar{g} = \sqrt{g_2^2 + g'^2}$. When *CP*-invariance is preserved the other degrees of freedom form two charged, one *CP*-odd and three *CP*-even Higgs states. The masses of the charged and *CP*-odd Higgs bosons are

$$m_{H^\pm}^2 = \frac{\sqrt{2}\lambda A_\lambda}{\sin 2\beta} s - \frac{\lambda^2}{2} v^2 + M_W^2 + \Delta_\pm,$$

$$m_A^2 \simeq \frac{\sqrt{2}\lambda A_\lambda}{\sin 2\beta} s + \Delta_A, \tag{26}$$

where $\Delta_\pm$ and Δ_A are the loop corrections.

The *CP*-even Higgs sector involves $\mathrm{Re}H_d^0$, $\mathrm{Re}H_u^0$ and $\mathrm{Re}S$. In the field space basis (h, H, N), rotated by an angle β with respect to the initial one,

$$\mathrm{Re}H_d^0 = (h\cos\beta - H\sin\beta + v_1)/\sqrt{2},$$

$$\mathrm{Re}H_u^0 = (h\sin\beta + H\cos\beta + v_2)/\sqrt{2}, \tag{27}$$

$$\mathrm{Re}S = (s + N)/\sqrt{2},$$

the mass matrix of the *CP*-even Higgs sector takes the form [38]

$$M^2 = \begin{pmatrix} \frac{\partial^2 V}{\partial v^2} & \frac{1}{v}\frac{\partial^2 V}{\partial v \partial \beta} & \frac{\partial^2 V}{\partial v \partial s} \\ \frac{1}{v}\frac{\partial^2 V}{\partial v \partial \beta} & \frac{1}{v^2}\frac{\partial^2 V}{\partial^2 \beta} & \frac{1}{v}\frac{\partial^2 V}{\partial s \partial \beta} \\ \frac{\partial^2 V}{\partial v \partial s} & \frac{1}{v}\frac{\partial^2 V}{\partial s \partial \beta} & \frac{\partial^2 V}{\partial^2 s} \end{pmatrix}$$

$$= \begin{pmatrix} M_{11}^2 & M_{12}^2 & M_{13}^2 \\ M_{21}^2 & M_{22}^2 & M_{23}^2 \\ M_{31}^2 & M_{32}^2 & M_{33}^2 \end{pmatrix}, \tag{28}$$

where

$$M_{11}^2 = \frac{\lambda^2}{2} v^2 \sin^2 2\beta + \frac{\bar{g}^2}{4} v^2 \cos^2 2\beta + g_1'^2 v^2 (\tilde{Q}_1 \cos^2\beta + \tilde{Q}_2 \sin^2\beta)^2 + \Delta_{11},$$

$$M_{12}^2 = M_{21}^2 = \left(\frac{\lambda^2}{4} - \frac{\bar{g}^2}{8}\right) v^2 \sin 4\beta + \frac{g_1'^2}{2} v^2 (\tilde{Q}_2 - \tilde{Q}_1) \times (\tilde{Q}_1 \cos^2\beta + \tilde{Q}_2 \sin^2\beta)\sin 2\beta + \Delta_{12},$$

$$M_{22}^2 = \frac{\sqrt{2}\lambda A_\lambda}{\sin 2\beta} s + \left(\frac{\bar{g}^2}{4} - \frac{\lambda^2}{2}\right) v^2 \sin^2 2\beta + \frac{g_1'^2}{4}(\tilde{Q}_2 - \tilde{Q}_1)^2 v^2 \sin^2 2\beta + \Delta_{22},$$

$$M_{23}^2 = M_{32}^2 = -\frac{\lambda A_\lambda}{\sqrt{2}} v\cos 2\beta + \frac{g_1'^2}{2}(\tilde{Q}_2 - \tilde{Q}_1)\tilde{Q}_S vs \sin 2\beta + \Delta_{23},$$

$$M_{13}^2 = M_{31}^2 = -\frac{\lambda A_\lambda}{\sqrt{2}} v\sin 2\beta + \lambda^2 vs + g_1'^2(\tilde{Q}_1 \cos^2\beta + \tilde{Q}_2 \sin^2\beta)\tilde{Q}_S vs + \Delta_{13}, \tag{29}$$

$$M_{33}^2 = \frac{\lambda A_\lambda}{2\sqrt{2}s} v^2 \sin 2\beta + M_{Z'}^2 + \Delta_{33}.$$

In Eq. (29) the Δ_{ij} represent the contributions from loop corrections which in the leading one-loop approximation are rather similar to the ones calculated in the NMSSM.[6] Since the minimal eigenvalue of the mass matrix (28) and (29) is always less than its smallest diagonal element, at least one Higgs scalar in the *CP*–even sector (approximately *h*) remains always light, i.e. $m_{h_1}^2 \lesssim M_{11}^2$. In the leading two-loop approximation the mass of the lightest Higgs boson in the E_6SSM does not exceed 150–155 GeV. When the SUSY breaking scale M_S and the VEV *s* of the singlet field are considerably larger than the EW scale, the mass matrix (28) and (29) has a

hierarchical structure and can be diagonalized using the perturbation theory [38–40]. In this case the masses of the heaviest Higgs bosons are closely approximated by the diagonal entries M_{22}^2 and M_{33}^2 [23]. As a result the mass of one *CP*-even Higgs boson (approximately given by *H*) is governed by m_A while the mass of another one (predominantly the *N* singlet field) is set by $M_{Z'}$. When $\lambda \gtrsim g_1'$, vacuum stability requires m_A to be considerably larger than $M_{Z'}$ and 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 [40,41]. In the considered limit the heaviest *CP*-even, *CP*-odd and charged states are almost degenerate around m_A and lie beyond the TeV range [23].

If all other Higgs states are much heavier than the lightest *CP*-even Higgs boson then the lightest Higgs state (approximately given by *h*) manifests itself in the interactions with gauge bosons and fermions as a SM-like Higgs

---

[6]Note that the explicit expressions for Δ_{ij}, $\Delta_\pm$ and Δ_A presented in the first paper in [38] contain a typo. In the corresponding formula μ is neither a parameter of the MSSM Lagrangian nor an effective μ-term in the NMSSM. It has to be associated with the renormalization scale.

boson. Since within the E$_6$SSM the mass of this state is predicted to be relatively low its production cross section at the LHC should be large enough so that it can be observed in the near future. In this context it is particularly interesting and important to analyze the decay modes of the lightest *CP*-even Higgs state. Furthermore we concentrate on the decays of the SM-like Higgs boson into the lightest and second lightest Inert neutralinos.

The couplings of the Higgs states to the Inert neutralinos originate from the interactions of H_u, H_d and S with the Inert Higgs superfields in the superpotential. Using Eqs. (27) one can express ReH_d^0, ReH_u^0 and ReS in terms of the components of the *CP*-even Higgs basis h, H and N. At the same time the components of the *CP*-even Higgs basis are related to the physical *CP*-even Higgs eigenstates by virtue of a unitary transformation:

$$\begin{pmatrix} h \\ H \\ N \end{pmatrix} = U^\dagger \begin{pmatrix} h_1 \\ h_2 \\ h_3 \end{pmatrix}. \tag{30}$$

Combining all these expressions together one obtains an effective Lagrangian that describes the interactions of the Inert neutralinos with the *CP*-even Higgs eigenstates

$$\mathcal{L}_{H\chi\chi} = \sum_{i,j,m} (-1)^{\theta_i+\theta_j} X_{ij}^{h_m} (\psi_i^{0T}(-i\gamma_5)^{\theta_i+\theta_j}\psi_j^0) h_m,$$

$$X_{ij}^{h_m} = -\frac{1}{\sqrt{2}} U_{Nh_m}^\dagger \Lambda_{ij} - \frac{1}{\sqrt{2}}(U_{hh_m}^\dagger \cos\beta - U_{Hh_m}^\dagger \sin\beta)F_{ij}$$
$$- \frac{1}{\sqrt{2}}(U_{hh_m}^\dagger \sin\beta + U_{Hh_m}^\dagger \cos\beta)\tilde{F}_{ij},$$

$$F_{ij} = f_{11}N_i^6 N_j^5 + f_{12}N_i^6 N_j^2 + f_{21}N_i^3 N_j^5 + f_{22}N_i^3 N_j^2,$$

$$\tilde{F}_{ij} = \tilde{f}_{11}N_i^6 N_j^4 + \tilde{f}_{12}N_i^6 N_j^1 + \tilde{f}_{21}N_i^3 N_j^4 + \tilde{f}_{22}N_i^3 N_j^1,$$

$$\Lambda_{ij} = \lambda_{11}N_i^4 N_j^5 + \lambda_{12}N_i^4 N_j^2 + \lambda_{21}N_i^1 N_j^5 + \lambda_{22}N_i^1 N_j^2,$$

(31)

where $i, j = 1, 2, \ldots 6$ and $m = 1, 2, 3$. In Eq. (31) $\psi_i^0 = (-i\gamma_5)^{\theta_i}\chi_i^0$ is the set of Inert neutralino eigenstates with positive eigenvalues, while θ_i equals 0 (1) if the eigenvalue corresponding to χ_i^0 is positive (negative). As before, the Inert neutralinos are labeled according to increasing

absolute value of mass, with ψ_1^0 being the lightest Inert neutralino and ψ_6^0 the heaviest.

