The Quark Mass Spectrum in the Universal Seesaw Model

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Abstract

In the context of a Universal Seesaw model implemented in a left-right symmetric theory, we show that, by allowing the two left-handed doublet Higgs fields to develop different vacuum-expectation-values (VEV's), it is possible to account for the observed structure of the quark mass spectrum without the need of any hierarchy among the Yukawa couplings. In this framework the top-quark mass is expected to be of the order of its present experimental lower bound, $m_t \simeq$ 90 to 100 GeV. Moreover, we find that, while one of the Higgs doublets gets essentially the "standard model" VEV of approximately 250 GeV, the second doublet is expected to have a much smaller VEV, of order 10 GeV. The identification of the large mass scale of the model with the Peccei-Quinn scale fixes the mass of the right-handed gauge bosons in the range 10^7 to 10^{10} GeV, far beyond the reach of present collider experiments. Also all FCNC processes are consequently suppressed, in agreement with the present bounds.

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The observed hierarchical structure of fermion masses does not have a satisfactory explanation in the Standard Model (SM) and in the conventional Grand Unified Theories (GUT's). In the quark sector, for example, which will be the subject of the present paper, the masses range from about 10 MeV for the u- and the d-quarks, to more than 100 GeV for the top-quark. In the context of the SM, one can account for this large mass difference only assuming a corresponding hierarchy in the Yukawa coupling constants, which should then range from $\simeq 10^{-5}$ to $\simeq 1$. Apart from the problem of understanding such a hierarchy, also the need of introducing unnaturally small ($\mathcal{O}(10^{-5})$) dimensionless parameters is quite unpleasant, especially if we compare them with, *e.g.*, the magnitude of the gauge couplings. In particular, since all masses, included those of the gauge bosons, arise from the Higgs-condensation mechanism as a consequence of the vacuum-expectation-value (VEV) $v_L \simeq 250$ GeV developed by the SM Higgs doublet, it would be more natural to expect all masses to be of order 100 GeV. In other words, it is not easy to understand the lightness of the first and second generation fermions with respect to the weak scale. We recall that an analogous problem is the one related to the extreme lightness of the neutrino ($m_{\nu_e} \leq 18$ eV), if compared for example to the electron mass. In this case, as is well known, a satisfactory explanation is furnished by the so-called "seesaw" mechanism [1], where the standard neutrinos get very suppressed masses as a consequence of the introduction of very heavy "singlet" fermions, the right-handed (RH) neutrinos ($\nu_R \sim {\nu^c}_L$). The extension of this idea to all fermions give rise to the so-called "Universal Seesaw" (US) models [2-6], extensively studied in the literature by many authors and with many variants. These models are especially addressed to the study of two particular problems which do not find a satisfactory explanation in the context of the SM, namely the fermion mass hierarchy and the understanding of the family replication which is observed in the fermionic sectors. Although the observation we wish to make in the present paper is rather model-independent, we shall mainly refer to the various versions [3,5] of the particular US model which was proposed by Davidson and Wali in ref.[2].

