Gausted Flavor, Supersymmetry and Grand Unification

Rabindra N. Mohapatra

Maryland Center for Fundamental Physics and Department of Physics, University of Maryland, College Park, MD 20742, USA

Abstract. I review a recent work on gauged flavor with left-right symmetry, where all masses and all Yukawa couplings owe their origin to spontaneous flavor symmetry breaking. This is suggested as a precursor to a full understanding of flavor of quarks and leptons. An essential ingredient of this approach is the existence of heavy vector-like fermions, which is the home of flavor, which subsequently gets transmitted to the familiar quarks and leptons via the seesaw mechanism. I then discuss implications of extending this idea to include supersymmetry and finally speculate on a possible grand unified model based on the gauge group SU(5)_L × SU(5)_R which provides a group theoretic origin for the vector-like fermions.

Keywords: Gauged flavor, vector-like quarks, grand unification

PACS: 12.15.Ff

Understanding the flavor of quarks and leptons is a major unsolved problem of the standard model (SM). It is generally believed that a full understanding of the flavor symmetry U(3)^5 that emerges in the limit of zero Yukawa couplings of SM and how it breaks may hold the key to this problem. Use of these symmetries also forms the basis for a recent surge of interest in the so-called minimal flavor violation hypothesis[1], which states that the reason why SM provides such a good account of observed flavor violation is that any beyond the standard model (BSM) physics that incorporates new Higgs doublets as a way to understand flavor, must have all Yukawa couplings transform as (3, 3, 1), (3, 1, 3) under the full quark sector chiral flavor group U(3)^3. If this hypothesis is taken literally, one is left with the choice to imagine that the Yukawa couplings of the SM are merely vev’s of a set of spurion scalar fields of a higher scale theory and that all Yukawa couplings are vevs of scalar fields of higher scale dynamics. In such an approach, to avoid Goldstone bosons from creating conflict with cosmological observations, one must assume that the flavor symmetry of SM is indeed a gauge symmetry. In this article, we explore this gauged flavor approach. The approach is however intrinsically different from the usual MFV models[1] in that there are no extra standard model Higgs doublets but rather heavy vector like fermions which carry all the flavor information.

A convenient implementation of gauge flavor within SM has been carried out recently[2] where it is assumed that quark masses may owe their origin to a seesaw mechanism, involving vector like quarks and leptons, similar to the neutrinos. The idea of using seesaw mechanism for quarks was already discussed in the literature[3]more than two decades ago. The interesting point made in ref. [2] is that the same quark-seesaw framework also allows for gauging of the flavor symmetry without any anomalies. The requirement that there be vector-like quarks at high scale- preferably in the TeV range also implies that the model can be probed in collider searches. It was further pointed out in [2] that consistent with current flavor changing neutral current constraints, part of the gauged flavor dynamics as well as parts of the vector-like quark spectrum can be probed in the colliders.

The work of [2] is however phenomenologically incomplete since the model did not address the issue of neutrino masses. It is however easy to see how a simple extensions to the lepton sector can be carried out. It requires that there be three right handed neutrinos, which is an interesting consequence of flavor gauging since it makes neutrino mass natural. For earlier examples of models where flavor gauging implies non-zero neutrino masses see for instance[4]. The problem then is that in minimal models of this type, tiny Yukawa couplings are needed to give small masses to the neutrinos and even if we accept that, the lepton mixing angles vanish, making the model unacceptable. Clearly some nontrivial extension is needed.

At the conceptual level, one finds that not all fermion mass terms in the theory are protected by a gauge symmetry and could therefore be arbitrary (even Planck scale, as e.g. the electron mass in QED). This is different from the SM where all fermion masses owe their origin to spontaneous symmetry breaking and their magnitude must therefore be limited by the gauge symmetry breaking scale. To solve both these problems, it was proposed in [5] that we extend...
the standard model gauge group to the left-right symmetric group based on \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \) [5]. There are
then no free mass parameters in the fermion sector of the theory and all masses arise out of spontaneous symmetry breaking like in the standard model. It was shown that indeed the lightest vector-like fermions as well the flavor gauge bosons do indeed remain in the TeV to sub-TeV range even in this extension. An additional advantage of the left-right version is that it provides a solution to the strong CP problem [6] without the need for an axion. In this talk I review this work and comment on possible extension of this idea to supersymmetry and grand unification. The latter may answer the question as to where the vectorlike quarks with the particular quantum numbers come from? It has been known for some time that if one considers a grand unified extension of the standard model based on \( SU(5)_L \times SU(5)_R \) [7], then the vector like \( SU(2)_{LR} \) singlet quarks and leptons are automatically part of the fermion spectrum. This is particularly suited to accomodate the left-right version of the model [5] rather the GRV version.  