The expressions for the couplings of the Higgs scalars to the Inert neutralinos (31) become much more simple in the case of the hierarchical structure of the Higgs spectrum. In this case U_{ij} is almost an identity matrix. As a consequence, the couplings of the SM-like Higgs boson to the lightest and second lightest Inert neutralino states are approximately given by

$$X_{\gamma\sigma}^{h_1} = -\frac{1}{\sqrt{2}}(F_{\gamma\sigma}\cos\beta + \tilde{F}_{\gamma\sigma}\sin\beta), \tag{32}$$

where γ, $\sigma = 1, 2$, labeling the two light, mostly Inert singlino states. 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, as defined in Sec. III A, and Inert neutralino mass matrix has a hierarchical structure (i.e. $\lambda_\alpha s \gg f_\alpha v$, $\tilde{f}_\alpha v$), one can use the expressions (18) for $N_{1,2}^a$ in order to derive the approximate analytical formula for $X_{\gamma\sigma}^{h_1}$. Substituting Eqs. (18) into (32) one obtains

$$X_{\gamma\sigma}^{h_1} \simeq \frac{|m_{\chi_\sigma^0}|}{v}\delta_{\gamma\sigma}. \tag{33}$$

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 E$_6$SSM. 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. We see that this is exactly what is found in the limit of $|m_{\chi_\sigma^0}|$ being small.

V. NOVEL HIGGS DECAYS AND DARK MATTER

A. Higgs decay widths

The interaction Lagrangian (31) gives rise to decays of the lightest Higgs boson into Inert neutralino pairs with partial widths given by

$$\Gamma(h_1 \to \chi_\alpha^0 \chi_\beta^0) = \frac{\Delta_{\alpha\beta}}{8\pi m_{h_1}}\left(X_{\alpha\beta}^{h_1} + X_{\beta\alpha}^{h_1}\right)^2 \left[m_{h_1}^2 - (|m_{\chi_\alpha^0}| + (-1)^{\theta_\alpha+\theta_\beta}|m_{\chi_\beta^0}|)^2\right]$$

$$\times \sqrt{\left(1 - \frac{|m_{\chi_\alpha^0}|^2}{m_{h_1}^2} - \frac{|m_{\chi_\beta^0}|^2}{m_{h_1}^2}\right)^2 - 4\frac{|m_{\chi_\alpha^0}|^2 |m_{\chi_\beta^0}|^2}{m_{h_1}^4}}, \tag{34}$$

where $\Delta_{\alpha\beta} = \frac{1}{2}$ (1) for $\alpha = \beta$ ($\alpha \neq \beta$).

The partial widths associated with the exotic decays of the SM-like Higgs boson (34) have to be compared with the Higgs decay rates into the SM particles. When the SM-like

Higgs state is relatively light ($m_{h_1} \lesssim 140$ GeV) it decays predominantly into *b*-quark and τ-lepton pairs. The partial decay width of the lightest *CP*-even Higgs boson into fermion pairs is given by (for recent review see [42])

$$\Gamma(h_1 \to f\bar{f}) = N_c \frac{g_2^2}{32\pi} \left(\frac{m_f}{M_W}\right)^2 g_{h_1 ff}^2 m_{h_1} \left(1 - \frac{4m_f^2}{m_{h_1}^2}\right)^{3/2}.$$ (35)

Equation (35) can be used for the calculation of the lightest Higgs decay rate into τ-lepton pairs. In this case the coupling of the lightest CP-even Higgs state to the τ-lepton normalized to the corresponding SM coupling, i.e. $g_{h_1 \tau\tau}$, is given by

$$g_{h_1 \tau\tau} = \frac{1}{\cos\beta} \left(U_{hh_1}^\dagger \cos\beta - U_{Hh_1}^\dagger \sin\beta\right).$$ (36)

For a final state that involves b-quarks one has to include the QCD corrections. In particular, the fermion mass in Eq. (35) should be associated with the running b-quark mass $\bar{m}_b(\mu)$. The bulk of the QCD corrections are absorbed by using the running b-quark mass defined at the appropriate renormalization scale, i.e. at the scale of the lightest Higgs boson mass ($\mu = m_{h_1}$) in the considered case. In addition to the corrections which are associated with the running b-quark mass there are other QCD corrections to the Higgs coupling to the b-quark that should be taken into account [43]. As a consequence, the partial decay width of the lightest CP-even Higgs boson into b-quark pairs can be calculated using Eq. (35) if one sets $N_c = 3$ and replaces

$$m_f \to \bar{m}_b(m_{h_1}),$$
$$g_{h_1 ff}^2 \to \frac{1}{\cos^2\beta} \left(U_{hh_1}^\dagger \cos\beta - U_{Hh_1}^\dagger \sin\beta\right)^2 \left[1 + \Delta_{bb} + \Delta_H\right],$$
$$\Delta_{bb} \simeq 5.67 \frac{\bar{\alpha}_s}{\pi} + (35.94 - 1.36 N_f) \frac{\bar{\alpha}_s^2}{\pi^2},$$
$$\Delta_H \simeq \frac{\bar{\alpha}_s^2}{\pi^2} \left(1.57 - \frac{2}{3} \log \frac{m_{h_1}^2}{m_t^2} + \frac{1}{9} \log^2 \frac{\bar{m}_b^2}{m_{h_1}^2}\right),$$ (37)

where $\bar{\alpha}_s = \alpha_s(m_{h_1}^2)$. Here we neglect radiative corrections that originate from loop diagrams that contain SUSY and exotic particles.[7]

From Eqs. (33)–(35) one can see that in the E$_6$SSM the branching ratios of the SM-like Higgs state into the lightest and second lightest Inert neutralinos depend rather strongly on the masses of these exotic particles. When the lightest Inert neutralino states are relatively heavy, i.e. $m_{\chi_1, \chi_2} \gtrsim \bar{m}_b(m_{h_1})$, the lightest Higgs boson decays predominantly into $\chi_\alpha \chi_\beta$ while the branching ratios for decays into SM particles are suppressed. On the other hand if the lightest Inert neutralino have masses which are considerably smaller than the masses of the b-quark and τ-lepton then

---

[7]Radiative corrections that are induced by SUSY particles can be very important particularly in the case of the bottom quark at high values of $\tan\beta$ (for a review, see [44]).

---

the branching ratios of the exotic decays of the SM-like Higgs state are small. In the E$_6$SSM the lightest and second lightest Inert neutralinos are expected to be heavier than a few MeV so that they would not contribute to the expansion rate prior to nucleosynthesis and thus not modify Big Bang Nucleosynthesis (BBN).

B. Dark matter

More stringent constraints on the masses of the lightest Inert neutralino can be obtained if we require that this exotic state accounts for all or some of the observed dark matter relic density which is measured to be $\Omega_{CDM}h^2 = 0.1099 \pm 0.0062$ [45]. 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 E$_6$SSM results in too large a density of dark matter. As we noted in Sec. III, $\tilde{\chi}_1^0$ is usually composed of Inert singlino and has a mass (Eq. (16)) which is inversely proportional to the charged Higgsino mass. Thus in this limit it is typically very light $|m_{\chi_\sigma^0}| \ll M_Z$. As a result the couplings of the lightest Inert neutralino to gauge bosons, the SM-like Higgs state, quarks and leptons are quite small leading to a relatively small annihilation cross section for $\tilde{\chi}_1^0 \tilde{\chi}_1^0 \to$ SM particles. Since the dark matter number density is inversely proportional to the annihilation cross section at the freeze-out temperature (see, for example [46]) the lightest Inert neutralino state gives rise to a relic density that is typically much larger than its measured value. Thus in the limit considered the bulk of the E$_6$SSM parameter space that leads to small masses of $\tilde{\chi}_1^0$ is ruled out.

The situation changes dramatically when the mass of the lightest Inert neutralino increases. In this case the Higgsino components of $\tilde{\chi}_1^0$ become larger and as a consequence the couplings of $\tilde{\chi}_1^0$ to the Z boson grow [26]. A reasonable density of dark matter can be obtained for $|m_{\chi_\sigma^0}| \sim M_Z/2$ when the lightest Inert neutralino states annihilate mainly through an s channel Z boson, via its Inert Higgsino doublet components which couple to the Z boson. It is worth noting that if $\tilde{\chi}_1^0$ was pure Inert Higgsino then the s channel Z boson annihilation would proceed with the full gauge coupling strength leaving the relic density too low to account for the observed dark matter. In the E$_6$SSM the LSP is mostly Inert singlino so that its coupling to the Z boson is typically suppressed, since it only couples through its Inert Higgsino admixture leading to an increased relic density. In practice, the appropriate value of $\Omega_{CDM}h^2$ can be achieved even if the coupling of $\tilde{\chi}_1^0$ to the Z boson is relatively small. This happens when $\tilde{\chi}_1^0$ annihilation proceeds through the Z boson resonance, i.e. $2|m_{\chi_\sigma^0}| \simeq M_Z$

[26,47]. Thus scenarios which result in a reasonable dark matter density correspond to lightest Inert neutralino masses that are much larger than $\bar{m}_b(m_{h_1})$, and hence the SM-like Higgs has very small branching ratios into SM particles.

