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May 9th 2007

**THE NEIGHBORHOOD OF THE STANDARD MODEL:  
MIXING ANGLES AND QUARK-LEPTON COMPLEMENTARITY  
FOR THREE GENERATIONS OF NON-DEGENERATE COUPLED FERMIONS**

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**Abstract:** We investigate the potential (small) deviations from the unitarity of the mixing matrix that are expected to occur, because of mass splittings, in the Quantum Field Theory of non-degenerate coupled systems. We extend our previous analysis concerning mixing angles, which led to a precise determination of the Cabibbo angle, to the case of three generations of fermions. We show that the same condition for neutral currents of *mass eigenstates*, *i.e.* that universality of diagonal currents is violated with the same strength as the absence of non-diagonal ones, is satisfied: on one hand, by the three CKM mixing angles with a precision higher than the experimental uncertainty; on the other hand, by a neutrino-like mixing pattern in which  $\theta_{23}$  is maximal, and  $\tan(2\theta_{12}) = 2$ . This last pattern turns out to satisfy exactly the “quark-lepton complementarity condition”  $\theta_c + \theta_{12} = \pi/4$ . Moreover, among all solutions, two values for the third neutrino mixing angle arise which satisfy the bound  $\sin^2(\theta_{13}) \leq 0.1$ :  $\theta_{13} = \pm 5.7 \cdot 10^{-3}$  and  $\theta_{13} = \pm 0.2717$ . The so-called “Neighborhood of the Standard Model” is thus confirmed to exhibit special patterns which presumably originate in physics “Beyond the Standard Model”.

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# 1 Introduction

Following the study of neutral kaons done in [1], we have shown in [2] and [3] that:

\* in Quantum Field Theory (QFT), mixing matrices linking flavour to mass eigenstates for non degenerate coupled systems should never be parametrized as unitary. Indeed, assuming that the effective renormalized quadratic Lagrangian is hermitian at any  $q^2$  and that flavour eigenstates form an orthonormal basis, different mass eigenstates, which correspond to different values of  $q^2$  (poles of the renormalized propagator) belong to different orthonormal bases <sup>1</sup>;

\* when it is so, the properties of universality for diagonal neutral currents and absence of flavor changing neutral currents (FCNC) which are systematically implemented, for the Standard Model (SM), in the space of flavour eigenstates, do not automatically translate anymore into equivalent properties in the space of mass eigenstates. In the case of two generations of fermions, imposing them for mass eigenstates yields two types of solutions for the mixing angles <sup>2</sup> of each doublet with identical electric charge: Cabibbo-like solutions <sup>3</sup> which reduce to a single unconstrained mixing angle, and a set of discrete solutions, unnoticed in the customary approach, including in particular the so-called maximal mixing  $\pi/4 \pm k\pi/2$ ;

\* for any of these solutions one recovers a unitary mixing matrix; but, as said above, very small deviations are expected due to mass splittings, which manifest themselves as a tiny departure from the exact two previous conditions. In particular, in the *neighborhood* of a Cabibbo-like solution, these deviations become of equal strength for a value of the mixing angle extremely close to the measured Cabibbo angle

$$\tan(2\theta_c) = 1/2. \quad (1)$$

This success was an encouragement to go further in this direction. We present below the outcome of our investigation of neutral current patterns in the case of three generations of fermions.

The intricate system of trigonometric equations has been solved by successive approximations, starting from configurations in which  $\theta_{13}$  is vanishing. We will see that this approximation, obviously inspired by the patterns of mixing angles determined from experimental measurements, turns out to be a very good one. Indeed, we precisely show, without exhibiting all the solutions of our equations, that the presently observed patterns of quarks as well as of neutrinos, do fulfill our criterion. While the three angles of the Cabibbo-Kobayashi-Maskawa (CKM) solution are ‘‘Cabibbo-like’’, the neutrino-like solution

$$\begin{aligned} \tan(2\theta_{12}) &= 2 \Leftrightarrow \theta_{12} \approx 31.7^\circ, \\ \theta_{23} &= \pi/4, \\ \theta_{13} &= \pm 5.7 \cdot 10^{-3} \text{ or } \theta_{13} = \pm 0.2717 \end{aligned} \quad (2)$$

is of a mixed type, where  $\theta_{23}$  is maximal while  $\theta_{12}$  and  $\theta_{13}$  are Cabibbo-like.