### GAUGED FLAVOR WITH LEFT-RIGHT SYMMETRY

In the SM, once the Yukawa couplings are set to zero, the maximal flavor symmetry group is \( SU(3)_{Q_L} \times SU(3)_{Q_R} \times SU(3)_{d_R} \times SU(3)_{t_L} \times SU(3)_{f_R} \). If the weak gauge group is extended to that of the left-right symmetric model, the flavor group becomes \( SU(3)_{Q_L} \times SU(3)_{Q_R} \times SU(3)_{t_L} \times SU(3)_{f_R} \), which is more economical and, unlike the SM, also simultaneously explains neutrino masses.

We will therefore start with the gauge group \( G_{LR} \equiv SU(3)_{c_L} \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_{Q_L} \times SU(3)_{Q_R} \times SU(3)_{t_L} \times SU(3)_{f_R} \), where \( SU(3)_{Q_L} \times SU(3)_{Q_R} \) represents the flavor gauge symmetries respectively in the left- and right-handed quark sector, and \( SU(3)_{t_L} \times SU(3)_{f_R} \), the corresponding ones for the lepton sector. The particle content and its transformation properties under fundamental representations of the group \( G_{LR} \) are as in table I below. It is easy to verify that this field content makes \( G_{LR} \) completely anomaly free. In fact the full anomaly free gauge group also contains chiral color \( SU(3)_{c_L} \times SU(3)_{c_R} \) and if this symmetry is broken at the TeV scale, it can give rise to near TeV mass axi-gluons [9] which is a class of particle being searched for at the LHC. Our detailed phenomenological considerations below do not depend on whether axigluons exist or not.

---

2 This is different from recent attempts to grand unify such models in [8]
The Yukawa couplings of the model are given by:

\[ \mathcal{L}_Y = \mathcal{L}_Y^{\text{kin}} - V(Y_u, Y_d, \chi_L, \chi_R) + \lambda_u(\hat{O}_L \hat{\chi}_L \psi_R^I + \hat{O}_R \hat{\chi}_R \psi_L^I) + \lambda_d(\hat{O}_L \chi_L \psi_R^I + \hat{O}_R \chi_R \psi_L^I) + \lambda_u \hat{\psi}_L^I Y_u \psi_R^I + \lambda_d \hat{\psi}_R^I Y_d \psi_R^I + \text{h.c.}, \]

We note at this point that, since under parity \( g_L \leftrightarrow g_R \) and \( \psi_L^I \leftrightarrow \psi_R^I \) (and similarly for \( \psi_{L,R}^I \)), parity symmetry requires \( Y_{u,d} \leftrightarrow Y_{d,u}^I \) and the \( \lambda_{u,d} \) as well as \( \lambda_{u,d}^I \) couplings to be real.

Concerning the breaking of the gauge groups, the flavor gauge group \( SU(3)_L \times SU(3)_R \) is broken spontaneously by the vevs of \( Y_u \) and \( Y_d \) while the group \( SU(2)_L \times SU(2)_R \) by the vevs of the Higgs doublets, \( \chi_{L,R} \), as already mentioned. In particular, we adopt the following vev normalization \( <\chi_L> = \left( \begin{array}{c} 0 \\ \nu_L \end{array} \right) \), \( <\chi_R> = \left( \begin{array}{c} 0 \\ \nu_R \end{array} \right) \), while diagonal \( Y \) vevs will be denoted henceforth as \( \langle \hat{Y}_{u,d} \rangle \). It is the \( <\hat{Y}_{u,d}> \)'s which are responsible for fermion masses and mixings. Thus all flavor originated from flavor breaking.

**Fermion masses**

From eq. (1) one can read off the up-type fermion mass Lagrangian to be \( \mathcal{L}_m = \hat{U}_LM_u U_R \), with \( U^T = (u, \psi^I) \), each of the \( u \) and \( \psi^I \) fields carrying a generation index. The mass matrix reads

\[
M_u = \begin{pmatrix}
0 & \lambda_u \nu_L I_{3 \times 3} \\
\lambda_u \nu_R I_{3 \times 3} & \lambda_u^I <\hat{Y}_u>
\end{pmatrix}, \quad M_d = \begin{pmatrix}
0 & \lambda_d \nu_L I_{3 \times 3} \\
\lambda_d \nu_R I_{3 \times 3} & \lambda_d^I <\hat{Y}_d>
\end{pmatrix}.
\]

For simplicity, let us work in the limit that the parameters \( \lambda_u \nu_L \) and \( \lambda_u \nu_R \) are much smaller than any of the \( \lambda_u^I <\hat{Y}_u> \).