VI. BENCHMARKS, CONSTRAINTS AND PREDICTIONS

In order to illustrate the features of the E_6SSM mentioned in the previous section, we shall specify a set of benchmark points (see Tables I and II). For each

TABLE I.   Benchmark scenarios for $m_{h_1} \approx 133$–135 GeV. 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, $\lambda = 0.6$, $A_\lambda = 1600$ GeV, $m_Q = m_U = M_S = 700$ GeV, $X_t = \sqrt{6}M_S$ that correspond to $m_{h_2} \simeq M_{Z'} \simeq 890$ GeV.

	i	ii	iii	iv
$\tan(\beta)$	1.5	1.5	1.7	1.564
$m_{H^\pm} \simeq m_A \simeq m_{h_3}$/GeV	1977	1977	2022	1990
m_{h_1}/GeV	135.4	135.4	133.1	134.8
λ_{22}	0.001	0.001	0.094	0.0001
λ_{21}	0.077	0.062	0	0.06
λ_{12}	0.077	0.062	0	0.06
λ_{11}	0.001	0.001	0.059	0.0001
f_{22}	0.001	0.001	0.53	0.001
f_{21}	0.61	0.61	0.05	0.476
f_{12}	0.6	0.6	0.05	0.466
f_{11}	0.001	0.001	0.53	0.001
$\tilde{f}_{22}$	0.001	0.001	0.53	0.001
$\tilde{f}_{21}$	0.426	0.426	0.05	0.4
$\tilde{f}_{12}$	0.436	0.436	0.05	0.408
$\tilde{f}_{11}$	0.001	0.001	0.53	0.001
$m_{\tilde{\chi}_1^0}$/GeV	41.91	47.33	33.62	-36.69
$m_{\tilde{\chi}_2^0}$/GeV	-42.31	-47.84	47.78	36.88
$m_{\tilde{\chi}_3^0}$/GeV	-129.1	-103.6	108.0	-103.11
$m_{\tilde{\chi}_4^0}$/GeV	132.4	107.0	-152.1	103.47
$m_{\tilde{\chi}_5^0}$/GeV	171.4	151.5	163.5	139.80
$m_{\tilde{\chi}_6^0}$/GeV	-174.4	-154.4	-200.8	-140.35
$m_{\tilde{\chi}_1^\pm}$/GeV	129.0	103.5	100.1	101.65
$m_{\tilde{\chi}_2^\pm}$/GeV	132.4	106.9	159.5	101.99
$\Omega_\chi h^2$	0.096	0.098	0.109	0.107
R_{Z11}	-0.0250	-0.0407	-0.144	-0.132
R_{Z12}	0.0040	0.0048	0.051	0.0043
R_{Z22}	-0.0257	-0.0429	-0.331	-0.133
$\Delta N_\nu^{\text{eff}}$	0.000090	0	0.0068	0.0073
D	2.011	2.000	2.85	2.91
$X_{11}^{h_1}$	0.137	0.147	0.110	-0.114
$X_{12}^{h_1} + X_{21}^{h_1}$	-1.9×10^{-6}	-3.4×10^{-6}	0.0136	1.15×10^{-6}
$X_{22}^{h_1}$	-0.138	-0.148	0.125	0.115
$\sigma_{\text{SI}}/10^{-44}$ cm^2	2.6–10.5	3.0–12.1	1.7–7.1	2.0–8.2
Br$(h \to \tilde{\chi}_1^0 \tilde{\chi}_1^0)$	49.5%	49.7%	57.8%	49.1%
Br$(h \to \tilde{\chi}_1^0 \tilde{\chi}_2^0)$	7.9×10^{-11}	2.5×10^{-10}	0.34%	49.2%
Br$(h \to \tilde{\chi}_2^0 \tilde{\chi}_2^0)$	49.0%	48.5%	39.8%	3.5×10^{-11}
Br$(h \to b\bar{b})$	1.36%	1.58%	1.87%	1.59%
Br$(h \to \tau\bar{\tau})$	0.142%	0.165%	0.196%	0.166%
$\Gamma(h \to \tilde{\chi}_1^0 \tilde{\chi}_1^0)$/MeV	98.3	85.1	81.7	82.9
Γ^{tot}/MeV	198.7	171.1	141.2	169.0

benchmark scenario we calculate the spectrum of the Inert neutralinos, Inert charginos and Higgs bosons as well as their couplings, the branching ratios of the decays of the lightest *CP*-even Higgs state and the dark matter relic density. In order to calculate the dark matter relic density

we use numerical methods. In particular, MICROMEGAS 2.2 [48] is used to numerically compute the present day density of dark matter. This includes the relevant (co-)annihilation channel cross sections and the LSP freeze-out temperature. MICROMEGAS achieves this by calculating all of the

TABLE II. Benchmark scenarios for $m_{h_1} \approx 114$–116 GeV. 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, $\lambda = g_1' = 0.468$, $A_\lambda = 600$ GeV, $m_Q = m_U = M_S = 700$ GeV, $X_t = \sqrt{6}M_S$ that correspond to $m_{h_2} \simeq M_{Z'} \simeq 890$ GeV.

	v	vi	vii	viii	ix
$\tan(\beta)$	1.5	1.7	1.5	1.5	1.5
$m_{H^\pm} \simeq m_A \simeq m_{h_3}$/GeV	1145	1165	1145	1145	1145
m_{h_1}/GeV	115.9	114.4	115.9	115.9	115.9
λ_{22}	0.004	0.104	0.094	0.001	0.468
λ_{21}	0.084	0	0	0.079	0.05
λ_{12}	0.084	0	0	0.080	0.05
λ_{11}	0.004	0.09	0.059	0.001	0.08
f_{22}	0.025	0.72	0.53	0.04	0.05
f_{21}	0.51	0.001	0.053	0.68	0.9
f_{12}	0.5	0.001	0.053	0.68	0.002
f_{11}	0.025	0.7	0.53	0.04	0.002
$\tilde{f}_{22}$	0.025	0.472	0.53	0.04	0.002
$\tilde{f}_{21}$	0.49	0.001	0.053	0.49	0.002
$\tilde{f}_{12}$	0.5	0.001	0.053	0.49	0.05
$\tilde{f}_{11}$	0.025	0.472	0.53	0.04	0.65
$m_{\tilde{\chi}_1^0}$/GeV	-35.76	41.20	35.42	-45.08	-46.24
$m_{\tilde{\chi}_2^0}$/GeV	39.63	44.21	51.77	55.34	46.60
$m_{\tilde{\chi}_3^0}$/GeV	-137.8	153.1	105.3	-133.3	171.1
$m_{\tilde{\chi}_4^0}$/GeV	151.7	176.7	-152.7	136.9	-171.4
$m_{\tilde{\chi}_5^0}$/GeV	173.6	-197.3	162.0	178.4	805.4
$m_{\tilde{\chi}_6^0}$/GeV	-191.3	-217.9	-201.7	-192.2	-805.4
$m_{\tilde{\chi}_1^\pm}$/GeV	135.8	152.7	100.1	133.0	125.0
$m_{\tilde{\chi}_2^\pm}$/GeV	149.3	176.5	159.5	136.8	805.0
$\Omega_\chi h^2$	0.102	0.108	0.107	0.0324	0.00005
R_{Z11}	-0.116	-0.0278	-0.115	-0.0217	-0.0224
R_{Z12}	0.0037	-0.00039	-0.045	-0.0020	-0.213
R_{Z22}	-0.118	-0.0455	-0.288	-0.0524	-0.0226
$\Delta N_\nu^{\mathrm{eff}}$	0.0049	0.00009	0.0034	1.57×10^{-6}	0
D	2.62	2.011	2.43	2.0002	2.0
$X_{11}^{h_1}$	-0.117	0.141	0.117	-0.147	-0.148
$X_{12}^{h_1} + X_{21}^{h_1}$	-0.000027	-0.00025	-0.0127	-0.0000140	-0.000031
$X_{22}^{h_1}$	0.130	0.147	0.141	0.174	0.149
$\sigma_{\mathrm{SI}}/10^{-44}$ cm^2	3.9–15.7	5.4–21.9	3.5–14.2	6.0–24.4	6.1–25.0
Br($h \to \tilde{\chi}_1^0 \tilde{\chi}_1^0$)	49.6%	53.5%	76.3%	83.4%	49.3%
Br($h \to \tilde{\chi}_1^0 \tilde{\chi}_2^0$)	2.1×10^{-8}	7.2×10^{-7}	0.26%	7.6×10^{-9}	3.0×10^{-8}
Br($h \to \tilde{\chi}_2^0 \tilde{\chi}_2^0$)	48.4%	44.2%	20.3%	12.3%	47.9%
Br($h \to b\bar{b}$)	1.87%	2.04%	2.83%	3.95%	2.58%
Br($h \to \tau\bar{\tau}$)	0.196%	0.21%	0.30%	0.41%	0.27%
$\Gamma(h \to \tilde{\chi}_1^0 \tilde{\chi}_1^0)$/MeV	61.5	60.1	62.6	49.0	44.4
Γ^{tot}/MeV	124.1	112.2	82.0	58.8	90.1

relevant tree-level Feynman diagrams using CALCHEP. The CALCHEP model files for the considered model are generated using LANHEP [49]. The MICROMEGAS relic density calculation assumes standard cosmology in which the LSP was in equilibrium with the photon at some time in the past.