Two significant features in these results must be stressed. First, the values for the third neutrino mixing angle  $\theta_{13}$  given in (2) are predictions which take into account present (loose) experimental constraints. Only two possibilities survive: an extremely small value  $\theta_{13} \sim V_{ub} \sim \text{a few } 10^{-3}$ , and a rather ‘‘large’’ one, at the opposite side of the allowed range. Secondly, our procedure yields in an exact, though quite simple way, the well-known ‘‘quark-lepton complementarity relation’’ [6] for 1-2 mixing:

$$\theta_{12} + \theta_c = \pi/4, \quad (3)$$

where  $\theta_{12}$  is the leptonic angle, and  $\theta_c$  the Cabibbo angle for quarks.

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<sup>1</sup>Since at any *given*  $q^2$ , the set of eigenstates of the renormalized quadratic Lagrangian form an orthonormal basis, the mixing matrix with all its elements evaluated at this  $q^2$  is unitary and the unitarity of the theory is never jeopardized.

<sup>2</sup>For two generations, one is led to introduce two mixing angles to parametrize each  $2 \times 2$  non-unitary mixing matrix.

<sup>3</sup>Cabibbo-like angles can only be fixed by imposing conditions on the *violation pattern* of the unitarity of the mixing matrix in its vicinity.

## 2 Constraints on neutral currents of mass eigenstates

### 2.1 The different basis of fermions

Three bases will appear throughout the paper:

\* flavour eigenstates, that we note  $(u_f, c_f, t_f)$  and  $(d_f, s_f, b_f)$  for quarks,  $(e_f, \mu_f, \tau_f)$  and  $(\nu_{ef}, \nu_{\mu f}, \nu_{\tau f})$  for leptons;

\* mass eigenstates that we note  $(u_m, c_m, t_m)$  and  $(d_m, s_m, b_m)$  for quarks,  $(e_m, \mu_m, \tau_m)$  and  $(\nu_{em}, \nu_{\mu m}, \nu_{\tau m})$  for leptons; they include in particular the charged leptons detected experimentally, since their identification proceeds through the measurement of their *charge/mass* ratio in a magnetic field;

\* the neutrinos that couple to mass eigenstates of charged leptons in charged weak currents. These are the usual "electronic", "muonic" and "tau" neutrinos  $\nu_e, \nu_\mu, \nu_\tau$  considered in SM textbooks [5]: they are indeed identified by the outgoing charged leptons that they produce through charged weak currents, and the latter are precisely mass eigenstates (see above). These states read (see Appendix D)

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = K_\ell^\dagger \begin{pmatrix} \nu_{ef} \\ \nu_{\mu f} \\ \nu_{\tau f} \end{pmatrix} = (K_\ell^\dagger K_\nu) \begin{pmatrix} \nu_{em} \\ \nu_{\mu m} \\ \nu_{\tau m} \end{pmatrix}, \quad (4)$$

where  $K_\ell$  and  $K_\nu$  are the mixing matrices respectively of charged leptons and of neutrinos (*i.e.* the matrices that connect their flavour to their mass eigenstates). Note that these neutrinos coincide with flavour eigenstates when the mixing matrix of *charged leptons* is taken equal to unity  $K_\ell = 1$ , *i.e.* when the mass and flavour eigenstates of charged leptons are aligned, which is often assumed without justification in the literature.

### 2.2 Neutral currents and mixing matrices; notations

Neutral currents in the basis of mass eigenstates are controlled by the product  $K^\dagger K$  of the  $3 \times 3$  mixing matrix  $K$  with its hermitian conjugate (see [2]); for example, the (left-handed) neutral currents for quarks with electric charge  $(-1/3)$  read

$$\overline{\begin{pmatrix} d_f \\ s_f \\ b_f \end{pmatrix}} \gamma_L^\mu \begin{pmatrix} d_f \\ s_f \\ b_f \end{pmatrix} = \overline{\begin{pmatrix} d_m \\ s_m \\ b_m \end{pmatrix}} \gamma_L^\mu K_d^\dagger K_d \begin{pmatrix} d_m \\ s_m \\ b_m \end{pmatrix}. \quad (5)$$

