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Phenomenology of the two Higgs doublet sector of a quark-lepton symmetric model

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In the simplest examples of models with a discrete quark-lepton symmetry, an electroweak symmetry-breaking sector with more than one Higgs doublet is necessary to obtain the correct mass relations between quarks and leptons. A two Higgs doublet model version has flavor-nonconserving Yukawa couplings, which are proportional to the masses of the quark-lepton symmetric partners of the fermions. We describe how flavor-changing leptonic decays can occur, with branching ratios not far beyond those currently measurable, enabling investigation of the phenomenology of such models.

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I. INTRODUCTION

The quark-lepton symmetric extension of the standard model involves the addition of leptonic color to the symmetry group, in order that the Lagrangian can be made to exhibit a discrete symmetry between quarks and leptons [1]. The symmetry group then becomes $G_{q\ell} \equiv \mathrm{SU}(3)_\ell \times \mathrm{SU}(3)_q \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_X$. Symmetry breaking occurs in two stages: first the leptonic color group is spontaneously broken, at a scale as low as about a few TeV, then the electroweak symmetry is broken. Electroweak symmetry breaking can lead to unsuitable fermion mass relations unless (i) more than one Higgs doublet is used, (ii) leptonic color is completely broken, or (iii) a left-right symmetry is incorporated (see Refs. [1(l)] and [1(i)], respectively].

This paper extends the analysis of Ref. [2] concerning an electroweak symmetry-breaking sector of two Higgs doublets, in particular model 1 of that paper which leads to the most interesting phenomenology. In this model, flavor-changing Yukawa couplings are proportional to the masses of the quark-lepton symmetric partners of the fermions, so leptonic processes will be of most interest, since they will depend on the heavier quark masses. Tree-level flavor-changing processes in both the charged and neutral sectors can occur, unlike in usual two Higgs doublet models (2HDM's) in which these interactions are deliberately suppressed [3].

Primarily processes involving muons will be investigated in Sec. III, as these give more opportunity for experimental investigation. First the neutral-Higgs-boson mediated decay $\mu^- \to e^- e^+ e^-$ will be investigated at the tree level, as will the charged-Higgs boson contribution to the decay $\mu^- \to \nu_\mu e^- \bar{\nu}_e$. Following that, a one-loop calculation of the process $\mu \to e\gamma$ will be performed. This will lead on to an examination of the effect on the muon anomalous magnetic moment, and the decay $b \to s\gamma$.

II. YUKAWA COUPLINGS IN A QUARK-LEPTON SYMMETRIC MODEL

Electroweak symmetry breaking in the two Higgs doublet version of the quark-lepton symmetric model gives rise to the same physical fields as in normal 2HDM's [4]. Thus there is a charged field $H^\pm$, a CP-odd neutral field η, and two CP-even neutral fields h_1 and h_2. There is also a neutral physical field left from the Higgs field used to break leptonic color and the discrete symmetry, but, following Ref. [2], we will neglect mixing between this field and the other neutral bosons, because the scale of quark-lepton symmetry breaking is expected to be much greater than the masses of these electroweak fields.

The Yukawa Lagrangians in Ref. [2] are written in terms of the gauge eigenstate fermion fields, e.g., u for the charge 2/3 quarks. The corresponding mass-eigenstate fields are defined as $U_{L,R} \equiv V^u_{L,R} u_{L,R}$, by introducing the unitary diagonalization matrices $V^u_{L,R}$. The diagonal mass matrix is then $M_u = V^u_L m_u V^{u\dagger}_R$. Analogous relations of course apply to the other kinds of fermions, d, e, and ν. (See Ref. [2] for a detailed introduction to the whole model.)