Without loss of generality we can assume that \(<Y_u>\) is diagonal and \(<Y_u> = V_{\text{CKM}} \hat{Y}_u V_{\text{CKM}}^T\) (With the subscript \( i \) in \(<\hat{Y}_{u(d)}>\), we shall henceforth label the diagonal entries of the flavon vev matrices.). We can then do a leading order diagonalization of quark fields and get quark masses by first changing to a basis where \( \psi^I_u = V_{\text{CKM}}^T \psi_u \) and \( u^I = V_{\text{CKM}}^T u \).

Then, to leading order in an expansion in the parameters, the masses of the up and down quarks can be written as:

\[ m_u = \frac{\lambda_u \nu_L}{\lambda_u^I <\hat{Y}_u>}, \quad m_d = \frac{\lambda_d \nu_L}{\lambda_d^I <\hat{Y}_d>}. \]

From this, the CKM matrices that govern the weak mixings in the light quark sector are inherited from off-diagonals in the flavon vevs \(<\hat{Y}_{u,d}>\). In the absence of exact parity, we will have a flavon vev pattern of the form \(<Y_u> = V_R^T \hat{Y}_u V_L >, <Y_d> = <\hat{Y}_d>, \) with \( V_{L,R} \) unitary. We summarize the salient points of the above discussion below:

(i) In the limit of \( \nu_R \ll <\hat{Y}_{u,d}> \), the elements of the diagonal \( <\hat{Y}_{u,d}> \) matrices follow an inverted hierarchy with respect to the quark masses [2, 3].

(ii) For a given value of \( \nu_R \) and of the \( \lambda^{(i)} \) couplings, or the corresponding exact expressions allow to fix the \( <\hat{Y}_{u,d}> \) entries. Since the \( Y_{u,d} \) vevs set also the mass scale for the flavor gauge bosons, the inverted hierarchy mentioned in item (i) implies a similar hierarchy in new flavor changing neutral current effects: the lighter the generations, the more suppressed the effects [2]. This is arguably one of the most attractive features of the model and has been quantitatively analyzed in ref.[5]. We summarize the results below.

(iii) The mass matrices \( M_{u,d} \) in the above discussion are hermitian, leading to \( \text{arg} \left\{ \det[M_{u,d}] \right\} = 0 \), and implying that the strong CP parameter at the tree level vanishes. The one loop calculation for a more general case of this type was carried out in Ref. [6]. Using this result, we conclude that the model solves the strong CP problem without the axion.

**Flavor gauge boson masses and phenomenology**

The masses of the \( SU(3)_L \times SU(3)_R \) gauge bosons \( G_{iL,R} \) (\( i = 1, \ldots, 8 \)) are obtained from the kinetic terms of \( Y_u \) and \( Y_d \) in the Lagrangian, \( \text{Tr} (\bar{D}^I Y_{u,d} D^I Y_{u,d}^T) \). In this subsections, we will discuss the various observables that are expected to provide a constraint (or else the possibility of a signal) for the model. Since in some cases – starting from the model spectrum – the model predictions vary in a wide range, we found it useful to explore these predictions with a flat scan of the model parameters in ref.[5] and the results are summarized below. This is done for both the cases with
and without TeV scale parity symmetry. In the other scenario where parity is not a good symmetry at the TeV scale, all the left vs. right couplings can be chosen as different from each other. Concerning the $SU(2)_{L,R}$ couplings, there are examples of scenarios where $g_R/g_L \neq 1$ for a UV complete theory which conserves parity. In ref.[5], we limited ourselves to the choice $g_R = 0.7 \cdot g_L$, which can be achieved in such models.

The flavor gauge bosons $G^{\alpha}_{aLR}$ couple to the currents $J^{\mu}_{LR} \equiv g_H \overline{Q}_{LR} \gamma^\mu \lambda^a \frac{2}{3} Q_{LR}$. Similarly as in Ref.[2], these interactions give rise to new, tree-level, contributions to the 4-fermion operators

\begin{equation}
Q^{\mu}_{1} = \left( \overline{q}^\dagger_{i} \gamma^\mu \gamma^\nu q^\nu_{j} \right) \left( \overline{q}^\dagger_{i} \gamma^\mu \gamma^\nu q^\nu_{j} \right),
\end{equation}

\begin{equation}
\tilde{Q}^{\mu}_{1} = \frac{\lambda^a}{c^a} Q^{\mu}_{1} |_{L \rightarrow R},
\end{equation}

\begin{equation}
Q^{\mu}_{2} = \left( q^\dagger_{i} \gamma^\mu \gamma^\nu q^\nu_{j} \right) \left( \overline{q}^\dagger_{i} \gamma^\mu \gamma^\nu q^\nu_{j} \right),
\end{equation}

with Latin and Greek indices on the quark fields denoting flavor and respectively color, and where $P_{LR} \equiv (1 \mp \gamma_5)/2$.

