We emphasise that even in the extreme chiral limit where only the top and bottom quarks acquire mass, quark mixing is physically meaningful. This implies that the natural value of $|V_{13}|^2 + |V_{23}|^2$ is of order one, which is to be compared to its experimental value of order $10^{-3}$. We show how this fine-tuning problem of the Standard Model can be solved through an extension of the Standard Model where vector-like quarks and a complex singlet are introduced, together with a flavour symmetry. The mixings of the light quarks are generated through the mixing of the vector-like quarks with the standard quarks.
1. Introduction

This talk is based on our previous work of Ref. [1]. Understanding the pattern of fermion masses and mixing remains one of the open fundamental questions in Particle Physics. The discovery of neutrino oscillations pointing towards non-vanishing neutrino masses and large leptonic mixing has rendered the above question even more challenging. In the quark sector, one may ask the following question: “In the framework of the Standard Model (SM), does the small quark mixing just reflect the strong hierarchy of quark masses?” In this talk, we will show that the answer to this question is a definite NO. Indeed we have shown [1] that even in the limit where only the third generation acquires mass, quark mixing is physically meaningful. More precisely, we have shown [2] that in the SM, the natural value of 

\[ |V_{13}|^2 + |V_{23}|^2 = O(1) \]  

(1.1)

This is to be compared with the experimental value:

\[ |V_{13}|^2 + |V_{23}|^2 \simeq 1.6 \times 10^{-3} \]  

(1.2)

It is clear that there is a novel fine-tuning problem in the SM. In order to display explicitly the origin of this fine-tuning, let us consider the extreme chiral limit (ECL) where the first two generations are massless:

\[ m_d = m_s = 0; \quad m_b \neq 0 \]  

(1.3)

\[ m_u = m_c = 0; \quad m_t \neq 0 \]  

(1.4)

In the EC limit, the general quark mass matrices can be written:

\[ M_d = U_{dL} U_{dR}^{\dagger} \text{diag}(0,0,m_b) U_{dL}^{\dagger}, \quad M_u = U_{uL} U_{uR}^{\dagger} \text{diag}(0,0,m_t) U_{uL}^{\dagger} \]  

(1.5)

where \( U_{dL,R} \) and \( U_{uL,R} \) are arbitrary unitary matrices. Note that the ordering of the eigenvalues in the diagonal matrices has no physical meaning, since a change of ordering can be included in the arbitrary matrices \( U_{dL,R} \) and \( U_{uL,R} \). Taking into account that in the EC limit the first two generations are massless, one can make an arbitrary redefinition of the light quark masses through a unitary transformation of the type

\[ W_{u,d} = \begin{pmatrix} X_{u,d} & 0 \\ 0 & 1 \end{pmatrix} \]  

(1.6)

where \( X_{u,d} \) are \( 2 \times 2 \) unitary matrices. Under this transformation, the quark mixing matrix \( V^0 \) transforms as:

\[ V^0 \rightarrow V' = W_u^{\dagger} V^0 W_d \]  

(1.7)

One can use the freedom to choose \( X_{u,d} \) at will, to diagonalise the upper left sector of \( V^0 \) leading to \( V'_{12} = V'_{21} = 0 \). So one has:

\[ V' = \begin{pmatrix} V'_{11} & 0 & V'_{13} \\ 0 & V'_{22} & V'_{23} \\ V'_{31} & V'_{32} & V'_{33} \end{pmatrix} \]  

(1.8)
Unitarity of $V'$ leads then to:

$$V'_{13} V'_{23} = 0 \quad \text{and} \quad V'_{31} V'_{32} = 0 \quad (1.9)$$

One can then choose, without loss of generality, $V'_{13} = V'_{31} = 0$ and the $V_{CKM}$ matrix becomes an orthogonal matrix:

$$V_{CKM} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & c\alpha & s\alpha \\ 0 & -s\alpha & c\alpha \end{bmatrix} \quad (1.10)$$

It is important to emphasize that mixing is meaningful even in the EC limit and it is arbitrary. In this limit the natural value for $\alpha$ is of order one, independently of the hierarchy of the quark masses. The apparent fine-tuning which we described above, provides motivation to introduce a symmetry which could provide a justification for the observed small mixing.