Here we shall briefly summarize the main features of these models, referring for the notation and for further details to the papers in refs.[2-5]. Essentially the idea which lies behind this type of models is, as anticipated above, the extension of the neutrino-seesaw mechanism to all fermions. The model, which is naturally implemented in the context of a left-right symmetric theory based on the gauge group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$, is obtained by enlarging the fermionic content of the theory. More precisely, for each left- and right-handed fermion ($f_{L(R)}$), is introduced a new heavy $SU(2)_L \otimes SU(2)_R$-singlet ($F_{L(R)}$), which mimics the role of the RH neutrino of the standard (neutrino) seesaw scenario. That is, each ordinary fermion f (included the RH neutrino) has a singlet heavy partner F, with the same electric and color charges. While the fermionic sector is extended in this way, the Higgs sector of the model, on the other hand, is kept as minimal as possible; in fact, at the one-generation level there are only one left- and one right-handed Higgs doublets[1], $\phi_L(2,1)_1$ and $\phi_R(1,2)_1$ (respectively with VEV's v_L and v_R), which couple the ordinary fermions to the heavy singlets. The latter get their large mass through the coupling with an extra odd-parity singlet Higgs field $\sigma(1,1)_0$ which, developing a very large VEV, χ, can also be used to induce the spontaneous breaking of the left-right symmetry à la Chang-Mohapatra-Parida [7]. In the multi-generation models, as we shall see below, in order to produce a mass hierarchy between the various fermion generations, one must introduce a suitable discrete symmetry broken spontaneously by $<\sigma>=\chi$. In the previous papers [2-5], this symmetry, which also plays the role of a global horizontal symmetry, was identified with the axial $U(1)_A$ introduced by Peccei-Quinn (PQ) in order to solve the strong-CP problem [8]. In this case the mass scale χ will be in the range 10^{10} to 10^{12} GeV [9]. Furthermore, as a consequence of the introduction of such an axial global symmetry, we have to double the number of the Higgs-doublets[2], resulting in two left- and two right-handed doublets, $\phi_{L1(2)}(2,1)_1$ and $\phi_{R1(2)}(1,2)_1$, distinguished by their opposite PQ-charge, and whose VEV's are denoted by $v_{L1(2)}$ and $v_{R1(2)}$. The important feature of the model is the absence of the "standard" (*i.e.*, common to all conventional left-right symmetric models [10]) bidoublet Higgs field $\phi(2,2)_0$, resulting in the vanishing, at the lowest order, of Dirac masses for the ordinary fermions. This is the reason for which in this type of models the fermion masses may only arise at higher order, and are therefore naturally suppressed. More explicitly, after the breaking of the gauge symmetry one gets, in a basis (f,F) (where $f_{L(R)}$ and $F_{L(R)}$ denote, respectively, some ordinary fermion and

[1] We adopt the standard notation, where $(i,j)_k$ label respectively $SU(2)_L$, $SU(2)_R$ representations, and k is the quantum number of $U(1)_{B-L}$.

[2] So that both the up- and the down-charge sectors get Yukawa couplings with the Higgs doublets [2,4].

its heavy singlet partner), the following mass matrix of "seesaw-type" [2-5]:

$$\mathcal{M} = \begin{pmatrix} 0 & M_L \\ M_R & M_\sigma \end{pmatrix} , \tag{1}$$

where M_L, M_R, M_σ represent mass terms proportional, respectively, to the VEV's v_{Li}, v_{Ri}, (i=1,2) and χ. The diagonalization of this type of mass matrix leads, in view of the hierarchy among the VEV's, $v_{Li} \ll v_{Ri} \ll \chi$ (which, as shown in ref.[3], corresponds to the minimum of the Higgs potential), to a very heavy singlet (F) of mass M_σ ($\propto \chi$), and a light "ordinary" fermion ($\simeq f$) with mass $\simeq M_L M_R / M_\sigma$ ($\ll v_L$), much smaller than the weak scale v_L without the need of unnaturally small Yukawa couplings. This is the essence of the explanation, in the context of the US models, of the relative lightness of the first two generation fermions. The diagrammatic origin of their mass in the present model is shown in Fig.1 of ref.[4].

