CP Violation in Seesaw Models of Quark Masses

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CP phenomenology in "seesaw" models of quark masses is shown to parallel that of the usual left-right-symmetric models with the additional advantage that it provides a natural solution to the strong *CP* problem. For the case where the third-generation mixing parameter V_{ub} is extremely small, the neutral-Higgs-boson interactions lead to $\epsilon'/\epsilon \simeq 10^{-3}$ and the electric dipole moment of the neutron $d_n \simeq 10^{-25}\, e\,\mathrm{cm}$. Smallness of the neutrino masses is understood as a two-loop effect.

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Understanding the smallness of the fermion masses compared to the scale of electroweak symmetry breaking (i.e., $m_{q,l} \ll m_{W_L}/g$) is one of the major puzzles of the standard model. If the neutrinos have mass, their small masses add another dimension to this puzzle. For Majorana neutrinos the "seesaw" mechanism[1] provides a way to understand the smallness of m_ν in terms of a new scale corresponding to the breaking of local[1] or global $B-L$ symmetry.[2] This mechanism, however, does not shed any light on the first question, i.e., why $m_{q,l} \ll m_W$? A new approach has recently been advocated[3-6] which addresses both these questions using the seesaw mechanism for the charged-fermion masses, rather than for the neutrinos. As a result, the Yukawa couplings needed to understand quark-lepton masses need not be smaller than 10^{-2}, alleviating the naturalness problem somewhat. The (Dirac) neutrino masses in this case arise[3,6] at the two-loop level, making them automatically small. In this Letter, we study the *CP*-violating consequences of this model. The result of our investigation is that, even though the model differs considerably in its structure of fermion masses from previous formulations of left-right models,[7] *CP*-violating phenomenology when the third generation is nearly decoupled closely parallels that of the earlier models with the additional advantage that the present model provides a natural solution to the strong *CP* problem. The effective $\bar\theta$ in this model arises only at the two-loop level. Furthermore, we find that the model leads naturally to an electric dipole moment of the neutron $d_n \simeq 10^{-25}\, e\,\mathrm{cm}$ and longitudinal muon polarization in $K_L^0 \to \mu\bar\mu$ at the level of 10^{-3}–10^{-2}. The main reason for the d_n to be considerably enhanced in our model over the standard model is that in our case it arises at the one-loop level from the existence of right-handed charged currents and Higgs-boson interactions whereas in the standard model d_n arises at higher loop level.

While we adopt a strategy similar to that of Refs. 4 and 5, there is an important difference in our approach: In Refs. 4 and 5, the neutrino mass $m_\nu \sim h^2 m_{W_L}^2/M$, where h is a typical Yukawa coupling and M is the scale of singlet masses. For $h \sim 1$–10^{-2}, we see that $M \geq 10^{12}$–10^8 GeV and since a typical charged-fermion mass $m_f \approx h^2 m_{W_L} m_{W_R}/M$, this implies that m_{W_R} is also in the 10^8–10^{12}-GeV range. Thus, effects such as *CP* violation associated with the new physics are negligible at low energies. In our case, we do not include a heavy neutral singlet lepton; as a result m_ν arises at the two-loop level,[3,6] and the singlet fermion as well as the W_R scale can be in the TeV range, making the new physics accessible to experiments in the near future. This is the crucial difference between our work and that of Refs. 4 and 5. In fact, if we envision $SU(5)_L \otimes SU(5)_R$ as the grand-unification group as in Ref. 4, the heavy singlet lepton is naturally absent.

The electroweak gauge group of the model is $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ with the known quarks and leptons assigned to the gauge group representations as follows (suppressing generation index i,j,k):

$$Q_L \equiv \begin{pmatrix} u \\ d \end{pmatrix}_L (2,1,\tfrac{1}{3}), \quad Q_R \equiv \begin{pmatrix} u \\ d \end{pmatrix}_R (1,2,\tfrac{1}{3}),$$
$$\Psi_L \equiv \begin{pmatrix} \nu \\ e^- \end{pmatrix}_L (2,1,-1), \quad \Psi_R \equiv \begin{pmatrix} \nu \\ e^- \end{pmatrix}_R (1,2,-1). \tag{1}$$

We introduce one set of heavy isosinglet fermions for each generation transforming as follows: singlet quarks $P(1,1,\tfrac{4}{3})$, $N(1,1,-\tfrac{2}{3})$; singlet lepton $E(1,1,-2)$. The Higgs sector of the model consists only of the following: $\chi_L \equiv (2,1,1)$; $\chi_R \equiv (1,2,1)$; $\sigma \equiv (1,1,0)$. As in the paper of Mohapatra,[3] we assume σ to be odd under parity transformation. It is taken to be charge-conjugation even, so that it is odd under *CP*.