We write each mixing matrix  $K$  as a product of three matrices, which reduce, in the unitarity limit, to the basic rotations by  $\theta_{12}, \theta_{23}$  and  $\theta_{13}$  (we are not concerned with *CP* violation)

$$K = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -\tilde{s}_{23} & \tilde{c}_{23} \end{pmatrix} \times \begin{pmatrix} c_{13} & 0 & s_{13} \\ 0 & 1 & 0 \\ -\tilde{s}_{13} & 0 & \tilde{c}_{13} \end{pmatrix} \times \begin{pmatrix} c_{12} & s_{12} & 0 \\ -\tilde{s}_{12} & \tilde{c}_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (6)$$

We parametrize each basic matrix, which is *a priori* non-unitary, with two angles, respectively  $(\theta_{12}, \tilde{\theta}_{12})$ ,  $(\theta_{23}, \tilde{\theta}_{23})$  and  $(\theta_{13}, \tilde{\theta}_{13})$ . We deal accordingly with six mixing angles, instead of three in the unitary case, where  $\tilde{\theta}_{ij} = \theta_{ij}$ . We will use throughout the paper the notations  $s_{ij} = \sin(\theta_{ij})$ ,  $\tilde{s}_{ij} = \sin(\tilde{\theta}_{ij})$ , and likewise, for the cosines,  $c_{ij} = \cos(\theta_{ij})$ ,  $\tilde{c}_{ij} = \cos(\tilde{\theta}_{ij})$ .

To lighten the text, the elements of  $K^\dagger K$  will be abbreviated by  $[ij]$ ,  $i, j = 1 \dots 3$  instead of  $(K^\dagger K)_{[ij]}$ , and the corresponding neutral current will be noted  $\{ij\}$ . So, in the quark case,  $\{12\}$  stands for  $\bar{u}_m \gamma_L^\mu c_m$

or  $\bar{d}_m \gamma_L^\mu s_m$ , and, in the neutrino case, for  $\bar{\nu}_{em} \gamma_L^\mu \nu_{\mu m}$  or  $\bar{e}_m \gamma_L^\mu \mu_m$ . Last, “channel  $(i, j)$ ” corresponds to two fermions  $i$  and  $j$  with identical electric charge; for example, “channel  $(2, 3)$ ” corresponds to  $(d, b)$ ,  $(c, t)$ ,  $(\mu^-, \tau^-)$  or  $(\nu_\mu, \nu_\tau)$ .

The general constraints that we investigate are five: three arise from the absence of non-diagonal neutral currents for mass eigenstates, and two from the universality of diagonal currents. Accordingly, one degree of freedom is expected to be unconstrained.

### 2.3 Absence of non-diagonal neutral currents of mass eigenstates

The three conditions read:

\* for the absence of  $\{13\}$  and  $\{31\}$  currents:

$$[13] = 0 = [31] \Leftrightarrow c_{12} [c_{13}s_{13} - \tilde{c}_{13}\tilde{s}_{13}(\tilde{c}_{23}^2 + s_{23}^2)] - \tilde{c}_{13}\tilde{s}_{12}(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23}) = 0; \quad (7)$$

\* for the absence of  $\{23\}$  and  $\{32\}$  currents:

$$[23] = 0 = [32] \Leftrightarrow s_{12} [c_{13}s_{13} - \tilde{c}_{13}\tilde{s}_{13}(\tilde{c}_{23}^2 + s_{23}^2)] + \tilde{c}_{13}\tilde{c}_{12}(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23}) = 0; \quad (8)$$

\* for the absence of  $\{12\}$  and  $\{21\}$  currents:

$$[12] = 0 = [21] \Leftrightarrow s_{12}c_{12}c_{13}^2 - \tilde{s}_{12}\tilde{c}_{12}(c_{23}^2 + \tilde{s}_{23}^2) + s_{12}c_{12}\tilde{s}_{13}^2(s_{23}^2 + \tilde{c}_{23}^2) + \tilde{s}_{13}(s_{12}\tilde{s}_{12} - c_{12}\tilde{c}_{12})(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23}) = 0. \quad (9)$$