A. Benchmark scenarios

In order to construct benchmark scenarios that are consistent with cosmological observations and collider constraints we restrict our considerations to low values of $\tan\beta \lesssim 2$. Figures 1 and 2 show that in principle the appropriate value of dark matter density can be obtained even when $\tan\beta > 2$. At the same time larger values of $\tan\beta$ lead to masses of the lightest and second lightest Inert neutralinos that are too small, as discussed in Sec. III. As a result larger couplings of the lightest Inert neutralinos to Z are required to reproduce the measured value of $\Omega_{\text{CDM}}h^2$. On the other hand according to Figs. 1 and 2, light Inert neutralinos with substantial couplings to Z boson give a considerable contribution to its invisible width leading to a conflict with LEP measurements (see discussion in Sec. VI B).

However, even for $\tan\beta \lesssim 2$ the lightest inert neutralino states can get appreciable masses only if either all or at least one of the Inert chargino mass eigenstates are light, i.e. $m_{\chi_1^\pm} \simeq 100\text{--}200$ GeV. As clarified in Sec. III and in [26], the masses of the lightest inert neutralino states
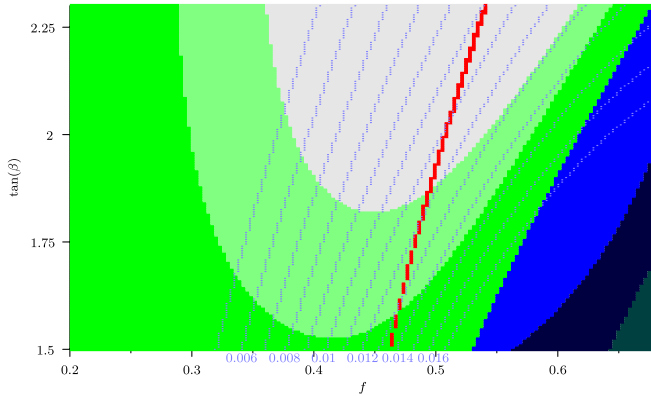

FIG. 1 (color online). Contour plot of $(X_{11}^{h_1})^2$ and relic density $\Omega_\chi h^2$ regions in the $(f, \tan\beta)$-plane with $s = 2400$ GeV, $f_{\alpha\alpha} = \tilde{f}_{\alpha\alpha} = \lambda_{\alpha\alpha} = 0$, $f_{12} = f$, $\tilde{f}_{12} = f_{12}/a$, $f_{21} = 1.02 \cdot f_{12}$, $\tilde{f}_{21} = 0.98 \cdot \tilde{f}_{12}$, $a = 0.75 + 0.25\tan\beta$ and $\lambda_{12} = \lambda_{21} = 0.06$ ($m_{\chi_{1,2}^\pm} = 101.8$ GeV). The red region is where the prediction for $\Omega_\chi h^2$ is consistent with the measured one σ range of $\Omega_{\text{CDM}}h^2 = 0.1099 \pm 0.0062$. The dark green region corresponds to $D < 3$ (D is defined in Sec. VI B) while the pale green region represents the part of the parameter space in which D varies from 3 to 4. The grey area indicates that $D > 4$. The blue region corresponds to $m_{\chi_1^0} > M_Z/2$, while the dark blue region to the right is ruled out by the requirement that perturbation theory remains valid up to the GUT scale.
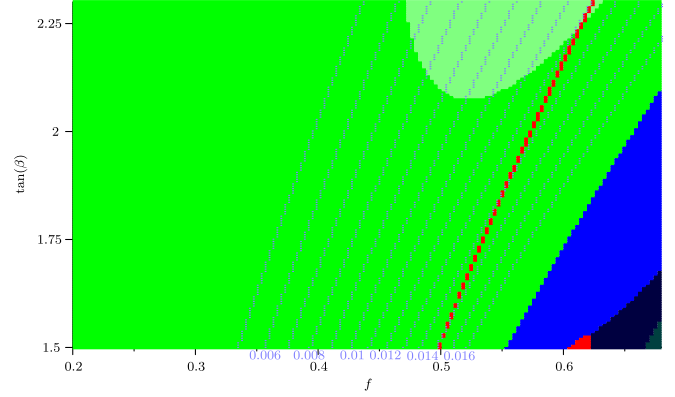

FIG. 2 (color online). Contour plot of $(X_{11}^{h_1})^2$ and relic density $\Omega_\chi h^2$ regions in the $(f, \tan\beta)$-plane with $s = 2400$ GeV, $f_{\alpha\alpha} = \tilde{f}_{\alpha\alpha} = \lambda_{\alpha\alpha} = 0$, $f_{12} = f$, $\tilde{f}_{12} = f_{12}/a$, $f_{21} = 1.02 \cdot f_{12}$, $\tilde{f}_{21} = 0.98 \cdot \tilde{f}_{12}$, $a = 0.5 + 0.5\tan\beta$ and $\lambda_{12} = \lambda_{21} = 0.06$ ($m_{\chi_{1,2}^\pm} = 101.8$ GeV). The red region is where the prediction for $\Omega_\chi h^2$ is consistent with the measured one σ range of $\Omega_{\text{CDM}}h^2 = 0.1099 \pm 0.0062$. The dark green region corresponds to $D < 3$ (D is defined in Sec. VI B) while the pale green region represents the part of the parameter space in which D varies from 3 to 4. The grey area indicates that $D > 4$. The blue region corresponds to $m_{\chi_1^0} > M_Z/2$, while the dark blue region to the right is ruled out by the requirement that perturbation theory remains valid up to the GUT scale.

decrease with increasing $m_{\chi_{1,2}^\pm}$ and it is therefore rather difficult to find benchmark scenarios consistent with cosmological observations for $m_{\chi_1^\pm} \gtrsim 200$ GeV. At the same time we demonstrate (see benchmark point (ix) in Table II) that one light Inert chargino mass eigenstate is enough to ensure that the lightest inert neutralino state gains a mass of the order of $M_Z/2$.

To obtain the kind of Inert neutralino and chargino spectrum discussed above one has to assume that some or all of the couplings $\lambda_{\alpha\beta}$ are rather small, e.g. they are expected to be much smaller than $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$. On the other hand in order to get $m_{\chi_1^0} \sim m_{\chi_2^0} \sim M_Z/2$ the Yukawa couplings $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$ need to be relatively close to their theoretical upper bounds which are caused by the requirement of the validity of perturbation theory up to the GUT scale. Since gauge coupling unification determines the RG flow and low energy value of g_1' the mass of the Z' gauge boson is set by the VEV of the singlet field s only. In our study we choose $s = 2400$ GeV so that the Z' mass is about 890 GeV. This value of the Z' boson mass is just above the present lower bound of 865 GeV set by CDF [50] and allows satisfaction of stringent limits on the Z' mass and $Z - Z'$ mixing that come from precision EW tests [51].

Since we restrict our analysis to low values of $\tan\beta \lesssim 2$ the mass of the SM-like Higgs boson is very sensitive to the choice of the coupling λ. Stringent LEP constraints require $\lambda(M_t)$ to be larger than the low energy value of $g_1' \simeq 0.47$. If we try to increase $\lambda(M_t)$ much further, then the

theoretical upper bounds on $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$ become substantially stronger. As a consequence, it is rather difficult to find solutions with $|m_{\chi_1^0}| \sim |m_{\chi_2^0}| \sim M_Z/2$. Therefore in our analysis we concentrate on values of $\lambda(M_t) \lesssim 0.6$. In addition, we set stop scalar masses to be equal to $m_Q = m_U = M_S = 700$ GeV and restrict our consideration to the so-called maximal mixing scenario when the stop mixing parameter $X_t = A_t - \lambda s/(\sqrt{2}\tan\beta)$ is equal to $X_t = \sqrt{6}M_S$. This choice of parameters limits the range of variations of the lightest *CP*-even Higgs mass. In the leading two-loop approximation the mass of the SM-like Higgs boson varies from 115 GeV($\lambda = g_1'$) to 136 GeV ($\lambda = 0.6$). From Tables I and II one can see that the large values of $\lambda \gtrsim g_1'$ that we choose in our analysis result in the extremely hierarchical structure of the Higgs spectrum, as pointed out in Sec. IV (see also [23]). In Tables I and II 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.

The set of the benchmark points that we specify demonstrates that one can get a reasonable dark matter density consistent with the recent observations if $|m_{\chi_1^0}| \sim |m_{\chi_2^0}| \sim M_Z/2$. Our benchmark scenarios also indicate that in this case the SM-like Higgs boson decays predominantly into the lightest inert neutralinos (χ_1 and χ_2) while the total branching ratio into SM particles varies from 2% to 4%.

The benchmark points (i), (ii), (iv), (v) and (viii) are motivated by a non-Abelian family symmetry Δ_{27} which describes well the observed hierarchy in the quark and lepton sectors. As was discussed in Sec. III these scenarios imply that all flavor diagonal Yukawa couplings $\lambda_{\alpha\alpha}$, $f_{\alpha\alpha}$ and $\tilde{f}_{\alpha\alpha}$ are rather small. Because of the approximate global $U(1)$ symmetry (21), that originates from the family symmetry Δ_{27}, the spectrum of Inert neutralinos involves a set of pseudo-Dirac states. When the masses of the lightest and second lightest Inert neutralinos are close or they form a Dirac state then the decays of h_1 into $\chi_\alpha\chi_\beta$ will not be observed at the LHC. Thus these decay channels give rise to a large invisible branching ratio of the SM-like Higgs boson.