The Yukawa coupling between the fermions and Higgs bosons are given by

$$\mathcal{L}_Y = h^u_{ab}\bar{Q}_{L,a}\tilde{\chi}_L P_{R,b} + h^d_{ab}\bar{Q}_{L,a}\chi_L N_{R,b} + h^e_{ab}\bar{\Psi}_{L,a}\chi_L E_{R,b} + (L \leftrightarrow R) + \mathrm{H.c.}$$
$$+ \bar{P}_a(m^u_{ab} + i f^u_{ab}\sigma\gamma_5)P_b + \bar{N}_a(m^d_{ab} + i f^d_{ab}\sigma\gamma_5)N_b + \bar{E}_a(m^e_{ab} + i f^e_{ab}\sigma\gamma_5)E_b. \tag{2}$$

Here, m^i and f^i are Hermitian matrices, whereas h^i's are arbitrary complex matrices. To get the correct symmetry-breaking pattern, it was shown in Ref. 3 that $\langle\sigma\rangle = \mu \neq 0$ leading to complex masses for P, N, and E. The gauge-symmetry breaking is achieved by the following vacuum expectation values for $\chi_{L,R}$ doublets which minimize the potential: $\langle\chi_L^0\rangle=\kappa_L$, $\langle\chi_R^0\rangle=\kappa_R$. In this model, the W_L-W_R mixing vanishes at the tree level and arises only at the one-loop level and is given by[3]

$$\zeta_{W_L\text{-}W_R}\simeq\frac{\alpha}{4\pi\sin^2\theta_W}\left(\frac{m_t m_b}{m_{W_R}^2}\right)\le 10^{-6}.$$

In the unitary gauge, this model has only three neutral Higgs bosons, σ_L, σ_R, and σ. They, however, mix at the tree level with σ_L-σ_R mixing of order $\lambda(\kappa_L/\kappa_R)$.

Before turning to discuss CP violation, let us briefly review the situation with neutrino mass in the model. It is clear from Eq. (2) that subsequent to symmetry breaking $m_\nu=0$. Nonzero contribution to the neutrino mass arises at the two-loop level involving W_L and W_R exchanges, leading to[6]

$$m_\nu\simeq\frac{1}{2}\left(\frac{\alpha}{4\pi\sin^2\theta_W}\right)^2\frac{M^2}{m_{W_R}^2}\left(\frac{m_t m_b}{m_{W_L}^2}\right)m_l \tag{3}$$

or $m_{\nu_i}\simeq 10^{-7}m_{l_i}$ for $M\simeq m_{W_R}$. (Here M is the singlet-fermion mass scale.) This not only explains the smallness of the neutrino masses, but is also consistent with the cosmological bounds.

The mass matrix for charged fermions (that includes the light and heavy sector) is given by (for the up-quark sector)

$$M^+=\begin{pmatrix}0 & \kappa_L h^u\\ \kappa_R h^{u\dagger} & M^u\end{pmatrix}, \tag{4}$$

where $M^u=m^u+i\langle\sigma\rangle f^u$; m^u and f^u are both Hermitian matrices but M^u is not. Similar mass matrices M^- and M^l arise for the down-quark and the charged-lepton sectors, respectively. To diagonalize the above mass matrices, define matrices

$$\rho_L^u=\kappa_L h^u M^{u-1},\quad \rho_R^{u\dagger}=\kappa_R M^{u-1}h^u. \tag{5}$$