2. Small Quark Mixing from a Flavour Symmetry

Let us introduce the following symmetry in the context of the SM:

$$Q^0_{Lj} \rightarrow e^{i\tau} Q^0_{Li}, \quad Q^0_{L3} \rightarrow e^{-i\tau} Q^0_{L3}, \quad d^0_{Rj} \rightarrow e^{-i\tau} d^0_{R1}, \quad d^0_{R2} \rightarrow e^{-i\tau} d^0_{R2}, \quad d^0_{R3} \rightarrow e^{-2i\tau} d^0_{R3},$$

$$u^0_{R1} \rightarrow e^{i\tau} u^0_{R1}, \quad u^0_{R2} \rightarrow e^{i\tau} u^0_{R3}, \quad u^0_{R3} \rightarrow u^0_{R3}, \quad \Phi \rightarrow e^{i\tau} \Phi$$

where the $Q^0_{Lj}$ are left-handed quark doublets, $d^0_{Rj}$ and $u^0_{Rj}$ are right-handed quark singlets and $\Phi$ denotes the Higgs doublet. The Yukawa interactions can be written:

$$\mathcal{L}_Y = \left[ -\overline{Q}^0_{Li} \Phi d^0_{Rj} - \overline{U}^0_{Li} \Phi u^0_{Rj} \right] + \text{h.c.}, \quad (2.1)$$

and this symmetry constrains the Yukawa couplings to be of the form:

$$Y_d = \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & \times \end{bmatrix}, \quad Y_u = \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & \times \end{bmatrix} \quad (2.2)$$

which generate a $V_{CKM}$ equal to the identity with only the third generation acquiring mass.

3. Generating a realistic quark mass spectrum and quark mixing

The generation of realistic quark masses and mixing will be obtained through the introduction of vector-like quarks (VLQ). Extensions of the SM with VLQ arise in a variety of scenarios [3–19] and can play an important role in generating a complex CKM matrix in models with spontaneous CP violation and they can also provide a solution of the strong CP problem without axions [20]. Note that there is experimental evidence for a complex CKM matrix even if one allows for the presence of New Physics [21].

Let us introduce three down ($D^0_{Li}, D^0_{Ri}$) and three up ($U^0_{Li}, U^0_{Ri}$) vector-like isosinglet quarks. The Yukawa interactions are given by:

$$\mathcal{L}_Y = \left[ -\overline{Q}^0_{Li} \Phi (Y_d)_{ia} d^0_{R\alpha} - \overline{U}^0_{Li} \Phi (Y_u)_{ib} u^0_{R\beta} \right] + \text{h.c.}, \quad (3.1)$$
where the index $i$ runs from 1 to 3, as in the SM, while the indices $\alpha$ and $\beta$ cover all right-handed quark singlets of the down and up sector, respectively. A generic bare mass term is also $SU(2) \times U(1)$ gauge invariant and therefore should be introduced:

$$
\mathcal{L}_{\text{b.m.}} = \left[ - \overline{D}_{L,i}^0 (\eta_d)_{j\alpha} d_{R\alpha}^0 - \overline{U}_{L,k}^0 (\eta_u)_{k\beta} u_{R\beta}^0 \right] + \text{h.c.} \quad (3.2)
$$

The indices $j$ and $k$ run over all left-handed vectorial quarks in each sector. As mentioned before, in all examples that follow $i$, $j$ and $k$ run from 1 to 3 and therefore $\alpha$ and $\beta$ run from 1 to 6 (obviously $D_{Ri}^0 \equiv d_{Ri+3}^0$ and $U_{Ri}^0 \equiv u_{Ri+3}^0$). In what follows we extend the discrete flavour symmetry of Eq. (2.1) and we introduce a complex scalar singlet $S$. This scalar singlet will couple to the quark singlets as:

$$
\mathcal{L}_g = \left[ - \overline{D}_{L,i}^0 [(g_d)_{ja} S + (g_u')_{ja} S^\dagger] d_{R\alpha}^0 - \overline{U}_{L,k}^0 [(g_u)_{k\beta} S + (g_u')_{k\beta} S^\dagger] u_{R\beta}^0 \right] + \text{h.c.} \quad (3.3)
$$

We assume that the order of magnitude of the modulus of the vacuum expectation value of the field $S$ is higher than the electroweak scale. After spontaneous symmetry breaking one generates:

$$
\mathcal{L}_M = \left[ - \frac{v}{\sqrt{2}} \overline{D}_{L,i}^0 (Y_d)_{i\alpha} d_{R\alpha}^0 - \frac{v}{\sqrt{2}} \overline{U}_{L,k}^0 (Y_u)_{k\beta} u_{R\beta}^0 - \overline{D}_{L,i}^0 (\mu_d)_{i\alpha} d_{R\alpha}^0 - \overline{U}_{L,k}^0 (\mu_u)_{k\beta} u_{R\beta}^0 \right] + \text{h.c.} \quad (3.4)
$$

