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Quark mixing from softly broken symmetries

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Abstract: Quark flavor mixing may originate in the soft breaking of horizontal symmetries. Those symmetries, which in the simplest case are three family U(1) groups, are obeyed only by the dimension-4 Yukawa couplings and lead, when unbroken, to the absence of mixing. Their breaking may arise from the dimension-3 mass terms of SU(2)-singlet vector-like quarks. Those gauge-singlet mass terms break the horizontal symmetries at a scale *much higher* than the Fermi scale, yet *softly*, leading to quark mixing while the quark masses remain unsuppressed.

Keywords: Quark Masses and SM Parameters, Beyond Standard Model.

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1. Introduction

The seesaw mechanism [1] is a very popular mechanism for explaining the smallness of the neutrino masses. Extending the lepton sector of the Standard Model (SM) by right-handed neutrino singlets ν_R, one obtains the full Majorana mass matrix

$$\begin{pmatrix} M_L & M_D^T \\ M_D & M_R \end{pmatrix} \tag{1.1}$$

for the neutral components of the left-handed SM lepton doublets, ν_L, together with the charge-conjugates of the neutrino singlets, $(\nu_R)^c$. If the scale inherent in M_R is much larger than the scales of M_L and M_D, then eq. (1.1) leads to an effective mass matrix

$$M_L - M_D^T M_R^{-1} M_D \tag{1.2}$$

for the light neutrinos.

On the one hand this mechanism provides a rationale for the smallness of the neutrino masses, yet on the other hand it complicates the situation with respect to lepton mixing. This has now four sources: the charged-lepton mass matrix, M_L, M_D, and M_R. In [2, 3] a scenario has been proposed where the sole source of lepton mixing is M_R. That scenario is based on the following assumptions:

1. The lepton-family numbers L_α ($\alpha = e, \mu, \tau$) are conserved by all terms in the lagrangian of dimension 4;

2. The L_α are broken *softly* by the ν_R Majorana mass terms in M_R, which have dimension 3.

Such a framework [4] leads to a renormalizable theory, and it has been shown [3] to be compatible with all experimental constraints. Within the above scenario several models for explaining maximal atmospheric neutrino mixing have been devised [2, 5, 6, 7].

In this letter we transfer the idea of obtaining lepton mixing from softly broken lepton-family symmetries to the quark sector. We add vector-like isosinglet quarks to the SM and suggest that it is the dimension-3 mass terms of those quarks which are responsible for mixing.

The idea that quark mixing is trivial at some high-energy scale is not new. For instance, in [8] non-trivial quark mixing is generated at the one-loop level from soft supersymmetry breaking and flavor-changing terms.

2. Extension of the SM quark sector

SU(2)-singlet (isosinglet) vector-like quarks may be added either to the SM down-type-quark sector, or to the up-type-quark sector, or to both simultaneously. For definiteness we shall first assume that they are added only to the down-type-quark sector; an extension of this sector through vector-like isosinglet quarks is motivated by the Grand Unified Theories based on the Lie algebra E_6 [9]. We add n left-handed and n right-handed quark SU(2) singlets, each of them with electric charge $-1/3$, to the SM [10, 11]. We then have 3 left-handed quark doublets $q_L = (p_L,\, n_L)^T$, n left-handed quark singlets N_L, and $3+n$ right-handed quark singlets N_R; there are also 3 right-handed quark fields p_R with electric charge $+2/3$ [12]. The gauge symmetry allows the terms

$$-\mathcal{L} = \cdots + \bar{q}_L \left(i\sigma_2 \phi^*\right) \Delta p_R + \bar{q}_L \phi\, \Gamma N_R + \bar{N}_L M N_R + \text{H.c.}\,, \tag{2.1}$$

where Δ is a dimensionless 3×3 matrix, Γ is a dimensionless $3\times(3+n)$ matrix, and M is a dimension-1 $n\times(3+n)$ matrix. For simplicity we have considered only one Higgs doublet, but this framework is independent of the number of Higgs doublets. Writing $M_p = v^*\Delta$ and $m = v\Gamma$, where v is the vacuum expectation value of ϕ^0, one obtains the mass terms

$$\bar{p}_L M_p p_R + \left(\, \bar{n}_L, N_L \,\right) \mathcal{M}\, N_R\,, \tag{2.2}$$

where the $(3+n)\times(3+n)$ matrix $\mathcal{M}$ is given by

$$\mathcal{M} = \begin{pmatrix} m \\ M \end{pmatrix}. \tag{2.3}$$

For the discussion of quark mixing we must consider

$$\mathcal{M}\mathcal{M}^\dagger = \begin{pmatrix} mm^\dagger & mM^\dagger \\ Mm^\dagger & MM^\dagger \end{pmatrix}. \tag{2.4}$$