If we assume that $\kappa_L h^u\ll\kappa_R h^u\ll M^u$ (which is reasonable since M^u is a gauge singlet mass), then typical elements of ρ_L and ρ_R are much smaller than 1. To second order in $\rho_{L,R}$, the matrix M^+ is block diagonalized by the transformation $D^+=U_L M^+ U_R^\dagger$, where

$$U_A=\begin{pmatrix}1-\rho_A\rho_A^\dagger/2 & -\rho_A\\ \rho_A^\dagger & 1-\rho_A^\dagger\rho_A/2\end{pmatrix},\quad A=L,R, \tag{6}$$

$$D^+=\begin{pmatrix}-h^u M^{u-1}h^{u\dagger}\kappa_L\kappa_R & 0\\ 0 & M^u\end{pmatrix}[1+O(\rho^2)]. \tag{7}$$

We can further diagonalize D^+ by the following two unitary matrices

$$K_{L,R}^+=\begin{pmatrix}K_{L,R}^u & 0\\ 0 & K_{L,R}^P\end{pmatrix} \tag{8}$$

via the biunitary transformation $M_{\rm diag}^+=K_L^+D^+K_R^{+\dagger}$. Two distinct cases arise depending on whether we assume (a) CP violation is intrinsic or (b) it is spontaneous. In case (a), the matrices h^u and M^u are arbitrary complex matrices; therefore K_L^u,K_L^P and K_R^u,K_R^P are totally unrelated to each other. We, therefore, have both mixing angles and phases arbitrary in the left and right sector. In this case, the weak Lagrangian in the mass-eigenstate basis can be written as (to order ρ)

$$\mathcal{L}_{\rm wk}^{\rm CC}=\frac{g}{\sqrt{2}}W_L^{\mu+}(\bar u_L\gamma_\mu V_L d_L+\bar u_L\gamma_\mu K_L^u\rho_L^d K_L^{N\dagger}N_L+\bar P_L\gamma_\mu K_L^P\rho_L^{u\dagger}K_L^{d\dagger}d_L)$$
$$+\frac{g}{\sqrt{2}}W_R^{\mu+}(\bar u_R\gamma_\mu V_R d_R+\bar u_R\gamma_\mu K_R^u\rho_R^d K_R^{N\dagger}N_R+\bar P_R\gamma_\mu K_R^P\rho_R^{u\dagger}K_R^{d\dagger}d_R)+\text{H.c.}+\text{leptonic part}. \tag{9}$$

Here $V_A=K_A^u K_A^{d\dagger}$, $A=L,R$ are the usual quark mixing matrices. The interactions involving the physical Higgs bosons $\sigma_L=\sqrt{2}\,\mathrm{Re}\chi_L^0$ and $\sigma_R=\sqrt{2}\,\mathrm{Re}\chi_R^0$ and the fermions can now be written as

$$\sqrt{2}\mathcal{L}^{\rm Higgs}=\bar u_L M_{\rm diag}^+u_R\left(\frac{\sigma_L}{\kappa_L}+\frac{\sigma_R}{\kappa_R}\right)+\bar u_L K_L^u h^u K_R^{P\dagger}P_R\sigma_L+\bar u_R K_R^u h^u K_L^{P\dagger}P_L\sigma_R+\text{H.c.} \tag{10}$$

and similarly for the down-quark and lepton sectors. Here we have neglected the couplings of the singlet σ connecting ordinary and heavy quarks, since these are suppressed at least by a factor $\kappa_R f^u M^{u-1}$. These formulas [Eqs. (9) and (10)] simplify once we turn to case (b), where CP is spontaneously violated by $\langle\sigma\rangle\neq 0$. In this case, the matrices M^i, $i=u,d,e$ are symmetric because of left-right symmetry. As a result, in Eq. (8), $K_{L,R}^i$ and $K_{L,R}^a$ satisfy the following relations: $K_L^i=K_R^{i*}$, $i=u,d$; $K_L^a=K_R^{a*}$, $a=P,N$. This implies that the rotation angles in L and R are the same, but the phases are different. In particular, for two generations, we have

$$V_R=\begin{pmatrix}e^{i\alpha} & 0\\ 0 & e^{-i\alpha}\end{pmatrix}V_L\begin{pmatrix}e^{i\beta} & 0\\ 0 & e^{-i\beta}\end{pmatrix}. \tag{11}$$

