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Unconventional superstring-derived E_6 models and neutrino phenomenology

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Conventional superstring-derived E_6 models can accommodate small neutrino masses if a discrete symmetry is imposed which forbids tree-level Dirac neutrino masses but allows for radiative mass generation. Since the only possible symmetries of this kind are known to be generation dependent, we explore the possibility that the three sets of light states in each generation do not have the same assignments with respect to the **27** of E_6, leading to nonuniversal gauge interactions under the additional U(1)′ factors for the known fermions. We argue that models realizing such a scenario are viable, with their structure being constrained mainly by the requirement of the absence of flavor-changing neutral currents in the Higgs sector. Moreover, in contrast with the standard case, rank 6 models are not disfavored with respect to rank 5. By requiring the number of light neutral states to be minimal, these models have an almost unique pattern of neutrino masses and mixings. We construct a model based on the unconventional assignment scenario in which (with a natural choice of the parameters) $m_{\nu_\tau} \sim 10$ eV is generated at one loop, m_{ν_μ} is generated at two loops and lies in a range interesting for the solar neutrino problem, and ν_e remains massless. In addition, since baryon and lepton number are conserved, there is no proton decay in the model. In order to illustrate the nonstandard phenomenology implied by our scheme we also discuss a second scenario in which an attempt for solving the solar neutrino puzzle with matter-enhanced oscillations and practically massless neutrinos can be formulated, and in which peculiar effects for the $\nu_\mu \to \nu_\tau$ conversion of the upward-going atmospheric neutrinos could arise as well.

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I. INTRODUCTION

It is generally believed that neutrinos possess very small but nonvanishing masses. While there is no fundamental reason for the neutrinos to be exactly massless, small ν masses are needed in any particle physics explanation of the solar neutrino problem and, at the same time, they imply several interesting phenomenological consequences. A very attractive way of generating naturally small neutrino masses is through the use of the seesaw mechanism [1]. In E_6 supersymmetric grand unified theories (GUT's) [2], as derived from superstring theories, the seesaw mechanism cannot be easily implemented since the Higgs representation necessary to generate a large Majorana mass for the right-handed neutrinos is absent. However, even in the absence of Majorana terms, small masses can be generated through radiative corrections in models in which at the lowest order $m_\nu=0$. As was pointed out by Campbell *et al.* [3] and Masiero, Nanopoulos, and Sanda [4] E_6 GUT's do offer the possibility of implementing this second mechanism.

The fermion content of models based on E_6 is enlarged with respect to the standard model (SM). In fact, two additional lepton SU(2) doublets, two SU(2)-singlet neutral states, and two color-triplet SU(2)-singlet d-type quarks are present in the fundamental representation of the group. In order to forbid neutrino masses at the tree level an appropriate discrete symmetry has to be imposed on the superpotential of the model. Branco and Geng (BG) [5] have shown that no generation-blind symmetry exists that forbids nonvanishing neutrino masses at the tree level, and at the same time allows for the radiative generation of the masses at one loop. As a result, in order to implement this mechanism a symmetry that does not act in the same way on the three generations is needed. In Sec. II we will briefly review the main features of the conventional E_6 models, and we will establish the notations.

Our work stems from the observation that once we choose to build a model based on a symmetry that does distinguish among the different generations, there is no reason, in principle, to expect that this symmetry will result in a set of *light* fermions (i.e., the known states) that will exactly replicate throughout the three generations. To state this idea more clearly, we wish to suggest the possibility that what we call "ν_τ" is actually assigned to an SU(2) doublet which has a different embedding in E_6 with respect to the doublet that contains what we call "ν_e." As a consequence the two neutrinos will have different E_6 gauge interactions. More drastically, we can envisage the possibility that the gauge interactions of the d quarks and leptons of one family (say the third one) are different from those of the corresponding states of the other two generations. Obviously, experimentally we know that the SU(2)×U(1) interactions of the fermions do respect universality with a high degree of precision; however, in the class of models that we want to investigate, one or two additional U(1)′ Abelian factors are always present, implying additional massive neutral gauge

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bosons possibly at energies ~ 1 TeV or less. The possibility that the U(1)′ interactions of the known fermions could violate universality then is indeed still phenomenologically viable.

In Sec. III we will develop a scenario that realizes this idea. Starting from the assumption of unconventional assignments (UA's) for the light neutrino of the third generation, we will show that the need for UA's is reflected in the d-quark sector as well, thus leading to a third generation of light fermions which is not a replica of the first two. In Sec. IV we will concentrate on the neutrino phenomenology, and we will describe the pattern of masses and mixings that is predicted by our scheme. We believe that the unconventional scenario that we are going to analyze here could be interesting in itself, since it is not *a priori* obvious that models in which the "low"-energy gauge interactions of the known fermions are not universal can be consistently constructed. However, it turns out that beyond being viable, these models also lead to an interesting phenomenology, especially in the neutrino sector, and as well imply some rather unusual consequences. In order to illustrate this, at the end of Sec. IV we will discuss a particular model in which a few peculiar effects in the propagation of the neutrinos through matter could arise. We will formulate an attempt to find a Mikheyev-Smirnov-Wolfenstein- (MSW)-like solution to the solar neutrino puzzle [6] with "almost massless" neutrinos (i.e., $m_\nu \ll 10^{-3}$ eV). We will also address the issue that these nonstandard effects could lead to a suppression of the "ν_μ"-"ν_τ" oscillations for the high-energy upward-going atmospheric neutrinos. Finally, in Sec. V we will summarize our results and draw the conclusions.

II. CONVENTIONAL E_6 MODELS

In E_6 GUT's, matter fields belong to the fundamental **27** representation of the group. E_6 contains $SO(10)\times U_\psi(1)$ as a maximal subalgebra, and the **27** branches to $\mathbf{1}+\mathbf{1}+\mathbf{16}$ of SO(10). In turn, SU(10) contains $SU(5)\times U_\chi(1)$.

The SO(10), SU(5), $U_\psi(1)$, and $U_\chi(1)$ assignments for the states in the **27** representation are listed in Table I. Usually the known particles of the three generations are assigned to the **16** representation of SO(10) that also contains a right-handed neutrino:

$$[\mathbf{16}]_i = \left[Q \equiv \begin{pmatrix} u \\ d \end{pmatrix}, u^c, e^c, d^c, L \equiv \begin{pmatrix} \nu \\ e \end{pmatrix}, \nu^c \right]_i , \quad i=1,2,3 , \tag{2.1}$$

while the **10** and the **1** of SO(10) contain the new fields

$$\begin{aligned} [\mathbf{10}]_i &= \left[H^c \equiv \begin{pmatrix} E^c \\ N^c \end{pmatrix}, h, H \equiv \begin{pmatrix} N \\ E \end{pmatrix}, h^c \right]_i , \\ [\mathbf{1}]_i &= [S^c]_i, \quad i=1,2,3 . \end{aligned} \tag{2.2}$$

As is clear from Table I there is an ambiguity in assigning the known states to the **27**, since under the SM gauge group

$$\mathcal{G}_{SM} \equiv SU(3)_c \times SU(2)_L \times U(1)_Y$$

the $\bar{\mathbf{5}}_{(10)}$ in the **10** of SO(10) has the same field content as the $\bar{\mathbf{5}}_{(16)}$ in the **16**. The same ambiguity is also present for the two $\mathcal{G}_{SM}$ singlets, namely $\mathbf{1}_{(1)}$ and $\mathbf{1}_{(16)}$.

Since E_6 is rank 6 as many as two additional neutral gauge bosons can be present, corresponding, for example, to some linear combinations of the $U_\chi(1)$ and $U_\psi(1)$ generators. The fermion interactions with these gauge bosons will depend on the specific assignments. The two additional neutral gauge bosons are usually parametrized as

$$\begin{aligned} Z'_\beta &= Z_\psi \sin\beta + Z_\chi \cos\beta , \\ Z''_\beta &= Z_\psi \cos\beta - Z_\chi \sin\beta , \end{aligned} \tag{2.3}$$

and in the following we will often collectively refer to them as Z_β bosons. In the presence of at least one "light" Z_β ($M_\beta \lesssim 1-2$ TeV), different assignments will lead to a different phenomenology that could be tested, e.g., at the CERN Large Hadron Collider (LHC) and Superconducting Super Collider (SSC) energies or possibly even at the CERN e^+e^- collider LEP II. In contrast, it is clear that in the limit $M_\beta \to \infty$ the choice of the assignment is irrelevant as far as we are only concerned with the gauge interactions. However, as we will show, even in this limit the requirement of $U_\beta(1)$ gauge invariance for the superpotential, together with the phenomenological constraints on the absence of flavor-changing neutral currents (FCNC's) in the Higgs sector, will have far reaching consequences for determining the structure of

TABLE I. SO(10), $U_\psi(1)$, SU(5), and $U_\chi(1)$ assignments for the left-handed fermions of the **27** fundamental representation of E_6. The SU(2) doublets H^c, H, L, and Q are explicitly written in components. The Abelian charges q_ψ and q_χ can be derived from the quantum numbers listed in the square brackets by dividing by $c_\psi = 6\sqrt{2/5}$ and $c_\chi = 6\sqrt{2/3}$, respectively. The charges are normalized to the hypercharge axis according to $\sum_{f=1}^{27}(q^f_{\psi,\chi})^2 = \sum_{f=1}^{27}(\frac{1}{2}Y^f)^2 = 5$.