In Tables I and II we presented a few benchmark scenarios (i), (ii), (iv)–(vi), (ix) with almost degenerate lightest and second lightest Inert neutralinos. In some of these benchmark points both lightest Inert neutralinos are lighter than $M_Z/2$. Thus the Z boson can decay into $\chi_\alpha\chi_\beta$ so that the lightest and second lightest Inert neutralino states contribute to the invisible Z-boson width. In other benchmark scenarios both of the lightest Inert neutralinos have masses above $M_Z/2$ and the decays $Z \rightarrow \chi_\alpha\chi_\beta$ are kinematically forbidden.

When the LSP and NLSP are close in mass, LSP-NLSP coannihilations may be an important factor in determining the dark matter relic density. If this is the case then the

LSP-NLSP mass splitting should be an important factor. Since annihilations of two like-neutralinos are p-wave suppressed, one should compare βR_{Z11} with R_{Z12} when trying to determine how important coannihilations are, where β is the relative speed of the incoming particles, approximately $1/6$. It is useful to consider the following situations. With the LSP and NLSP almost degenerate and with equal self-annihilation cross sections, but a negligible coannihilation cross section, the relic density of dark matter would be twice what it would have been if the NLSP had not been present. If, alternatively, the coannihilation cross section was equal to the self-annihilation cross sections, the existence of this extra channel would lead to a lower relic density. In this case it would in fact be equal to the relic density calculated in the absence of the NLSP. In this way, in such a scenario where coannihilations and self-annihilations are about as important as each other, the relic density is largely independent of the LSP-NLSP mass splitting.

For the benchmark scenarios (i) and (ii) this latter situation is approximately the case and the LSP-NLSP mass splitting turns out not to be an important factor. The mass splitting is in fact small, about half a GeV, but if it were larger and the NLSPs were made to have frozen-out much earlier, the relic density would only be decreased slightly (by about a tenth). In benchmark scenario (iv), even though the LSP and NLSP are close in mass, coannihilations are unimportant due to the small value of R_{Z12}. In this case increasing the NLSP mass substantially while keeping everything else fixed would lead to an approximate halving of the predicted relic density, since the NLSPs would have decoupled much earlier than, rather than at the same time as, the LSPs. The only other benchmark scenario where the LSP and NLSP are close enough in mass for coannihilations to be potentially important is scenario (ix). Here coannihilation is in fact the dominant process and changing the LSP-NLSP mass splitting would have a large effect on the predicted relic density. In fact, in this scenario, if the NLSP were not present the predicted relic density would be within the measured range.

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 χ_2 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 $|m_{\chi_2^0}| - |m_{\chi_1^0}| \gtrsim 20$ 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 χ_2 produced through a Higgs decay. However, it might be

possible to detect some lepton-antilepton pairs that come from the decays $h_1 \to \chi_2\chi_\alpha$. In particular, we hope that $\mu^+\mu^-$ pairs that come from the exotic decays of the lightest *CP*-even Higgs state mentioned above can be observed at the LHC.

In Tables I and II benchmark scenarios (iii), (vii), (viii) 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 $l^+l^- + X$ might play an essential role in Higgs searches.

In addition to the exotic Higgs decays, 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. Since 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.

B. Neutralino and chargino collider limits

The remarkable signatures discussed above raise serious concerns that they could have already been observed at the Tevatron and/or even earlier at LEP. For example, the light Inert neutralino and chargino states could be produced at the Tevatron [52]. Recently, the CDF and D0 collaborations set a stringent lower bound on chargino masses using searches for SUSY with a trilepton final state (i.e. trilepton signal) [53]. These searches ruled out chargino masses below 164 GeV. However this lower bound on the chargino mass was obtained by assuming that the corresponding chargino and neutralino states decay predominantly into the LSP and a pair of leptons. In our case, however, the Inert neutralino and chargino states are expected to decay via virtual Z and W exchange, i.e. they decay predominantly into the LSP and a pair of quarks. As a consequence the lower limit on the mass of charginos that is set by the Tevatron is not directly applicable to the benchmark scenarios that we consider here. Instead in our study we use the 95%% C.L. lower limit on the chargino mass of about 100 GeV that was set by LEP II [54].

In principle LEP experiments also set constraints on the masses and couplings of neutral particles that interact with the Z boson. As mentioned above when the masses of χ_1 and χ_2 are below $M_Z/2$ they are almost degenerate and thus the decays of Z into $\chi_\alpha\chi_\beta$ contribute to the invisible width of the Z boson changing the effective number of neutrino species $N_\nu^{\rm eff}$. The contribution of χ_1 and χ_2 ($\Delta N_\nu^{\rm eff}$) to $N_\nu^{\rm eff}$ is given by

$$\Delta N_\nu^{\rm eff} = \delta_{11} + 2\delta_{12} + \delta_{22}, \tag{38}$$

where

$$\delta_{\alpha\beta} = R_{Z\alpha\beta}^2\left[1 - \frac{|m_{\chi_\alpha^0}|^2 + |m_{\chi_\beta^0}|^2}{2M_Z^2} - 3(-1)^{\theta_\alpha+\theta_\beta}\right.$$
$$\times \frac{|m_{\chi_\alpha^0}||m_{\chi_\beta^0}|}{M_Z^2} - \left.\frac{(|m_{\chi_\alpha^0}|^2 - |m_{\chi_\beta^0}|^2)^2}{2M_Z^4}\right]$$
$$\times \sqrt{\left(1 - \frac{|m_{\chi_\alpha^0}|^2 + |m_{\chi_\beta^0}|^2}{M_Z^2}\right)^2 - 4\frac{|m_{\chi_\alpha^0}|^2|m_{\chi_\beta^0}|^2}{M_Z^4}}. \tag{39}$$

All three terms in Eq. (38) contribute to $N_\nu^{\rm eff}$ only if $2|m_{\chi_2^0}| < M_Z$. In the case where only the Z boson decays into $\chi_1^0\chi_1^0$ are kinematically allowed the values of δ_{12} and δ_{22} should be set to zero. If $|m_{\chi_1^0}| + |m_{\chi_2^0}| < M_Z$ while $2|m_{\chi_2^0}| > M_Z$ then only δ_{11} and δ_{12} need to be taken into account.

In order to compare the measured value of N_ν with the effective number of neutrino species in the E_6SSM, i.e. $N_\nu^{\rm eff} = 3 + \Delta N_\nu^{\rm eff}$, it is convenient to define the variable

$$D = \frac{N_\nu^{\rm eff} - N_\nu^{\rm exp}}{\sigma^{\rm exp}}, \tag{40}$$

where $N_\nu^{\rm exp} = 2.984$ and $\sigma^{\rm exp} = 0.008$ [55]. The value of D represents the deviation between the predicted and measured effective number of neutrinos contributing to the Z-boson invisible width. It is worth pointing out that in the SM $D = 2$. In the benchmark scenarios presented in Tables I and II the value of D is always less than 3. Figures 1 and 2 also demonstrate that there is a substantial part of the E_6SSM parameter space where $m_{\chi_{1,2}^0} < M_Z/2$ and $D < 3$. This indicates that the relatively light Inert neutralinos with masses below $M_Z/2$ are not ruled out by different constraints on the effective number of neutrinos set by LEP experiments (see, for example [55,56]). Indeed, as argued in Sec. III the Yukawa couplings $f_{\alpha\beta}$ and $\tilde{f}_{\alpha\beta}$ can be chosen such that the $R_{Z\alpha\beta}$ are very small. The couplings of the lightest and second lightest Inert neutralinos to the Z boson are relatively small anyway because of the Inert singlino admixture in these states. Nevertheless Figs. 1 and 2 show that the scenarios with light Inert neutralinos which have masses below $M_Z/2$ and relatively small couplings to the Z boson can lead to the appropriate dark matter density consistent with the recent observations.

LEP has set limits on the cross section of $e^+e^- \to \chi_2^0\chi_1^0(\chi_1^+\chi_1^-)$ in the case when $\chi_2^0 \to q\bar{q}\chi_1^0(\chi_1^\pm \to q\bar{q}'\chi_1^0)$ is predominant[57]. Unfortunately, the bounds are not directly applicable for our study because OPAL limits were set for a relatively heavy χ_2^0 ($\chi_1^\pm$) only ($|m_{\chi_2^0}| \gtrsim 60$ GeV). Nevertheless, these bounds demonstrate that it

was difficult to observe light neutralinos with $|m_{\chi_{2,1}^0}| \lesssim 100$ GeV if their production cross section $\sigma(e^+e^- \to \chi_\alpha^0 \chi_\beta^0) \lesssim 0.1{-}0.3$ pb^{-1}. Since at LEP energies of the cross sections of colorless particle production through the s-channel γ/Z exchange are typically a few picobarns, the lightest and second lightest Inert neutralino states in the E$_6$SSM could escape detection at LEP if their couplings $R_{Z\alpha\beta} \lesssim 0.1{-}0.3$.