This corresponds to the pseudomanifest left-right symmetry.[7-9] Here V_L is the usual real two-generation Cabibbo rotation matrix.[10] Obviously, the exchange of left-handed W_L bosons does not lead to any CP-violating effect. On the other hand, the CP-violating effects induced by the right-handed W_R exchange satisfy the "isoconjugate" relations[7] leading to $\epsilon'=0$ and the magnitude of ϵ being proportional to $(m_{W_L}/m_{W_R})^2$. These results have been extensively discussed before and we do not repeat them here except to note that they are applicable in our case as well. There are, of course, new contributions to $\Delta S=2$ CP-violating Hamiltonian from box graphs involving heavy-single-quark intermediate states. They can arise from $W_L W_R$ and $W_R W_R$ exchanges, the former contribution being the dominant one. To estimate these contributions, we need to know the orders of magnitude of the singlet-fermion masses M, and the coupling h^i_{ab}. The values of these parameters are constrained by the seesaw formula for the quark masses. Ignoring the effects of mixing, we expect $m^u_a \simeq h^{u^2}_{aa}\kappa_L\kappa_R/M_P$ (and similarly for down quarks and leptons). For m_{W_R} in the TeV range, we expect the singlet-quark masses ($M_{P,N}$) in the 10 TeV range, leading to h^d_{11}, h^u_{11} of order 10^{-2}, h^u_{22} and h^d_{22} of order 10^{-1}, etc. From Eq. (9) we conclude that the dominant singlet-quark mediated W_LW_R contribution to $\Delta S=2$ amplitude is given by[11] (in the limit $m_{W_L} \ll m_{W_R} \ll M_{P_i}$, assuming the ith heavy quark to dominate)

$$H_{\Delta S=2} \simeq \frac{G_F}{\sqrt{2}}\frac{2\alpha}{\pi\sin^2\theta_W}\frac{m^2_{W_L}}{M^2_{P_i}}\lambda_{L_i}\lambda_{R_i}\left[1+\ln\left(\frac{m^2_{W_L}}{M^2_{P_i}}\right)-2\ln\left(\frac{m^2_{W_L}}{m^2_{W_R}}\right)\right], \tag{12}$$

$$\lambda_{A_i} = (K^P_A \rho^{u\dagger}_A K^{d\dagger}_A)^*_{is}(K^P_A\rho^{u\dagger}_A K^{d\dagger}_A)_{id}, \quad A=L,R. \tag{13}$$

We estimate the strength of the above $\Delta S=2$ amplitude (for $m_{W_R}\simeq 2$ TeV, $M_P \simeq 10$ TeV) to be

$$\langle K^0 | H^{\text{singlet}}_{\Delta S=2} | \bar{K}^0\rangle \simeq \langle K^0 | H^{LL}_{\Delta S=2} | \bar{K}^0\rangle 10^3 \lambda_{L_i}\lambda_{R_i}. \tag{14}$$

This implies that $\lambda_{L_i}\lambda_{R_i} < 10^{-3}$. Additional constraints on $\rho^{L,R}_{ij}$ arise from a tree-level flavor-changing neutral-current contribution, which provide much stronger bounds on ρ. For this purpose, let us write down the flavor-changing neutral-current interactions of Z_1 and Z_2 in the down-quark sector in this model[12]:

$$\mathcal{L}^{\text{FCNC}} = \frac{e}{\sin 2\theta_W} Z^\mu_1 \sum_{i\neq j}\left[\left(\cos\beta+\sin\beta\frac{\sin^2\theta_W}{(\cos 2\theta_W)^{1/2}}\right)F^q_{ij,L}\bar{q}_{i,L}\gamma_\mu q_{j,L} + 2\sin\beta\tan^2\theta_W(\cos 2\theta_W)^{1/2}F^q_{ij,R}\bar{q}_{i,R}\gamma_\mu q_{j,R}\right], \tag{15}$$