In the quark mass eigenstates basis, the Wilson coefficients of the above operators read

\begin{equation}
C^{q_{1}}_{1} = \frac{g_{H}}{8} (M_{q}^{2})_{a,b}^{-1} (V_{L}^{a} \lambda^{b} V_{L}^{\dagger})_{ij} (V_{L}^{a} \lambda^{b} V_{L}^{\dagger})_{ij},
\end{equation}

\begin{equation}
C^{q_{1}}_{2} = -\frac{g_{H}}{8} (M_{q}^{2})^{-1}_{a,b},
\end{equation}

\begin{equation}
C^{q_{1}}_{5} = \frac{g_{H}}{2} (M_{q}^{2})^{-1}_{a,b},
\end{equation}

where $q$ can be $u$ or $d$, and a sum over $a$ and $b$ in the range $1, \ldots, 8$ is understood. Updated bounds on the Wilson coefficients in eq. (6) have been reported by the UTfit collaboration [10] and usefully tabulated in their table 4 for the different meson-antimeson mixing processes. The contributions, predicted in our model, to the above coefficients have been explored by the random scan mentioned above. As previously anticipated, these contributions are well within the existing bounds in the bulk of the explored parameter space.

For the exact TeV scale parity, meaning $g_R/g_L = 1$, we find the lower bounds on the masses of the lightest vectorlike fermion (the top partner) and the lowest allowed mass for the lightest gauge boson to be $5b$ TeV and $10$ TeV’s respectively and for the case of no TeV-scale parity (where we assume, as mentioned $g_R/g_L = 0.7$ ), both those values come down to the sub-TeV range and hence accessible at the LHC. For details see, [5]. The current ATLAS and CMS bound on the vector-like top partners are 760 GeV and 475 GeV respectively.

Similarly for the mixings with vector-like quarks, the first and second generation quark partners have very small mixings with $u, d, c, s$ quarks whereas for the third generation and the right handed top partner, the mixings can be of order one. This has several interesting consequences for LHC search of vector-like quarks.

- In the pp collision, we could expect a large cross section for the production of a pair of $\psi, \bar{\psi}$, and each vector like quark through its mixing will decay to $\psi, t \rightarrow t + H$ with $t \rightarrow b + W$ and $H \rightarrow b \bar{b}$ and similarly for the $\bar{\psi}$ leading to a spectacular signature of six b-quarks in the final state.
- One would expect large FCNC effects in $t$ decays e.g. $t \rightarrow c + g$ is much enhanced over the SM prediction. With the definition $\mathcal{L}_{eff} = \kappa \sigma_{\mu \nu \gamma} G_{\mu \nu}$ in SM we expect $\kappa \sim 10^{-5}$ (TeV)$^{-1}$ whereas in our model we expect $\kappa \sim 10^{-3}$ (TeV)$^{-1}$. The current D0 limit on this is 0.018 TeV$^{-1}$.

It has also been pointed out that the presence of the light flavor gauge bosons allows a reconciliation of the "brewing" $\varepsilon_{K} - \sin 2\beta$ anomaly[11].

The same mechanism can be replicated in the lepton sector and the neutrinos now have mass. The vector-like fermions in the lepton sector include three vectorlike charged fermions ($E_{\ell,LR}$) and heavy neutral leptons ($N_{i,LR}$). In the presence of these fermions, the flavor gauge group becomes $SU(3)_{\ell,L} \times SU(3)_{\ell,R}$ under which the lepton doublets of the LR model as well as the $E, N$ transform as triplets. The gauge group is then anomaly free. Like in the quark case, there are flavon fields $Y_{\nu}$ which carry lepton flavor in their vevs. The mass matrices have similar forms as in the quark case. Without any additional Higgs fields, the neutrino are Dirac fermions.
FLAVOR PATTERN FROM SYMMETRY BREAKING

What distinguishes this approach to flavor from other ones in the literature is that all flavor originates from the vev of the flavon fields $Y_{u,d}$. It is therefore necessary to say a few words about this. In an unpublished work, we have looked at the minimization of the flavon potential in this approach. The flavon potential at the renormalizable level can be written as for the $Y_u$ and $Y_d$ as $V_u + V_d + V_{ud}$

$$V_u = -m_u^2 Tr Y_u^3 Y_u + \lambda_1 Tr (Y_u^3 Y_u)^2 + \lambda_2 Tr (Y_u^4 Y_u Y_u^2 Y_u + Det Y_u)$$