It is natural to assume that m is of the order of the electroweak (Fermi) scale while M is of a much larger scale. Then the effective mass-squared matrix for the light down quarks is [13]

$$D_{\text{light}} = mm^\dagger - mM^\dagger \left(MM^\dagger\right)^{-1} Mm^\dagger\,. \tag{2.5}$$

Notice the similarity of the transition from eq. (2.4) to eq. (2.5) to the well-known seesaw transition from eq. (1.1) to eq. (1.2). However, eq. (2.5) does *not* involve a true seesaw mechanism; in general, no suppression of the magnitude of the quark masses is produced by eq. (2.5).

3. Softly broken quark-family numbers

Equation (2.5) can be used [14] to generate *unsuppressed* CP violation at low energies from the spontaneous breaking of that symmetry in a heavy sector. In this note we point out that quark *mixing* can also be generated from eq. (2.5), if some horizontal symmetry is *softly* broken in the matrix M.

As a first example, let us assume that $n = 3$; then there are six N_R, which we divide in two sets: three N_{R1} and three N_{R2}. We furthermore assume that there are three family numbers F_α ($\alpha = u, c, t$) such that the fields $q_{L\alpha}$, $N_{L\alpha}$, $p_{R\alpha}$, $N_{R1\alpha}$, and $N_{R2\alpha}$ all have $F_\alpha = 1$, and $F_\beta = 0$ for $\beta \neq \alpha$. We require family-number conservation in all terms of dimension 4 in the lagrangian. Then,

$$m = (m_1,\, m_2)\,, \tag{3.1}$$

where m_1 and m_2 are diagonal 3×3 matrices. On the other hand,

$$M = (M_1,\, M_2)\,, \tag{3.2}$$

and we allow the 3×3 matrices M_1 and M_2 to be non-diagonal, since they correspond to dimension-3 (soft) terms in the lagrangian. Equation (2.5) now reads

$$D_{\text{light}} = m_1 m_1^\dagger + m_2 m_2^\dagger - \left(m_1 M_1^\dagger + m_2 M_2^\dagger\right)\left(M_1 M_1^\dagger + M_2 M_2^\dagger\right)^{-1}\left(M_1 m_1^\dagger + M_2 m_2^\dagger\right), \tag{3.3}$$

and D_{light} will in general be non-diagonal. On the other hand, the Yukawa-coupling matrix Δ in eq. (2.1) is diagonal, and so is the up-type-quark mass matrix:

$$M_p = \text{diag}\left(m_u e^{i\phi_u},\, m_c e^{i\phi_c},\, m_t e^{i\phi_t}\right), \tag{3.4}$$

where $m_{u,c,t}$ are the up-type-quark masses; the $\phi_{u,c,t}$ are arbitrary phases, which should be removed through rephasings of the p_R.

Both M_1 and M_2 must be present for non-trivial mixing to be obtained. For instance, if $M_1 = 0$ then eq. (3.3) gives

$$D_{\text{light}} = m_1 m_1^\dagger + m_2 m_2^\dagger - m_2 M_2^\dagger \left(M_2 M_2^\dagger\right)^{-1} M_2 m_2^\dagger = m_1 m_1^\dagger, \tag{3.5}$$

which is diagonal.