	S^c	$\begin{pmatrix} E^c \\ N^c \end{pmatrix}$ h	$\begin{pmatrix} N \\ E \end{pmatrix}$ h^c	ν^c	$\begin{pmatrix} \nu \\ e \end{pmatrix}$ d^c	e^c u^c $\begin{pmatrix} u \\ d \end{pmatrix}$
SO(10) $(c_\psi q_\psi)$	**1**(4)	**10**(−2)		**16**(1)		
SU(5) $(c_\chi q_\chi)$	**1**(0)	**5**(2)	$\bar{\mathbf{5}}$(−2)	**1**(−5)	$\bar{\mathbf{5}}$(3)	**10**(−1)

the viable models.

The most general renormalizable superpotential arising from the coupling of the three **27**'s in Table I and invariant under the SM gauge group is [7]

$$W = W_1 + W_2 + W_3 + W_4 ,$$

where

$$\begin{aligned} W_1 &= \lambda^{(1)} H^c Q u^c + \lambda^{(2)} H Q d^c + \lambda^{(3)} H L e^c + \lambda^{(4)} S^c h h^c , \\ W_2 &= \lambda^{(5)} h u^c e^c + \lambda^{(6)} L Q h^c + \lambda^{(7)} \nu^c h d^c , \\ W_3 &= \lambda^{(8)} h Q Q + \lambda^{(9)} h^c u^c d^c , \\ W_4 &= \lambda^{(10)} H^c L \nu^c + \lambda^{(11)} H^c H S^c . \end{aligned} \tag{2.4}$$

The Yukawa couplings in (2.4) are three index tensors in generation space, e.g.,

$$\lambda^{(1)} H^c Q u^c \equiv \lambda^{(1)}_{ijk} H^c_i Q_j u^c_k$$

with $i,j,k=1,2,3$ generation indices, and, in general, they are not constrained by the E_6 Clebsch-Gordan relations [8]. The presence of W_4 would produce tree-level Dirac masses for all the neutral states in the model. In particular, the ν's would acquire a Dirac mass $m_\nu = \lambda^{(10)} \langle \tilde{N}^c \rangle$. An unnatural tuning of the $\lambda^{(10)}$ Yukawa couplings is then required to make these masses small. If $\lambda^{(10)}$ were absent, then ν and ν^c would be massless at the tree level. Furthermore, if at the same time the couplings $\lambda^{(6)}$ and $\lambda^{(7)}$ in W_2 were nonvanishing, naturally small Dirac masses would be produced at the one-loop level through diagrams such as the one depicted in Fig. 1 [3,4]. However, a problem arises due to the fact that the simultaneous presence of $W_2(\supset \lambda^{(6)}, \lambda^{(7)})$ and W_3 induces fast proton decay. The vanishing of W_3 can cure this problem still allowing for radiative neutrino masses. The conclusion is that in the conventional schemes the vanishing of W_3 and $\lambda^{(10)}$ together with nonvanishing $\lambda^{(6)}$ and $\lambda^{(7)}$ couplings is required in order to have an interesting neutrino phenomenology and not to conflict with the limits on the proton lifetime. As was discussed in detail by BG [5], the correct pattern of vanishing Yukawa couplings leading to small ν_e, ν_μ, and ν_τ masses can be realized only by means of a generation-dependent discrete symmetry, i.e., a symmetry under which the fields transform with a generation-dependent phase $\psi_j \to e^{i\alpha_{\psi_j}} \psi_j$ where j is a generation index.

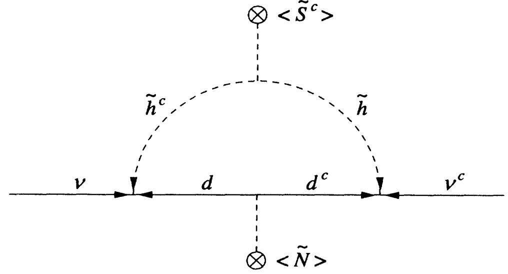

FIG. 1. A typical diagram contributing to the neutrino Dirac masses at the one-loop level.

III. THE UNCONVENTIONAL ASSIGNMENT SCENARIO

Once we give up the assumption that, from the point of view of the symmetry transformation properties, the three generations are exact replicas of one another, we may also abandon the assumption that the known particles of the three generations should be identified with the same states in the three different **27** representations. We will now explore the possibility of constructing a consistent model in which the assignments of the known fermions to the **27** are different for the different generations. Models of these kind turn out to be phenomenologically viable and they clearly imply a few unusual phenomenological consequences. For example, the known fermions will have generation-dependent neutral current (NC) gauge interactions induced by Z_β exchange, due to the difference in the $U_\psi(1)$ and $U_\chi(1)$ charges.

A. The assignments for the leptons

As a starting point for investigating E_6 models with UA we will assume that what we call "ν_τ" is, in fact, the N_3 weak doublet neutral state belonging to the $\bar{5}_{10}$, while ν_e and ν_μ are still assigned as usual to the $\bar{5}_{16}$. We will henceforth use quotation marks to denote the known states with their conventional labels, since they might not correspond to the entries in Table I. Labels not enclosed within quotation marks will always refer to the fields listed in the table. We will also keep the same assignments as well as the same transformation properties under the discrete symmetry group for the *known* states of the first and second generations (other assignments leading to different models are trivially obtained by interchanging the generation labels, e.g., $1 \leftrightarrow 3$). Accordingly, when we refer to the first generation it is understood that the same applies to the second generation as well. Latin indices $i,j,\ldots$ run from 1 to 3, while the Greek indices $\alpha,\beta,\ldots=1,2$ will refer only to the first two generations.

With the notation given in (2.1) and (2.2), and referring to the **10** and **16** representations of SO(10), our starting assumption for the assignments of the three SU(2)-doublet light neutrinos reads

$$\begin{aligned} \text{“}\nu_\alpha\text{”} &\in L_\alpha \in \mathbf{16}, \quad \alpha = 1,2 , \\ \text{“}\nu_\tau\text{”} &\in H_3 \in \mathbf{10} . \end{aligned} \tag{3.1}$$

In order to realize this scenario we first have to require that the three-level masses for ν_α and N_3 vanish. This can be achieved by setting

$$\lambda^{(10)}_{\langle i \rangle \alpha j}(H^c_i L_\alpha \nu^c_j) = 0, \quad \lambda^{(11)}_{\langle i \rangle 3 j}(H^c_i H_3 S^c_j) = 0 . \tag{3.2}$$

For the sake of clarity we have enclosed inside angular brackets the indices labeling the particular vacuum expectation values (VEV's) which are relevant for the actual discussion. From the LEP measurement of the number of weak-doublet neutrinos we know that all the remaining SU(2)-doublet neutral states N_α, ν_3, and N^c_i must be heavy ($\gtrsim 50$ GeV). This in turn implies that the following terms must be nonvanishing:

$$\lambda^{(10)}_{i3\langle\beta\rangle}(H^c_i L_3 \nu^c_\beta)\neq 0, \quad \lambda^{(11)}_{i\alpha\langle\beta\rangle}(H^c_i H_\alpha S^c_\beta)\neq 0 \ . \qquad (3.3)$$

If any two of the scalar components of the three neutral fields in the trilinear terms $\lambda^{(10)}$ and $\lambda^{(11)}$ acquire a VEV, then the VEV of the third scalar field must also be nonvanishing. Therefore, almost all of the neutral scalars in W_4 (doublets and singlets) will eventually acquire a VEV. In particular, it is not difficult to see that in order to have all the H^c_i heavy, none of their scalar components can be prevented from eventually acquiring a VEV. This is the reason why we have forbidden all the couplings between the massless neutrinos and the H^c_i fields in (3.2). We note that at the same time the conditions (3.2) allow for $\langle\tilde{\nu}_\alpha\rangle=\langle\tilde{N}_3\rangle=0$ which, as we will discuss, is necessary if we want to prevent spontaneous violation of lepton number.