C. Dark matter direct detection

Another constraint on the couplings of the lightest Inert neutralino comes from experiments for the direct detection of dark matter. Recently the CDMS-II and XENON100 collaborations have set upper limits on the weakly interacting massive particle (WIMP)-nucleon elastic-scattering spin-independent cross section [15,58]. The XENON100 Collaboration claims a limit on the spin-independent cross section of 3.4×10^{-44} cm^2 for a 55 GeV WIMP. This limit remains fairly constant for lower WIMP masses and does not increase above about 4×10^{-44} cm^2 even for the lowest LSP masses that are consistent with our thermal freeze-out scenario. Since in the E$_6$SSM the couplings of the lightest Inert neutralino to quarks (leptons) and squarks (sleptons) are suppressed, the χ_1^0-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 χ_1^0-nucleon cross section in the E$_6$SSM takes the form [59,60]

$$\sigma_{\rm SI} = \frac{4m_r^2 m_N^2}{\pi v^2 m_{h_1}^4} |X_{11}^{h_1} F^N|^2, \qquad m_r = \frac{m_{\chi_1^0} m_N}{m_{\chi_1^0} + m_N},$$

$$F^N = \sum_{q=u,d,s} f_{Tq}^N + \frac{2}{27} \sum_{Q=c,b,t} f_{TQ}^N, \tag{41}$$

where

$$m_N f_{Tq}^N = \langle N|m_q \bar{q}q|N \rangle, \qquad f_{TQ}^N = 1 - \sum_{q=u,d,s} f_{Tq}^N.$$

Here for simplicity we assume that the lightest Higgs state has the same couplings as the Higgs boson in the SM and ignore all contributions induced by heavy Higgs and squark exchange.[8] Because of the hierarchical structure of the particle spectrum and the approximate Z_2^H symmetry this approximation works very well. Using the

---

[8]The presence of almost degenerate lightest and second lightest Inert neutralinos could result in the inelastic scattering of χ_1^0 on nuclei (A), i.e. $\chi_1^0 + A \to \chi_2^0 + A$, that could affect the direct detection of χ_1^0 at the experiment. However such processes may take place only if the mass splitting between χ_1^0 and χ_2^0 is less than 100 KeV [61]. Since in all of the benchmark scenarios considered here the corresponding mass splitting is substantially larger the inelastic scattering of χ_1^0 does not play any significant role.

experimental limits set on $\sigma_{\rm SI}$ and Eqs. (41) one can obtain upper bounds on $X_{11}^{h_1}$ [62].

In Tables I and II we specify the interval of variations of $\sigma_{\rm SI}$ for each benchmark scenario. As one can see from Eq. (41) the value of $\sigma_{\rm SI}$ depends rather strongly on the hadronic matrix elements, i.e. the coefficients f_{Tq}^N, that are related to the π-nucleon σ term and the spin content of the nucleon. The hadronic uncertainties in the elastic scattering cross section of dark matter particles on nucleons were considered in [59,63]. In particular, it was pointed out that f_{Ts}^N could vary over a wide range. In Tables I and II the lower limit on $\sigma_{\rm SI}$ corresponds to $f_{Ts}^N = 0$ while the upper limit implies that $f_{Ts}^N = 0.36$ (see [60]). From Tables I and II and Eq. (41) it also becomes clear that $\sigma_{\rm SI}$ decreases substantially when m_{h_1} grows.

Since in all of the benchmark scenarios presented in Tables I and II 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 χ_1^0 to the lightest *CP*-even Higgs state is always large giving rise to a χ_1^0-nucleon spin-independent cross section which is of the order of or larger than the experimental upper bound. However it is worth keeping in mind that the obtained experimental limits on $\sigma_{\rm SI}$ are not very robust [64]. Moreover, CDMS-II and XENON100 quote 90% C.L. upper bounds while the 95% confidence level bounds are larger by a factor of 1.3. By the same token the 99% C.L. and 99.9% C.L. upper bounds, which are associated with 2.6 and 3.3 standard deviations, are expected to be 2 and 3 times larger than the 90% C.L. bounds, respectively. Following these estimates it is clear that the benchmark scenarios presented in Tables I and II cannot yet be ruled out by either XENON100 or CDMS-II. However in the near future the expected new analysis from XENON100 may either confirm or refute our scenario.

VII. SUMMARY AND CONCLUSIONS

In this paper we have considered novel decays of the SM-like Higgs boson which can occur within a particular dark matter motivated scenario of the Exceptional Supersymmetric Standard Model (E$_6$SSM). This model implies that at high energies the E_6 GUT gauge group is broken to the SM gauge group together with an additional $U(1)_N$ gauge group under which right-handed neutrinos have zero charge. To ensure anomaly cancellation and gauge coupling unification, the low energy matter content of the E$_6$SSM includes three 27 representations of E_6 and a pair of $SU(2)$ doublets from an additional 27$'$ and $\overline{27}'$. Thus the E$_6$SSM involves extra exotic matter beyond that of the MSSM that includes two families of Inert Higgs doublet superfields H_α^u and H_α^d and two Inert SM-singlet superfields S_α that carry $U(1)_N$ charges. The fermion components of these superfields form Inert neutralino and chargino states.

To satisfy LEP constraints we restricted our consideration to scenarios with relatively heavy Inert chargino states, i.e. $m_{\chi^\pm_{1,2}} \gtrsim 100$ GeV. In our analysis we also required the validity of perturbation theory up to the GUT scale which sets stringent constraints on the values of the Yukawa couplings at low energies. Using these restrictions we argued that the lightest and the second lightest Inert neutralinos (χ^0_1 and χ^0_2) are always light, viz. they typically have masses below 60–65 GeV. These neutralinos are mixtures of Inert Higgsinos and singlinos. In our model χ^0_1 tends to be the LSP and can play the role of dark matter, while χ^0_2 tends to be the NLSP. The masses of χ^0_1 and χ^0_2 can be induced even if only one family of the Inert Higgsinos couples to the two SM singlinos. The masses of χ^0_1 and χ^0_2 decrease with increasing $\tan\beta$ and Inert chargino masses.

An important requirement of this paper is that the lightest Inert neutralino account for all or most of the observed dark matter relic density. This sets another stringent constraint on the masses and couplings of χ^0_1. Indeed, because the lightest Inert neutralino states are almost Inert singlinos, their couplings to the gauge bosons, Higgs states, quarks (squarks) and leptons (sleptons) are rather small resulting in a relatively small annihilation cross section of $\tilde{\chi}^0_1\tilde{\chi}^0_1 \to$ SM particles and the possibility of an unacceptably large dark matter density. In the limit when all non-SM states except the Inert neutralinos and charginos are heavy ($\gtrsim$ TeV) a reasonable density of dark matter can be obtained for $|m_{\chi^0_{1,2}}| \sim M_Z/2$ where the Inert LSPs annihilate mainly through Z in the s channel [26]. If $\tilde{\chi}^0_1$ annihilation proceeds through the Z-boson resonance, i.e. $2|m_{\chi^0_\sigma}| \approx M_Z$, then an appropriate value of $\Omega_{\rm CDM}h^2$ can be achieved even for a relatively small coupling of $\tilde{\chi}^0_1$ to Z.

The above scenario naturally emerges when a Δ_{27} family symmetry is included in the E$_6$SSM [33]. The family symmetry was not introduced for this purpose, instead it was introduced earlier to provide an explanation of the Z^H_2 symmetry and to account for the quark and lepton masses and mixings, including tribimaximal neutrino mixing. It is therefore encouraging to find that the same symmetry leads to a spectrum of inert pseudo-Dirac neutralinos which allows for a successful dark matter relic abundance, and also predicts novel Higgs decays. The Δ_{27} family symmetry also implies two almost degenerate families of D-fermion states [33] and in addition may have interesting consequences for B physics [65]. As discussed in Sec. III B this symmetry leads to a cancellation of different contributions to the off-diagonal couplings of the LSP and NLSP. In addition, due to the singlino component of the lightest Inert neutralino states, the diagonal couplings of χ^0_1 and χ^0_2 to the Z boson can also be rather small. Therefore these states could have escaped detection at LEP.

The main point we make in this paper is that, within the above dark matter motivated scenario, although the lightest

and the second lightest Inert neutralinos might have very small couplings to the Z boson, their couplings to the SM-like Higgs state h_1 are always large. Indeed, we argued that in the first approximation the couplings of χ^0_1 and χ^0_2 to the lightest CP-even Higgs boson are proportional to $|m_{\chi^0_{1,2}}|/v$. Since $|m_{\chi^0_{1,2}}| \sim M_Z/2$ these couplings are much larger than the corresponding b-quark coupling. Thus the SM-like Higgs boson decays predominantly into the lightest inert neutralino states and has very small branching ratios (2%–4%) for decays into SM particles. We have illustrated this, together with the other phenomenological aspects of the dark matter motivated scenario considered in this paper, by presenting a set of benchmark points in Tables I and II. If the masses of the lightest and second lightest Inert neutralinos are very close then the decays of h_1 into $\chi_\alpha\chi_\beta$ will not be observed at the LHC giving rise to a large invisible branching ratio of the SM-like Higgs boson. When the mass difference between the second lightest and the lightest Inert neutralinos is larger than 10 GeV the invisible branching ratio remains dominant but some of the decay products of χ_2 might be observed at the LHC. In particular, there is a chance that $\mu^+\mu^-$ pairs could be detected. Since the branching ratios of h_1 into SM particles are extremely suppressed, the decays of the SM-like Higgs boson into $l^+l^- + X$ could be important for Higgs searches.