Let us consider the mass terms involving the fields n_L. Since m_1 and m_2 are diagonal, there are three separate expressions (one for each α)

$$\bar{n}_{L\alpha}\left[(m_1)_{\alpha\alpha} N_{R1\alpha} + (m_2)_{\alpha\alpha} N_{R2\alpha}\right]. \tag{3.6}$$

We may rotate each pair $(N_{R1\alpha}, N_{R2\alpha})^T$ by a unitary 2×2 matrix such that $(m_2)_{\alpha\alpha}$ becomes zero. Then $m_2 = 0$ and eq. (3.3) reads

$$D_{\text{light}} = m_1 m_1^\dagger - m_1 M_1^\dagger \left(M_1 M_1^\dagger + M_2 M_2^\dagger \right)^{-1} M_1 m_1^\dagger \,. \tag{3.7}$$

One may further perform unitary transformations of the N_L and N_{R2} in such a way that M_2 becomes diagonal and real non-negative (bi-diagonalization of M_2).

It is interesting to notice that a genuine seesaw mechanism for Dirac fermions may be obtained from eq. (3.7) if M_1 happens to be of a much larger order of magnitude than M_2; we stress, though, that the inequality $M_1 M_1^\dagger \gg M_2 M_2^\dagger$ must hold in the basis where $m_2 = 0$. Then we obtain the mass terms in leading order as [15]

$$-\bar{n}_L \left(m_1 M_1^{-1} M_2 \right) N_{R2} + \bar{N}_L M_1 N_{R1} \,. \tag{3.8}$$

The diagonal elements of M_1 and M_2 conserve the quark-family numbers. It is natural to assume that the off-diagonal elements of those matrices are small compared to the diagonal ones.[1] This might give a rationale for rather small quark mixing, in contrast to large lepton mixing. Let us then assume that

$m_2 = 0$;

both m_1 and M_2 are diagonal and real;

$M_1 = \hat{M}_1 + A$, where $\hat{M}_1$ is diagonal and real while A has zeros in the diagonal and its off-diagonal elements are small compared to the diagonal elements of both $\hat{M}_1$ and M_2.

Then, to first order in A one obtains

$$D_{\text{light}} = \frac{m_1^2 M_2^2}{\hat{M}_1^2 + M_2^2} - \frac{m_1 \hat{M}_1}{\hat{M}_1^2 + M_2^2} A \frac{m_1 M_2^2}{\hat{M}_1^2 + M_2^2} - \frac{m_1 M_2^2}{\hat{M}_1^2 + M_2^2} A^\dagger \frac{m_1 \hat{M}_1}{\hat{M}_1^2 + M_2^2} \,. \tag{3.9}$$

4. Less than 3 isosinglet quarks

In order to simplify the above scheme one may allow for $n < 3$. There will still be three $\mathrm{U}(1)_\alpha$ groups ($\alpha = u, c, t$), but one (if $n = 2$) or two (if $n = 1$) of these groups act trivially on the N_{R2}. For instance, if $n = 1$ then M_1 is a row vector with three elements and M_2 is just a number; m_2 is a column vector with three elements.[2] As discussed before, with a change of basis we achieve $m_2 = 0$. Then, with the notation $m_1 = \mathrm{diag}\,(a_1, a_2, a_3)$, $S^2 = M_1 M_1^\dagger + M_2 M_2^\dagger$ (now this is simply a number), and $\ell_j = (M_1)_j \big/ S$, from eq. (3.7) one obtains

$$(D_{\text{light}})_{ij} = a_i a_j^* \left(\delta_{ij} - \ell_i^* \ell_j \right) . \tag{4.1}$$

[1] This assumption may be misleading, though, since the down-type-quark masses have different orders of magnitude, and therefore the diagonal elements of M_1 and M_2 are also likely to be widely different.

[2] The position of its non-zero element is determined by the transformation property of N_{R2} under the $\mathrm{U}(1)_\alpha$ groups.

From this equation we see that

$$(D_{\mathrm{light}})_{12}\,(D_{\mathrm{light}})_{23}\,(D_{\mathrm{light}})_{31} = -\prod_{j=1}^{3} |a_j|^2\,|\ell_j|^2 \tag{4.2}$$

is real and negative. Since the up-type-quark mass matrix in eq. (3.4) is diagonal in this basis, $D_{\mathrm{light}} = V M_d^2 V^\dagger$, where M_d^2 is the diagonal matrix of the squared down-type-quark masses and V is the Cabibbo–Kobayashi–Maskawa (CKM) matrix. Thus,