So, by transforming to the mass-eigenstate fields [and using the notation $V^{fg}_{L,R} \equiv V^f_{L,R} V^{g\dagger}_{L,R}$, so that for instance the Cabibbo-Kobayashi-Maskawa (CKM) quark mixing matrix is V^{ud}_L], we obtain, for the flavor-changing interactions,

$$\mathcal{L}^+_{\mathrm{Yuk}} = \frac{1}{u}\bar{U}\left(V^{u\nu}_L M_\nu V^{d\nu\dagger}_R \gamma_+ + V^{ue}_R M_e V^{de\dagger}_L \gamma_-\right) D H^+ + \frac{1}{u}\bar{N}\left(V^{u\nu\dagger}_L M_u V^{ue}_R \gamma_+ + V^{d\nu\dagger}_R M_d V^{de}_L \gamma_-\right) E H^+ + \mathrm{H.c.}, \tag{1}$$

$$\begin{aligned}\mathcal{L}^\eta_{\mathrm{Yuk}} = {} & \frac{i}{\sqrt{2}u}\bar{U}\left(V^{ue}_L M_e V^{ue\dagger}_R \gamma_+\right) U\eta + \frac{i}{\sqrt{2}u}\bar{E}\left(V^{ue\dagger}_L M_u V^{ue}_R \gamma_+\right) E\eta \\ & + \frac{i}{\sqrt{2}u}\bar{D}\left(V^{d\nu}_L M_\nu V^{d\nu\dagger}_R \gamma_+\right) D\eta + \frac{i}{\sqrt{2}u}\bar{N}\left(V^{d\nu\dagger}_L M_d V^{d\nu}_R \gamma_+\right) N\eta + \mathrm{H.c.},\end{aligned} \tag{2}$$

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$$\mathcal{L}^h_{\rm Yuk} = \frac{1}{\sqrt{2}u}\bar{U}\Big(V_L^{ue}M_eV_R^{ue\dagger}\gamma_+\Big)U\Big[-\sin(\omega-\varphi)h_1+\cos(\omega-\varphi)h_2\Big]$$
$$+\frac{1}{\sqrt{2}u}\bar{E}\Big(V_L^{ue\dagger}M_uV_R^{ue}\gamma_+\Big)E\Big[\sin(\omega-\varphi)h_1-\cos(\omega-\varphi)h_2\Big]$$
$$+\frac{1}{\sqrt{2}u}\bar{D}\Big(V_L^{d\nu}M_\nu V_R^{d\nu\dagger}\gamma_+\Big)D\Big[\sin(\omega-\varphi)h_1-\cos(\omega-\varphi)h_2\Big]$$
$$+\frac{1}{\sqrt{2}u}\bar{N}\Big(V_L^{d\nu\dagger}M_dV_R^{d\nu}\gamma_+\Big)N\Big[-\sin(\omega-\varphi)h_1+\cos(\omega-\varphi)h_2\Big]+\text{H.c.}, \quad (3)$$

where $u \equiv \sqrt{u_1^2+u_2^2}$, with u_1 and u_2 being the VEV's of the two electroweak doublets. Also in the above, $\tan\omega \equiv u_2/u_1$, and φ is a mixing angle relating the mass-eigenstate fields h_1 and h_2 to the CP-even parts of the original fields.

As discussed in Ref. [2], the down-quark sector neutral flavor-changing vertices are proportional to neutrino Dirac masses, and so are highly suppressed. However, processes in the charged Higgs sector are proportional to the masses of the charged leptons, and as such are not as suppressed.

In the rest of this paper, the following notation will be used:

$$(u_1,u_2,u_3)\equiv(u,c,t);\ (d_1,d_2,d_3)\equiv(d,s,b),$$
$$(e_1,e_2,e_3)\equiv(e,\mu,\tau);\ (\nu_1,\nu_2,\nu_3)\equiv(\nu_e,\nu_\mu,\nu_\tau),$$
$$\frac{1}{m_\pm^2}\equiv\frac{\sin^2(\omega-\varphi)}{m_{h_1}^2}+\frac{\cos^2(\omega-\varphi)}{m_{h_2}^2}\pm\frac{1}{m_\eta^2}, \quad (4)$$

$$y_f\equiv\frac{m_f^2}{m_H^2};\ z_f\equiv\frac{m_f^2}{m_\eta^2};\ w_f^{(1,2)}\equiv\frac{m_f^2}{m_{h_{1,2}}^2};\ w_f^\pm\equiv\frac{m_f^2}{m_\pm^2}. \quad (5)$$