$$V_d = -m_d^2 Tr Y_d^3 Y_d + \lambda_1 Tr (Y_d^3 Y_d)^2 + \lambda_2 Tr (Y_d^4 Y_d Y_d^2 Y_d + Det Y_d)$$

$$V_{ud} = + m_{ud}^2 Tr (Y_u^3 Y_d) + \sum_{i,j,k,l} \lambda_{ijkl} Tr (Y_u^i Y_d^j Y_d^k Y_d^l) + \sum_{i,j,k,l} \lambda^{ij}_{ijkl} Tr(Y_d^i Y_d^j) Tr(Y_d^k Y_d^l) + \epsilon^{ijkl} Y_i Y_j Y_k + h.c.$$ (10)

where $i,j,k,l$ in the third line go over $u,d$ with the understanding that all $u$ and all $d$ terms are omitted. The minimum of this potential corresponds to

$$< Y_{u,d} > = \begin{pmatrix} M_{u,d} & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}$$ (11)

Thus this gives rise to the leading flavon vev, that corresponding to $Y_{u,d,11}$. Once we include $d=6$ terms in the potential, it induces a smaller vev in the 22 entries and the Det terms then induce the 33 entries. With further higher order terms, we can also induce the off diagonal terms, although to get the hierarchical pattern, we need to do fine tuning. One might contemplate generating the higher dimensional terms from a radiative symmetry breaking scheme.

LEPTON SECTOR

We now briefly discuss the lepton sector of the model. Before proceeding to the lepton in the left-right symmetric gauged flavor model, let us discuss the situation in the model of ref.[2]. While the ref.[2] does not discuss the lepton sector, a possible extension of this model to include the leptons is straight forward and would be to introduce the leptonic flavor gauge groups, $G_{L,H} \equiv SU(3)_L \times SU(3)_E \times SU(3)_{u,H}$. This group becomes anomaly free if in addition to SM leptons, we add two $SU(3)_L$ triplet but SM singlet fermions $\psi_{E_R}, \psi_{N_6}$ and an $\psi_{E_L}$ which is a triplet under $SU(3)_{E_R,H}$ group. We also include the flavor Higgs field $Y_3(3,\bar{3})$ under $G_{L,H}$. One can then write down the full gauge invariant Yukawa interaction for leptons and the mass terms to be:

$$\mathcal{L}_L = h_L L \bar{H} \psi_{E_R} + h_\nu L H \psi_{N_6} + \psi_{E_R} Y_\ell \psi_{E_L}$$ (12)

It is now clear that the charged fermion masses arise in this model from the seesaw mechanism via the vev of the $Y_\ell$ field where the neutrino mass is simply $m_\nu = h_\nu < H >$.

As far as neutrinos go, several points are worth noting:

- Anomaly freedom requires the existence of three right handed neutrinos and hence massive neutrinos (unlike the standard model where the right handed neutrino has to be added by hand).
- In the minimal version of the model, the neutrino is a Dirac fermion. However to get small masses for them, we need to have $h_\nu \sim 10^{-12}$ and less. While phenomenologically, there is nothing wrong with this, such tiny Yukawa coupling needs some explanation. and is generally considered undesirable.
- Finally, since $< Y_\ell >$ is the only source for flavor mixing, in the minimal version, by a choice of basis, $< Y_\ell >$ can be diagonalized without affecting any other term in the Lagrangian (since $h_\nu$ matrix is a unit matrix). As a result, there is no mixing among the neutrinos. Thus the minimal version of the CRV model in the lepton sector is not phenomenologically viable. As we see below, extension of the electroweak sector of the model to make it left-right symmetric, cures this problem.
- The flavor mixing can be generated by extending the Higgs sector to include a $SU(3)_c$ sextet scalar which gives a heavy Majorana mass to the right handed neutrinos and hence the mass to the light neutrinos via the seesaw mechanism[15]. In this case, the $< Y_\ell >$ is inversely proportional to the observed light neutrino mass matrix.
An implication of the above model is that since the right handed neutrino \( N_R \) and the SM leptonic doublet transform as fundamental representation of the same horizontal group, in a left-handed neutrino interaction with another neutrino, one can produce right handed sterile neutrinos. Such interactions will then provide a new drain on the energy on the supernova explosion. Considerations similar to the discussion of right handed neutrinos, imply that the leptonic gauge flavor scale in this model must be at least 20 TeV if the gauge flavor coupling is of the order of the weak gauge coupling[16]. We will see below that these constraints can be avoided in the left-right symmetric gauge flavor models.

