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NATURAL R PARITY CONSERVATION WITH HORIZONTAL SYMMETRIES. A FOUR GENERATION MODEL

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Abstract

The absence of R parity violating operators can be naturally ensured in the presence of a $SU(N)_H$ ($N = 4, 6, \dots$) horizontal gauge symmetry, independently of the vertical gauge group. We study an extension of the supersymmetric standard model with four families and gauged $SU(4)_H$. Beyond preserving R parity, the model gives rise to the realistic fermion mass matrices which naturally ensure the heaviness of the fourth family fermions b', t', τ', ν' . Their masses are all in the 100 GeV range. Phenomenological implications of the model are discussed. It is shown that cosmological constraints on the lifetime of the fourth family quarks translate into a lower bound on the τ -neutrino mass of a few eV. Hence this neutrino can provide a hot component of the cosmological dark matter, while in our R-parity conserving model a cold component can be naturally provided by the stable lightest supersymmetric particle.

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1 Introduction

In the Standard Model (SM), Baryon (B) and Lepton (L) numbers are conserved as a result of accidental global $U(1)_B$ and $U(1)_L$ symmetries that follow from the requirement of gauge invariance and renormalizability.¹ In the Supersymmetric version of the Standard Model (SSM) this is not true anymore. Consider in fact the quark and lepton left-handed chiral superfields, which transform under $SU(3) \times SU(2) \times U(1)$ as follows:

$$\begin{aligned} q = \begin{pmatrix} u \\ d \end{pmatrix} &\sim (3, 2, 1/6), & u^c &\sim (\bar{3}, 1, -2/3), & d^c &\sim (\bar{3}, 1, 1/3), \\ l = \begin{pmatrix} \nu \\ e \end{pmatrix} &\sim (1, 2, -1/2), & e^c &\sim (1, 1, 1). \end{aligned} \quad (1)$$

The two Higgs superfields $\Phi_{1,2}$ transform as $(1, 2, \mp 1/2)$ respectively. As the scalar components of $\Phi_{1,2}$ acquire nonzero vacuum expectation values (VEVs), the fermions acquire a mass through the superpotential terms $\Phi_1 q d^c$, $\Phi_2 q u^c$ and $\Phi_1 l e^c$. Since l and Φ_1 have the same transformation properties under the gauge group, the L and B violating terms obtained by substituting $\Phi_1 \rightarrow l$ are also allowed, as well as an additional term involving three quark superfields. Explicitly these three terms read

$$l q d^c, \quad l l e^c, \quad u^c d^c d^c. \quad (2)$$

If both the first and third of these terms are present, their combination would lead to catastrophically fast proton decay mediated by d^c -type squark exchange.

The relevant symmetry that ensures the B and L conservation in the SSM is called R parity, which is defined as $R \equiv (-1)^{2J+3B+L}$, where J is the spin of the particle and B(L) its baryon (lepton) number [2]. R parity does not commute with supersymmetry. On the other hand, it is an automatic consequence of a Z_2 matter parity under which the fermion superfields change the sign while the ‘Higgs’ ones $\Phi_{1,2}$ remain invariant. It is a well known fact that an unsatisfactory feature of the SSM is that the Z_2 (or equivalently R) parity conservation has to be imposed by hand.

In the context of Grand Unification Theories (GUT) based on the gauge group $SU(5)$, the fermion superfields are assigned to the $10 + \bar{5}$ representation of $SU(5)$, while Φ_1 and Φ_2 belong respectively to the $\bar{5}_{\Phi_1}$ and 5_{Φ_2} . The down-quark and lepton masses are generated through the coupling $10 \bar{5} \bar{5}_{\Phi_1}$, and the gauge invariant term obtained by $\bar{5}_{\Phi_1} \rightarrow \bar{5}$, namely $10 \bar{5} \bar{5}$, leads again to the set of B and L violating couplings in (2). Thus, with respect to automatic R-parity conservation the supersymmetric (SUSY) $SU(5)$ model does not differ much from the SSM. The $SO(10)$ model offers an elegant solution to this problem, since the fermion superfields are in the spinor representation **16** whereas the Higgs ones are generally assigned to vector representations as **10**, **45**, **54**, **126** etc.. The masses of the fermions, including the neutrinos, can be generated through the gauge invariant couplings $16 \bar{16} 10_{\Phi}$ and $16 \bar{16} 126_{\Phi}$ [3], while the term $(16)^3$ is forbidden since it does not contain an $SO(10)$ singlet. In other words, as long as all the $SO(10)$ invariant couplings allow

¹These symmetries can be broken only by higher order non-renormalizable operators, cutoff at the Planck scale, which can arise from non-perturbative quantum gravity or string effects [1]. For example, the lepton number violating term $(1/M_{Pl})ll\Phi\Phi$ provides the neutrino Majorana masses of about 10^{-5} eV, which could be relevant for the solar neutrino oscillations. However, analogous terms violating baryon number are very small to cause any observable effect.

only pairs of 16-plets, the theory has an automatic Z_2 matter parity under which 16-plets change the sign whereas the superfields in vector representations remain invariant. This is not true anymore for the $SO(10)$ models in which the symmetry breaking is triggered also by the scalar components of superfields belonging to the $16_\Phi + \overline{16}_\Phi$. (Examples of SUSY models in which these representations play the role of the standard $126 + \overline{126}$ can be found in ref. [4]). In fact, after substituting the VEV $\langle 16_\Phi \rangle$, the couplings $\frac{1}{M}16^316_\Phi$ which are allowed by the gauge symmetry lead again to R-parity violating terms. Since in these models the right handed neutrino masses are generated by operators $\sim \frac{1}{M}16^216_\Phi^2$ which have the same structure, the ratio $\langle 16_\Phi \rangle/M$ cannot be very small, implying in turn that the magnitude of the resulting R-parity violating terms is again in conflict with the limits on the proton lifetime. We conclude that R-parity conservation is not automatic anymore for the $SO(10)$ models with Higgs fields belonging to the 16_Φ : in this case some additional discrete symmetry has to be imposed by hand in order to distinguish the fermion 16-plets from the Higgs ones.

In this paper we wish to put forward the idea that R (or equivalently Z_2) parity conservation can be naturally ensured in models based on *gauged* horizontal symmetries. Such models are particularly interesting since they can explain the observed pattern of fermion masses and mixing. Namely, the structure of the fermion mass matrices can be related to the horizontal symmetry breaking pattern, while the mass hierarchy between families is due to certain hierarchy in this breaking (see for example the models [5, 6] based on $SU(3)_H$ horizontal symmetry). In this kind of models it is natural to assume that the horizontal group G_H acts only on the quark-lepton superfields, while the Higgses $\Phi_{1,2}$ responsible for the electroweak symmetry breaking are G_H -singlets. Hence, independently of the choice of the vertical gauge group and/or of the particular superfield assignments to its representations, the Higgs and fermion superfields can be always distinguished since they carry different horizontal quantum numbers. This leads to the possibility of allowing the necessary mass terms which are bilinear in the fermion superfields, while forbidding the B and L violating trilinear couplings in (2).