$$F_{ij,A} = (K_A\rho_A\rho^\dagger_A K^\dagger_A)_{ij}. \tag{16}$$

The interaction with Z_2 is obtained from the above by replacing $\sin\beta$ by $\cos\beta$ and $\cos\beta$ by $-\sin\beta$. Here $\sin\beta$ is the mixing angle in the neutral-gauge-boson sector which is of order κ^2_L/κ^2_R.[3] The parameters $\rho^{L,R}$ will be constrained by flavor-changing neutral-current processes such as $K^0_L \to \mu\bar{\mu}$ and $K_L - K_S$ mass difference. The former process leads to $\rho_{L,ij} \lesssim 10^{-2}$ and $\rho_{R,ij} \lesssim 10^{-2}(m_{Z_2}/m_{Z_1}) \simeq 10^{-1}$, in which case the latter constraint is satisfied automatically. This will also contribute to the $\Delta S=2$ CP-violating Hamiltonian via Z_1 and Z_2 exchange. They have strengths (we omit the CP-violating phase, but note that it is different from the phase that appears in the study of ϵ parameter)

$$f_{\Delta S=2,Z_1} \simeq G_F\rho^4_L \simeq G_F \times 10^{-8}\sin^2\theta_C, \quad f_{\Delta S=2,Z_2} \simeq G_F\rho^4_R(m_{Z_1}/m_{Z_2})^2. \tag{17}$$

Since, in this model, $m^2_{Z_2} \simeq m^2_{W_R}\cos^2\theta_W \sec 2\theta_W$, $f_{\Delta S=2,Z_2} \simeq G_F \times 10^{-11}$. This should be compared with the usual CP-violating contribution of order $G^2_F m^2_c/4\pi^2 \times 10^{-3} \simeq G_F \times 10^{-9}$. Therefore, these contributions are negligible.

To calculate ϵ' in this model we need the direct CP-violating $\Delta S=1$ nonleptonic Hamiltonian. They cannot come[9] from W_R exchange since W_R exchange contributions satisfy the isoconjugate relation and therefore lead to $\mathrm{Im}A_2/\mathrm{Re}A_2 = \mathrm{Im}A_0/\mathrm{Re}A_0$. Two possible sources of ϵ' in our model are (a) tree-level flavor-changing neutral currents and (b) penguin graphs involving singlet-fermion intermediate states.

The tree-level flavor-changing contributions to CP-violating nonleptonic Hamiltonian can be induced by Z_1 or Z_2 exchanges. Their strengths are given by $G_F\rho^2_L$ and $G_F\rho^2_R(m_{Z_1}/m_{Z_2})^2$. Using the bounds for ρ_L and ρ_R derived in the previous section, it is easily seen that their contribution to ϵ'/ϵ is of order 10^{-6} and hence negligible.

Turning to case (b), the contribution of the penguin graphs can be either of chirality-flip or of chirality-conserving type. In the case of chirality flip, the strength of the penguin graph involving σ_L-σ_R mixing is

$$H^{P,f}_{\Delta S=1} \simeq \zeta_{\sigma_L-\sigma_R}(K^d_L h^d K^{N\dagger}_R)_{1i}(K^d_R h^d K^{N\dagger}_L)^*_{2i}\frac{\alpha_s}{4\pi M_N Q^2}\bar{d}_L\sigma_{\mu\nu}\lambda^a s_R Q^\nu \sum_i \bar{q}_i\gamma_\mu\lambda^a q_i + L \leftrightarrow R + \text{H.c.}, \tag{18}$$

whereas the chirality-conserving Hamiltonian is given by

$$H^{P,c}_{\Delta S=1} \simeq (K^d_L h^d K^{N\dagger}_R)_{1i}(K^d_L h^d K^{N\dagger}_R)^*_{2i}\frac{\alpha_s}{4\pi M^2_N}\bar{d}_L\gamma_\mu\lambda^a s_L \sum_i \bar{q}_i\gamma_\mu\lambda^a q_i + L \leftrightarrow R + \text{H.c}. \tag{19}$$

In order to estimate the relative contribution of the above two operators [Eqs. (18) and (19)] we need an estimate of the $K \to 2\pi$ matrix element of the hadronic operators. We use the estimate of the chirality-preserving operator from Ref. 13 and the chirality-flipping one from Ref. 14 and find that the chirality-flipping contribution dominates in the CP-violating $K \to 2\pi$ transition by a factor of 100 leading to a phase $\xi_0 \simeq 10^{-4}$, which in turn implies $\epsilon'/\epsilon \simeq 10^{-3}$. This is in agreement with the recent NA-31[15] observation. The gauge-boson contribution to the above amplitude with heavy internal quark is of the same order as Eq. (19) and hence negligible. It is worth suggesting that the above estimate requires the singlet masses to be less than about 10 TeV.