### 2.4 Universality of diagonal neutral currents of mass eigenstates

The two independent conditions read:

\* equality of  $\{11\}$  and  $\{22\}$  currents:

$$[11] - [22] = 0 \Leftrightarrow (c_{12}^2 - s_{12}^2) [c_{13}^2 + \tilde{s}_{13}^2(s_{23}^2 + \tilde{c}_{23}^2)] - (\tilde{c}_{12}^2 - \tilde{s}_{12}^2)(c_{23}^2 + \tilde{s}_{23}^2) + 2\tilde{s}_{13}(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23})(c_{12}\tilde{s}_{12} + s_{12}\tilde{c}_{12}) = 0; \quad (10)$$

\* equality of  $\{22\}$  and  $\{33\}$  currents:

$$[22] - [33] = 0 \Leftrightarrow s_{12}^2 + \tilde{c}_{12}^2(c_{23}^2 + \tilde{s}_{23}^2) - (s_{23}^2 + \tilde{c}_{23}^2) + (1 + s_{12}^2) [\tilde{s}_{13}^2(s_{23}^2 + \tilde{c}_{23}^2) - s_{13}^2] + 2s_{12}\tilde{s}_{13}\tilde{c}_{12}(c_{23}\tilde{s}_{23} - c_{23}s_{23}) = 0. \quad (11)$$

The equality of  $\{11\}$  and  $\{33\}$  currents is of course not an independent condition.

### 2.5 Solutions for $\theta_{13} = 0 = \tilde{\theta}_{13}$

In a first step, to ease solving the system of trigonometric equations, we shall study the configuration in which one of the two angles parametrizing the 1-3 mixing vanishes<sup>4</sup>, which is very close to what is observed experimentally in the quark sector, and likely in the neutrino sector. It turns out, as demonstrated

<sup>4</sup>By doing so, we exploit the possibility to fix one degree of freedom left *a priori* unconstrained by the five equations; see subsection 2.2.

in Appendix A, that the second mixing angle vanishes simultaneously. We accordingly work in the approximation (the sensitivity of the solutions to a small variation of  $\theta_{13}, \tilde{\theta}_{13}$  will be studied afterwards)

$$\theta_{13} = 0 = \tilde{\theta}_{13}. \quad (12)$$

Eqs. (7), (8), (9), (10) and (11), reduce in this limit to

$$-\tilde{s}_{12}(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23}) = 0, \quad (13a)$$

$$\tilde{c}_{12}(c_{23}s_{23} - \tilde{c}_{23}\tilde{s}_{23}) = 0, \quad (13b)$$

$$s_{12}c_{12} - \tilde{s}_{12}\tilde{c}_{12}(c_{23}^2 + \tilde{s}_{23}^2) = 0, \quad (13c)$$

$$(c_{12}^2 - s_{12}^2) - (\tilde{c}_{12}^2 - \tilde{s}_{12}^2)(c_{23}^2 + \tilde{s}_{23}^2) = 0, \quad (13d)$$

$$s_{12}^2 + \tilde{c}_{12}^2(c_{23}^2 + \tilde{s}_{23}^2) - (s_{23}^2 + \tilde{c}_{23}^2) = 0. \quad (13e)$$

It is shown in Appendix B that the only solutions are  $\theta_{12}$  and  $\theta_{23}$  Cabibbo-like ( $\tilde{\theta}_{12,23} = \theta_{12,23} + k\pi$ ) or maximal ( $\theta_{12,23} = \pi/4 + n\pi/2$ ,  $\tilde{\theta}_{12,23} = \pi/4 + m\pi/2$ ).

Accordingly, the two following sections will respectively start from:

- \*  $\theta_{12}$  and  $\theta_{23}$  Cabibbo-like (and, in a first step, vanishing  $\theta_{13}$ ), which finally leads to a mixing pattern similar to what is observed for quarks;
- \*  $\theta_{23}$  maximal and  $\theta_{12}$  Cabibbo like (and, in a first step, vanishing  $\theta_{13}$ ), which finally leads to a mixing pattern similar to the one observed for neutrinos.