However, if in the SM the problem is the understanding of the smallness of some of the fermion masses, in the US models the problem is why, for example the top quark, is so heavy, with a mass unsuppressed with respect to the weak scale. This means that one must find a way of protecting the heavy fermions from getting a seesaw-suppression mechanism. At a naive single-generation level, this protection may be obtained by setting to zero in eq.(1) the mass term M_σ for the corresponding singlet partner. In this case, in fact, the mass matrix in eq.(1) yields a heavy singlet with mass $\simeq M_R$ and an ordinary fermion with an unsuppressed mass of order M_L ($\simeq M_{W_L}$), in agreement with our expectations for m_t. The absence of the mass term M_σ for the singlets may be imposed by introducing a suitable discrete symmetry, which, as shown in refs.[2-5], may be identified with the axial $U(1)_A$ global symmetry of Peccei-Quinn. In a more realistic multi-generation model, this protection from the US mechanism may be obtained by imposing the mass matrix for the heavy singlets to be singular ($\mathrm{Det}(M_\sigma) = 0$), and with a rank equal to the number of generations with masses much smaller than M_{W_L}. This is very important, because it means that in the framework of the US models, it may be possible to explain not only the relative lightness of some of the fermions, but also their hierarchical structure, without the need of assuming a corresponding hierarchy in the Yukawa couplings.

In the previous papers [2,5], however, where for simplicity the VEV's of the two left- and the two right-handed Higgs doublets were assumed to be equal, *i.e*, $v_{L1} = v_{L2}$

and $v_{r1} = v_{R2}$, it was possible to generate a mass hierarchy only between the first two generations and the third one ($m_{1,2} \simeq v_L v_R/\chi\,, \quad m_3 \simeq v_L$), without accounting for the mass difference between the first two. In ref.[5], for example, the mass splitting between the electron and the muon was obtained by fine-tuning the Yukawa couplings. In the present paper, on the other hand, in a model-independent way[3], we shall show that by relaxing the above assumption on the VEV's, and allowing in particular $v_{L1} \neq v_{L2}$, it is possible to explain the full structure of the quark mass spectrum, also obtaining a prediction for the top-quark mass. We recall that the only constraint on the two left-handed (LH) Higgs doublet VEV's is that $({v_{L1}}^2 + {v_{L2}}^2)^{1/2} \equiv v_L \simeq 250$ GeV, in order to reproduce the correct mass for the SM gauge bosons.

Now, let us analyse in detail the structure of the quark mass spectrum. Of course, as usual when one studies the structure of mass matrices, it will be understood throughout the paper that all masses are actually the "running" masses evaluated at a fixed scale, say $\mu = 1$ GeV. The *Review of Particle Data 1992* [11] and the results given by Leutwyler and Gasser in Ref.[12] then suggest the following structure:

$$\begin{gathered}
m_u = 5.6 \pm 1.1\,\mathrm{MeV}\,, \qquad\qquad m_d = 9.9 \pm 1.1\,\mathrm{MeV}\,,\\
m_s = 199 \pm 33\,\mathrm{MeV}\,,\\
m_c = 1.35 \pm 0.05\,\mathrm{GeV}\,, \qquad\qquad m_b = 5.3 \pm 0.1\,\mathrm{GeV}\,,\\
m_t \geq 150\,\mathrm{GeV}\,,
\end{gathered} \tag{2}$$

where we have conveniently put on the same line the quarks characterized approximately by the same mass scale[4], namely the u and d, as well as the c and b. According to eq.(2), all quark masses may essentially be arranged in four distinct classes. The first two (including the u, d, and s quarks) require a seesaw suppression, while the last two (c, b, and t quarks) should get unsuppressed masses. In particular, the top-quark mass being of the order of M_{W_L}, it may be assumed to arise from the contribution of the LH Higgs doublet whose VEV, v_{L2}, is of the order of the SM vacuum-expectation-value, $v_L \simeq 250$

[3] Our analysis is model-independent in the sense that it does not depend, for example, on the particular discrete symmetry chosen to protect the heavy fermions from the seesaw suppression mechanism.

[4] The lower bound of 150 GeV for $m_t(1\text{ GeV})$ corresponds to a physical top-quark mass of about 91 GeV [11,12].