Our task is now to find and classify the theories in which the horizontal gauge group G_H naturally forbids the terms in (2) due to gauge principles, or in other words in which R parity (or equivalently Z_2 matter parity) appears as an automatic consequence of the horizontal gauge symmetry and of the field content of the model. We demand that the models we are interested in should satisfy the following list of basic requirements:

(i) In order to ensure a straightforward definition of the horizontal gauge symmetry, all the fermion superfields q, l, u^c, d^c and e^c should be assigned to the same representation of the horizontal group G_H . In other words, we forbid G_H singlet families. For example, for the group $G_H = SU(N)$ each fermion superfield can be assigned either to the N or to the \overline{N} representation of the group.

(ii) We require that the couplings in (2) are forbidden by gauge invariance. Therefore $SO(3)_H$ and chiral $SU(3)_H$, which are the only simple groups containing three dimensional representations, are excluded since in both cases $\mathbf{3}^3$ contains a gauge singlet. However we immediately notice that for $SU(N)_H$ ($N > 3$) the term N^3 does not contain gauge singlets. Hence these groups represent a class of interesting candidates.

To have phenomenologically realistic theories, the following additional constraints should be also imposed:

(iii) In order to avoid the proliferation of Higgs doublets with masses at the electroweak scale, the standard Higgs superfields $\Phi_{1,2}$ must be G_H singlets. The presence of several Higgs doublets (as implied by non-singlet horizontal representations) would in fact spoil the natural suppression of Flavor Changing Neutral Currents (FCNC) [7]. It would also destroy gauge coupling unification, thus preventing any attempt to embed the model in some vertical GUTs.

(iv) The fermion masses should arise from effective operators with the structure

$$\mathcal{O}_{eff} \sim \frac{\mathcal{P}^{(n)}(\xi_k)}{M^n} f f^c \Phi_{1,2}, \quad (3)$$

where f, f^c are the fermion superfields in eq. (1), $\mathcal{P}^{(n)}$ represents some n -order polynomial of the scalars ξ_k responsible for the breaking of G_H , and M is some cutoff mass scale.² Clearly, in order to ensure G_H invariance, $\mathcal{P}^{(n)}$ should transform as the conjugate of the tensor product of f and f^c . We must also require that the operators (3) should produce a realistic pattern of fermion masses and mixings. In this respect, models based on chiral $SU(3)_H$ [5, 6] have proven to be quite effective, since the fermion mass hierarchy can be related to the hierarchy in the $SU(3)_H$ breaking, that is to the hierarchy among the horizontal VEVs $\langle \xi_k \rangle$. Even if chiral $SU(3)_H$ fails to satisfy condition (ii), its success in accounting for the pattern of fermion masses and mixings points again towards chiral $SU(N)_H$ ($N > 3$) as possible interesting candidates.³

(v) A final strong condition is that R-parity breaking terms should not appear even after G_H breaking. More precisely, all effective operators of the form

$$\frac{\mathcal{P}^{(n)}(\xi_k)}{M^n} f f f^c, \quad \frac{\mathcal{P}^{(n)}(\xi_k)}{M^n} f^c f^c f^c, \quad (4)$$

should be forbidden by the G_H symmetry, since after the horizontal symmetry breaking $\xi_k \rightarrow \langle \xi_k \rangle$ these terms would generate again the R parity violating couplings (2).

In particular this last condition restrict the viable $SU(N)_H$ models to the cases when N is even. Consider in fact $SU(N)_H$ with the f and f^c fermion superfields assigned to the fundamental N dimensional representation. The mass terms transform as $N \times N$ and thus belong to two-index (symmetric and antisymmetric) representations. In order to construct horizontal gauge invariant mass terms, the horizontal Higgses ξ_k can be also taken in two-index representations. Then for $N=4,6,\dots$ terms of the form $N \times N \times N \times \mathcal{P}^{(n)}$ (that is the R-parity violating effective operators (4)) cannot arise, since it is impossible to saturate all the indices and construct horizontal gauge invariants. In contrast, for $SU(N)_H$ with N odd the totally antisymmetric ϵ tensor allows to rewrite some combinations of Higgs

²The non-renormalizable couplings (3) with $M \sim M_{Pl}$ could appear due to quantum gravity effects. Alternatively, these operators with arbitrary M can be effectively generated through the exchange of some superheavy fields with $O(M)$ masses [8].

³We note that vectorlike $SU(3)_H$ with q, l transforming as $\mathbf{3}$ and u^c, d^c, e^c as $\bar{\mathbf{3}}$, forbids the first two terms in (2), which is enough to ensure the proton stability. However, this case can be hardly regarded as realistic, since it allows the $SU(3)_H$ invariant terms $f f^c \Phi_{1,2}$. Then the mass splitting between different families can be achieved by means of the additional effective operators (3) only in a very unnatural way, at the price of *many* fine tunings. In addition vectorlike $SU(3)_H$ would also impede the unification of the fermions (1) within one irreducible GUT multiplet.

fields with an even number of free indices as tensors with an odd number of free indices which are suitable for generating gauge invariants, when matched with the $N \times N \times N$ term. Indeed, after the horizontal symmetry breaking ($\xi \rightarrow \langle \xi \rangle$) operators of the form

$$f_\alpha^{(c)} f_\beta^{(c)} f_\gamma^{(c)} \frac{(\xi_1 \xi_2 \cdots \xi_n)_{\delta \dots \sigma}}{M^n} \epsilon_{\alpha \beta \gamma \delta \dots \sigma} \quad (5)$$

(which can be constructed also when the ξ 's belong to the symmetric part of $N \times N$) will again spoil R-parity.

This first brief analysis suggests that natural conservation of R-parity could be achieved in models based on chiral horizontal symmetries $SU(N)_H$ with N even, under which the quark and lepton superfields transform as fundamental N -plets. Then unwanted terms transforming as N^3 are automatically forbidden by horizontal gauge invariance. Clearly this implies that the number of families must be extended to $N_f = 4, 6, \dots$. As is well known, the possibility of extra families with a light neutrino is ruled out by the results of the Mark II and LEP collaborations [9]. However, they do not exclude sequential generations with heavy neutrinos ($m > M_Z/2$). On the other hand, detailed studies [10] of the effects of radiative corrections due to additional families show that precise electroweak data are not incompatible with a fourth family, while six families (which would be our next interesting case) are ruled out [10]. In addition, a dedicated analysis showing the viability of supersymmetric models with four families with respect to gauge coupling unification was presented in ref. [11]. These results are relevant for our analysis, since the condition (iii) ensures that the field content in our $SU(4)_H$ model is the same than that of the four family SSM of ref. [11], up to some large energy scale where the horizontal symmetry breaks down (see Sect. 2). Hence, we conclude that if natural R-parity conservation has to be achieved by means of some horizontal gauge symmetry, then theoretical and phenomenological constraints hint to models based on the $SU(4)_H$ group, on which we will concentrate in the rest of the paper.