Because of our choice of the light states and the vanishing of the couplings in (3.2) the following VEV's can be generated: $\langle\tilde{\nu}^c_\alpha\rangle$, $\langle\tilde{S}^c_\alpha\rangle$, $\langle\tilde{N}_\alpha\rangle$, $\langle\tilde{\nu}_3\rangle$, and $\langle\tilde{N}^c_i\rangle$. It was argued in Ref. [3] that in the conventional E_6 models it might be difficult to achieve $\langle\tilde{\nu}^c\rangle\neq 0$ since the Yukawa couplings $\lambda^{(7)}$ and/or $\lambda^{(10)}$ that are needed for driving $m^2_{\tilde{\nu}^c}$ negative through the renormalization group, are constrained to be either vanishing or too small to generate this VEV. This implies that the set of VEV's present in the conventional models does not allow for lowering the rank of the gauge group by more than one, and since the SM is rank 4 it is probably not possible to construct a dynamical model based on rank 6. In contrast, we will see that in the present scheme some of the $\lambda^{(7)}$ are not constrained to be particularly small, and indeed some of the $\lambda^{(10)}$ couplings [those in (3.3)] are expected to be rather large. We can conclude that rank 6 models are indeed viable in our UA scenario since $\langle\tilde{\nu}^c\rangle\neq 0$ can be easily achieved.

Now, in order to allow for radiatively generated Dirac masses, we need massless right-handed neutrinos as well. For the sake of simplicity, we will require a *minimum* number of light neutral SU(2) singlets. In (3.3) we have already assumed that the couplings involving ν^c_3 and S^c_3 are forbidden, thus preventing their fermionic component from acquiring a mass at tree level.

Another consequence of (3.3) regards the charged lepton mass matrix. In fact, it is clear that the e_α fields and the left-handed "τ" lepton E_3 have to acquire their mass from the $\lambda^{(3)}$ Yukawa coupling, since the $\lambda^{(10)}$ and $\lambda^{(11)}$ couplings for these states are forbidden. Then the "τ" lepton mass term $m_\tau E_3 e^c_3$ must be generated from the $\tilde{L}_3$ scalar doublet, while m_e and m_μ are generated from the VEV of one of the $\tilde{H}_\alpha$ Higgs multiplet (for example, $\tilde{H}_2$). As a consequence of this, it is true that, in general, all the right-handed leptons e^c_j will couple to both $\tilde{L}_3$ and $\tilde{H}_2$ through the couplings $\lambda^{(3)}_{3\langle 3\rangle j}$ and $\lambda^{(3)}_{\langle 2\rangle\beta j}$. It is well known that this situation can give rise to dangerous lepton-flavor-violating (LFV) couplings between the fermions and the Higgs fields [9] since the rotation that diagonalizes the lepton mass matrix does not diagonalize the fermion couplings to the Higgs fields. In this respect the couplings with $\tilde{H}_1$ are also dangerous since its neutral scalar component will eventually acquire a VEV as well. In addition, nonzero mass terms connecting $E_\alpha - e^c_j$ ($e_3 - e^c_j$) which can be generated by nonvanishing $\lambda^{(3)}_{\alpha\langle 3\rangle j}$ ($\lambda^{(3)}_{\langle\alpha\rangle 3j}$) couplings will induce an isospin-violating ($\Delta I=\frac{1}{2}$) light-heavy mixing between the e^c and the E^c fields. It is well known [3,10] that a mixing of this kind can give rise to tree-level LFV processes mediated by Z_0 exchange. Therefore, in order not to conflict with the tight limits on LFV processes such as $\mu\rightarrow eee$, μ-e conversion in muonic atoms, etc., we have to require all the $\lambda^{(3)}$ couplings to be absent, with the exception of $\lambda^{(3)}_{\langle 2\rangle\alpha\beta}$ and $\lambda^{(3)}_{3\langle 3\rangle 3}$ which are needed to generate masses for the light leptons. Together with the conditions in (3.2), this additional requirement ensures that all the light-heavy lepton mixings are absent. In addition, the resulting mass matrix for the light leptons turns out to be block diagonal, with

$$[m_l]_{\alpha\beta}=\lambda^{(3)}_{\langle 2\rangle\alpha\beta}\langle\tilde{N}_2\rangle, \quad [m_l]_{33}=\lambda^{(3)}_{3\langle 3\rangle 3}\langle\tilde{\nu}_3\rangle \ . \qquad (3.4)$$

We note that an important consequence of the constraints just discussed is that any possible mixing of the third generation neutrino can only arise in the neutrino sector.

B. The assignments for the quarks

At this stage three SU(2)-doublet and two SU(2)-singlet neutral states are massless, namely, ν_α, N_3 and ν^c_3, S^c_3. Dirac masses for these states cannot be generated via loops of leptons, since this would require some of the couplings in (3.2), but they can indeed be induced by loops involving quarks through a set of diagrams that are analogous to the one depicted in Fig. 1.

The relevant couplings for generating these diagrams are $\lambda^{(2)}HQd^c$, $\lambda^{(4)}S^c hh^c$, $\lambda^{(6)}LQh^c$, and $\lambda^{(7)}\nu^c hd^c$. In addition to appearing as vertices for the external states, each of these couplings will also provide a mass insertion for the quarks running inside the loop. We have recast these couplings into the schematic partial diagrams A and B depicted in Fig. 2. In these diagrams it is understood that one of the two neutral fields is external while the other one corresponds to a VEV insertion. We will label A_1 and B_1 the partial diagrams A and B in which the VEV insertions correspond, respectively, to the SU(2) singlets $\langle\tilde{S}^c_\alpha\rangle$ and $\langle\tilde{\nu}^c_\alpha\rangle$, and A_2 and B_2 are those diagrams in which the VEV insertion corresponds to the doublets $\langle\tilde{\nu}_3\rangle$ and $\langle\tilde{N}_\alpha\rangle$. By gluing two partial diagrams together, we can generate the following entries in the neutrino Dirac mass matrix:

$$(\nu_\alpha \ \ N_3)\begin{bmatrix}[A_1A_2] & [A_1B_2]\\ [B_1A_2] & [B_1B_2]\end{bmatrix}\begin{bmatrix}S^c_3\\ \nu^c_3\end{bmatrix} \ . \qquad (3.5)$$

Obviously, not all the couplings $\lambda^{(m)}_{ijk}$ ($m=2,4,6,7$) are allowed, and we will now proceed to select the couplings that must be forbidden. In the first place, we note that in order to generate a diagram that will provide a radiative mass for N_3, at least one of the couplings $\lambda^{(2)}_{3ij}H_3Q_id^c_j$ in B_1 must be nonvanishing. However, $\lambda^{(2)}$ is precisely the Yukawa that, in the conventional models, is needed to give mass to the down-type quarks, for example,

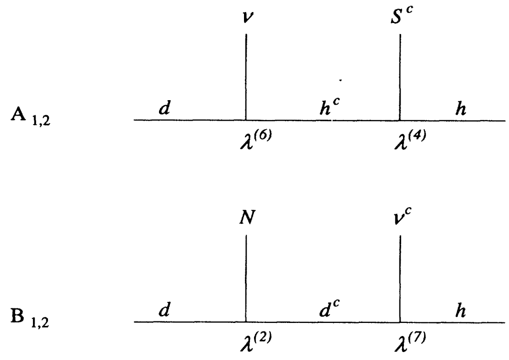

FIG. 2. Schematic partial diagrams for generating one-loop mass entries in the Dirac neutrino mass matrix. Both in A and B it is understood that one of the two neutral fields is external, while the other one represents a VEV insertion that provides a mass for the quark running inside the loop. A_1 and B_1 correspond, respectively, to the diagrams where the mass insertions are provided by the VEV's $\langle \tilde{S}^c_\alpha \rangle$ and $\langle \tilde{\nu}^c_\alpha \rangle$ while ν_α and N_3 are, respectively, external. A_2 and B_2 correspond to insertions of the doublets VEV's $\langle \tilde{\nu}_3 \rangle$ and $\langle \tilde{N}_\alpha \rangle$ while S^c_3 and ν^c_3 are external.