GeV; so that $m_t = Y\, v_{L2} \equiv L_2$, where Y is a "characteristic" Yukawa coupling constant. Then, we may give smaller masses (but still unsuppressed) to c and b by assuming their coupling to the other Higgs doublet, whose VEV, v_{L1}, is taken to be much smaller than v_{L2} (so that $v_{L2} \simeq v_L$); so that $m_{c,b} \simeq Y v_{L1} \equiv L_1$. As a consequence, the ratio of the two LH Higgs doublet VEV's may be fixed by choosing a specific value for the top-quark mass. In particular, if m_t is taken to be of the order of its present lower bound, we get $\tan\beta \equiv v_{L1}/v_{L2} \simeq m_{c,b}/m_t \simeq 1/30$. Now, we notice that this is also approximately equal to the ratio of the mass scales characterizing the first two classes, $m_{u,d}/m_s$. Therefore, we may be able to fit all masses, by saying that the three smaller quark masses are all suppressed, but while u and d couple to the Higgs doublet whose VEV is v_{L1}, the s-quark couples to the other one, with the larger VEV equal to v_{L2}. More explicitly, we shall set $m_{u,d} \simeq Y v_{L1} v_R/\chi \equiv L_1 R/K$, and $m_s \simeq Y v_{L2} v_R/\chi \equiv L_2 R/K$. Since in the present model the fermion mass hierarchy does not arise from large differences of the Yukawa couplings, we are essentially assuming that they are all within one order of magnitude[5]. The above structure of the spectrum can easily be implemented in an *ansatz* for the up- and the down-quark mass matrices, written respectively in a basis (u_i, U_j) and (d_i, D_j), where i,j=1,2,3, are generation indices. Referring to the notation introduced in eq.(1), where now M_L, M_R, and M_σ are of course 3×3 matrices, we have[6]:

$$M_{Lu} = \begin{pmatrix} L_1 & \cdot & \cdot \\ \cdot & L_1 & \cdot \\ \cdot & \cdot & L_2 \end{pmatrix}, \tag{3a}$$

$$M_{\sigma u} = \mathrm{diag}(K, 0, 0), \tag{3b}$$

for the up-quark sector, and:

$$M_{Ld} = \begin{pmatrix} L_1 & \cdot & \cdot \\ \cdot & L_2 & \cdot \\ \cdot & \cdot & L_1 \end{pmatrix}, \tag{4a}$$

[5] Actually, we are assuming that also the Yukawa couplings in the RH and in the singlet sectors are similar to those characterizing the couplings with the LH doublets, *i.e.*, $R \simeq Y v_R$, and $K \simeq Y\chi$.

[6] For simplicity, we assume all mass matrices to be real, disregarding the possibility of spontaneous CP violation.

$$M_{\sigma d} = \mathrm{diag}(K, K, 0)\,, \tag{4b}$$

for the down-type quarks. In both cases the lower-left sub-matrix M_R may be any arbitrary non-singular 3×3 matrix with elements proportional to the VEV of the RH Higgs doublets, $v_{R1} = v_{R2} \equiv v_R$. In the upper-right sub-matrices M_L we have fixed only the diagonal elements, with $L_1 \equiv Y v_{L1}$ and $L_2 \equiv Y v_{L2}$; the off-diagonal *"dots"* standing for (relatively small) unspecified terms responsible for the observed quark mixing. Since here we are mainly interested in understanding the structure of the mass spectrum, leaving aside the problem of explaining the observed structure of the Cabibbo-Kobayashi-Maskawa (CKM) mixing, we shall not try to fix these off-diagonal elements. It is interesting to notice, from eqs.(3b) and (4b), the different rank of the singlet mass matrices M_σ (where, we have set $K \equiv Y\chi$) in the two sectors, corresponding to the seesaw-suppression of two generations for the down-quarks, and of only one generation for the up-quarks.