III. TREE-LEVEL AND ONE-LOOP PROCESSES

A. $\mu^-\to e^-e^+e^-$

At the tree level, the flavor-changing decay $\mu^-\to e^-e^+e^-$ can be mediated by the neutral Higgs particles, η, h_1, and h_2. The calculation of the branching ratio involves many unknown parameters in the Higgs boson masses and the mixing matrices. In order to get an indication of how large $B(\mu\to e\bar{e}e)$ might be, we will suppose that all of the mixing matrices display a hierarchical structure similar to that of the CKM matrix, which can be written in the qualitative form

$$V\sim\begin{pmatrix}1&\epsilon&\epsilon^3\\ \epsilon&1&\epsilon^2\\ \epsilon^3&\epsilon^2&1\end{pmatrix}, \quad (6)$$

where, for the CKM matrix, for instance, $\epsilon\sim 0.22$. In using this substitution for the mixing matrix of course, some detail is lost, such as features of the unitarity of the matrices, and the possibilites for accidental cancellation.

The branching ratio is then calculated to be

$$B(\mu\to e\bar{e}e)\sim\frac{3}{4}m_c^2\epsilon^2(m_u+m_c\epsilon^2)^2$$
$$\times\Big[\frac{1}{2m_-^4}+\frac{1}{3m_+^4}+\frac{m_e}{m_\mu}\frac{4}{m_-^2m_+^2}\Big]. \quad (7)$$

There are three extreme possibilities for the relations between the masses of the various neutral scalars, each of which leads to a slightly different relation between the mass and the ϵ parameter: if $m_\eta\approx m_{h_1}\approx m_{h_2}$ then $m_-^2\gg m_+^2\equiv\frac{1}{2}m_\phi^2$; if $m_\eta^2\ll m_{h_1}^2,m_{h_2}^2$ then $m_+^2\approx -m_-^2\approx m_\eta^2$; and if $m_\eta^2\gg m_{h_1}^2,m_{h_2}^2$ then $m_-^2\approx m_+^2\equiv m_h^2$.

The dependence of the branching ratio on ϵ for various choices of the Higgs boson masses is shown for these three cases in Fig. 1. The minimum value of the mass used here is 50 GeV, but for a more accurate lower limit those obtained for the standard model [5] give a good indication, as they use Higgs couplings to the Z bosons, which are quite similar in our model. Only two lines for each ϵ value are visible, because the plots for m_η and m_h are inseparable on the scale used. It can be seen from this that an increase in the precision of the measurement of the branching ratio by only a couple of orders of magnitude opens up for investigation vast new regions of the parameter space. Indeed, a considerable range is already excluded, since for example for $\epsilon=0.22$, corresponding to the CKM matrix, the Higgs boson mass has to be greater than about 330 GeV.

Similar calculations can also be performed for the tauon as well, in such processes as $\tau^-\to\mu^-\mu^+\mu^-$ and $\tau^-\to e^-e^+e^-$. The heavy top quark mass contributes to these calculations, and using Higgs boson masses of the