$$[m_d]_{ij} = \lambda^{(2)}_{\langle 1\rangle ij} \langle \tilde{N}_1 \rangle d_i d^c_j \ .$$

Then it is clear that N_3 cannot couple to a pair of d-d^c fields that acquire a mass through $\lambda^{(2)}$, otherwise it will necessarily have to transform in the same way as N_1 does under the discrete symmetry, and N_3 will be forced to acquire a large tree-level mass through the same mechanism that makes N_1 heavy. We conclude that the requirement of generating a one-loop mass for N_3 implies that some of the "d" quarks cannot acquire their mass through $\lambda^{(2)}$. Assuming that all the "d" quarks acquire a mass at the tree level implies that we need to flip the assignments for some of the light right-handed $Q=-\frac{1}{3}$ fields as well. More precisely, if $d_i d^c_j$ are the fields coupled to N_3, the (light) SU(2)-doublet field d_i has to acquire its mass through the term $\lambda^{(6)}_{\langle 3\rangle ik} L_3 Q_i h^c_k$ while the (heavy) SU(2)-singlet field d^c_j will acquire its mass through a singlet VEV from the term $\lambda^{(7)}_{\langle \alpha\rangle kj} \nu^c_\alpha h_k d^c_j$. The same argument implies that we cannot flip the assignments for all the "d" quarks. In fact, the need for generating a radiative mass entry for ν_α implies that some $\lambda^{(6)}_{\alpha ij} L_\alpha Q_i h^c_j$ vertices in A_1 must be nonvanishing. At the same time the quark fields entering this vertex cannot acquire a $\lambda^{(6)}\langle \tilde{\nu}_3 \rangle$ mass, otherwise we could not prevent ν_α from acquiring a large mass as ν_3 does.

All these requirements are satisfied, for example, by the assignments

$$\text{"}d^c_\alpha\text{"} \equiv h^c_\alpha \in \mathbf{10}, \quad \alpha = 1,2 \ ,$$
$$\text{"}d^c_3\text{"} \equiv d^c_3 \in \mathbf{16} \ , \tag{3.6}$$

meaning that the massive states corresponding to the known (light) $Q=-\frac{1}{3}$ quarks having SU(2) chiral interactions are $(d_\alpha h^c_\beta)$ and $(d_3 d^c_3)$ with their mass generated, respectively, by $\langle \tilde{\nu}_3 \rangle$ and $\langle \tilde{N}_1 \rangle$, while the heavy vectorlike SU(2) singlets $(h_i d^c_\beta)$ and $(h_i h^c_3)$ acquire a mass through the singlet VEV's $\langle \tilde{\nu}^c_\alpha \rangle$ and $\langle \tilde{S}^c_\alpha \rangle$ ($\alpha=1$ and/or 2), respectively. In order to realize this scenario, we are once again faced with a problem of FCNC that must be confronted with the tight experimental limits derived mainly from analyses of the K and B systems. A large number of Yukawa couplings must be forbidden in order to avoid an excessive tuning for those parameters responsible for the FCNC processes. In order to avoid the Higgs-mediated FCNC in the light "d" quarks sector, which are a direct consequence of the asymmetric assignments among the three families, we must require $\lambda^{(2)}_{\langle i\rangle \alpha 3} = \lambda^{(6)}_{\langle 3\rangle 3\alpha} = 0$ and we must forbid all the couplings between $(d_\alpha h^c_\beta)$ and H_2 as well. Moreover, as in the charged lepton case, the $\Delta I = \frac{1}{2}$ light-heavy mixing between the d_i and h_j states will induce Z_0 mediated FCNC [3,10]. In addition, in the present case a new source of FCNC is represented by the $\Delta I = 0$ light-heavy mixings among h^c_α, d^c_3 and d^c_β, h^c_3. Unlike the $\Delta I = \frac{1}{2}$ case, these mixings are not suppressed by any small doublet-VEV–to–singlet-VEV ratio, and are then expected to be large [11]. However, since they do not violate weak isospin, no FCNC processes can be induced by Z_0 exchange. Nevertheless, these mixings do still affect the Z_β couplings, and could indeed constitute an additional dangerous source of FCNC in the presence of a Z_β with mass below ~ 1 TeV [11]. Both these additional sources of FCNC can be avoided by setting $\lambda^{(2)}_{\langle \alpha\rangle j\beta} = \lambda^{(6)}_{\langle 3\rangle j3} = 0$ and $\lambda^{(4)}_{\langle \alpha\rangle j\beta} = \lambda^{(7)}_{\langle \alpha\rangle j3} = 0$. In particular, we note that the second condition is also needed for the sake of keeping a well-defined meaning to our UA since, in principle, there is no reason to expect that the $\lambda^{(4)}$ and $\lambda^{(7)}$ couplings generating the h_j-h^c_β and h_j-h^c_3 mixing mass terms should be much smaller than the $\lambda^{(6)}$ and $\lambda^{(2)}$ Yukawa couplings responsible, in our scheme, for the "d," "s," and "b" masses.

After all these conditions are implemented, there are no light-heavy mixings in the whole quark sector. The mass matrix for the light down-quarks reads

$$[m_{\mathcal{D}}]_{\alpha\beta} = \lambda^{(6)}_{\langle 3\rangle \alpha\beta} \langle \tilde{\nu}_3 \rangle \ , \tag{3.7}$$

$$[m_{\mathcal{D}}]_{33} = \lambda^{(2)}_{\langle 1\rangle 33} \langle \tilde{N}_1 \rangle \ . \tag{3.8}$$

The remaining SU(2)-singlets $Q=-\frac{1}{3}$ quark states h_i, d^c_β, and h^c_3 are vectorlike and acquire (large) masses through VEV's of singlets:

$$[M_{\mathcal{D}}]_{i\beta} = \lambda^{(7)}_{\langle \alpha\rangle i\beta} \langle \tilde{\nu}^c_\alpha \rangle \ , \tag{3.9}$$

$$[M_{\mathcal{D}}]_{i3} = \lambda^{(4)}_{\langle \alpha\rangle i3} \langle \tilde{S}^c_\alpha \rangle \ . \tag{3.10}$$

From (3.7) and (3.8) we see that our starting assumption about the flipped assignments for the "τ" neutrino has had the far reaching consequence that the down-quark mass matrix is block diagonal. Then all the mixing of the third family can be generated only in the up-quark sector (this might as well suggest a mechanism for explaining the smallness of these mixings relative to the Cabibbo mixing). As a result, in order to have a Cabibbo-Kobayashi-Maskawa (CKM) matrix without zero entries, the up-quark mass matrix must be truly 3×3. We note that the Q_3 doublet cannot transform under the discrete symmetry like the Q_α doublets, otherwise it would not be possible to forbid the d_α-d_3^c mass terms and simultaneously allow for a nonzero d_3-d_3^c mass. At the same time, in order to allow for a 3×3 up-quark mass matrix without zero entries all the u_i^c fields appearing in $\lambda^{(1)}$ have to transform with the same phase. Then, since the bilinears $Q_\alpha u_j^c$ and $Q_3 u_j^c$ have to transform with an overall different discrete phase, in order to construct trilinear invariants they must be coupled to different H^c Higgs fields. As a conclusion, we see that Higgs-mediated FCNC cannot be completely avoided in the UA scenario. However, in the scheme we are analyzing here they appear only in the up-quark sector. Since there are no experimental data on FCNC involving the t quark, the only existing constraints are for the u-c transitions. The strongest bounds on these FCNC come from the limits on D^0-$\bar{D}^0$ oscillations that receive contributions from the $c\bar{u}\to\bar{c}u$ amplitude. Other rare processes, as the rare decays $D\to\pi\pi$, KK that could be induced by the $c\bar{u}\to\bar{u}u$ amplitude, do not give additional constraints.

A bound $\lambda^{(1)}\sim\lambda'_{uc}<2\times10^{-4}$ on the off-diagonal u-c coupling was obtained in Ref. [3] from the limits on D^0-$\bar{D}^0$ oscillations and assuming $M_{H^c}\simeq100$ GeV. Since the $\lambda^{(1)}$ couplings are also responsible for generating the up-quark mass ($\sim$few MeV) from VEV's$\sim$100 GeV, we indeed expect some of the $\lambda^{(1)}$ to be of order 10^{-4} or less. We can conclude that the previous bound does not constitute a serious constraint for the UA scheme, since it does not require a particular tuning of the FCNC parameters. However, a definite prediction of the present model is the existence of an amount of FCNC in the up-quark sector which is larger than in the SM.