2 Horizontal symmetry $SU(4)_H$

Let us now consider the standard $SU(3) \times SU(2) \times U(1)$ vertical gauge group, with local chiral $SU(4)_H$ horizontal symmetry acting on four families of left chiral superfields

$$\begin{aligned} f_\alpha : \quad q_\alpha &= \begin{pmatrix} u \\ d \end{pmatrix}_\alpha \sim (3, 2, 1/6, 4), & l_\alpha &= \begin{pmatrix} \nu \\ e \end{pmatrix}_\alpha \sim (1, 2, -1/2, 4) \\ f_\alpha^c : \quad u_\alpha^c &\sim (\bar{3}, 1, -2/3, 4), & d_\alpha^c &\sim (\bar{3}, 2, 1/3, 4), & e_\alpha^c &\sim (1, 1, 1, 4) \end{aligned} \quad (6)$$

where each superfield is assigned to the fundamental $\mathbf{4}$ representation ($\alpha = 1, \dots, 4$ is the $SU(4)_H$ index). With this field content the horizontal $SU(4)_H$ is anomalous. In order to cancel the horizontal anomaly we introduce the following superfields which are vectorlike with respect to $SU(3) \times SU(2) \times U(1)$ and belong to the $\bar{\mathbf{4}}$ of $SU(4)_H$:

$$\begin{aligned} F^\alpha : \quad U^\alpha &\sim (3, 1, 2/3, \bar{4}), & D^\alpha &\sim (3, 1, -1/3, \bar{4}), & E^\alpha &\sim (1, 1, -1, \bar{4}) \\ F_c^\alpha : \quad U_c^\alpha &\sim (\bar{3}, 1, -2/3, \bar{4}), & D_c^\alpha &\sim (\bar{3}, 1, 1/3, \bar{4}), & E_c^\alpha &\sim (1, 1, 1, \bar{4}), & N_c^\alpha &\sim (1, 1, 0, \bar{4}) \end{aligned} \quad (7)$$

As we will see in short, these superfields turn out to be necessary also for providing masses to the known fermions.

In the Higgs sector, we choose the standard Higgs doublet superfields $\Phi_{1,2}$ to be singlets under $SU(4)_H$. The additional Higgs scalars needed for the breaking of the horizontal symmetry at some large scale cannot couple to the standard $SU(2) \times U(1)$ gauge bosons, and thus must be singlets under the electroweak group. In order to break completely the horizontal symmetry and to generate realistic mass matrices for the fermions, we introduce a set of $SU(3) \times SU(2) \times U(1)$ singlet ‘horizontal’ superfields, transforming either as the symmetric $\mathbf{10}$ ($\xi_{\{\alpha\beta\}}$) or as the antisymmetric $\mathbf{6}$ ($\chi_{[\alpha\beta]}$) representations of $SU(4)_H$. Let us first consider the case with the fields ξ^k , $k = 1, 2, \dots$ in symmetric representation. Additional superfields $\bar{\xi}^k$ transforming as $\overline{\mathbf{10}}$ are then needed to render the Higgsino sector free from chiral anomalies. However, these additional scalars do not contribute to the fermion masses.⁴

What remains now to show, is that within the framework of the gauge horizontal symmetry $SU(4)_H$ one can obtain a realistic mass pattern for three families, ensuring at the same time that all fermions of the fourth family are naturally heavy, say in the 100 GeV range. Although we have started our considerations by a general analysis which included also non-renormalizable operators as in (3) and (4), in building the model we will restrict ourselves to consider only renormalizable interactions.

The most general Yukawa superpotential for the down (up) quark and for the lepton superfields allowed by gauge invariance reads

$$W_F = g_f f_\alpha F_c^\alpha \Phi_{1(2)} + \sum_k h_F^k F^\alpha F_c^\beta \xi_{\alpha\beta}^k + \mu_f F^\alpha f_\alpha^c \quad (8)$$

with f and F respectively from (6) and (7). The analogous couplings for the neutrinos have the form

$$W_N = g_\nu l_\alpha N_c^\alpha \Phi_2 + \sum_k h_N^k N_c^\alpha N_c^\beta \xi_{\alpha\beta}^k. \quad (9)$$

Here the g 's and h 's are Yukawa couplings which we assume to be $O(1)$. The last term in eq. (8) is a gauge invariant bilinear, and the μ_f 's are gauge invariant large mass parameters. As already stated, no terms trilinear in the quark and lepton superfields are allowed by the $SU(4)_H$ gauge symmetry, ensuring naturally the absence of the B and L violating couplings $lq d^c$, $ll e^c$ and $u^c d^c d^c$. We are facing here a situation analogous to the $SO(10)$ model, since R-parity does not have to be imposed by hand, but appears as an accidental symmetry that follows from the requirement of horizontal gauge invariance. Indeed, the superpotential is invariant with respect to the Z_2 transformation under which the fermion superfields f, f^c, F and F_c (which have an odd number of $SU(4)_H$ indices) change sign, while the Higgs superfields $\Phi_{1,2}$ and ξ (with an even number of $SU(4)_H$ indices) stay invariant. More in general, the superpotential $W_F + W_N$ has an automatic global symmetry $U(1)_H$ under the following transformations:

$$f, f^c \rightarrow e^{i\omega} f, f^c, \quad F, F^c \rightarrow e^{-i\omega} F, F^c, \quad \xi^k \rightarrow e^{2i\omega} \xi^k, \quad \Phi_{1,2} \rightarrow \Phi_{1,2} \quad (\bar{\xi}^k \rightarrow e^{-2i\omega} \bar{\xi}^k) \quad (10)$$

⁴Let us note that $\bar{\xi}^{\{\alpha\beta\}}$ in the $\overline{\mathbf{10}}$ cannot couple in renormalizable way to the heavy vectorlike ‘matter’ fields in the $\overline{\mathbf{4}}$, and being an $SU(2)$ singlet, neither it can couple to quarks and leptons. However, it is still possible to introduce a direct non-renormalizable terms cutoff by the Planck scale $(1/M_{Pl}) f_\alpha f_\beta^c \Phi_{1,2} \bar{\xi}^{\alpha\beta}$. In this case, in order to reproduce the observed values of the fermion masses, the $SU(4)_H$ symmetry should be broken at a scale very close to M_{Pl} . On the other hand, as we will see in Sect. 3, the phenomenology of the model requires a horizontal symmetry breaking scale substantially smaller than M_{Pl} . As a consequence, these non-renormalizable terms would provide only negligible contribution to the fermion masses.

where a Z_2 subgroup ($\omega = \pi$) remains unbroken even when the scalars ξ get non-zero VEVs. This Z_2 matter parity ensures R parity conservation and hence proton stability.