In conclusion, the E$_6$SSM predicts three Higgs families plus three Higgs singlets, where one family develop VEVs, while the remaining two which do not are called Inert. This pattern of Higgs VEVs is due to a broken Z^H_2 symmetry whose origin can be understood from a Δ_{27} family symmetry. The model can account for the dark matter relic abundance if the two lightest Inert neutralinos, identified as the LSP and NLSP, have masses close to half the Z mass, with a pseudo-Dirac structure as predicted by the Δ_{27} family symmetry. Within this scenario we find that the usual SM-like Higgs boson decays more that 95% of the time into either LSPs or NLSPs, with the latter case producing a final state containing two soft leptons l^+l^- with an invariant mass less than or about 10 GeV. We have illustrated this with a set of benchmark points satisfying phenomenological constraints and the WMAP dark matter relic abundance. This scenario also predicts other light Inert chargino and neutralino states below 200 GeV, and large LSP direct detection cross sections close to current limits and observable soon at XENON100.

[1] S. Chang, R. Dermisek, J. F. Gunion, and N. Weiner, Annu. Rev. Nucl. Part. Sci. **58**, 75 (2008); A. Djouadi and R. M. Godbole, arXiv:0901.2030; R. Dermisek, Mod. Phys. Lett. A **24**, 1631 (2009).

[2] R. E. Shrock and M. Suzuki, Phys. Lett. B **110**, 250 (1982); L. F. Li, Y. Liu, and L. Wolfenstein, Phys. Lett. B **159**, 45 (1985); A. S. Joshipura and S. D. Rindani, Phys. Rev. Lett. **69**, 3269 (1992); A. S. Joshipura and J. W. F. Valle, Nucl. Phys. **B397**, 105 (1993); T. Binoth and J. J. van der Bij, Z. Phys. C **75**, 17 (1997); A. Datta and A. Raychaudhuri, Phys. Rev. D **57**, 2940 (1998); C. P. Burgess, M. Pospelov, and T. ter Veldhuis, Nucl. Phys. **B619**, 709 (2001); M. Hirsch, J. C. Romao, J. W. F. Valle, and A. Villanova del Moral, Phys. Rev. D **73**, 055007 (2006); B. Patt and F. Wilczek, arXiv:hep-ph/0605188; J. J. van der Bij, Phys. Lett. B **636**, 56 (2006); D. G. Cerdeno, A. Dedes, and T. E. J. Underwood, J. High Energy Phys. 09 (2006) 067; V. Barger, P. Langacker, M. McCaskey, M. J. Ramsey-Musolf, and G. Shaughnessy, Phys. Rev. D **77**, 035005 (2008); H. Sung Cheon, S. K. Kang, and C. S. Kim, J. Cosmol. Astropart. Phys. 05 (2008) 004; A. Dedes, T. Figy, S. Hoche, F. Krauss, and T. E. J. Underwood, J. High Energy Phys. 11 (2008) 036; H. Sung Cheon, S. K. Kang, and C. S. Kim, Phys. Lett. B **675**, 203 (2009); V. Barger, P. Langacker, M. McCaskey, M. Ramsey-Musolf, and G. Shaughnessy, Phys. Rev. D **79**, 015018 (2009); C. S. Kim, S. C. Park, K. Wang, and G. Zhu, Phys. Rev. D **81**, 054004 (2010); M. Farina, D. Pappadopulo, and A. Strumia, Phys. Lett. B **688**, 329 (2010); I. M. Shoemaker, K. Petraki, and A. Kusenko, J. High Energy Phys. 09 (2010) 060.

[3] O. J. P. Eboli, M. C. Gonzalez-Garcia, A. Lopez-Fernandez, S. F. Novaes, and J. W. F. Valle, Nucl. Phys. **B421**, 65 (1994); F. de Campos, O. J. P. Eboli, J. Rosiek, and J. W. F. Valle, Phys. Rev. D **55**, 1316 (1997).

[4] D. Choudhury and D. P. Roy, Phys. Lett. B **322**, 368 (1994).

[5] S. P. Martin and J. D. Wells, Phys. Rev. D **60**, 035006 (1999).

[6] R. Schabinger and J. D. Wells, Phys. Rev. D **72**, 093007 (2005); M. J. Strassler and K. M. Zurek, Phys. Lett. B **661**, 263 (2008); S. Gopalakrishna, S. J. Lee, and J. D. Wells, Phys. Lett. B **680**, 88 (2009); S. Gopalakrishna, in *SUSY09: 7th International Conference on Supersymmetry and the Unification of Fundamental Interactions*, edited by G. Alverson, P. Nath, and B. Nelson, AIP Conf. Proc. No. 1200 (AIP, New York, 2010).

[7] K. Belotsky, D. Fargion, M. Khlopov, R. Konoplich, and K. Shibaev, Phys. Rev. D **68**, 054027 (2003).

[8] K. Belotsky, V. A. Khoze, A. D. Martin, and M. G. Ryskin, Eur. Phys. J. C **36**, 503 (2004).

[9] N. Arkani-Hamed, S. Dimopoulos, G. R. Dvali, and J. March-Russell, Phys. Rev. D **65**, 024032 (2001); G. F. Giudice, R. Rattazzi, and J. D. Wells, Nucl. Phys. **B595**, 250 (2001); N. G. Deshpande and D. K. Ghosh, Phys. Lett. B **567**, 235 (2003); A. Datta, K. Huitu, J. Laamanen, and B. Mukhopadhyaya, Phys. Rev. D **70**, 075003 (2004); D. Dominici and J. F. Gunion, Phys. Rev. D **80**, 115006 (2009).

[10] M. Battaglia, D. Dominici, J. F. Gunion, and J. D. Wells, The Invisible Higgs Decay Width in the Add Model at the LHC, (to be published).

[11] M. Asano, S. Matsumoto, N. Okada, and Y. Okada, Phys. Rev. D **75**, 063506 (2007); R. S. Hundi, B. Mukhopadhyaya, and A. Nyffeler, Phys. Lett. B **649**, 280 (2007); L. Wang and J. M. Yang, Phys. Rev. D **79**, 055013 (2009).

[12] S. Kanemura, S. Matsumoto, T. Nabeshima, and N. Okada, Phys. Rev. D **82**, 055026 (2010).

[13] G. Bertone, D. Hooper, and J. Silk, Phys. Rep. **405**, 279 (2005).

[14] H. Baer, D. Dicus, M. Drees, and X. Tata, Phys. Rev. D **36**, 1363 (1987); J. F. Gunion and H. E. Haber, Nucl. Phys. **B307**, 445 (1988); **B402**, 569(E) (1993); K. Griest and H. E. Haber, Phys. Rev. D **37**, 719 (1988); J. L. Lopez, D. V. Nanopoulos, H. Pois, X. Wang, and A. Zichichi, Phys. Rev. D **48**, 4062 (1993); A. Djouadi, J. Kalinowski, and P. M. Zerwas, Z. Phys. C **57**, 569 (1993); A. Djouadi, P. Janot, J. Kalinowski, and P. M. Zerwas, Phys. Lett. B **376**, 220 (1996); G. Belanger, F. Boudjema, F. Donato, R. Godbole, and S. Rosier-Lees, Nucl. Phys. **B581**, 3 (2000); G. Belanger, F. Boudjema, A. Cottrant, R. M. Godbole, and A. Semenov, Phys. Lett. B **519**, 93 (2001).

[15] Z. Ahmed *et al.* (The CDMS-II Collaboration), Science **327**, 1619 (2010).

[16] A. Djouadi, Int. J. Mod. Phys. A **10**, 1 (1995).

[17] LEP Higgs Working for Higgs boson searches and ALEPH Collaboration and DELPHI Collaboration and CERN-L3 Collaboration and OPAL Collaboration, arXiv:hep-ex/0107032.

[18] S. G. Frederiksen, N. Johnson, G. L. Kane, and J. Reid, Phys. Rev. D **50**, R4244 (1994); R. M. Godbole, M. Guchait, K. Mazumdar, S. Moretti, and D. P. Roy, Phys. Lett. B **571**, 184 (2003).

[19] H. Davoudiasl, T. Han, and H. E. Logan, Phys. Rev. D **71**, 115007 (2005).

[20] J. F. Gunion, Phys. Rev. Lett. **72**, 199 (1994); B. P. Kersevan, M. Malawski, and E. Richter-Was, Eur. Phys. J. C **29**, 541 (2003).

[21] E. E. Boos, S. V. Demidov, and D. S. Gorbunov, arXiv:1010.5373.

[22] O. J. P. Eboli and D. Zeppenfeld, Phys. Lett. B **495**, 147 (2000).

[23] S. F. King, S. Moretti, and R. Nevzorov, Phys. Rev. D **73**, 035009 (2006).

[24] S. F. King, S. Moretti, and R. Nevzorov, Phys. Lett. B **634**, 278 (2006).

[25] S. F. King, R. Luo, D. J. Miller, and R. Nevzorov, J. High Energy Phys. 12 (2008) 042.

[26] J. P. Hall and S. F. King, J. High Energy Phys. 08 (2009) 088.

[27] E. Keith and E. Ma, Phys. Rev. D **56**, 7155 (1997).

[28] S. F. King, S. Moretti, and R. Nevzorov, Phys. Lett. B **650**, 57 (2007).