Up to this point we have analyzed the requirements which the superpotential (2.4) must satisfy in order to realize the UA scenario and produce an acceptable and possibly interesting phenomenology. We have carried out a general discussion without referring to any particular discrete symmetry. However, since we have required the set of Yukawa couplings in the superpotential (2.4) to satisfy a rather large number of constraints, it could well be that one particular transformation for one field, which is needed in order to forbid a dangerous coupling, at the same time implies the vanishing of another coupling that we want to be nonzero. A general proof that our set of constraints is self-consistent would be lengthy and cumbersome. However, it is enough to prove that there is at least one set of discrete transformations for the fields that satisfies all our constraints, and this will automatically insure that our set of constraints is not self-contradictory.

We have found that our scheme can be implemented by imposing on the superpotential (2.4) a simple $Z_2\times Z_3$ symmetry. The transformation properties of the fields in the three **27** representations are listed in Table II. Beyond satisfying all our requirements it is easily seen that this symmetry also ensures that there is no fast proton decay, since it forbids the terms $\lambda^{(8)}$ and $\lambda^{(9)}$ all together (note that in spite of the UA these still represent the dangerous terms, being both invariant under the exchange $d^c\leftrightarrow h^c$).

In more generality it can be shown that proton stability is just a consequence of additional symmetries which are implied by the discrete symmetry in Table II. In fact, the terms in the superpotential which are invariant under this discrete symmetry are invariant under two global U(1) symmetries as well. The first one acts only on the color-triplet fields for which the global $U_B(1)$ charges are, respectively, $B(Q_i)=B(h_i)=+\frac{1}{3}$ and $B(u_i^c)=B(d_i^c)=B(h_i^c)=-\frac{1}{3}$. $U_B(1)$ can be identified with the baryon number. Under the second global $U_L(1)$ the color-singlet fields transform with $L(L_\alpha)=L(H_3)=+1$, $L(\nu_3^c)=L(S_3^c)=-1$, and $L=0$ for the remaining fields L_3, H_α, H_i^c, ν_α^c, and S_α^c. For the color triplets the L charges are $L(h_i)=+1$, $L(d_\alpha^c)=L(h_3^c)=-1$, and $L(Q_i)=L(d_3^c)=L(h_\alpha^c)=0$. $U_L(1)$ can be identified with lepton number. The h_i, d_α^c, and h_3^c heavy states which carry both baryon and lepton numbers are leptoquarks. B and L conservation in turn imply that R parity is unbroken, and then the model predicts a stable lightest supersymmetric particle (LPS).

From the assignments in Table II it is clear that most of the fields acquiring a VEV transform nontrivially under the $Z_2\times Z_3$ symmetry. This is indeed unavoidable for any discrete symmetry suitable for implementing the scheme which we have been discussing. As a consequence, when the neutral components of the scalar fields acquire a VEV, the discrete symmetry is spontaneously broken. In the early Universe, when a phase transition occurs during the expansion, symmetry breaking takes place independently in different causally disconnected regions that are filled with different discrete phases, and separated from one another by domain walls. In the standard hot universe theory, domain walls cause cosmological problems since they would dominate the energy density of the Universe, as well as astrophysical problems, since they would lead to a considerable anisotropy in the primordial background radiation [12]. Thus cosmological arguments would suggest that we have to renounce

TABLE II. Transformations of the fields in the three **27** representations of E_6 under the discrete $Z_2\times Z_3$ symmetry. The index i ranges from 1 to 3, while $\alpha=1,2$ refers to the first two generations.

Z_2		Z_3		
$+$	$-$	1	$e^{-i(2\pi/3)}$	$e^{i(2\pi/3)}$
$[Q,u^c,H^c]_i$	$[e^c,h]_i$	$[u^c,h]_i$	$[d^c,h^c]_i$	
$[H,h^c,\nu^c,S^c]_\alpha$	$[d^c,L]_\alpha$	$[\nu^c]_2$	$[H^c]_\alpha[H^c,S^c]_2$	$[H,\nu^c,S^c]_1$
$[d^c,L]_3$	$[H,\nu^c,h^c,S^c]_3$	$[Q,L,e^c,H^c]_3$		$[H,\nu^c,S^c]_3$

the kind of models with spontaneously broken discrete symmetry discussed here. However, in an inflationary universe scenario it is possible to get rid of this problem since inflation ensures that each region containing a different phase becomes exponentially large, up to the point that there would not be a single domain wall in the observable part of the Universe. In order for this mechanism to be effective, inflation has to go on long enough after the phase transition, and the reheating temperature of the Universe after inflation should be low enough in order to ensure that the symmetry is not restored when the ordinary adiabatic expansion of the standard cosmology begins. Viable cosmological models in which inflation takes place at the electroweak scale, and that satisfy at these requirements exist and, for example, have been recently discussed in Ref. [13].

IV. PHENOMENOLOGY

We will now concentrate on the pattern of masses and mixings allowed, in our scheme, for the light neutrinos. We will first discuss the neutrino mass matrix in the one-loop approximation, and then we will briefly describe the effects of the additional contributions that arise from two-loop diagrams.

A. One-loop neutrino masses

In the previous section the general form (3.5) of the one-loop neutrino mass matrix was derived. The mass terms for the light and heavy $Q=-\frac{1}{3}$ quarks given in Eqs. (3.7)–(3.10) were also worked out, according to the choice of the assignments in (3.6). Now from Fig. 2, we see that in order to generate at one-loop the ν_α-ν_3^c mass term $[A_1B_2]$ in (3.5), the two mass insertions $[M_{\mathcal{D}}]_{i3}=\lambda^{(4)}_{\langle\alpha\rangle i3}\langle\tilde{S}^c_\alpha\rangle$ and $[m_{\mathcal{D}}]_{33}=\lambda^{(2)}_{\langle 1\rangle 33}\langle\tilde{N}_1\rangle$ are needed. This fixes d_3 and h_3^c as the quarks that couple to the external ν_α, implying that the vertex $\lambda^{(6)}_{\alpha 33}\nu_\alpha Q_3 h_3^c$ must be simultaneously nonvanishing. On the other hand, in order to generate the ν_α-S_3^c mass term $[A_1A_2]$, we need the mass insertion $[M_{\mathcal{D}}]_{i3}$ together with $[m_{\mathcal{D}}]_{\beta\gamma}=\lambda^{(6)}_{\langle 3\rangle\beta\gamma}\langle\tilde{\nu}_3\rangle$. This in turn implies the nonvanishing of the $\lambda^{(6)}_{\alpha\beta 3}\nu_\alpha Q_\beta h_3^c$ vertex. However, the two $\lambda^{(6)}$ vertices cannot be simultaneously nonvanishing. In fact, this would require Q_β to transform like Q_3 under the discrete symmetry, since they both couple to the same bilinear $\nu_\alpha h_3^c$. The result is that to this order only one of the two possible radiative mass terms for the ν_α is allowed. A similar argument implies that also for the N_3 light state only one of the two N_3-ν_3^c and N_3-S_3^c one-loop mass terms is allowed, corresponding to only one of the two $\lambda^{(2)}_{33\beta}N_3Q_3d_\beta^c$ or $\lambda^{(2)}_{3\alpha\beta}N_3Q_\alpha d_\beta^c$ being nonvanishing.

By requiring that our scheme should allow for nonzero ν_α-N_3 mixings, we are left with the two choices

$$\text{(i)}\ \lambda^{(4)}_{\alpha\beta 3}\nu_\alpha Q_\beta h_3^c=\lambda^{(2)}_{3\alpha\beta}N_3Q_\alpha d_\beta^c=0$$

or

$$\text{(ii)}\ \lambda^{(4)}_{\alpha 33}\nu_\alpha Q_3 h_3^c=\lambda^{(2)}_{33\beta}N_3Q_3d_\beta^c=0\ .$$

The transformation properties listed in Table II correspond to the first choice and lead to nonvanishing mass terms with ν_3^c while, at this order, S_3^c does not couple to any doublet neutrino and remains massless. As a result, at one loop the Dirac mass matrix for the light neutrinos acquires the very simple form

$$(\nu_1\nu_2N_3)\mathcal{M}_1\begin{bmatrix}0\\S_3^c\\\nu_3^c\end{bmatrix},\quad \mathcal{M}_1=\begin{bmatrix}0&0&a_1\\0&0&a_2\\0&0&a_3\end{bmatrix}, \tag{4.1}$$

where for convenience we have added one dummy entry in the vector of the right-handed neutrinos. From (4.1) it is apparent that two mass eigenstates n_1 and n_2 will be massless, while the third one n_3 will acquire a Dirac mass

$$\mu_\nu=\sqrt{a_1^2+a_2^2+a_3^2}\ .$$

Clearly, here the SU(2) singlet ν_3^c plays the role of the right-handed component of the massive SU(2)-doublet neutrino. As we will see in the following, the degeneracy between n_1 and n_2 will be effectively removed due to two-loop corrections; however, for the moment we will discuss the results implied by the mass matrix $\mathcal{M}_1$ in the one-loop approximation.