The Yukawa couplings (8) lead to the so called "universal seesaw" mechanism [12] for the fermion mass generation, which for the case of neutrinos reduces to the ordinary seesaw mechanism [13]. Indeed, after the horizontal scalars ξ_k develop non-zero VEVs, the extra fermions F and F_c of eq. (7) acquire large masses through the second term in eq. (8). Then the first and third terms cause a "seesaw" mixing of the ordinary quarks and leptons f, f^c with the heavy ones. As a result, in the base (f, F) (f^c, F_c), the 8×8 mass matrix for the down (up) type charged fermions $f = e, d, (u)$ reads ⁵

$$\mathcal{M}_f = \begin{pmatrix} 0 & g_f v_{1(2)} \\ \mu_f & \hat{M}^F \end{pmatrix}, \quad \hat{M}_{\alpha\beta}^F = \sum_k h_F^k \langle \xi_{\alpha\beta}^k \rangle \quad (11)$$

where $v_{1,2} = \langle \tilde{\Phi}_{1,2} \rangle$ are the VEVs of the two electroweak Higgs doublets. As for the neutrinos, in the base (ν, N_c) the 8×8 Majorana mass matrix has the form

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & g_\nu v_2 \\ g_\nu v_2 & \hat{M}^N \end{pmatrix}, \quad \hat{M}_{\alpha\beta}^N = \sum_k h_N^k \langle \xi_{\alpha\beta}^k \rangle. \quad (12)$$

The universal seesaw picture provides a natural possibility to obtain three light families, while the fourth one is heavy, say with masses of the order of the electroweak scale. Indeed, let us assume that the 4×4 mass matrices $\hat{M}_{F(N)}$ for the heavy fermions are rank-3 matrices of the following form

$$\hat{M}_F = \begin{pmatrix} M_F^{(3)} & 0 \\ 0 & 0 \end{pmatrix}, \quad F = U, D, E, N \quad (13)$$

where the 3×3 blocks $M_F^{(3)}$ contain non-zero entries. In other words, we assume that all the VEVs of the type $\langle \xi_{\alpha 4}^k \rangle$ are vanishing, so that a diagonal $\tilde{U}(1)$ subgroup of $SU(4)_H \times U(1)_H$, given by the generator $\tilde{T} = \text{diag}(0, 0, 0, 1)$, is left unbroken. In this case there is no seesaw mechanism for the fermions of the fourth family: the right-handed components of the fields $f_4 = b', t', \tau', \nu'$ are actually the F_c^4 states, whereas the f_4^c form with the F_4 superheavy particles of mass μ_f .

From eqs. (11) and (12) we obtain for the fourth family fermions

$$\begin{aligned} m_{b'} &= g_d v \cos \beta & m_{t'} &= g_u v \sin \beta \\ m_{\tau'} &= g_e v \cos \beta & m_{\nu'} &= g_\nu v \sin \beta \end{aligned} \quad (14)$$

where $v = 174 \text{ GeV}$ is the electroweak breaking scale and $\tan \beta = v_2/v_1$. Since all the Yukawa couplings are assumed to be $O(1)$, for moderate values of $\tan \beta$ all the masses in (14) are of the order $\sim 100 \text{ GeV}$. On the experimental side, the firmest constraints on the masses of any new sequential fermion, quark or lepton, have been set at LEP: $m_f \gtrsim M_Z/2$. This indeed represent the best constraint on $m_{\tau'}$ and $m_{\nu'}$. Searches for new quarks at

⁵The scheme considered here is a direct $SU(4)_H$ extension of the model [14] based on the horizontal symmetry $SU(3)_H$. In the case when the vertical gauge symmetry is extended to the left-right symmetric model $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, the values of μ 's are given by the scale of the $SU(2)_R$ breaking [14].

the TEVATRON collider could in principle give much better bounds for $m_{t'}$ and $m_{b'}$ [16]. However, let us note that the structure of the heavy mass matrix (13) implies that the fourth family is unmixed with the three lighter ones. Hence the usual signatures, as for example $b' \rightarrow c, u$, that have been used to set the limits on new sequential quarks [16] do not occur in our case. In the absence of a detailed experimental analysis of the unmixed case, the only reliable limit is again the LEP one also for the new quarks. Hence we can safely conclude that the predictions in (14) are by no means in conflict with the existing experimental limits. However, it is clear that for the masses of the fourth family fermions not much room is left. The allowed parameter space is in fact strongly constrained by the CDF measurement of the top mass, $m_t = 174 \pm 10 \pm 13$ GeV [15], by the precision tests of the SM which do not leave much space for additional sizeable radiative corrections as would be induced by a too large $m_{t'}-m_{b'}$ splitting, and by renormalization group (RG) analysis of the Yukawa couplings, much in the spirit of ref. [11].

In particular, while the general analysis in [11] do allow for the possibilities $m_{t'} < m_t$ or $m_{t'} < m_{b'}$, in our model these pattern of masses are not allowed. In fact the universal seesaw mechanism outlined before implies $m_{t'} \geq m_t$, and most likely $m_{t'} \geq m_{b'}$. Then, according to [11], for $m_{t'} \geq m_t > 150$ GeV the consistency of the model implies not too large values for the masses of the other fermions in fourth family. Namely, for the low values of $\tan \beta$ we are interested in (e.g. $\tan \beta \sim 2$), the maximal values allowed are about $m_{b'} \sim 100$ GeV and $m_{\tau', \nu'} \sim 50$ GeV, that is within the reach of LEP II.

Let us now consider the mass matrices for the first three families. In this case the seesaw mechanism is indeed effective for suppressing the fermion masses from the electroweak scale down to the observed values. By assuming $M_F^{(3)} > \mu_f$, it is apparent from (11) that the fermions of the first three families will acquire their masses through a mixing with the superheavy F fermions. Namely, after decoupling the heavy states, the 3×3 mass matrices of the light charged down (up) type fermions are ⁶

$$m_f^{(3)} = g_f \mu_f (M_F^{(3)})^{-1} v_{1(2)}, \quad f = d, e, (u) \quad (15)$$

while the 3×3 Majorana mass matrix for the light neutrinos obtained from (12) reads

$$m_\nu^{(3)} = (M_N^{(3)})^{-1} (g_\nu v_2)^2. \quad (16)$$

In contrast to the SM and to most GUT models, in our picture the fermion mass hierarchy is not generated by an *ad hoc* choice of the Yukawa coupling constants. In fact, in our scheme all the Yukawas are assumed to be of the same order of magnitude, for example $O(1)$ or close to the size of the gauge couplings. As long as the off-diagonal blocks in eqs. (11) and (12) are flavour blind (unit) matrices, all the informations on the fermion mass and mixing pattern is contained in the heavy fermion mass matrices $M_F^{(3)}$. Since the structure of the latters is determined by the different VEVs $\langle \xi^k \rangle$ (modulo differences in the Yukawa constants h_F^k), the observed hierarchy of the light fermion masses is ultimately determined by the hierarchy in the VEVs which break the horizontal symmetry. In other

⁶After decoupling the heavy states at the horizontal symmetry scale V_H , our model simply reduces to the SSM with four families. In fact, eqs. (14), (15) and (16) define the fermion running masses at $\mu = V_H$. In order to deduce the fermion physical masses the RG running has to be taken into account.

words, the VEV pattern should provide the step-by-step breaking of the chiral horizontal symmetry

$$SU(4)_H \times U(1)_H \xrightarrow{V_1} SU(3)_H \times U(1)'_H \xrightarrow{V_2} SU(2)_H \times U(1)''_H \xrightarrow{V_3} \tilde{U}(1) \quad (17)$$

so that the first breaking (at the scale $V_1 \sim \langle \xi_{11} \rangle$) defines the mass terms for the first heavy family F_1 , the second breaking (at $V_2 \sim \langle \xi_{12} \rangle, \langle \xi_{22} \rangle$) for the second family F_2 etc. Through the seesaw mechanism this horizontal VEV hierarchy is reflected in the observed pattern of fermion masses. Namely, from (15) and (16) it is clear that the hierarchy among the light families is inversely proportional to the one between the heavies [14] (see also [17]), while the unbroken global symmetry $\tilde{U}(1)$ ensures the natural heaviness of the fourth family fermions.