### 3 The quark sector; constraining the three CKM angles

Mass splittings entail that the previous general conditions, which, when exactly satisfied, correspond *de facto* to unitary mixing matrices, cannot be exactly fulfilled. We investigate the vicinity of their solutions, and show that the same violation pattern that led to an accurate determination of the Cabibbo angle in the case of two generations, is also satisfied by the CKM angles in the case of three generations.

#### 3.1 The simplified case $\theta_{13} = 0 = \tilde{\theta}_{13}$

In the neighborhood of the solution with both  $\theta_{12}$  and  $\theta_{23}$  Cabibbo-like, we write

$$\begin{aligned} \tilde{\theta}_{12} &= \theta_{12} + \epsilon, \\ \tilde{\theta}_{23} &= \theta_{23} + \eta. \end{aligned} \quad (14)$$

The pattern ( $\theta_{13} = 0 = \tilde{\theta}_{13}$ ) can be reasonably considered to be close to the experimental situation, at least close enough for trusting not only the relations involving the first and second generation, but also the third one.

Like in [3], we impose that the absence of  $\{12\}, \{21\}$  neutral currents is violated with the same strength as the universality of  $\{11\}$  and  $\{22\}$  currents. It reads

$$|2\eta s_{12}c_{12}s_{23}c_{23} + \epsilon(c_{12}^2 - s_{12}^2)| = |-2\eta s_{23}c_{23}(c_{12}^2 - s_{12}^2) + 4\epsilon s_{12}c_{12}|. \quad (15)$$

We choose the “+” sign for both sides, such that, for two generations only, the Cabibbo angle satisfies  $\tan(2\theta_{12}) = +1/2$ . (15) yields the ratio  $\eta/\epsilon$ , that we then plug into the condition equivalent to (15) for the (2, 3) channel.

$$|\eta c_{12}(c_{23}^2 - s_{23}^2)| = |2\eta s_{23}c_{23}(1 + c_{12}^2) - 2\epsilon s_{12}c_{12}|. \quad (16)$$

(15) and (16) yield

$$\tan(2\theta_{23}) = \frac{c_{12}}{1 + c_{12}^2 - 2s_{12}c_{12}} \frac{(s_{12}c_{12} + c_{12}^2 - s_{12}^2)}{4s_{12}c_{12} - (c_{12}^2 - s_{12}^2)} \approx \frac{c_{12}}{2 - \frac{5}{4} \frac{s_{12}c_{12}}{\tan(2\theta_{12}) - \frac{1}{2}}}. \quad (17)$$

In the r.h.s. of (17), we have assumed that  $\theta_{12}$  is close to its Cabibbo value  $\tan(2\theta_{12}) \approx 1/2$ .  $\theta_{23}$  is seen to vanish with  $[\tan(2\theta_{23}) - 1/2]$ . The predicted value for  $\theta_{23}$  is plotted in Fig. 1 as a function of  $\theta_{12}$ , together with the experimental intervals for  $\theta_{23}$  and  $\theta_{12}$ . There are two [7] for  $\theta_{12}$ ; the first comes from the measures of  $V_{ud}$  (in black on Fig. 1)

$$V_{ud} \in [0.9735, 0.9740] \Rightarrow \theta_{12} \in [0.2285, 0.2307], \quad (18)$$

and the second from the measures of  $V_{us}$  (in purple on Fig. 1)

$$V_{us} \in [0.2236, 0.2278] \Rightarrow \theta_{12} \in [0.2255, 0.2298]. \quad (19)$$

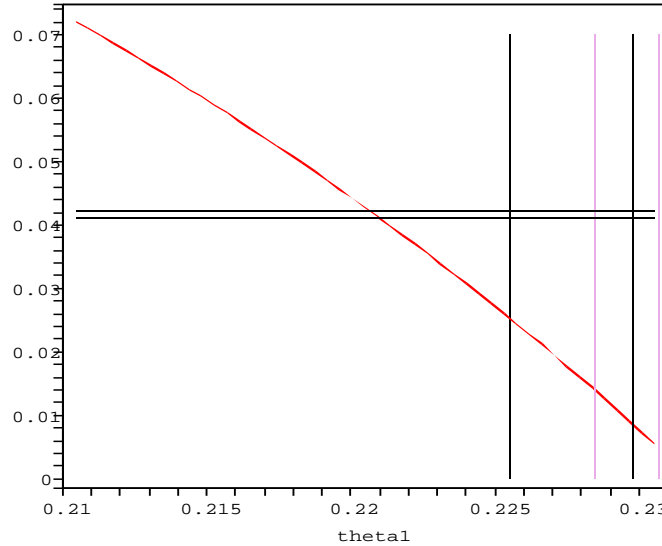


Fig. 1:  $\theta_{23}$  for quarks as a function of  $\theta_{12}$ ; simplified case  $\theta_{13} = 0 = \tilde{\theta}_{13}$