The unitarity transformation that relates the flavor to the mass eigenstates is

$$\begin{bmatrix}\nu_1\\\nu_2\\N_3\end{bmatrix}=\mathcal{R}\begin{bmatrix}n_1\\n_2\\n_3\end{bmatrix}$$

with (4.2)

$$\mathcal{R}=\begin{bmatrix}c_\alpha&0&s_\alpha\\-s_\alpha s_\gamma&c_\gamma&c_\alpha s_\gamma\\-s_\alpha c_\gamma&-s_\gamma&c_\alpha c_\gamma\end{bmatrix},$$

where

$$s_\alpha=a_1/\mu_\nu,\quad c_\alpha s_\gamma=a_2/\mu_\nu,\quad c_\alpha c_\gamma=a_3/\mu_\nu\ .$$

In (4.2) we have made use of the freedom in rotating the degenerate massless states n_1 and n_2 in such a way that n_2 does not couple to the electron (we have implicitly assumed that the rotation of the e and μ fields needed to diagonalize their mass matrix has already been absorbed in the definition of ν_1 and ν_2).

Following Ref. [3], in first approximation the mass entries in (4.1) can be estimated to be

$$a_\alpha\sim\frac{1}{32\pi^2}\lambda^{(6)}_{\alpha 33}\lambda^{(7)}_{3j3}m_b,\quad \alpha=1,2 \tag{4.3}$$

$$a_3\sim\frac{1}{32\pi^2}\lambda^{(2)}_{33\beta}\lambda^{(7)}_{3j3}m_b\ . \tag{4.4}$$

Since all the three entries a_1, a_2, and a_3 in the mass matrix are proportional to the "b"-quark mass (which is the only light quark allowed to appear inside the corresponding loops), in principle, there is no reason to expect any hierarchy among the a_i, and the mixing of the third generation neutrino can then be large. However, since m_d

and m_s are both proportional to $\lambda^{(6)}\langle\tilde{L}_3\rangle$ and at the same time the "τ" mass is given by $\lambda^{(3)}\langle\tilde{L}_3\rangle$, we would expect the $\lambda^{(6)}$ couplings, rather than $\langle\tilde{L}_3\rangle$, to be small. On the other hand, $m_b\sim\lambda^{(2)}\langle\tilde{H}_1\rangle$ so that if we assume $\langle\tilde{L}_3\rangle\sim\langle\tilde{H}_1\rangle$ and that excessively large hierarchies inside each set of $\lambda^{(m)}$ couplings are absent, it is reasonable to expect that $\lambda^{(2)}>\lambda^{(6)}$. In turn, this implies $a_3>a_\alpha$ meaning that the light massive neutral state will be mainly the third generation "ν_τ" neutrino with $\mu_\nu\sim a_3$.

We can now question if μ_ν can be small enough to lie in the range of values required for solving the solar neutrino problem via matter-enhanced "ν_e"-"ν_τ" oscillations. As is well known a MSW solution would require values of μ_ν as small as $\lesssim 10^{-2}$ eV. In order to estimate how small μ_ν can be for natural values of the parameters, we will assume the doublet VEV's to be ~ 100 GeV, so that in (4.4) $\lambda^{(2)}\sim m_d/100\text{ GeV}\simeq 10^{-4}$. Clearly we need a value for $\lambda^{(7)}_{3j3}$ as well, since all the entries in $\mathcal{M}_1$ are proportional to this coupling. The set of $\lambda^{(7)}$ is responsible for generating the masses for the new heavy states h_j-d^c_α through $\langle\tilde{\nu}^c_\alpha\rangle$, thus their order of magnitude is, in principle, unknown. However, if we assume that the new fermions and the new gauge bosons are as light as allowed by the present phenomenological constraints (thus implying that they will be detected with the next generation of colliders) we can still work out an estimate for $\lambda^{(7)}$. According to the present limits from direct searches at colliders, the heavy fermions cannot be much lighter than ~ 100 GeV, so that $\lambda^{(7)}\gtrsim 100\text{ GeV}/\langle\tilde{\nu}^c_\alpha\rangle$. On the other hand, we can argue that the lowest value for this VEV is still expected to be ~ 10 TeV. By confronting the data on the light element abundances with the standard nucleosynthesis calculations, a limit of 3.6 relativistic neutrinos in thermal equilibrium at the time of nucleosynthesis can be derived [14]. This implies that not even one additional neutrino can remain in equilibrium in addition to the three known light states. Though these are singlets under $SU(2)\times U(1)$, the two additional light neutrinos present in our scheme do have $U(1)'$ interactions. Therefore, we have to require this interaction to be weak enough to allow for the decoupling of both the SU(2) singlets ν^c_3 and S^c_3 at a sufficiently early time (for example, before the QCD phase transition) so that their number density can be safely diluted. This argument implies that the mass of the lightest additional gauge boson M_β should be at least of the order of ~ 1–2 TeV [15]. Such a large mass will be mainly generated by the singlet VEV's giving

$$M_\beta\sim g_\beta[(Q^{\nu^c}_\beta\langle\tilde{\nu}^c_\alpha\rangle)^2+(Q^{S^c}_\beta\langle\tilde{S}^c_\alpha\rangle)^2]^{1/2}\ .$$

Since the coupling constant $g_\beta\sim g_1\sim 0.16$ and the Q_β charges $\lesssim 1$, we see that indeed VEV's ~ 10 TeV are required, implying as a result $\lambda^{(7)}\sim 10^{-2}$. We also note that since the only nonvanishing Yukawa coupling for ν^c_3 is precisely $\lambda^{(7)}_{3j3}$ we could have also attempted to estimate it directly by requiring that the exchange of scalar quarks of a typical supersymmetry (SUSY) mass $\sim$few 100 GeV should not be able to keep this particular species in thermal equilibrium. As a result of a rough computation we have found that scalar quark masses below 1 TeV are indeed consistent with $\lambda^{(7)}_{3j3}\sim 10^{-2}$.

Now, according to (4.3) the order of magnitude of the "ν_τ" mass is $\mu_\nu\sim 10$ eV. This value is indeed too large to play any role in the solar (or atmospheric) neutrino problem; however, it is not in conflict with the cosmological limit $\mu_\nu\lesssim 92\Omega h^2$ eV implied by requiring the Universe not to be overclosed (here $\Omega=\rho/\rho_c\sim 1$ is the ratio of the energy density of the Universe to the critical density, and $h\sim 0.4$–1 is the Hubble parameter in units of 100 km s^{-1} Mpc^{-1}. Since this neutrino is effectively stable it could be a natural candidate for the hot component of the dark matter (DM). We note as well that our scheme is also consistent with a certain amount of cold DM, since R parity is unbroken and the LPS is stable.

On the other hand, if $\lambda^{(7)}_{3j3}$ were about two orders of magnitude smaller, then the n_3 mass would fall in the right range of values for a possible explanation of the atmospheric neutrino deficit via "ν_μ"-"ν_τ" oscillations. We will briefly discuss in the following a possible scheme in which the fact that the two flavor states do have different NC interactions could play an interesting role for these oscillations.

B. Two-loop neutrino masses

The previous discussion does not imply that the scenario that we are analyzing cannot offer a solution to the solar neutrino problem. In fact, even if at the one-loop level both "ν_e" and "ν_μ" are massless, nonzero ν_α-S^c_3 entries can be generated at the two-loop level due to the presence of the $\lambda^{(5)}hu^ce^c$ couplings in W_2 (2.4). A typical two-loop diagram is depicted in Fig. 3. We note that since the set of couplings needed to generate this diagram is indeed allowed by the assignments in Table II, the generation of two-loop mass entries is not in conflict with the other constraints on the superpotential discussed in the previous section.