Let us now briefly analyse the issue of horizontal symmetry breaking. The simplest way to achieve non-zero VEVs for the scalars ξ is to introduce a set of superfields S_a and Σ_a , respectively in $SU(4)_H$ singlet and adjoint representations ($a = 1, 2, \dots$), and to consider, according to our ‘renormalizability’ paradigm, the general superpotential

$$W_H = M_k \xi_k \bar{\xi}_k + \lambda_{akl} S_a \xi_k \bar{\xi}_l + \lambda'_{akl} \xi_k \Sigma_a \bar{\xi}_l + P(S, \Sigma) \quad (18)$$

where $P(S, \Sigma)$ is a general 3^{rd} order polynomial of the S_a and Σ_a fields (containing linear, bilinear and trilinear terms). Notice that this superpotential automatically respects the $U(1)_H$ invariance (10), but has no additional accidental global symmetries.

The superpotential (18) in itself does not break SUSY. Moreover, in the exact SUSY limit the vacuum state is highly degenerated – there are several zero-energy vacua with different configurations of horizontal VEVs. It would be a difficult task to provide an exhaustive analysis of all the possible vacua in the general case, that is to decide which configuration of VEVs is chosen as the true vacuum once the soft SUSY breaking terms are included. However, taking into account that after SUSY breaking the potential of the horizontal scalars has to a large extent the general structure of usual (non-SUSY) Higgs polynomial, one can argue that for a certain choice of parameters it is possible to obtain the needed pattern of VEVs (see for example the analysis in refs. [6, 18] for the case of $SU(3)_H$ symmetry).

Indeed, let us consider a first case with *only* one pair of $\xi_1 + \bar{\xi}_1$ superfields, the ones which have the largest VEV (V_1) in the exact SUSY limit. The constraint from the D -term tells us that in this case $\langle \xi_1 \rangle = \langle \bar{\xi}_1 \rangle$. Then it is easy to show that after SUSY breaking, for a proper choice of the range of values for the relevant parameters, the true vacuum can have the configuration $\langle \xi_1 \rangle = V_1 \text{diag}(1, 0, 0, 0)$ which breaks $SU(4)_H \times U(1)_H$ down to $SU(3)_H \times U(1)'_H$. Therefore, at this stage only the first family of F fermions gets a mass through the couplings (8) while the others, being protected by the residual chiral symmetry $SU(3)_H \times U(1)'_H$, remain massless.⁷

In analysing the scalar potential of the fields $\xi_2 + \bar{\xi}_2$ with next largest VEV (V_2), we have to take into the account that after ‘decoupling’ ξ_1 (i.e. substituting $\xi_1 \rightarrow \langle \xi_1 \rangle$), the symmetry group is reduced to $SU(3)_H \times U(1)'_H$, under which ξ_2 branches as **10=6+3+1**.

⁷Alternatively, for the complementary choice of the parameter range, one would have the vacuum $\langle \xi_1 \rangle \propto \text{diag}(1, 1, 1, 1)$ which breaks $SU(4)_H \times U(1)_H$ down to $SO(4)_H$. This pattern, however, does not maintain chirality and leads to degenerate fermion masses.

The VEVs which give masses to the second heavy generation belong to the **6** and **3**, and can be chosen in the form $\langle \xi_{22} \rangle$ and $\langle \xi_{12} \rangle$ respectively. These break $SU(3)_H \times U(1)'_H$ down to $SU(2)_H \times U(1)''_H$, thus respecting a residual chiral symmetry for the third and fourth family of heavy fermions which at this stage remain massless. Finally, yet another pair $\xi_3 + \bar{\xi}_3$ can develop VEVs in the $\alpha 3$ directions ($\alpha = 1, 2, 3$) thus breaking $SU(2)_H \times U(1)''_H$ down to $\tilde{U}(1)$, which acts only on the fourth family.

One could try to avoid introducing the adjoint representations Σ_a and keep in the superpotential (18) only the singlet fields S_a . Then, in the exact SUSY limit the vacuum state would have a continuous degeneration – there will be vacuum valleys. The reason for this is that in this case the superpotential would have an accidental global symmetry $SU(10)$ larger than local $SU(4)_H$. In the SUSY limit these valleys correspond to massless Goldstone modes given by certain components of the horizontal superfields. When SUSY breaking is taken into account, the radiative corrections explicitly break the extra global symmetry, lifting the vacuum degeneracy and providing ~ 100 GeV masses to the horizontal Goldstone modes, which would then become pseudo-Goldstone, massive familon-like scalars. In general these states would have diagonal as well as non-diagonal Yukawa couplings with the fermions, and in particular the strength of the couplings to the light fermions would be suppressed by a factor $\sim v/V_H$.

It is a very difficult task to provide a full analysis of the VEV pattern in this case and deduce which configuration of VEVs is fixed as the true vacuum state after SUSY breaking. Namely, already for two pairs of $\xi + \bar{\xi}$ the general VEV structure of vacuum valleys *completely* breaks the $SU(4)_H$ symmetry, not maintaining the $\tilde{U}(1)$ subgroup. One can still argue that for a proper choice of the relevant parameters the needed pattern of VEVs can be obtained. Though this possibility can be interesting from the phenomenological point of view, it deserves a special investigation. Therefore, in the following we will assume that there are no light familon-like scalars and that all the horizontal fields have $O(V_H)$ masses.

As we have mentioned at the beginning of this section, in order to generate masses for the heavy states F and F_c it is also possible to introduce horizontal Higgs fields $\chi_{[\alpha\beta]}$ in the antisymmetric **6** representations of $SU(4)_H$. The $\langle \chi_k \rangle$ VEVs would then contribute to the mass matrices of the heavy *charged* fermions through the term $\sum_k h_F^k F^\alpha F_c^\beta \chi_{[\alpha\beta]}^k$, while the corresponding term for the Majorana mass matrix of the heavy neutral states N^c is forbidden due to the antisymmetry of the representation. However, in this case the appearance of the terms like $M'_k \chi_k \chi_k$ in the superpotential for the horizontal fields would break explicitly the global $U(1)_H$ in eq. (10) and hence the residual $\tilde{U}(1)$ invariance, thus rendering unnatural the degenerate structure of the heavy matrices \hat{M}_F in eq. (13). Although for $\langle \chi_k \rangle \leq \mu_f$ the heaviness of the fourth family charged fermions would still be guaranteed, we would lose a natural explanation for a heavy ν' . In fact, through additional terms like $\xi \Sigma \chi$ sizable VEVs in the $\alpha 4$ directions would be induced also for the ξ fields. If, as it seems natural to occur, the induced VEVs are larger than the electroweak scale, the fourth family neutrino mass will also result from a seesaw giving $m_{\nu'} \sim (g_\nu v_2)^2 / \langle \xi \rangle_{induced}$. That is, the ν' will also be light, thus rendering the model phenomenologically unacceptable.