$$\sum_{i,j,k=d,s,b} m_i^2 m_j^2 m_k^2 V_{ui} V_{ci}^* V_{cj} V_{tj}^* V_{tk} V_{uk}^* \tag{4.3}$$

is real and negative in this model. This means that there is no CP violation in the CKM matrix, and thus V may be set real. Moreover, given the known hierarchy of the down-type-quark masses and of the CKM matrix elements, the quantity in eq. (4.3) is approximately equal to

$$m_b^4 V_{tb}^2 V_{ub} V_{cb} \left(m_b^2 V_{ub} V_{cb} + m_s^2 V_{us} V_{cs}\right), \tag{4.4}$$

which must be negative. This implies that the CKM-matrix Wolfenstein parameter ρ [16] is negative in this model. A negative ρ is experimentally disfavored in the SM — see for instance [12] — but here things are more complicated, since this model has flavor-changing quark currents coupling to the Z^0 [10, 11], and, moreover, it must be endowed with some extra structure, like for instance several Higgs doublets, in order to account for CP violation.

5. Non-abelian symmetries

Another possibility consists in assuming that the horizontal symmetry is non-abelian, instead of being the product of three $\mathrm{U}(1)_\alpha$ groups. This is interesting since m_1 and/or m_2 may then become more constrained, in particular two or three of their diagonal elements may become equal. Relevant discrete groups might for instance be the dihedral groups [17], S_4, or A_4, which have three-dimensional irreducible representations (irreps), or S_3, or D_4, which have two-dimensional irreps.

For instance, the group of the permutations of four objects, S_4, has a three-dimensional faithful irrep $\underline{3}$ generated by

$$(13) \to \begin{pmatrix} 0 & 0 & -1 \\ 0 & 1 & 0 \\ -1 & 0 & 0 \end{pmatrix}, \quad (24) \to \begin{pmatrix} 0 & 0 & 1 \\ 0 & 1 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \quad (12) \to \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & -1 & 0 \end{pmatrix}, \tag{5.1}$$

where (13), (24), and (12) are transpositions. The product $\underline{3} \otimes \underline{3}$ contains only one invariant of S_4; if ψ_j $(j = 1, 2, 3)$ and ψ_j' $(j = 1, 2, 3)$ are two triplets of S_4, then $\psi_1 \psi_1' + \psi_2 \psi_2' + \psi_3 \psi_3'$ transforms trivially under S_4. One may then postulate that the three $q_{L\alpha}$, the three $N_{R1\alpha}$, and the three $N_{R2\alpha}$ are all triplets of S_4; if the Higgs doublets are invariant under S_4, this immediately makes m_1 and m_2 proportional to the unit matrix, and thus diagonal, even in the absence of any $\mathrm{U}(1)_\alpha$ symmetries. As for M_1 and M_2, they are arbitrary since we allow S_4 to be softly broken.

Another example is the symmetry D_4, which has a two-dimensional faithful irrep $\underline{2}$ given by the matrices

$$\begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \tag{5.2}$$

and all the products thereof. The product $\underline{2} \otimes \underline{2}$ contains only one invariant $\underline{1}$ of D_4; if (ψ_1, ψ_2) and (ψ'_1, ψ'_2) are two doublets of D_4, then $\psi_1\psi'_1 + \psi_2\psi'_2$ is invariant under D_4. Thus, if the three $q_{L\alpha}$ transform as $\underline{2} \oplus \underline{1}$ and if $n = 1$ with the four N_R transforming as $\underline{2} \oplus \underline{1} \oplus \underline{1}$, then, choosing a basis where $m_2 = 0$, $\mathcal{M}$ will be as described before in eq. (4.1) but with $a_1 = a_2$.

A disadvantage of non-abelian symmetries, though, is that one must allow for isosinglet vector-like quarks in both the down-type- and up-type-quark sectors. Indeed, since the $q_{L\alpha}$ transform as a multi-dimensional representation of some discrete (or continuous) non-abelian group, the absence of soft breaking of that group in one of the sectors leads to degenerate quarks in that sector. In order to avoid having two quarks with equal masses, the discrete group must be (either softly or spontaneously) broken in both the up-type-quark and down-type-quark sectors.