At the next order two additional entries are generated in the neutrino Dirac mass matrix $\mathcal{M}_2$, namely, $[\mathcal{M}_2]_{12}\equiv b_1$ and $[\mathcal{M}_2]_{22}\equiv b_2$. They can be roughly estimated to be [3]

$$b_\alpha\sim\lambda^{(1)}_{ik\delta}\lambda^{(3)}_{2\alpha\beta}\lambda^{(4)}_{3j\gamma}\lambda^{(5)}_{jk\beta}\frac{m_s}{(16\pi^2)^2}\ . \tag{4.5}$$

We note that in the present case the "s" quark ($d_2h^c_2$) is the heaviest one allowed to run inside the loop. Now if we take $\lambda^{(3)}_{21\beta}\sim m_e/100$ GeV, $\lambda^{(3)}_{22\beta}\sim m_\mu/100$ GeV and the representative value $\lambda^{(1)}\sim m_c/100$ GeV, we obtain

$$b_2\sim 10^{-1}\lambda^{(4)}\lambda^{(5)}\text{ eV}\gg b_1$$

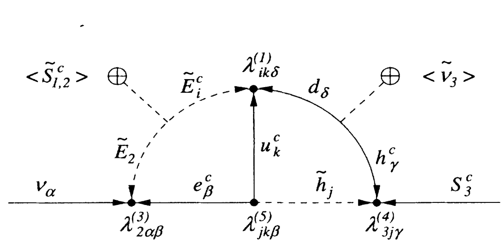

FIG. 3. A two-loop diagram giving rise to ν_α-S^c_3 entries in the neutrino Dirac mass matrix (4.2).

showing that values of the "ν_μ" mass, interesting for a MSW [6] solution of the solar neutrino problem, can indeed be accommodated for natural values of the remaining Yukawas.

To summarize our results, in this model we have a massive n_3 neutrino, mainly "ν_τ", with a mass that can easily fall in the range $\sim 0.1-10$ eV interesting for providing a hot DM candidate or for the oscillations of the "ν_μ" atmospheric neutrinos. A second neutrino n_2, mainly "ν_μ," acquires a much smaller mass at the two-loop level and can be relevant for matter-enhanced "ν_e"-"ν_μ" oscillations in the Sun and finally, due to the absence in our minimal scheme of a helicity partner, n_1 remains massless.

C. A scheme with a light Z_β boson

Before concluding this section we want to illustrate a different scheme in which an attempt for an unconventional solution to the solar neutrino problem can be formulated, and in which unusual effects for the "ν_μ"-"ν_τ" oscillations of the atmospheric neutrinos could arise as well. The following analysis is mainly intended as an example of the unusual phenomenology that can be implied by UA models.

From the previous discussion it should be clear that if we insist on trying to achieve a very small one-loop μ_ν mass, we must require $\langle \tilde{\nu}^c_\alpha \rangle$ to be very large in order to allow for $\lambda^{(7)} \ll 10^{-2}$ while still keeping the vectorlike quarks as heavy as $\gtrsim 100$ GeV. Then let us assume that $\langle \tilde{\nu}^c_\alpha \rangle$ is much larger than all the other VEV's including $\langle \tilde{S}^c_\alpha \rangle$. With this assumption the Z_ψ-Z_χ mixing angle is

$$\tan 2\beta = \frac{2\sum_j Q^j_\psi Q^j_\chi \langle \phi_j \rangle^2}{\sum_j (Q^j_\psi)^2 \langle \phi_j \rangle^2 - \sum_j (Q^j_\chi)^2 \langle \phi_j \rangle^2} \simeq -\frac{\sqrt{15}}{7}, \tag{4.6}$$

where the $\langle \phi_j \rangle$'s represent the various singlet and doublet VEV's occurring in the model. Then the first gauge boson in (2.3)

$$Z'_\beta = -\tfrac{1}{4} Z_\psi + (\sqrt{15}/4) Z_\chi$$

is very heavy. The second one Z''_β, which corresponds to the orthogonal combination of generators, does not couple to the ν^c states. A major consequence of this fact is that, as long as the remaining VEV's which contribute to its mass are not very large, the Z''_β boson will be light. A second consequence is that the gauge interactions cannot keep the right-handed light ν^c_3 neutrino in thermal equilibrium in the early Universe, since the ν^c's are effectively singlets with respect to all the "light" gauge bosons. Furthermore, since now $\lambda^{(7)}_{3j3} \ll 10^{-2}$, even the exchange of scalar quarks as light as ~ 100 GeV cannot help in thermalizing these light states, so that the ν^c_3 degree of freedom will not be populated at the time of nucleosynthesis. Clearly, the other singlet S^c_3 will still be coupled to the light Z''_β since there is no possible choice for the angle β for which both the states are decoupled. However, in contrast with the previous case in which the presence of a light S^c_3 was needed in order to generate two-loop masses, we will now assume that S^c_3 transforms under the discrete symmetry like one of the S^c_α so that it will acquire a large mass. We obtain here a truly "minimal" scheme with only one light singlet neutrino. However, in contrast with a similar scheme first proposed in Ref. [16] there are no Majorana entries in the mass matrix. We stress that in this second scheme the requirement of allowing for a very small "ν_τ" mass $\mu_\nu \ll 1$ eV automatically allows for a new "light" neutral gauge boson as well.

Since the nucleosynthesis constraints on the mass of the lightest additional gauge boson Z''_β are evaded, this boson could be as light as allowed by the present limits from direct searches at colliders [17] and from the analysis of Z' indirect effects [18], resulting in both cases in $M_\beta \gtrsim 200$ GeV. One might object that these limits cannot be straightforwardly applied to the present situation, since they are derived from analyses based on the conventional scheme, while in the present case a large number of fermion couplings are clearly different. However, since we have no reason to expect that in the UA schemes the bounds could be greatly strengthened or relaxed, we will assume that the quoted limit still holds here also.

The presence of a light Z_β is crucial for the following discussion. In fact, as we have already stressed, the UA scheme implies that the "τ" neutrino does not have the same U(1)$'$ interactions with respect to "ν_e" and "ν_μ." It is then interesting to study the implications of having the different neutrino flavors interacting differently with matter through NC. For the present discussion we will restrict ourselves to the two flavor case ν_α ($\alpha=1$ or 2) and N_3.

The propagation of the two neutrino flavor eigenstates through matter is governed by the Schrödinger-like time evolution equation [6]

$$i\frac{d}{dx}\begin{bmatrix}\nu_\alpha \\ N_3\end{bmatrix} = \frac{1}{2E}\mathcal{H}\begin{bmatrix}\nu_\alpha \\ N_3\end{bmatrix}. \tag{4.7}$$

The effective Hamiltonian in (4.7) relevant to the present case is

$$\mathcal{H} = \frac{1}{2}\mathcal{R}_\alpha \begin{bmatrix}-\mu_\nu^2 & 0 \\ 0 & \mu_\nu^2\end{bmatrix}\mathcal{R}^\dagger_\alpha + 2E\begin{bmatrix}\mathcal{A}_{CC}-\Delta_{NC} & 0 \\ 0 & -\mathcal{A}_{CC}+\Delta_{NC}\end{bmatrix}. \tag{4.8}$$

In (4.8) μ_ν is the n_3 mass while $\mathcal{R}_\alpha$ is the relevant 2×2 $\nu_\alpha - N_3$ vacuum mixing matrix. $\mathcal{A}_{CC}$ represents the coherent neutrino forward-scattering due to the charged current (CC) interaction. Then, in the case of the solar electron neutrinos ($\alpha=1$) $\mathcal{A}_{CC}=\sqrt{2}G_F\mathcal{N}_e$ with $\mathcal{N}_e$ the electron density in the Sun, while for the upward-going atmospheric "ν_μ" neutrinos propagating through the earth ($\alpha=2$) $\mathcal{A}_{CC}=0$. Finally, Δ_{NC} represents the difference in the forward scattering between ν_α and N_3 that is due to the difference in their NC interactions. Clearly, because of universality, this term vanishes exactly in the SM. However, in the present case it is nonzero due to the additional contribution from Z''_β exchange. The additional term can be written as

$$\Delta_{\rm NC}=2\sqrt{2}G_F\frac{M_{Z_0}^2}{M_\beta^2}\frac{g_\beta^2}{g_0^2}\mathcal{F}(Q_\beta)\ ,$$
$$\mathcal{F}(Q_\beta)=(Q_\beta^{\nu_\alpha}-Q_\beta^{N_3})(Q_\beta^e+Q_\beta^p+Y_nQ_\beta^n)\mathcal{N}_e\ , \tag{4.9}$$