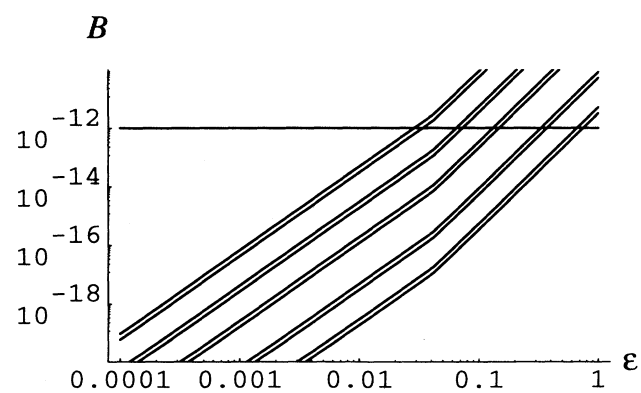

FIG. 1. Branching ratio of $\mu\to e\bar{e}e$ against the mixing matrix parameter ϵ, for a range of masses of the neutral Higgs particles. The region $\epsilon\sim 1$ is to be disregarded, however, because there our approximations fail. The upper line in each pair corresponds to the mass m_ϕ, the lower line to m_η and m_h, and the masses are, from left, 50 GeV, 100 GeV, 200 GeV, 500 GeV, and 1 TeV. The experimentally obtained upper limit on the branching ratio of 1.0×10^{-12} is indicated by the horizontal line.

order $\sim$ 100 GeV, and for the ϵ parameter $\epsilon \sim 0.1$, the branching ratios are of the respective orders 10^{-8} and 10^{-14}. Experimentally, these branching ratios are less than 1.7×10^{-5} and 2.7×10^{-5}, respectively; therefore, these processes do not constrain our model to nearly the same extent as does $\mu^- \to e^- e^+ e^-$.

B. $\mu^- \to \nu_\mu e^- \bar{\nu}_e$

At the tree level in the standard model, the decay $\mu^- \to \nu_\mu e^- \bar{\nu}_e$ is mediated by the W boson. In our model it can also be mediated by a charged scalar $H^\pm$. This can give a constraint on the various parameters (i.e., mass and mixing matrices), since the experimentally measured muon decay agrees so well with standard model predictions. That is, the charged-Higgs-boson contribution must be of the order of the uncertainty in the decay rate, $\Gamma(\mu \to \nu_\mu e \bar{\nu}_e) = (2.995\,92 \pm 0.000\,05) \times 10^{-16}$MeV [5].

Again with the approximation $m_e^2 \ll m_\mu^2$, and using hierarchical mixing matrices as in Eq. (6), the total decay rate, including both W and Higgs boson contributions is

$$\begin{aligned}\Gamma &= \Gamma_W + \Gamma_H + \Gamma_{HW} \\ &\sim \Gamma_W \Big[1 + \frac{1}{4}\epsilon^2 (y_c + y_t \epsilon^4)^2 \\ &\quad + 2\frac{m_e}{m_\mu}\frac{1}{m_H^2}(m_u + 2m_c\epsilon^2 + 3m_t\epsilon^6)(m_c + 2m_t\epsilon^4)\Big], \end{aligned} \tag{8}$$

where

$$\Gamma_W = \frac{G_F^2 m_\mu^5}{192\pi^2}. \tag{9}$$

In the calculation of Eq. (8), the matrix $V_L^{e\nu}$, essentially the leptonic equivalent of the CKM matrix, had to be used in the W boson couplings.

The magnitude of the branching ratio drops off very quickly as the Higgs boson mass is increased and the mixing parameter ϵ is decreased, but for values of $m_H =$ 45.3 GeV and $\epsilon = 0.22$, the total contribution from the terms Γ_H and Γ_{HW} is $(\Gamma_H + \Gamma_{HW})/\Gamma_W \approx 1.1 \times 10^{-5}$, which is only just less than the experimental uncertainty, $\Delta\Gamma/\Gamma_W \approx 1.8 \times 10^{-5}$.

At this point it should be noted that the minimum mass for the charged Higgs boson used in this paper is 45.3 GeV, as obtained by the ALEPH Collaboration [6] with the assumption that the branching ratio $B(H^+ \to \tau^+\nu) = 1$. In this particular 2HDM, with its quark-lepton symmetry, leptonic decays are proportional to the square of quark masses (in particular the top quark mass for $\tau^+\nu$), whereas hadronic decays are proportional to the squares of lepton masses. Thus here we can take $B(H^+ \to \tau^+\nu) = 1$, and so use the limit $m_H > 45.3$ GeV.