Since these 6×6 mass matrices, being of the form shown in eq.(1), are non-symmetric and therefore non-hermitian, in order to evaluate the corresponding (real and positive) eigen-masses, one must consider the hermitian matrices $\mathcal{M}_u {\mathcal{M}_u}^T$ and $\mathcal{M}_d {\mathcal{M}_d}^T$, whose eigenvalues give in fact the fermion masses squared. We have explicitly studied the leading behaviour of the invariants ($\Delta_1 \equiv$Trace$, ..., \Delta_6 \equiv$ Determinant) of these matrices, for an arbitrary form of M_R (also including small off-diagonal terms in M_L which generate the mixing among the different generations). In particular, using the hierarchy $K \gg R \gg L_2 \gg L_1$, and also taking into account the changing of the rank of $\mathcal{M}_u$ and $\mathcal{M}_d$ in the limits $v_{L1(2)}/v_R \to 0$ and $v_R/\chi \to 0$, we get the following relations for the eigenvalues $\lambda_{iu(d)}\,(i = 1, \ldots, 6)$ of $\mathcal{M}_u {\mathcal{M}_u}^T$ and $\mathcal{M}_d {\mathcal{M}_d}^T$:

$$\begin{aligned} &\lambda_{5d} + \lambda_{6d} \simeq \Delta_{1d}\,, && \lambda_{5d}\;\lambda_{6d} \simeq \Delta_{2d}\,, && \lambda_{4d} \simeq \Delta_{3d}/\Delta_{2d}\,, \\ &\lambda_{3d} \simeq \Delta_{4d}/\Delta_{3d}\,, && \lambda_{1d} + \lambda_{2d} \simeq \Delta_{5d}/\Delta_{4d}\,, && \lambda_{1d}\;\lambda_{2d} \simeq \Delta_{6d}/\Delta_{4d}\,, \end{aligned} \tag{5a}$$

and:

$$\begin{aligned} &\lambda_{6u} \simeq \Delta_{1u}\,, && \lambda_{4u} + \lambda_{5u} \simeq \Delta_{2u}/\Delta_{1u}\,, && \lambda_{4u}\;\lambda_{5u} \simeq \Delta_{3u}/\Delta_{1u}\,, \\ &\lambda_{2u} + \lambda_{3u} \simeq \Delta_{4u}/\Delta_{3u}\,, && \lambda_{2u}\;\lambda_{3u} \simeq \Delta_{5u}/\Delta_{3u}\,, && \lambda_{1d} \simeq \Delta_{6u}/\Delta_{5u}\,, \end{aligned} \tag{5b}$$

where, in leading order:

$$\begin{aligned} &\Delta_{1d} \simeq 2K^2\,, && \Delta_{2d} \simeq K^4\,, && \Delta_{3d} \simeq 2K^4R^2\,, \\ &\Delta_{4d} \simeq K^4R^2{L_1}^2\,, && \Delta_{5d} \simeq K^2R^4{L_1}^2{L_2}^2\,, && \Delta_{6d} \simeq R^6{L_1}^4{L_2}^2\,, \end{aligned} \tag{6a}$$

and:

$$\begin{aligned} &\Delta_{1u} \simeq K^2\,, && \Delta_{2u} \simeq 2K^2R^2\,, && \Delta_{3u} \simeq K^2R^4\,, \\ &\Delta_{4u} \simeq K^2R^4{L_2}^2\,, && \Delta_{5u} \simeq K^2R^4{L_1}^2{L_2}^2\,, && \Delta_{6u} \simeq R^6{L_1}^4{L_2}^2\,. \end{aligned} \tag{6b}$$

Since the $\lambda_{iu(d)}$'s are, for i=1,2,3, just the masses squared of the three ordinary up- (down-) quarks, from eqs.(5,6) we get the following formulæ:

$$\begin{aligned} &m_d \simeq L_1\frac{R}{K}\,, && m_s \simeq L_2\frac{R}{K}\,, && m_b \simeq L_1\,, \\ &m_u \simeq L_1\frac{R}{K}\,, && m_c \simeq L_1\,, && m_t \simeq L_2\,, \end{aligned} \tag{7}$$