In a more compact form we can write:

$$
\mathcal{L}_M = - \left( \overline{D}_L^0 D_R^0 \right) \mathcal{M}_d \left( \begin{array}{c} d_R^0 \\ U_R^0 \end{array} \right) - \left( \overline{U}_L^0 U_R^0 \right) \mathcal{M}_u \left( \begin{array}{c} u_R^0 \\ \bar{u}_L^0 \end{array} \right) \quad (3.5)
$$

where $\mathcal{M}_d$ and $\mathcal{M}_u$, are $6 \times 6$ matrices denoted as:

$$
\mathcal{M}_d = \begin{pmatrix} m_d & \omega_d \\ X_d & M_d \end{pmatrix} \quad \mathcal{M}_u = \begin{pmatrix} m_u & \omega_u \\ X_u & M_u \end{pmatrix} \quad (3.6)
$$

### 3.1 Extension of the symmetry to the full Lagrangian

We have introduced three down-type and three up-type vector-like quarks. In the scalar sector, in addition to the standard Higgs, we have introduced a complex scalar $S$ and extend the symmetry to the full Lagrangian, with the new fields transforming in the following way under the family symmetry:

$$
\begin{align*}
D_{R1}^0 &\rightarrow e^{-3i\tau} D_{R1}^0 & D_{L2}^0 &\rightarrow e^{-2i\tau} D_{L2}^0 & D_{L3}^0 &\rightarrow e^{-i\tau} D_{L3}^0 \\
D_{R1}^0 &\rightarrow e^{-2i\tau} D_{R1}^0 & D_{R2}^0 &\rightarrow e^{-3i\tau} D_{R2}^0 & D_{R3}^0 &\rightarrow D_{R3}^0 \\
U_{L1}^0 &\rightarrow e^{-i\tau} U_{L1}^0 & U_{L2}^0 &\rightarrow U_{L2}^0 & U_{L3}^0 &\rightarrow e^{i\tau} U_{L3}^0 \\
U_{R1}^0 &\rightarrow U_{R1}^0 & U_{R2}^0 &\rightarrow e^{-i\tau} U_{R2}^0 & U_{R3}^0 &\rightarrow e^{2i\tau} U_{R3}^0; \quad S \rightarrow e^{i\tau} S; \quad \tau = \frac{2\pi}{6}
\end{align*}
$$

(3.7)

Together with the transformations for the standard-like quarks given in Eq. (2.1). The singlet scalar $S$ is introduced in order to be able to obtain realistic quark masses and mixing, without breaking the symmetry in the Yukawa couplings.

Table 1 and Table 2 summarise the information on the combination of the different fermionic charges and allow to see what is the pattern of the mass matrices. The first three rows come from
Yukawa terms of the form given by Eq. (3.1) and therefore are only allowed when the fermionic charge cancels the one coming from the scalar doublet. In these cases we write this charge explicitly. In the forbidden terms we put a bullet sign. The last three rows come from bare mass terms of the form given by Eq. (3.2) or else from couplings to the field $S$. We denote with 1 the entries which correspond to allowed bare mass terms and by the fermionic charges those terms that allow coupling to either $S$ or $S^*$. The introduction of this singlet scalar field provides a rationale for the choice of terms that would otherwise softly break the symmetry in the Yukawa couplings and would seem arbitrary.

Table 1: Down sector, summary of transformation properties. In the forbidden terms we put a bullet sign. We denote by 1 the entries corresponding to allowed bare mass terms. The fermionic charges are given for those terms that are allowed through couplings to scalar fields: $\Phi$, $S$ or $S^*$, to which we assign appropriate charges.

|                | $d_R^0$ | $d_R^0$ | $d_R^0$ | $D_R^0$ | $D_R^0$ | $D_R^0$ |
|----------------|---------|---------|---------|---------|---------|---------|
| $(Q_L^1)_1$    | $-\tau$| $-\tau$| $-2\tau$| $-2\tau$| $-3\tau$| $0$     |
| $(Q_L^2)_2$    | $\bullet$| $\bullet$| $\bullet$| $\bullet$| $-\tau$| $\bullet$|
| $(Q_L^3)_\tau$| $\bullet$| $\bullet$| $-\tau$| $-\tau$| $\bullet$| $\bullet$|
| $(D_L^0)_1$    | $\bullet$| $\bullet$| $\tau$| $\tau$| $1$     | $\bullet$|
| $(D_L^0)_2$    | $\tau$| $\tau$| $1$     | $1$     | $-\tau$| $\bullet$|
| $(D_L^0)_3$    | $1$    | $1$    | $-\tau$| $-\tau$| $\bullet$| $\tau$|