[29] S. F. King, S. Moretti, and R. Nevzorov, arXiv:hep-ph/0601269; S. Kraml et al.CERN–2006–009; S. F. King, S. Moretti, and R. Nevzorov, in Cairo International Conference on High Energy Physics (CICHEP II), edited by S. Khalil, AIP Conf. Proc. No. 881 (AIP, New York, 2007); R. Howl and S. F. King, J. High Energy Phys. 01 (2008) 030; S. F. King, S. Moretti, and R. Nevzorov, in Proceedings of the 33rd International Conference, edited by A. Sissakian, G. Kozlov, and E. Kolganova (World Scientific Publishing Co., Singapore, 2007, pp. 1125–1128; P. Athron, J. P. Hall, R. Howl, S. F. King, D. J. Miller, S. Moretti, and R. Nevzorov, Nucl. Phys. B, Proc. Suppl. **200–202**, 120 (2010).

[30] P. Athron, S. F. King, D. J. Miller, S. Moretti, and R. Nevzorov, Phys. Rev. D **80**, 035009 (2009); P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, and R. Nevzorov, Phys. Lett. B **681**, 448 (2009); P. Athron, S. F. King, D. J. Miller, S. Moretti, R. Nevzorov, arXiv:0810.0617.

[31] J. Rich, M. Spiro, and J. Lloyd-Owen, Phys. Rep. **151**, 239 (1987); P. F. Smith, Contemp. Phys. **29**, 159 (1988); T. K. Hemmick et al., Phys. Rev. D **41**, 2074 (1990).

[32] S. Hesselbach, D. J. Miller, G. Moortgat-Pick, R. Nevzorov, and M. Trusov, Phys. Lett. B **662**, 199 (2008); S. Hesselbach, D. J. Miller, G. Moortgat-Pick, R. Nevzorov, and M. Trusov, The lightest Neutralino in the MNSSM (to be published).

[33] R. Howl and S. F. King, Phys. Lett. B **687**, 355 (2010).

[34] I. de Medeiros Varzielas, S. F. King, and G. G. Ross, Phys. Lett. B **648**, 201 (2007); E. Ma, Mod. Phys. Lett. A **21**, 1917 (2006).

[35] S. Nussinov, Phys. Lett. B **165**, 55 (1985); S. M. Barr, R. S. Chivukula, and E. Farhi, Phys. Lett. B **241**, 387 (1990); D. B. Kaplan, Phys. Rev. Lett. **68**, 741 (1992); D. Hooper, J. March-Russell, and S. M. West, Phys. Lett. B **605**, 228 (2005); S. B. Gudnason, C. Kouvaris, and F. Sannino, Phys. Rev. D **73**, 115003 (2006); T. A. Ryttov and F. Sannino, Phys. Rev. D **78**, 115010 (2008); D. E. Kaplan, M. A. Luty, and K. M. Zurek, Phys. Rev. D **79**, 115016 (2009); G. D. Kribs, T. S. Roy, J. Terning, and K. M. Zurek, Phys. Rev. D **81**, 095001 (2010); H. An, S. L. Chen, R. N. Mohapatra, and Y. Zhang, J. High Energy Phys. 03 (2010) 124.

[36] K. Griest and D. Seckel, Nucl. Phys. **B283**, 681 (1987); **B296**, 1034(E) (1988).

[37] M. T. Frandsen and S. Sarkar, Phys. Rev. Lett. **105**, 011301 (2010).

[38] P. A. Kovalenko, R. B. Nevzorov, and K. A. Ter-Martirosian, Yad. Fiz. **61**, 898 (1998) [Phys. At. Nucl. **61**, 812 (1998)].R. B. Nevzorov, K. A. Ter-Martirosyan, and M. A. Trusov, Yad. Fiz. **65**, 311 (2002) [Phys. At. Nucl. **65**, 285 (2002)].

[39] R. B. Nevzorov and M. A. Trusov, Zh. Eksp. Teor. Fiz. **91**, 1251 (2000) [J. Exp. Theor. Phys. **91**, 1079 (2000)].

[40] D. J. Miller, R. Nevzorov, and P. M. Zerwas, Nucl. Phys. **B681**, 3 (2004); R. Nevzorov and D. J. Miller, Proceedings to the 7th Workshop "What Comes beyond the Standard Model", edited by N. S. Mankoc-Borstnik, H. B. Nielsen, C. D. Froggatt, D. Lukman (DMFA-Zaloznistvo, Ljubljana, 2004), p. 107.

[41] C. Panagiotakopoulos and A. Pilaftsis, Phys. Rev. D **63**, 055003 (2001); D. J. Miller, S. Moretti, and R. Nevzorov, Proceedings to the 18th International Workshop on High-Energy Physics and Quantum Field Theory (QFTHEP 2004), edited by M. N. Dubinin and V. I. Savrin (Moscow State University, Moscow, 2004), p. 212.

[42] A. Djouadi, Phys. Rep. **459**, 1 (2008).

[43] S. G. Gorishnii, A. L. Kataev, S. A. Larin, and L. R. Surguladze, Mod. Phys. Lett. A **5**, 2703 (1990); Phys. Rev. D **43**, 1633 (1991); K. G. Chetyrkin and A. Kwiatkowski, Nucl. Phys. **B461**, 3 (1996); S. A. Larin, T. van Ritbergen, and J. A. M. Vermaseren, Phys. Lett. B **362**, 134 (1995).

[44] S. Heinemeyer, Int. J. Mod. Phys. A **21**, 2659 (2006).

[45] J. Dunkley et al. (WMAP Collaboration), Astrophys. J. Suppl. Ser. **180**, 306 (2009).

[46] J. D. Wells, Perspectives on Supersymmetry, edited by G. Kane (World Scientific, Singapore (1997), pp. 276–292.

[47] V. Barger, C. Kao, P. Langacker, and H. S. Lee, Phys. Lett. B **600**, 104 (2004); A. Menon, D. E. Morrissey, and C. E. M. Wagner, Phys. Rev. D **70**, 035005 (2004); V. Barger, P. Langacker, and H. S. Lee, Phys. Lett. B **630**, 85 (2005); S. V. Demidov and D. S. Gorbunov, J. High Energy Phys. 02 (2007) 055; V. Barger, P. Langacker, I. Lewis, M. McCaskey, G. Shaughnessy, and B. Yencho, Phys. Rev. D **75**, 115002 (2007).

[48] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **180**, 747 (2009); Comput. Phys. Commun. **176**, 367 (2007); **174**, 577 (2006); Comput. Phys. Commun. **149**, 103 (2002).

[49] A. Semenov, Comput. Phys. Commun. **180**, 431 (2009); A. V. Semenov, arXiv:hep-ph/0208011.

[50] J. Erler, P. Langacker, S. Munir, and E. rojas, arXiv:1010.3097; C. Hays, A. Kotwal, and O. Stelzer-Chilton, "A Search for Dimuon Resonances with CDF in Run II", http://www-cdf.fnal.gov/physics/exotic/r2a/20080710.dimuon_resonance/.

[51] J. Erler, P. Langacker, S. Munir, and E. R. Pena, J. High Energy Phys. 08 (2009) 017.

[52] H. Baer and X. Tata, Phys. Rev. D **47**, 2739 (1993); H. Baer, C. Kao, and X. Tata, Phys. Rev. D **48**, 5175 (1993).

[53] O. Mundal (D0 Collaboration), Search for Supersymmetry in Trilepton Final States with the DO Detector (to be published); T. Aaltonen et al. (CDF Collaboration), Phys. Rev. Lett. **101**, 251801 (2008); V. M. Abazov et al. (D0 Collaboration), Phys. Lett. B **680**, 34 (2009);

J. Strologas (CDF Collaboration), *SUSY09: 7th International Conference on Supersymmetry and the Unification of Fundamental Interactions*, edited by G. Alverson, P. Nath, and B. Nelson, AIP Conf. Proc. No. 1200 (AIP, New York, 2010).

[54] A. C. Kraan, arXiv:hep-ex/0505002.

[55] ALEPH Collaboration and DELPHI Collaboration and L3 Collaboration, Phys. Rep. **427**, 257 (2006); C. Amsler *et al.* (Particle Data Group), Phys. Lett. B **667**, 1 (2008).

[56] J. Abdallah *et al.* (DELPHI Collaboration), Eur. Phys. J. C **38**, 395 (2005).

[57] G. Abbiendi *et al.* (OPAL Collaboration), Eur. Phys. J. C **35**, 1 (2004).

[58] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. Lett. **105**, 131302 (2010).

[59] J. R. Ellis, K. A. Olive, and C. Savage, Phys. Rev. D **77**, 065026 (2008).

[60] J. Kalinowski, S. F. King, and J. P. Roberts, J. High Energy Phys. 01 (2009) 066.

[61] D. Tucker-Smith and N. Weiner, Phys. Rev. D **64**, 043502 (2001).

[62] K. Cheung and T. C. Yuan, Phys. Lett. B **685**, 182 (2010); K. Cheung, K. H. Tsao, and T. C. Yuan, arXiv:1003.4611.

[63] A. Bottino, F. Donato, N. Fornengo, and S. Scopel, Astropart. Phys. **13**, 215 (2000); J. R. Ellis, A. Ferstl, and K. A. Olive, Phys. Lett. B **481**, 304 (2000); J. R. Ellis, A. Ferstl, and K. A. Olive, Phys. Rev. D **63**, 065016 (2001); A. Bottino, F. Donato, N. Fornengo, and S. Scopel, Astropart. Phys. **18**, 205 (2002).

[64] J. I. Collar, arXiv:1010.5187.

[65] S. F. King, J. High Energy Phys. 09 (2010) 114.