3 Phenomenological consequences of the model

As was discussed above, the quark and charged lepton masses at the scale V_H are given by eqs. (14) and (15). By assumption, the heavy fermion mass matrices $M_F^{(3)}$ are non-degenerate, and thereby have three massive eigenstates, with mass hierarchy reflecting the $SU(4)_H$ symmetry breaking pattern. The weak mixing angles are determined by the structure of these matrices, whereas the quark and lepton masses are inversely proportional to the masses of their heavy partners. For the down quark and charged lepton masses we have

$$m_{d,s,(b)} \simeq \eta_{d,(b)} m_{b'} \frac{\mu_d}{M_{D,S,(B)}} < m_{b'} , \quad m_{e,\mu,(\tau)} \simeq \eta_{e,(\tau)} m_{\tau'} \frac{\mu_e}{M_{E,\mathcal{M},(\mathcal{T})}} < m_{\tau'} \quad (19)$$

where D, S, B ($E, \mathcal{M}, \mathcal{T}$) are the mass eigenstates of $M_{D(E)}^{(3)}$, and the factors η account for the differences in the RG running of masses from the horizontal scale to lower energies. The fact that the b and τ masses are of order a few GeV, implies that the masses of the corresponding heavy states B and \mathcal{T} are not much larger (say, within one or two orders of magnitude) than the mass scale $\mu_{d,e}$. As for the top quark, the value of its mass $m_t \gtrsim 150$ GeV requires $M_T \sim \mu_u$. In this case corrections to the seesaw formula (15) should be taken into account in relating m_t to the heavy scales (see e.g. ref. [17]).

As a result of the seesaw mechanism for the fermion masses generation, the light charged states correspond to some superposition of the (f, f^c) and (F, F^c) states. It is well known that a mixing between the light $SU(2)$ doublet states f and the heavy $SU(2)$ singlets F could induce FCNC in the electroweak interactions and will also alter the flavor diagonal couplings of the light states [19]. However, in our case such a mixing is suppressed as the ratio $v_{1,(2)}/\hat{M}_F$ and thus negligibly small when compared with the present experimental bounds [20, 21]. On the other hand, the mixing between the $SU(2)$ singlets f^c and F^c can be large, since it is controlled by the ratio of the two mass scales μ_f and \hat{M}_F which, as we have seen, can be as large as $\sim 10^{-1}$ or even close to unity in the case of the t quark. However, this kind of mixing between states transforming in the same way under $SU(2) \times U(1)$ cannot affect the electroweak quantities, and is essentially unobservable.

Let us discuss now the neutrino masses. As we have already stated, three neutrinos ν_e, ν_μ, ν_τ are light Majorana particles. Their running masses at $\mu = V_H$ are determined by the heavy 'right-handed' neutrino eigenstates N_e, N_μ, N_τ at their decoupling, according to the seesaw formula (16). As for the fourth neutrino ν' , it appears to be a heavy Dirac particle with mass ~ 100 GeV. Then for the neutrino physical masses we have

$$m_{\nu_e, \nu_\mu, \nu_\tau} = \eta_\nu \frac{(m_{\nu'})^2}{M_{N_e, N_\mu, N_\tau}} , \quad m_{\nu'} \geq \frac{M_Z}{2} \quad (20)$$

where the factor η_ν accounts for the different RG running of Majorana and Dirac masses from the $SU(4)_H$ breaking scale to lower energies (for the RG running of Majorana neutrino masses see e.g. [22]). Therefore, modulo the different Yukawa couplings h_F , the neutrino mass hierarchy is expected to be qualitatively the same as the hierarchy between the quarks or the charged leptons:

$$m_{\nu_e} : m_{\nu_\mu} : m_{\nu_\tau} \sim m_u : m_c : m_t \quad \text{or} \quad m_e : m_\mu : m_\tau \quad (21)$$

Below the scale V_H our theory is just the SSM, and all FCNC phenomena related with the horizontal symmetry are strongly suppressed. Therefore, all neutrinos are effectively stable on a cosmological scale. In order to respect the cosmological upper bound [23] on the light stable neutrino masses $m_{\nu_\tau} \leq 92 \Omega h^2 \text{ eV}$ (here $\Omega = \rho/\rho_c \sim 1$ is the ratio of the energy density of the Universe to the critical density, and $h = 0.4 - 1$ is the Hubble parameter in units of $H_0 = 100 \text{ Km s}^{-1} \text{ Mpc}^{-1}$) the following lower bound on the mass of N_τ must be respected :

$$M_{N_\tau} = h_N V_H \geq \frac{\eta_\nu}{4\Omega h^2} \left(\frac{m_Z^2}{92 \text{ eV}} \right) \gtrsim 10^{11} \text{ GeV}. \quad (22)$$

In deriving this limit we have taken into account that the present age of Universe $t_0 \gtrsim 12 \cdot 10^9 \text{ yr}$ requires, for $\Omega = 1$, the Hubble parameter $h \simeq 0.5$, and we have assumed that $\eta_\nu \sim 1$. As long as the Yukawa constants h_N are $O(1)$, this bound translates into a lower bound on the smallest scale $V_H = V_3$ in the $SU(4)_H$ symmetry breaking chain (17).

Let us now address some phenomenological issues regarding the fourth family fermions. The structure of the heavy mass matrix (13) which leaves unbroken the diagonal $\tilde{U}(1)$ subgroup of $SU(4)_H \times U(1)_H$, also ensures that the fourth family is unmixed with the three lighter ones. We assume that the lightest member of the fourth generation is the neutral one ν' , as is also suggested by the analysis of ref. [11]. for simplicity we also assume that $m_{\nu'} > m_{b'}$. Then b' and ν' are stable with respect to electroweak interactions.

The presence of stable neutrinos ν' with mass in the 100 GeV range is phenomenologically and in particular cosmologically acceptable, since their contribution to the cosmological energy density is vanishingly small. Only in the presence of a sizeable $\nu' - \bar{\nu}'$ primordial asymmetry the stable relics ν' would contribute to the present cosmological density, and this contribution would still be acceptable as long as their present number density $n_{\nu'}$ does not exceed the baryon number density n_B . However, as we will argue in the following, in our model no sizeable asymmetry has to be expected for the fourth family fermions.

In contrast, the existence of stable heavy quarks carrying colour and electric charge would constitute a potential problem for the model, since it will conflict with the constraints arising from superheavy element searches, as well as with other cosmological and astrophysical constraints [24, 25]. Indeed, the stable b' would behave essentially as d quarks, hadronising into heavy ‘protons’ and giving rise to heavy hydrogen-like ‘isotopes’ with masses $\sim 100 \text{ GeV}$. The existing experimental limits on this kind of isotopes are extremely tight. For example for masses $m_{b'} < 1 \text{ TeV}$ the limit on their abundance relative to normal hydrogen is $n_{b'}/n_B < 10^{-28}$ [26].