The measured value for  $\theta_{23}$  is seen on Fig. 1 to correspond to  $\theta_{12} \approx 0.221$ , that is  $\cos(\theta_{12}) \approx 0.9757$ . Our prediction for  $\cos(\theta_{12})$  is accordingly  $1.7 \cdot 10^{-3}$  away from the upper limit of the present upper bound for  $V_{ud} \equiv c_{12}c_{13}$  [8] [7]; it corresponds to twice the experimental uncertainty. It also corresponds to  $\sin(\theta_{12}) = 0.2192$ , while  $V_{us} \equiv s_{12}c_{13}$  is measured to be  $0.2247(19)$  [9] [7]; there, the discrepancy is  $2/100$ , only slightly above the  $1.8/100$  relative width of the experimental interval.

The approximation which sets  $\theta_{13} = 0 = \tilde{\theta}_{13}$  is accordingly reasonable, though it yields results slightly away from experimental bounds. We show in the next subsection that relaxing this approximation gives results in excellent agreement with present experiments.

### 3.2 Going to ( $\theta_{13} \neq 0, \tilde{\theta}_{13} \neq 0$ )

Considering all angles to be Cabibbo-like with, in addition to (14)

$$\tilde{\theta}_{13} = \theta_{13} + \rho, \quad (20)$$

the l.h.s.'s of eqs. (7),(8),(9), (10), (11) and the sum (10 + 11) depart respectively from zero by

$$\eta c_{13} [s_{12}(c_{23}^2 - s_{23}^2) + 2s_{13}c_{12}c_{23}s_{23}] - \rho c_{12}(c_{13}^2 - s_{13}^2); \quad (21a)$$

$$\eta c_{13} [-c_{12}(c_{23}^2 - s_{23}^2) + 2s_{13}s_{12}c_{23}s_{23}] - \rho s_{12}(c_{13}^2 - s_{13}^2); \quad (21b)$$

$$-\epsilon(c_{12}^2 - s_{12}^2) + \eta [s_{13}(c_{23}^2 - s_{23}^2)(c_{12}^2 - s_{12}^2) - 2c_{23}s_{23}c_{12}s_{12}(1 + s_{13}^2)] + 2\rho c_{13}s_{13}c_{12}s_{12}; \quad (21c)$$

$$4\epsilon c_{12}s_{12} + \eta [-4s_{13}s_{12}c_{12}(c_{23}^2 - s_{23}^2) - 2c_{23}s_{23}(c_{12}^2 - s_{12}^2)(1 + s_{13}^2)] + 2\rho c_{13}s_{13}(c_{12}^2 - s_{12}^2); \quad (21d)$$

$$-2\epsilon s_{12}c_{12} + \eta [2s_{13}c_{12}s_{12}(c_{23}^2 - s_{23}^2) + 2c_{23}s_{23}((c_{12}^2 - s_{12}^2) + c_{13}^2(1 + s_{12}^2))] + 2\rho c_{13}s_{13}(1 + s_{12}^2); \quad (21e)$$

$$2\epsilon s_{12}c_{12} + \eta [-2s_{13}c_{12}s_{12}(c_{23}^2 - s_{23}^2) + 2c_{23}s_{23}(c_{13}^2(1 + c_{12}^2) - (c_{12}^2 - s_{12}^2))] + 2\rho c_{13}s_{13}(1 + c_{12}^2). \quad (21f)$$

We have added (21f), which is not an independent relation, but the sum of (21d) and (21e); it expresses the violation in the universality of diagonal {11} and {33} currents.