Next, we discuss the electric dipole moment of the neutron (d_n) in this model. It arises from the σ_L-σ_R exchange graph (Fig. 1) and using Eq. (10) leads to the electric dipole moment of the d quark (and a similar expression for the u quark),

$$d_e^d \simeq \frac{1}{64\pi^2 M_N}\,\frac{1}{3}\,e\zeta_{\sigma_L\text{-}\sigma_R}\,\mathrm{Im}[(K_L^d h^d K_R^{N\dagger})_{1i}(K_R^d h^d K_L^{N\dagger})^*_{1i}] \simeq 10^{-26}\text{–}10^{-25}\,e\,\mathrm{cm}\,, \tag{20}$$

for the same range of parameters used earlier. This implies that $d_n \simeq 10^{-26}$–10^{-25} e cm. Note that this prediction can distinguish this model from the standard Kobayashi-Maskawa model.[16] It is perhaps worth noting that the dominant contribution to ϵ'/ϵ and d_n arises from the Feynman graph involving the exchange of σ-Higgs bosons; in this sense, ϵ'/ϵ and d_n are related to each other as in the usual left-right models, where they arise from W_L-W_R mixing.

Another interesting aspect of the model is a sizable electric dipole moment of the electron. Since the neutrino masses vanish at the tree level, there is no mixing between lepton generations in the weak current. However, there are interactions involving the doublet charged leptons (e,μ,τ) and E_1, E_2, E_3 that are CP-violating. They would contribute to the electric dipole moment of the electron via diagrams similar to Fig. 1. Since $h^e \simeq 6\times 10^{-3}$ for the first generation, we expect $d_e \simeq 10^{-27}$ e cm.

If we parametrize the $K_L \to \mu\bar{\mu}$ decay Hamiltonian as[17]

$$H_{\mathrm{wk}} = \frac{G_F}{\sqrt{2}} i\bar{d}\gamma_5 s(a\bar{\mu}\gamma_5\mu + ib\bar{\mu}\mu)\,, \tag{21}$$

the longitudinal muon polarization is $\mathcal{P} \simeq 1.8\times 10^6 b$. In our model, the dominant contribution to b comes from box graphs involving σ_L-σ_R exchange and heavy-singlet fermions. Making the conservative assumption that the vectorlike lepton dominates, we get

$$b \simeq \frac{G_F^{-1}}{8\sqrt{2}\pi^2 M_E^2} h_{22}^{d2} h_{22}^{e2} \sin\theta_C \sin\alpha \ln(M_E^2/M_{\sigma_L}^2)\,, \tag{22}$$

where h^d and h^e stand for the Yukawa coupling matrices, θ_C is the Cabibbo angle, and α is the CP phase. Using previous estimates that $h_{22}^d \simeq h_{22}^e \simeq \frac{1}{20}$ and choosing $M_E \simeq 10$ TeV, we obtain as a conservative bound $\mathcal{P} \geq 2\times 10^{-3}$.

Finally we comment on strong CP violation in this model. The effective strong CP parameter is defined as $\bar{\theta} = \theta + \mathrm{Arg\,Det}(M_u M_d)$. Parity invariance of the theory implies that $\theta = 0$ and a look at the mass matrices [Eq. (4)] immediately shows that $\mathrm{Arg\,Det}(M^{\pm}) = 0$ so that at the tree level $\bar{\theta} = 0$. It is easy to see that the one-loop contribution to $M^{\pm}$ arises from σ_L-σ_R exchange and leads to $\mathrm{Arg\,Det}(M^{\pm} + \delta M^{\pm}) = 0$, implying the vanishing of $\bar{\theta}_{\text{one loop}}$. Thus, the lowest order in which $\bar{\theta}$ may arise is the two-loop order.

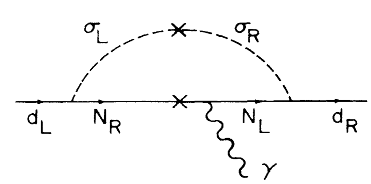

FIG. 1. Typical one-loop graph contributing to electric dipole moment of d quark and hence the neutron.

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