However, the exchange of the $SU(4)_H$ gauge bosons Z_H would allow the heavy quark to decay, dominantly through the channel $b' \rightarrow b \bar{\nu}_\tau \nu'$, with a lifetime ⁸

$$\tau_{b'} \sim \left(\frac{V_H}{v} \right)^4 \left(\frac{m_\mu}{m_{b'}} \right)^5 \tau_\mu = \left(\frac{V_H}{10^{12} \text{ GeV}} \right)^4 \left(\frac{150 \text{ GeV}}{m_{b'}} \right)^5 \cdot 4 \cdot 10^{17} \text{ s} \quad (23)$$

where $V_H = V_3$ is the lowest scale in the horizontal symmetry breaking (see eq. (17)), v is the electroweak scale and $\tau_\mu = \tau(\mu \rightarrow e \bar{\nu}_e \nu_\mu) \simeq 2.2 \times 10^{-6} \text{ s}$. is the muon lifetime. We can

⁸Much faster decay $b' \rightarrow b + J$ can occur if there are pseudo-Goldstone familon-like scalars J , arising from the breaking of accidental global $SU(10)$ symmetry of the superpotential (see the discussion in Sect. 2). However, here we will not consider such a possibility.

use cosmological arguments, together with the experimental limits on searches of heavy isotopes, to put an upper bound on $\tau_{b'}$, which in turn will translate in an upper limit on V_H . Indeed, taking into account the finite lifetime of the heavy quarks, their present number abundance relative to baryons is $n_{b'}/n_B = r_0 \exp(-t_0/\tau_{b'}) \leq 10^{-28}$, where $r_0 = (n_{b'}/n_B)_0$ represents the relic abundance for stable b' . From this equation we get the upper limit on the b' lifetime

$$\tau_{b'} \leq 3 \cdot 10^{15} h^{-1} (1 + 0.036 \lg r_0)^{-1} \text{ s.} \quad (24)$$

One cannot say definitely what is the value of r_0 , due to many theoretical uncertainties related to the actual annihilation cross section for the b' . However, an estimate of the relic abundance of heavy stable d -type quarks has been given in [24]. Under the assumption that there is no cosmological baryon asymmetry between the b' and \bar{b}' , it was found that for $m_{b'} \sim 150 \text{ GeV}$ the energy density of these relics, relative to critical density (namely $\Omega_{b'} h^2$) could range from 10^{-9} to 10^{-4} (smaller values are obtained for lighter b' masses). The lower limit corresponds to the case when the relic density is determined by the annihilation after the QCD phase transition, and it was obtained by taking as an upper bound on the annihilation cross section the geometrical cross section ($\sigma_0 \sim 100 \text{ mb}$). The upper limit was obtained under the opposite assumption, namely that annihilation after confinement is negligible, and that the relic density is essentially determined by the QCD annihilation cross section for free quarks. Then the ratio of the b' to baryon number densities $r_0 = (n_{b'}/n_B)_0 = \Omega_{b'}/\Omega_B \cdot m_B/m_{b'}$ lies in the range $3 \cdot 10^{-10} < r_0 < 3 \cdot 10^{-5}$ where we have taken $\Omega_B \sim 0.02$ as suggested by nucleosynthesis estimates. As is discussed in [24], the most reasonable assumption is that the relevant annihilation process happens after confinement, however with a cross section much smaller than the geometrical one, giving $r_0 \sim 10^{-7} - 10^{-8}$. Clearly, in the presence of a sizeable Baryon asymmetry in the fourth family sector, the relic abundance of the heavy b' quarks would be some orders of magnitude larger than the quoted estimates.

As we see, the bound (24) very weakly depends on the initial b' abundance. Even if we allow for a large primordial asymmetry for the b' , and let r_0 range between $1 - 10^{-10}$, by taking $h = 0.5$ we obtain $\tau_{b'} \leq 6 \cdot 10^{15} - 10^{16} \text{ s}$. On the other hand, according to eq. (23), this bound translates into an extremely strong upper limit

$$V_H \leq (1.1 - 1.3) \cdot 10^{11} \cdot \left(\frac{m_{b'}}{150 \text{ GeV}} \right)^{5/4} \text{ GeV.} \quad (25)$$

Thus, the scale V_H cannot exceed $2 \cdot 10^{11} \text{ GeV}$ unless $m_{b'} \gg 200 \text{ GeV}$.

More stringent limit on r_0 and $\tau_{b'}$ can be derived by considering that the late decay of the b' can cause a significant contribution to observed cosmic ray fluxes, in particular to the isotropic diffuse gamma-ray background [27]. Indeed, at the moment of decay, the b' quarks are bounded within colorless hadrons like $b'ud$ or $\bar{b}'u$ [24]. Then in the decay $b' \rightarrow b\bar{\nu}_\tau\nu'$ an unstable hadronic state emerges with the excitation energy $E_0 \simeq \frac{1}{3}m_{b'}$. This will essentially appear as a hadronic jet with the b quark being the leading particle. The fragmentation of this jet produces π^0 , η etc., with the subsequent radiative decay resulting in a specific photon spectrum. Obviously, the amount of produced photons is directly proportional to r_0 . In order to estimate their flux in the present Universe, the redshift of their energies has to be taken into account as well. As long as the decay happens at the matter dominated epoch, and the small amount of relativistic decay products does

not affect sensibly the Universe expansion rate, we have $1 + z = (t_0/\tau_{b'})^{2/3} \sim 10 - 20$ for the values of $\tau_{b'}$ estimated above. We also need to know what fraction of the jet energy E_0 is taken by the photons and what is the energy spectrum. These issues were studied in ref. [28], where the photon spectra produced at jet hadronization were computed for different leading particles using a Monte Carlo simulation program [29]. It was shown that these spectra exhibit a remarkable scaling property in terms of the variable $x = E_\gamma/E_0$, and in the case of leading particle being a b quark, the photons carry away about 25 percent of the initial jet energy. Using the results of ref. [28] we have computed the value of the isotropic cosmological gamma-flux $d\Phi_\gamma/dE_\gamma$ and we have compared it with the existing observational limits (see [30] and references therein). For example, for $E_\gamma = 100$ MeV the experimental upper bound on the γ -flux is of about $10^{-7} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ MeV}^{-1}$. We have obtained that the cosmic gamma-flux produced due to b' decay at $z = 10 - 20$, saturates the above bound for $r_0 \sim 10^{-7}$ which is close to, but still not in conflict with our estimate of the b' relic abundance in the $b'\text{-}\bar{b}'$ symmetric case. Substantially larger r_0 would require much larger redshift, and hence much smaller $\tau_{b'}$. On the other hand, the lower bound (22) on the horizontal symmetry breaking scale V_H already excludes much smaller lifetimes.