### 3.2.1 A guiding calculation

Before doing the calculation in full generality, and to make a clearer difference with the neutrino case, we first do it in the limit where one neglects terms which are quadratic in the small quantities  $\theta_{13}$  and  $\rho$ . By providing simple intermediate formulæ, it enables in particular to suitably choose the signs which occur in equating the moduli of two quantities. Eqs.(21) become

$$\eta [s_{12}(c_{23}^2 - s_{23}^2) + 2s_{13}c_{12}c_{23}s_{23}] - \rho c_{12}; \quad (22a)$$

$$\eta [-c_{12}(c_{23}^2 - s_{23}^2) + 2s_{13}s_{12}c_{23}s_{23}] - \rho s_{12}; \quad (22b)$$

$$-\epsilon(c_{12}^2 - s_{12}^2) + \eta [s_{13}(c_{23}^2 - s_{23}^2)(c_{12}^2 - s_{12}^2) - 2c_{23}s_{23}c_{12}s_{12}]; \quad (22c)$$

$$4\epsilon c_{12}s_{12} - 2\eta [2s_{13}s_{12}c_{12}(c_{23}^2 - s_{23}^2) + c_{23}s_{23}(c_{12}^2 - s_{12}^2)]; \quad (22d)$$

$$-2\epsilon s_{12}c_{12} + 2\eta [s_{13}c_{12}s_{12}(c_{23}^2 - s_{23}^2) + c_{23}s_{23}(1 + c_{12}^2)]; \quad (22e)$$

$$2\epsilon s_{12}c_{12} + 2\eta [-s_{13}c_{12}s_{12}(c_{23}^2 - s_{23}^2) + c_{23}s_{23}(1 + s_{12}^2)]. \quad (22f)$$

The principle of the method is the same as before. From (22c) = -(22d)<sup>5</sup>, which expresses that the absence of non-diagonal {12} current is violated with the same strength as the universality of {11} and

<sup>5</sup>The (-) signs ensures that  $\tan(2\theta_{12}) \approx (+)1/2$ .



{22} currents, one gets  $\epsilon/\eta$  as a function of  $\theta_{12}, \theta_{23}, \theta_{13}$ <sup>6</sup>. This expression is plugged in the relation (22b) = (-)(22e)<sup>7</sup>, which expresses the same condition for the (2, 3) channel; from this, one extracts  $\rho/\eta$  as a function of  $\theta_{12}, \theta_{23}, \theta_{13}$ <sup>8</sup>. The expressions that have been obtained for  $\epsilon/\eta$  and  $\rho/\eta$  are then inserted into the third relation, |(22a)| = |(22f)|, which now corresponds to the (1, 3) channel. This last step yields a relation  $F_0(\theta_{12}, \theta_{23}, \theta_{13}) = 1$  between the three angles  $\theta_{12}, \theta_{23}, \theta_{13}$ .

It turns out that  $\frac{\partial F_0(\theta_{12}, \theta_{23}, \theta_{13})}{\partial \theta_{13}} = 0$ , such that, in this case, a condition between  $\theta_{12}$  and  $\theta_{23}$  alone eventually fulfills the three relations under concern

$$1 = \left| \frac{\text{viol}([11] = [22])}{\text{viol}([12] = 0 = [21])} \right| = \left| \frac{\text{viol}([22] = [33])}{\text{viol}([23] = 0 = [32])} \right| = \left| \frac{\text{viol}([11] = [33])}{\text{viol}([13] = 0 = [31])} \right| \Leftrightarrow \tilde{F}_0(\theta_{12}, \theta_{23}) = 1. \quad (25)$$