C. $\mu \to e\gamma$

One-loop processes such as $\mu \to e\gamma$ and $b \to s\gamma$ have been extensively studied, both in terms of the standard model [7] and 2HDM's [8]. We have calculated the amplitude of $\mu \to e\gamma$ exactly to one-loop, but only specific interesting limiting cases will be presented here. If the neutral scalars are all much heavier than the charged scalars, then the branching ratio can be written [assuming the mixing matrices follow the hierarchical structure of Eq. (6), neutrino masses are negligible, and that $m_e^2 \ll m_\mu^2$],

$$B(\mu \to e\gamma) \sim \frac{\alpha}{96\pi}\frac{1}{m_H^4}\left| m_c^2 V_{R21}^{ue*} V_{R22}^{ue} + m_t^2 V_{R31}^{ue*} V_{R32}^{ue}\right|^2. \tag{10}$$

Alternatively, the neutral bosons could be much lighter; then we have the same three extremes of neutral boson masses as used in the discussion of $B(\mu \to e\bar{e}e)$, represented by m_ϕ, m_η, and m_h:

$$B(\mu \to e\gamma) \sim \frac{\alpha}{24\pi}\frac{1}{m_\phi^4}\left[\left| m_c^2 V_{L21}^{ue*} V_{L22}^{ue} + m_t^2 V_{L31}^{ue*} V_{L32}^{ue}\right|^2 + \left| L \leftrightarrow R\right|^2\right], \tag{11}$$

$$\begin{aligned} B(\mu \to e\gamma) \sim \frac{\alpha}{96\pi}\frac{1}{m_\eta^4}\Big[\Big| & m_c^2 (V_{L21}^{ue*} V_{L22}^{ue} - 9 V_{L21}^{ue*} V_{R22}^{ue} V_{L22}^{ue*} V_{R22}^{ue}) \\ & + m_t^2 (V_{L31}^{ue*} V_{L32}^{ue} - 9\frac{m_\tau}{m_\mu} V_{L31}^{ue*} V_{R33}^{ue} V_{L33}^{ue*} V_{R32}^{ue})\Big|^2 + \left| L \leftrightarrow R\right|^2\Big], \end{aligned} \tag{12}$$

$$\begin{aligned} B(\mu \to e\gamma) \sim \frac{\alpha}{96\pi}\frac{1}{m_h^4}\Big[\Big| & m_c^2 (V_{L21}^{ue*} V_{L22}^{ue} + 9 V_{L21}^{ue*} V_{R22}^{ue} V_{L22}^{ue*} V_{R22}^{ue}) \\ & + m_t^2 (V_{L31}^{ue*} V_{L32}^{ue} + 9\frac{m_\tau}{m_\mu} V_{L31}^{ue*} V_{R33}^{ue} V_{L33}^{ue*} V_{R32}^{ue})\Big|^2 + \left| L \leftrightarrow R\right|^2\Big]. \end{aligned} \tag{13}$$

Experimentally, $B(\mu \to e\gamma) < 4.9 \times 10^{-11}$ [5]. Parametrizing as per Eq. (6), the branching ratio has been plotted in Fig. 2 as a function of ϵ for various values of the Higgs boson mass for the case in which the charged Higgs boson contribution is dominant. The corresponding plots for the other cases are similar. In all these plots the experimental limit is only a couple of orders of magnitude greater than the branching ratios corresponding to that region of parameter space of greatest interest. As for current limits, an assumed value of $\epsilon = 0.22$ gives the bound $m_H > 92$ GeV on the mass of the charged Higgs boson.

Again, these calculations can be performed for the tauon as well as the muon. Once more using the ϵ parametrization, we obtain the branching ratios

$$B(\tau \to \mu\gamma) \sim B_\tau \frac{\alpha}{96\pi} m_t^4 \epsilon^4 \Big[\Big(\frac{m_\mu}{m_\tau}\frac{1}{m_H^2} - \frac{1}{m_+^2} - \frac{m_\mu}{m_\tau}\frac{1}{m_+^2} + 9\frac{1}{m_-^2}\Big)^2 + \Big(\frac{1}{m_H^2} - \frac{1}{m_+^2} - \frac{m_\mu}{m_\tau}\frac{1}{m_+^2} + 9\frac{1}{m_-^2}\Big)^2\Big], \tag{14}$$

$$B(\tau \to e\gamma) \sim B_\tau \frac{\alpha}{96\pi} m_t^4 \epsilon^6 \left[\left(\frac{1}{m_+^2} - \frac{9}{m_-^2} \right)^2 + \left(\frac{1}{m_H^2} - \frac{1}{m_+^2} + \frac{9}{m_-^2} \right)^2 \right]. \quad (15)$$