consistent with the observed structure of the mass spectrum, as given in eq.(2). Of course, these formulæ give the correct masses up to small (of order $\mathcal{O}(1)$) numerical factors, which depend on the particular form chosen for M_R and for the (small) off-diagonal elements in M_L. Eqs.(7) are our main results. By using the actual quark masses, they allow us to predict the various parameters of the model. First of all, we see that in our framework the not yet observed top quark is predicted to have a mass (at μ=1 GeV) of order:

$$m_t \simeq m_b\left(\frac{m_s}{m_{u,d}}\right) \simeq 150\ \text{GeV}\,, \tag{8}$$

equivalent to a physical mass of the order of its present lower bound, $\simeq$ 90 to 100 GeV. Then, as we have already shown above, we predict $\tan\beta \equiv v_{L1}/v_{L2} \simeq m_{u,d}/m_s \simeq 1/30$, which corresponds to a VEV for the second LH Higgs doublet of the order $\simeq$ 10 GeV. Moreover, since $v_{L1} \ll v_{L2} \simeq v_L \simeq 250$ GeV, the order of magnitude of the Yukawa coupling for all quarks is expected to be of order $Y \simeq m_t/v_{L2} \simeq 0.6$. Such a magnitude for this dimensionless parameter is certainly much more natural than the value needed in the SM to accomodate the first generation fermions ($Y_{SM} \simeq \mathcal{O}(10^{-5})$). Of course, since in the present model we may only account for the four different mass scales occuring in the quark spectrum according to eq.(2), the observed splitting between m_c and m_b, and between m_u and m_d, may be explained through the differences in their actual Yukawa coupling constants; the important feature is that in this model all Yukawa's are all of the same order, and do not need to be chosen unnaturally small. Furthermore, the three lighter quarks allow us to predict the size of the "seesaw" suppression factor; *e.g.*:

$$\frac{R}{K} \simeq \frac{v_R}{\chi} \simeq \frac{m_d}{m_b} \simeq \frac{1}{600}\, ; \tag{9}$$

(if we had used the mass ratios m_u/m_c and m_t/m_s, we had obtained, respectively, 1/300 and 1/750, all within about a factor of two). In order to estimate the actual scale of the $SU(2)_R \otimes U(1)_{B-L}$ breaking down to $U(1)_Y$, we need to fix the large mass scale $\chi \equiv < \sigma >$. In view of the identification [2-5] of the singlet Higgs field $\sigma(1,1)_0$ with the scalar responsible for breaking the Peccei-Quinn $U(1)_A$ symmetry à la Dine-Fishler-Srednicki (DFS) [14], we may use the astrophysical and cosmological constraints [9] on the axion to restrict χ in the range 10^{10} to 10^{12} GeV. Consequently, from eq.(9) we may conclude that in the present model the mass scale of the RH gauge bosons W_R, Z', is very large, say 10^7 to 10^{10} GeV, far beyond the reach of the present terrestrial experiments. Incidently, this also implies that not only the right-handed current interactions are suppressed, but also all flavour-changing neutral current (FCNC) processes, which as shown in ref.[4] are always proportional to the ratio ${v_L}^4/(v_R \chi)^2$, are well consistent with the present experimental bounds.

In conclusion, we have considered a Universal Seesaw model which, allowing the two LH Higgs doublets to develop different VEV's, may account for the observed structure of the quark mass spectrum as given in eq.(2), without requiring any hierarchy or any unnaturally small value for the Yukawa coupling constants. In this framework the top quark mass is expected to be in the region of its present lower bound. Furthermore, we have been able to predict that, while one of the LH Higgs doublets gets a VEV essentially equal to the Standard Model one ($v_{L2} \simeq v_L \simeq 250$ GeV), the other doublet, responsible for the mass of the b- and the c-quarks (but also for the seesaw-suppressed masses of u and d), has a much smaller VEV of order 10 GeV. In this paper we have not considered explicitly the quark mixing and CP violation. A more detailed model, also addressed to the understanding of the structure of the CKM mixing matrix is in progress, and will be presented elsewhere [14].

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