3.2 Effective Hermitian squared mass matrix

The $6 \times 6$ mass matrices $\mathcal{M}_d, \mathcal{M}_u$ are diagonalised through the bi-unitary transformations:

$$
\mathcal{U}^\dagger_L \mathcal{M}_d \mathcal{U}_L^d = \mathcal{D}_d \equiv \text{diag}(d_d, D_d)
$$

(3.8)

where $d_d \equiv \text{diag}(m_d, m_s, m_b)$, $D_d \equiv \text{diag}(M_{D1}, M_{D2}, M_{D3})$ and with $M_{Di}$ denoting the masses of the heavy quarks of charge $-1/3$. A similar equation can be written for $\mathcal{M}_u$. In order to have an idea of the main physical features involved, it is useful to perform an approximate evaluation of $\mathcal{U}^\dagger_L, \mathcal{U}^u_L$ and of the quark mass eigenvalues. For this purpose, we write $\mathcal{U}^d_L, \mathcal{U}^u_L$ in block form:

$$
\mathcal{U}_L = \begin{pmatrix} K & R \\ S & T \end{pmatrix}
$$

(3.9)
Table 2: Up sector, summary of transformation properties. In the forbidden terms we put a bullet sign. We denote by 1 the entries corresponding to allowed bare mass terms. The fermionic charges are given for those terms that are allowed through couplings to scalar fields: Φ, S or S∗, to which we assign appropriate charges.

| (U^0_{R1})/τ | (U^0_{R2})/τ | (U^0_{R3})/0 | (U^0_{R1})/0 | (U^0_{R2})/−τ | (U^0_{R3})/2τ |
|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| (Q^0_{L1})/τ   | •               | •               | •               | •               | •               |
| (Q^0_{L2})/2τ  | •               | •               | •               | τ               | •               |
| (Q^0_{L3})/τ   | •               | •               | τ               | τ               | •               |
| (U^0_{L1})/τ   | •               | •               | τ               | τ               | 1               | •               |
| (U^0_{L2})/0    | τ               | τ               | 1               | 1               | −τ             | •               |
| (U^0_{L3})/−τ  | 1               | 1               | −τ             | −τ             | •               | τ               |

where \( K, R, S, T \) are \( 3 \times 3 \) matrices. For simplicity, we drop the indices \( d \) and \( u \). It can be shown that the deviations of the unitarity of the matrix \( K \) are naturally small, of order \( m^2/M^2 \). From unitarity of \( U_L \) one has:

\[
KK^\dagger = 1 - RR^\dagger
\]  

(3.10)

with

\[
R \approx \frac{(mX^\dagger + \omega M^\dagger)T}{D^2} \approx (m/M)
\]  

(3.11)

and

\[
K^\dagger K = 1 - S^\dagger S
\]  

(3.12)

with

\[
S \approx \left( \frac{X(m^\dagger + M\omega^\dagger)}{XX^\dagger + MM^\dagger} \right) K \approx (m/M)
\]  

(3.13)

The matrices \( K_d, K_u \) can be evaluated from an effective Hermitian squared matrix \( \mathcal{H}_{\text{eff}} \) through:

\[
K^{-1}\mathcal{H}_{\text{eff}}K = d^2
\]  

(3.14)

with

\[
\mathcal{H}_{\text{eff}} = (mm^\dagger + \omega \omega^\dagger) - (mX^\dagger + \omega M^\dagger)(XX^\dagger + MM^\dagger)^{-1}(Xm^\dagger + M\omega^\dagger)
\]  

(3.15)

It can be shown that a realistic spectrum for the standard quarks can be generated and a realistic \( 3 \times 3 \) \( V_{\text{CKM}} \) matrix can be obtained, using the \( \mathcal{H}_{\text{eff}} \) of Eq. (3.15).
4. Conclusions

We have shown that there is a fine-tuning problem in the SM, related to the experimental fact that $|V_{13}|^2 + |V_{23}|^2 \approx 1.6 \times 10^{-3}$. We describe a possible solution which involves the introduction of a flavour symmetry, together with vector-like quarks of charge $(-1/3)$ and charge $(2/3)$, as well as a complex singlet scalar. In the absence of vector-like quarks only the bottom and the top quarks acquire mass and $V_{CKM} = 1$. In the presence of the vector-like quarks which mix with the standard quarks, a realistic quark mass spectrum can be obtained and a correct CKM matrix can be generated. These results are obtained in a framework where the imposed symmetry is an exact symmetry of the Lagrangian, only softly broken in the scalar potential. In Ref. [1] we have presented a detailed analysis of some of the salient phenomenological implications of this class of models, in particular the structure of FCNC, loop FCNC constraints, including $\Delta F = 2$ and $\Delta F = 1$ constraints, and we have also studied the heavy vector-like quark decay channels.

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