where g_β^2/g_0^2 in the first equation is the ratio of the squared $U_\beta(1)$ and SU(2) gauge coupling constants. In the expression for $\mathcal{F}(Q_\beta)$, $Q_\beta^f\equiv Q_\beta(f)-Q_\beta(f^c)$ is the vector coupling of the $f=e,p,n$ fermion to the Z_β'' boson. We have taken for the proton density $\mathcal{N}_p=\mathcal{N}_e$ corresponding to electrically neutral matter and, finally, Y_n is the ratio of the neutron to electron density. We note that in contrast to the SM, here the NC forward scattering off electrons and off protons do not cancel. Since both the u and u^c quarks belong to the **10** of SU(5) (see Table I) the Q_β^u vector charge vanishes, and the contribution of the scattering off nucleons is determined only by the "d" quark density. We also note that for the present case, corresponding to $\sin\beta=-\frac{1}{4}$, the light Z_β'' is mainly a Z_ψ boson. Had we chosen for the "d" quarks instead of the assignments (3.6) the alternative assignments "d_α^c"$\equiv d_\alpha^c\in\mathbf{16}$ and "d_3^c"$\equiv h_3^c\in\mathbf{10}$, the Q_ψ^d vector charges of the d quark would have been zero, as is the case for Q_ψ^e, and the $\Delta_{\rm NC}$ term would have been suppressed accordingly.

The presence of interactions implies that the "in matter" mass eigenstates are different from the vacuum mass eigenstates [6]. They can be obtained by diagonalizing $\mathcal{H}$ in (4.8). The two eigenvalues $\pm\delta$ and the matter mixing angle α_m are given, respectively, by

$$\delta^2=\left[\frac{4E}{\mu_\nu^2}(\mathcal{A}_{\rm CC}-\Delta_{\rm NC})-\cos2\alpha\right]^2+\sin^22\alpha\ ,$$
$$\sin^22\alpha_m=\sin^22\alpha/\delta^2\ . \tag{4.10}$$

The second equation in (4.10) shows that if the vacuum mixing angle is close to maximal ($\sin2\alpha\simeq1$), the effect of the additional interactions would be that of *reducing* the mixing in matter by the factor $4E(\mathcal{A}_{\rm CC}-\Delta_{\rm NC})/\mu_\nu^2$, thus suppressing the oscillation of the high-energy neutrinos with respect to the low-energy ones. If, in contrast, $\sin2\alpha$ is small, the mixing in matter will be maximal in the resonance region defined by

$$\frac{\mu_\nu^2}{4E}\cos2\alpha=\mathcal{A}_{\rm CC}-\Delta_{\rm NC}\ . \tag{4.11}$$

We see that as long as the second term on the right-hand side (RHS) is not completely negligible with respect to the first one, the allowed regions in the $(\Delta m^2,\sin^22\theta)$ plane would be different than in the standard case.

Now by confronting (4.9) with the electron-neutrino CC forward scattering, we obtain that the RHS in (4.11) would vanish for

$$\frac{\Delta_{\rm NC}}{\mathcal{A}_{\rm CC}}=2\mathcal{F}(Q_\beta)\frac{M_{Z_0}^2}{M_\beta^2}\frac{g_\beta^2}{g_0^2}=1\ , \tag{4.12}$$

i.e., when the difference between the ν_1 and N_3 NC interactions compensates in full the ν_1 CC interaction. Clearly, in this case we would have found the possibility of a ν_1-N_3 resonant conversion even in the case of practically massless neutrinos ($\mu_\nu\ll10^{-3}$). In the present case, in which the requirement of a large $\langle\nu_\alpha^c\rangle$ and a ν_3^c decoupled from the light Z_β'' selects $\sin\beta=-\frac{1}{4}$ and by taking $Y_n=0.5$ as the maximum value of the neutron density at the center of the Sun, the term $\mathcal{F}(Q_\beta)$ in (4.12) gives an enhancement factor ~2. However, with the normalizations given in Table I, the ratio of the squared $U_\beta(1)$ and SU(2) coupling constants is of the order of the electroweak mixing angle $\sin^2\theta_W\sim\frac{1}{4}$ and then we see that the massless neutrino case is indeed ruled out, since it would require $M_\beta\sim M_{Z_0}$.

In the case of "ν_μ" propagation through matter, the same mechanism could affect the rate of ν_2-N_3 conversion for the upward-going atmospheric muon neutrinos. The Kamiokande II [19] and IMB [20] Collaborations have observed an anomaly in the ratio of muons to electrons events induced by atmospheric neutrinos of energies of a few hundred MeV, and a possible explanation of the effects has been given in terms of $\nu_\mu\to\nu_x$ oscillations where $\nu_x=\nu_e$, ν_τ, or a sterile neutrino (ν_s). In order to explain the data, the ν_μ-ν_x mixing angle is required to be close to maximal ($\sin^22\theta>0.5$; see; e.g., Ref. [21]). However, the IMB [22] and Baksan [23] experiments have observed no reduction for the ν_μ flux of upward-going neutrinos with $E\gtrsim1-2$ GeV, and these data have been used to set stringent limits on the allowed region in the $(\Delta m^2,\sin^22\theta)$ plane [24].

As we have already said, for large $\sin2\theta$ the effect of possible additional interactions with matter would, in general, be that of shifting the matter mixing angle away from maximal, thus *suppressing* the rate of conversion. The equations describing this case would still be (4.10) and (4.11) with the correct values of the electron, proton, and neutron densities in the earth and with $\mathcal{A}_{\rm CC}=0$ for $\nu_x\neq\nu_e$. While negligible at low energy, the effect of the interaction with matter could become particularly relevant for high-energy neutrinos, thus helping to explain the data.

Such a mechanism has been investigated in Ref. [25] for the case of ν_μ oscillating into electron or sterile neutrinos. In particular, for ν_μ-ν_s oscillations the difference in the interaction strength of the two neutrino species is about $\frac{1}{2}$ that of the standard ν_e CC interaction, and the analysis in Ref. [25] shows that in this case matter effects are indeed important. In our case, due to the difference in the $U_\beta(1)$ charges, similar effects could arise for "ν_μ"-"ν_τ" oscillations as well. For the propagation through the Earth, $Y_n\simeq1$ gives $\mathcal{F}(Q_\beta)\sim2.7$, and we can assume, consistent with the present direct [17] and indirect [18] limits, $M_\beta\sim200$ GeV. Then from the RHS of (4.12) we see that the effective strength of the new interaction relative to the standard CC interaction is ~0.27, showing that in this case sizable effects could be present as well.

V. CONCLUSIONS

In conclusion we have examined the possibility of constructing consistent models in which the known fermions of the three different generations do not have the same gauge interactions under possible additional U(1)′ factors. We have carried out our analysis in the frame of the superstring-inspired E_6 models, taking as a guideline for constructing our scheme the requirement of having interesting neutrino phenomenology with naturally small radiatively generated Dirac masses. We have shown that models based in this scheme are indeed viable. They can be realized by imposing a family-nonblind discrete symmetry on the superpotential. We have discussed in some detail a minimal model, in which only two additional light SU(2)-singlet neutrinos are present thus leaving one doublet neutrino massless. Clearly, other models based on the same scenario but with a more rich structure in the neutrino sector can also be constructed.

We have shown that, in our model, values of the neutrino masses in interesting ranges for explaining the solar and atmospheric neutrino anomalies, or possibly for providing a hot component of the DM, can be obtained with a natural choice of the parameters. In addition, since baryon and lepton numbers are both conserved, the proton is effectively stable. Because of the presence of FCNC's in the up-quark sector, a rate for D^0-$\bar{D}^0$ oscillations larger than in the SM, but still consistent with the present limits, is predicted. However, there are no other dangerous sources of FCNC's in the model.

In order to illustrate some unusual consequences of our scheme, we have also investigated a different scenario in which only one SU(2)-singlet neutrino is light, and an additional neutral gauge boson is allowed at energies as low as ~200 GeV. Since the additional neutrino naturally decouples from the new light gauge boson, there is no conflict with the nucleosynthesis constraints on the number of neutrino species. We have shown that the generation-dependent NC interaction mediated by this gauge boson, though probably not relevant in the case of the propagation of the solar electron-neutrinos through matter, could however be of some importance in the case of ν_μ-ν_τ oscillations for the upward-going atmospheric neutrinos.

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