The lepton sector of the left-right gauged flavor model is specified by the fermion assignment of Table I. Again, the neutrinos are Dirac fermions in the minimal version of the model. Following the same procedure as for the quarks, we see that mass matrix for the Dirac neutrinos has the form:

\[
M_\nu = \left( \begin{array}{cc}
\bar{\nu}_L & \bar{N}_L \\
0 & \lambda_{V_L} v_L
\end{array} \right) \left( \begin{array}{c}
\nu_R \\
\bar{N}_R
\end{array} \right)
\]

The neutrino mass is given by: \( m_\nu \sim \frac{\lambda_{V_L}^2 v_L}{v^2} \). The Yukawa coupling can now be in more reasonable range depending on the ratio \( \frac{v}{M_{\text{susy}}} \). For instance, for \( \frac{v}{M_{\text{susy}}} \sim 10^{-4} \), we find the largest \( \lambda \sim 10^{-4} \).

As far as the supernova bound on the flavor scale is concerned, due to the heavy mass of the right handed the channel that is open for the GRV model is now blocked. There is now of course constraints on the right handed \( W_R \) boson. If parity symmetry is not exact at the weak scale, we can dial down the right handed gauge coupling so that the \( W_R \) goes down can easily be in the 5-10 TeV range. This also satisfies the BBN constraints on the new interactions.

**SUPERSYMMETRIC GENERALIZATION**

The model is easily generalized to accomodate TeV scale supersymmetry. In this case, all fields in the Table I become superfields and each flavon field and Higgs field e.g. \( \chi_{L,R} \) are accompanied by their conjugate fields \( \bar{\chi}_{L,R} \) so that the model remains anomaly free.

An immediate issue with susy gauged flavor with TeV susy breaking is that the soft susy breaking mass terms for the fields \( Y_{a,d} \) and \( Y_{a,d} \) are in general different and therefore their vevs differ by order TeV mass. This has the consequence that the D-terms of the theory induce large mass differences between different squark flavors which in turn will lead to large flavor changing neutral current effects. In our model we can choose gauge mediated origin for the susy breaking using flavor blind messenger fields so that the \( Y_{a,d} \) and \( Y_{a,d} \) soft masses differ only at the three loop level. As a result, the induced squark flavor mass difference is of order \( \sim \alpha M^2_{\text{susy}} \) thereby keeping the FCNC effect under control. The detailed implications of this approach is currently under study.

An interesting point of the SUSY embedding of gauged flavor in our model is that in general it restricts the form of the R-parity violating terms in the superpotential. Since the leptonic and the quark flavor symmetries are separate, the only R-parity violating term at the renormalizable level is:

\[
W_{RPV} = \lambda''_{L,R} \psi''_L \psi''_d + \lambda''_{L,R} \psi_d \psi_d + \lambda_{L,R} \psi''_L \psi''_e + \lambda_{L,R} \psi''_L \psi''_e
\]

Note that both the leptonic and the quark couplings in Eq.14 are antisymmetric in the family indices, since they must be gauged flavor invariant. Focusing on the quark sector, we note that prior to symmetry breaking the interactions in 14 only connects the heavy quarks. Once symmetry is completely broken, the heavy quarks will mix with the light quarks and lead to effective R-P breaking terms of type \( u'd'd' \) in the superpotential involving the mass eigenstate quark super-fields. The strength of these interactions are given by \( \lambda''_{L,R} \left( \frac{v^2}{\xi_{\frac{1}{2},\frac{3}{2}}} \right) \). Note that the dominant R-parity violation comes from the \( \lambda_R \) term involving SM singlet fields and secondly for the first two generations this contribution is highly suppressed. For example a typical expectation for the coupling \( \lambda_{321} \sim 10^{-8} \). This leads to \( \Delta B = 2 \) neutron-anti-neutron oscillation via the diagram in Fig.1. The expected strength of \( N - \bar{N} \) oscillation is \( \sim \frac{10^{-32}}{M_{\text{Planck}}} \approx 10^{-37} \), which puts it beyond the reach of contemplated experiments. This property of the gauged flavor models is very similar to the case of MFV models[12].
The flavor group is a chiral group and it breaks down to QCD at some lower scale. We therefore use left and right flavon fields in terms of the form

$$\text{and top quarks, similarly for down (D) and charged leptons (E).}$$

Vector like quarks and leptons. It is clear that this will give same mass to all three flavor partners (U, C, T) of up, charm and top quarks, similarly for down (D) and charged leptons (E).

First point to note is that the quark and charged lepton fields denoted by $L, R$ transform as

$$SU(2)_L \times SU(2)_R \times SU(3)_C / \Sigma_5, \Sigma_10$$ under which both the left-handed as well as the right handed fields transform as three dimensional representations. In order to give mass to the neutrinos, we can add $N_{i,L,R}$ as $SU(3)$ triplets and $SU(5)$ singlets.

Turning now to gauge symmetry breaking and fermion masses, we first note that at the GUT scale, the color $SU(3)$ group is a chiral group and it breaks down to QCD at some lower scale.