6. A particular case

Suppose for instance that $n = 2$ and that there is a softly broken symmetry S_4 such that

$$\mathcal{M} = \begin{pmatrix} m & 0 \\ M & N \end{pmatrix}, \tag{6.1}$$

where m is proportional to the unit matrix and (without loss of generality) real, while M is an arbitrary 2×3 matrix and N an arbitrary 2×2 matrix. Then,

$$D_{\text{light}} = m^2 \left[\mathbb{1} - M^\dagger \left(M M^\dagger + N N^\dagger \right)^{-1} M \right]. \tag{6.2}$$

We assume that the two row vectors in M are linearly independent. Then there is a column vector v_3 — unique up to a phase — such that $Mv_3 = 0$. It is clear that $D_{\text{light}} v_3 = m^2 v_3$. There is a 2×2 unitary matrix U and a 3×3 unitary matrix W such that

$$M = U \begin{pmatrix} \mu_1 & 0 & 0 \\ 0 & \mu_2 & 0 \end{pmatrix} W^\dagger, \tag{6.3}$$

with real and positive $\mu_{1,2}$. Then, with $\hat{\mu} = \text{diag}(\mu_1, \mu_2)$,

$$M M^\dagger = U \hat{\mu}^2 U^\dagger. \tag{6.4}$$

It is easy to see that

$$M^\dagger \left(M M^\dagger + N N^\dagger \right)^{-1} M = W \begin{pmatrix} 1 & 0 \\ 0 & 1 \\ 0 & 0 \end{pmatrix} (\mathbb{1} + B)^{-1} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \end{pmatrix} W^\dagger, \tag{6.5}$$

where

$$B = \hat{\mu}^{-1} U^\dagger N N^\dagger U \hat{\mu}^{-1}. \tag{6.6}$$

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Let x_1 and x_2 be the eigenvectors of the 2×2 hermitean matrix B, corresponding to the (real and positive) eigenvalues λ_1 and λ_2, respectively. Write $W = (v_1, v_2, v_3)$, where $\{v_1, v_2, v_3\}$ is an orthonormal system of 3-vectors. (Obviously, v_3 is the vector satisfying $Mv_3 = 0$.) We define

$$y_j = (x_j)_1 v_1 + (x_j)_2 v_2 \quad \text{for } j = 1, 2\,. \tag{6.7}$$

Then,

$$D_{\text{light}} y_j = \frac{m^2 \lambda_j}{1 + \lambda_j} y_j \quad \text{for } j = 1, 2\,. \tag{6.8}$$

We have thus found that y_1, y_2, and v_3 are the eigenvectors of D_{light} corresponding to the eigenvalues $m^2\lambda_1/(1+\lambda_1)$, $m^2\lambda_2/(1+\lambda_2)$, and m^2, respectively.

Since N is arbitrary, $NN^\dagger$ is an arbitrary hermitean matrix, and B too can be taken as a general hermitean positive 2×2 matrix. Therefore, λ_1 and λ_2 are arbitrary. Also, m^2 is unconstrained by the symmetry S_4. Finally, since M is arbitrary, W too is an arbitrary 3×3 unitary matrix, which means that y_1, y_2, and v_3 are arbitrary. We thus conclude that the matrix D_{light} in eq. (6.2) is a completely general 3×3 mass-squared matrix. This means that our simple case with $n = 2$ and softly broken S_4 symmetry leaves D_{light} unconstrained.

7. Conclusions

In this letter we have suggested that the Yukawa couplings of an extension of the SM with isosinglet vector-like quarks may respect a horizontal symmetry which would lead to trivial quark mixing. However, that horizontal symmetry is *softly* broken *at a high scale* (much above the Fermi scale) by the gauge-invariant mass terms of the vector-like quarks. We have stressed that this leads to *unsuppressed* quark mixing at low energies, even though that mixing arises from the heavy sector of the theory. The number of vector-like quarks is largely arbitrary, and the horizontal symmetry may either consist in the product of three U(1) quark-family groups or it may be a discrete or continuous non-abelian group implying flavor-diagonal Yukawa couplings. Our model leads to an understanding of the proximity of the CKM matrix to the identity matrix; if it also has the potential to understand the actual values of the off-diagonal elements of the CKM matrix needs further investigation.

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