This analysis implies that r_0 should be rather small, so that any sizeable cosmological baryon asymmetry between b' and \bar{b}' is excluded. This severely constrains the possible baryogenesis mechanisms applicable to our model. The appearance of baryon asymmetry in the fourth family in itself is hardly expected, since it is unmixed with the other three families and hence it has no source of CP violation. However, the sphaleron effects [31, 32] would immediately redistribute the baryon asymmetry produced within the first three families to the fourth family fermions. Therefore, no mechanism is acceptable which generates the baryon asymmetry before the sphaleron effects are switched off, that is before the electroweak phase transition.⁹ In the context of our model the most appealing possibility is to assume that no baryon asymmetry is produced before the electroweak epoch, and baryogenesis takes place at the electroweak (first order) phase transition. Such a baryogenesis mechanism is associated with the walls of the expanding bubbles of the broken phase [35]. Outside of the bubbles electroweak symmetry is unbroken, quarks are massless and the rate of the fermion number violation due to sphaleron transitions greatly exceeds the Universe expansion rate. Inside the bubbles the quarks are massive due to non-zero VEVs of the Higgs fields, while the sphaleron processes are strongly suppressed and fermion number is effectively conserved. Then baryon asymmetry inside the bubbles could be produced (and maintained) due to CP violating effects, as a difference between the quark and anti-quark fluxes penetrating the walls from the unbroken phase to the

⁹In principle, in our model the baryogenesis with non-zero B-L could occur due to CP violation effects in out-of-equilibrium decays $N^c \rightarrow l + \Phi$ of the heavy right-handed neutrino [33] (for the viability of this mechanism in the SUSY case see ref. [34]), or in the decays of $SU(4)_H$ gauge or scalar bosons. Then sphaleron effects would immediately transfer the produced net lepton number into a baryon asymmetry also in the fourth family sector. Fortunately, our model naturally avoids the possibility of such a leptobaryogenesis. As it was shown in ref. [34], the large scale density fluctuations hinted by the COBE measurements require rather low inflationary reheat temperature ($T_R \sim 10^8$ GeV) and correspondingly low inflaton mass ($m_\eta \sim 10^{11}$ GeV). On the other hand, the lower bound (22) on M_N (and respectively on V_H) tells us that masses of the right-handed neutrinos and horizontal bosons should exceed 10^{11} GeV, and therefore they are not produced after inflation.

broken one. Obviously, this concerns only the first three family fermions. Since the fourth family is unmixed, has no CP violation, and moreover all the fermions are very heavy, no baryon excess is expected in this sector. Although the viability of such a baryogenesis in the SM is still disputed in the literature [36], in the context of SSM it could be more effective and sufficient for providing the observed baryon asymmetry. Clearly this topic deserves additional special considerations.

According to eq. (20), the upper limit on the horizontal symmetry breaking scale $V_H \lesssim 2 \cdot 10^{11}$ GeV together with the experimental limit $m_{\nu'} \geq M_Z/2$ translates into a lower bound on the τ -neutrino mass:

$$m_{\nu_\tau} \geq \frac{\eta_\nu}{h_N} \frac{m_Z^2}{4V_H} \gtrsim (1 - 10) \text{ eV} \quad (26)$$

where in the numerical estimate we have taken into account the $O(1)$ uncertainties in the relative renormalization factor η_ν and in the Yukawa coupling h_N (for perturbativity we have to assume $h_N < 3$ at $\mu = V_H$). A ν_τ with mass in the range $1 - 10$ eV will give a sizeable contribution to the cosmological energy density as a hot dark matter (HDM) component, while according to (21) ν_μ and ν_e are expected to have much smaller masses. We remind here that the COBE measurements of the cosmic microwave background anisotropy, together with other data on the density distribution of the Universe at all distance scales (galaxy-galaxy angular correlations, correlations of galactic clusters, etc.), can all be fit by assuming some HDM admixture to the dominant CDM component [37]. The best fits hint to a neutrino mass $m_{\nu_\tau} \sim 5 - 7$ eV [38] which does appear naturally in our model. As for the CDM itself, in our R parity conserving SUSY model it is naturally provided by the lightest supersymmetric particle (LSP), presumably a neutralino.

As we commented earlier, the neutrino mass hierarchy should be qualitatively the same as that for the charged quarks and leptons. However, the spread in the Yukawa coupling constants h_F does not allow to put severe limits on the other neutrino masses. For example, by taking $m_{\nu_\mu}/m_{\nu_\tau} \sim m_c/m_t$, as is suggested by the first estimate in eq. (21), one obtains $m_{\nu_\mu} \sim (2 - 5) \cdot 10^{-3}$. This range corresponds to the Mikheyev-Smirnov-Wolfenstein (MSW) solution of the solar neutrino problem [39] via $\nu_e \rightarrow \nu_\mu$ oscillations. Alternatively, if we had to attempt an explanation of the deficit of the atmospheric ν_μ via $\nu_\mu \rightarrow \nu_e$ oscillations, then we would need $m_{\nu_\mu} \sim 0.1$ eV [40] which is compatible with the second estimate in eq. (21). Obviously the MSW explanation to the solar neutrino deficit would not be viable in this latter case.

4 Conclusions

In this paper we have put forward the idea that natural conservation of R parity in SUSY models can be guaranteed in the presence of some suitable horizontal gauge symmetries. We have shown how these symmetries can indeed forbid all the dangerous terms in the superpotential, which are trilinear in the fermion superfields, and how an accidental Z_2 matter parity (equivalent to R parity) then follows in a quite satisfactory way only due to gauge invariance and to the field content of the model. On theoretical and phenomenological grounds, we have uniquely identified $SU(4)_H$ as the only viable horizontal gauge

group. As a consequence, our scheme requires a fourth generation of superfields in addition to the three known families. Hence we have focused our analysis on a four generation SUSY model based on the SM vertical gauge group $SU(3) \times SU(2) \times U(1)$ and equipped with an $SU(4)_H$ anomaly free horizontal gauge symmetry. We have discussed in some details the structure of the fermion mass matrices arising in the model, as well as some possible patterns for the breaking of the horizontal symmetry. We have shown that the simplest symmetry breaking scheme which ensures that all the horizontal modes acquire large masses, can also lead to a particular form for the fermion mass matrices which ensures that the masses for the fourth generation fermions are naturally close to the electroweak scale. A recent RG analysis of SUSY models with four generations [11] does apply straightforwardly to our case, and suggests that if the hypothesis of unification of the vertical gauge group is correct, then at least the new leptons should be well in the reach of LEP II. As regards the light masses of the first three families, our model leads to a seesaw suppression of their magnitude from the electroweak scale down to the observed values. In particular, this is achieved without the need of any tuning for the Yukawa couplings, which can be assumed to be all $O(1)$ or close to the typical values of the gauge couplings. By means of cosmological and astrophysical arguments, we have managed to constrain rather precisely the scale V_H at which the horizontal gauge symmetry is completely broken, obtaining a very narrow window around 10^{11} GeV. Below this scale, our model is essentially the SSM with four generations. In turn, the upper bound on the scale V_H feeds back into the neutrino mass matrix, implying a mass for the τ -neutrino not much lighter than a few eV. A neutrino mass in this range will then give a sizeable contribution to the present energy density of the Universe. Thus, our model naturally provides cosmological HDM in the form of ν_τ 's and, due to R parity conservation, also CDM in the form of stable LSPs. Since in our scheme conservation of R-parity is ensured by the horizontal gauge symmetry independently of the particular choice for the vertical gauge group, it would be interesting to extend the present analysis to phenomenologically appealing GUT models, such as $SU(5)$ or E_6 , for which R-parity conservation is not automatic.

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