We can envision the symmetry breaking as follows:

(i) $(24,1)\oplus(1,24)$ to break $SU(5)_L \times SU(5)_R$ down to $SU(3)_{c,L} \times SU(2)_L \times U(1)_{Y,L} \times SU(3)_{c,R} \times SU(2)_R \times U(1)_{Y,R}$;

(ii) Use $(5,5) + (10,10)$ (denoted by $\Sigma_5, \Sigma_{10}$) vevs break $SU(3)_{c,L} \times SU(3)_{c,R}$ to QCD and also give mass to the vector like quarks and lepton. It is clear that this will give same mass to all three flavor partners (U, C, T) of up, charm and top quarks, similarly for down (D) and charged leptons (E).

As far as the light heavy mixed masses are concerned, since the mass term is of the form $F_i T_b$ that generates mass terms of the form $dD^c$ and $F_i F_b$ of the form $uU^c$ (a, b being the flavor or horizontal quantum numbers), they are flavor nonsinglet. We therefore use left and right flavon fields $Y_{ab}$ which are singlets under $SU(5) \times SU(5)$ but sextets under the flavor group $SU(3)_Y$. These masses arise from Yukawa couplings of the form

$$\mathcal{L}_Y = h_1 F_L T_H \Sigma_5 + h_2 T_L T_R \Sigma_{10} +$$

$$\left( h_1 F_i T_j H_L + h_4 T_i T_j H_L \right) \frac{Y_{ij}}{M} + L \rightarrow R + h.c$$

Here $H_{L,R}$ transform as $(5,1) \oplus (1,5)$ under $SU(5)_L \times SU(5)_R$. After electroweak and right handed symmetry breaking, the quark-vectorlike quark mass matrices take the form:

$$M_{ab} = \left( \begin{array}{cc} 0 & h_3 Y_{ab} \\ h_{3,\nu_{ab}} & h_1 < \Sigma_5 > \end{array} \right)$$

and similarly for the up quarks and the charged leptons. Here $y_{ab} = \frac{<Y_{ab}>}{M}$. The mass formula for the light down quarks is then given by

$$M_{ab} \approx \frac{h_3^2 y_{ab}}{h_1 < \Sigma_5 >} (y^T)_{ab}$$

GRAND UNIFICATION POSSIBILITIES

The model is based on the left-right symmetric $SU(5)_L \times SU(5)_R$ gauge group with fermions assigned in a left-right symmetric manner to the $5 \oplus 10$ of each group. They are given below for one of them below and the other follows by replacing L by R:

$$M_L = \begin{pmatrix} D_1^c \\ D_2^c \\ D_3^c \\ e^- \\ \nu \end{pmatrix}_{L} ; T_L = \begin{pmatrix} 0 & U_3^c & -U_2^c & u_1 & d_1 \\ -U_3^c & 0 & U_1^c & u_2 & d_2 \\ U_2^c & -U_1^c & 0 & u_3 & d_3 \\ -u_1 & -u_2 & -u_3 & 0 & E^{+} \\ -d_1 & -d_2 & -d_3 & -E^{-} & 0 \end{pmatrix}_{L}$$

FIGURE 1. The tree level diagram for neutron-anti-neutron oscillation due to R-parity violating interaction.
As in [5], we will need two flavon multiplets to have nonzero quark mixings since one flavon vev can always be diagonalized by an $SU(3)_V$ transformation.

As far as coupling unification is concerned, it is worth pointing out that the GUT scale value of $\sin^2 \theta_W(M_U) = \frac{3}{8(1+\alpha_L/\alpha_R)}$ as noted in the third paper of [7]. This is to be contrasted with the $\sin^2 \theta_W(M_U)$ values for simple GUT theories e.g. $SU(5)$ or $SO(10)$ where it is equal to $\frac{3}{8}$. This implies that in order to get the weak scale value of $\sin^2 \theta_W(M_U)$, we must have $\alpha_R \gg \alpha_L$ i.e. parity must be broken before right handed gauge symmetry breaks. Typically this requires that the unification scale be much lower than the canonical $10^{15} - 10^{16}$ GeV. This raises the question as to whether the model is consistent with current proton life time bounds. The answer to this question is "yes" since the tree level gauge exchange generates typical baryon number violating operators of the form:

$$O_B = \bar{\psi}_d C^{-1} \gamma^\mu L \cdot \bar{\psi}_u \gamma_\mu Q / M_U$$

(19)

In order to generate the proton decay operator, one must use two heavy light mixing factors and with each one being very small, this gives rise to the strength of proton decay operator which is consistent with current proton decay life time bounds.