These branching ratios are $\sim 10^{-8}$ and $\sim 10^{-10}$, respectively (with masses ~ 100 GeV and $\epsilon \sim 0.1$), compared to the experimental upper bounds 5.5×10^{-4} and 2.0×10^{-4}.

D. Other one-loop processes

The calculation of the anomalous magnetic moment of the muon to one loop is very similar to the above calculation of $\mu \to e\gamma$. The total contribution of the various Higgs particles is

$$\Delta a_\mu \approx \frac{G_F m_\mu^2}{24\sqrt{2}\pi} \left[-(y_c + y_t \epsilon^4)^2 + (w_c^+ + w_t^+ \epsilon^4) - 9 \left(w_c^- + \frac{m_\tau}{m_\mu} w_t^- \epsilon^4 \right) \right]. \quad (16)$$

If typical values of the masses and ϵ are substituted into Eq. (16), the anomalous magnetic moment is found to be less than $\sim 10^{-11}$. Experimentally it is known [5] that $-13 \times 10^{-9} < \Delta a_\mu < 21 \times 10^{-9}$, so this constraint is easily satisfied.

The nonleptonic decay $b \to s\gamma$ is currently of interest due to the recent detection of this process by the CLEO Collaboration [9], and is as yet not in disagreement with standard model predictions. Unlike other 2HDM's which can give a substantial contribution to this process, [8], the Yukawa couplings of the quarks are proportional to the lepton masses, and the resulting contribution is at most proportional to m_τ^2, whereas the standard model part is proportional to m_t^2. Thus the quark-lepton symmetric contribution is greatly suppressed, and should give no detectable contribution to the decay.

The same restraints also apply to the decays $B_s \to \ell\bar{\ell}$ [10] at the one-loop level. At the tree level, mediated by a neutral Higgs boson, they are proportional to neutrino Dirac masses, and as such are even more highly suppressed.

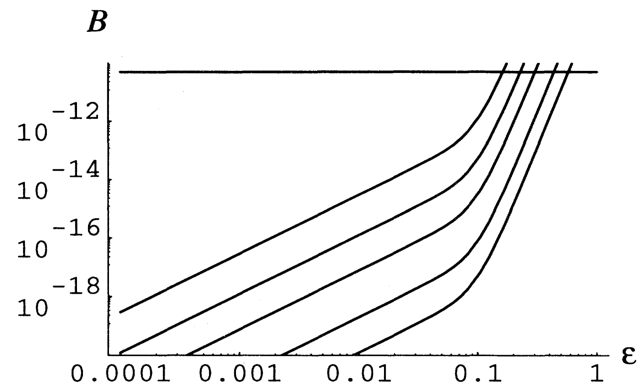

FIG. 2. Branching ratio $B(\mu \to e\gamma)$ against ϵ for various values of the mass of the charged Higgs boson H^+: from left, 45.3 GeV, 100 GeV, 200 GeV, 500 GeV, and 1 TeV. The horizontal line represents the current experimental upper bound of 4.9×10^{-11}. Again, our approximations fail in the region $\epsilon \sim 1$.

IV. CONCLUSION

If the two Higgs doublet version of the quark-lepton symmetric model is correct, then as we have seen, its contribution to various leptonic decays should be measurable for a wide range of parameters, with suitable improvements in experimental search capacity. As, however, the Higgs boson masses and the mixing between lepton families are unknown, it is possible that extreme values (i.e., very diagonal mixing matrices, or very heavy scalars), could lead to extremely small branching ratios. Measurement of the above decays is at least one way to obtain information about these parameters.

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