**SUMMARY AND OUTLOOK**

In summary, we have discussed a new approach to the fermion flavor problem where the introduction of TEV scale vector-like quarks to the standard model have made it possible to gauge the flavor symmetry. In this framework, all Yukawa couplings arise from flavor symmetry breaking while leaving some new particles with masses in the TeV range to sub-TeV range. The left-right version of this theory allows a solution of the strong CP problem without the axion. We also discuss an extension of the model to include supersymmetry and a possible grand unification is outlined. In the works on the subject to date, the flavon vevs are chosen by hand. Although there is some preliminary work on how to generate them from a complete theory, it will be interesting to generate the vevs either from radiative corrections in an UV complete theory or higher dimensional terms so that new insight into the flavor problem can emerge.

**ACKNOWLEDGMENTS**

This work is supported by National Science Foundation grant No. PHY-0968854. I would like to thank my collaborators on the project, Diego Guadagnoli and Ilmo Sung for many discussions and insights.

**REFERENCES**

1. R. S. Chivukula and H. Georgi, Phys. Lett. B 188, 99 (1987); G. D’Ambrosio, G. F. Giudice, G. Isidori and A. Strumia, Nucl. Phys. B 645, 155 (2002) [hep-ph/0207036]; A. J. Buras, M. V. Carlucci, S. Gori and G. Isidori, JHEP 1010, 009 (2010) [arXiv:1005.5310 [hep-ph]]; For an early MFV model, see G. C. Branco, W. Grimus and L. Lavoura, Phys. Lett. B 380, 119 (1996) [hep-ph/9601383].

2. B. Grinstein, M. Redi and G. Villadoro, JHEP 1011, 067 (2010) [arXiv:1009.2049 [hep-ph]].

3. Z. G. Berezhiani, Phys. Lett. B 150, 177 (1985); Z. G. Berezhiani and J. L. Chkareuli, Sov. Phys. Usp. 28, 104 (1985) [Usp. Fiz. Nauk 145, 137 (1985)]. A. Davidson and K. C. Wali, Phys. Rev. Lett. 60, 1813 (1988); K. S. Babu and R. N. Mohapatra, Phys. Rev. Lett. 62, 1079 (1989); Phys. Rev. D41, 1286 (1990).

4. R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. D 66, 051301 (2002) [hep-ph/0207110].

5. D. Guadagnoli, R. N. Mohapatra, I. Sung, JHEP 1104, 093 (2011); [arXiv:1103.4170 [hep-ph]].

6. K. S. Babu and R. N. Mohapatra, Phys. Rev. D 41, 1286 (1990).

7. A. Davidson, K. C. Wali, Phys. Rev. Lett. 58, 2623 (1987); P. L. Cho, Phys. Rev. D48, 3531-3541 (1993); R. N. Mohapatra, Phys. Rev. D54, 5728-5733 (1996); Phys.Lett. B379 (1996) 115-120; D. Emmanuel-Costa, E. T. Franco, R. Gonzalez Felipe, [arXiv:1104.2046 [hep-ph]].

8. T. Feldmann, JHEP 1104, 043 (2011); [arXiv:1010.2116 [hep-ph]].

9. P. H. Frampton and S. L. Glashow, Phys. Lett. B 190, 157 (1987).

10. M. Bona et al. [UTfit Collaboration], JHEP 0803, 049 (2008) [arXiv:0707.0636 [hep-ph]].

11. A. J. Buras, M. V. Carlucci, L. Merlo and E. Stamou, arXiv:1112.4477 [hep-ph].

12. E. Nikolaidakis and C. Smith, Phys. Rev. D 77, 015021 (2008) [arXiv:0710.3129 [hep-ph]]; C. Csaki, Y. Grossman and B. Heidenreich, arXiv:1111.1239 [hep-ph]; G. Arcadi, L. Di Luzio and M. Nardecchia, arXiv:1111.3941 [hep-ph].
13. R. Alonso, M. B. Gavela, L. Merlo and S. Rigolin, JHEP 1107, 012 (2011) [arXiv:1103.2915 [hep-ph]].
14. D. Guadagnoli, R. N. Mohapatra and I. Sung, presented by RNM at the Planck 2011 conference, Lisbon, June (2011).
15. P. Minkowski, Phys. Lett. B67, 421 (1977); T. Yanagida in Workshop on Unified Theories, KEK Report 79-18, p. 95 (1979);
M. Gell-Mann, P. Ramond and R. Slansky, Supergravity, p. 315; Amsterdam: North Holland (1979); S. L. Glashow, 1979 Cargese Summer Institute on Quarks and Leptons, p. 687; New York: Plenum (1980); R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
16. R. Barbieri and R. N. Mohapatra, Phys. Rev. D 39, 1229 (1989).