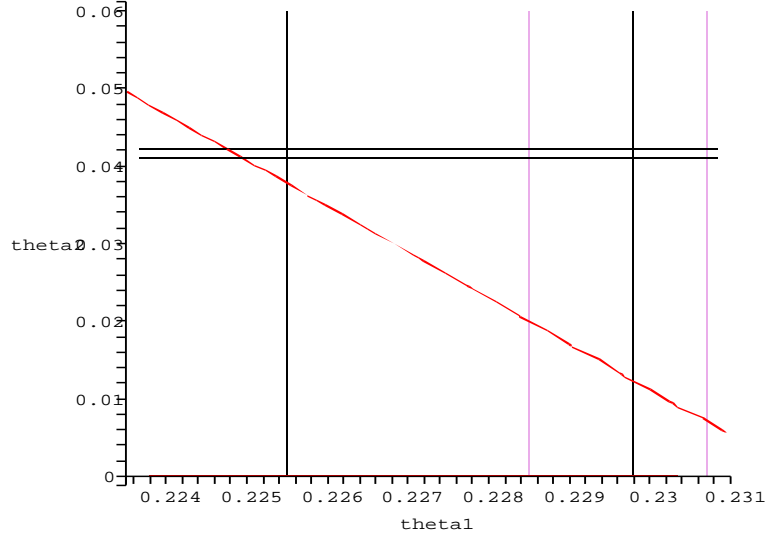


Fig. 2:  $\theta_{23}$  for quarks as a function of  $\theta_{12}$ ; neglecting terms quadratic in  $\theta_{13}$

$\theta_{23}$  is plotted on Fig. 2 as a function of  $\theta_{12}$ , together with the experimental intervals for  $\theta_{23}$  and  $\theta_{12}$  (the intervals for  $\theta_{12}$  come respectively from  $V_{ud}$  (eq. (18)) and  $V_{us}$  (eq. (19))).

The precision obtained is much better than in Fig. 1 since, in particular, for  $\theta_{23}$  within its experimental range, the discrepancy between the predicted  $\theta_{12}$  and its lower experimental limit coming from  $V_{us}$  is smaller than the two experimental intervals, and even smaller than their intersection.

### 3.2.2 The general solution

The principle for solving the general equations (21) is the same as above. One first uses the relation (21c) = (-) (21d) to determine  $\rho/\epsilon$  in terms of  $\eta/\epsilon$ . The result is plugged in the relation (21b) = (-) (21e), which

<sup>6</sup>

$$\frac{\epsilon}{\eta} = s_{13}(c_{23}^2 - s_{23}^2) + 2s_{23}c_{23} \frac{s_{12}c_{12} + c_{12}^2 - s_{12}^2}{4c_{12}s_{12} - (c_{12}^2 - s_{12}^2)}; \quad (23)$$

$\epsilon/\eta$  has a pole at  $\tan(2\theta_{12}) = 1/2$ , the predicted value of the Cabibbo angle for two generations.

<sup>7</sup>

<sup>8</sup>

There, again, the (-) sign has to be chosen so as to recover approximately (17).

$$\frac{\rho}{\eta} = 2c_{23}s_{23} \left[ s_{13} - c_{12} \left( 2 \frac{(c_{12}s_{12} + c_{12}^2 - s_{12}^2)}{4s_{12}c_{12} - (c_{12}^2 - s_{12}^2)} - \frac{1 + c_{12}^2}{c_{12}s_{12}} + \frac{1}{s_{12}} \frac{c_{23}^2 - s_{23}^2}{2s_{23}c_{23}} \right) \right]. \quad (24)$$

$\rho/\eta$  has a pole at  $\tan(2\theta_{12}) = 1/2$  and, for  $\theta_{13} = 0$ , it vanishes, as expected, when  $\theta_{12}$  and  $\theta_{23}$  satisfy the relation (17), which has been deduced for  $\tilde{\theta}_{13}(\equiv \theta_{13} + \rho) = 0 = \theta_{13}$ .

fixes  $\eta/\epsilon$ , and thus  $\rho/\epsilon$  as functions of  $(\theta_{12}, \theta_{23}, \theta_{13})$ . These expressions for  $\eta/\epsilon$  and  $\rho/\epsilon$  are finally plugged in the relation  $|(21a)| = |(21f)|$ , which provides a condition  $F(\theta_{12}, \theta_{23}, \theta_{13}) = 1$ . When it is fulfilled, the universality of each pair of diagonal neutral currents of mass eigenstates and the absence of the corresponding non-diagonal currents are violated with the same strength, in the three channels (1, 2), (2, 3) and (1, 3).

The results are displayed in Fig. 3;  $\theta_{23}$  is plotted as a function of  $\theta_{12}$  for  $\theta_{13} = 0.004$  and  $0.01$ . The present experimental interval is [7]

$$V_{ub} = \sin(\theta_{13}) \approx \theta_{13} \in [4 \cdot 10^{-3}, 4.6 \cdot 10^{-3}]. \quad (26)$$

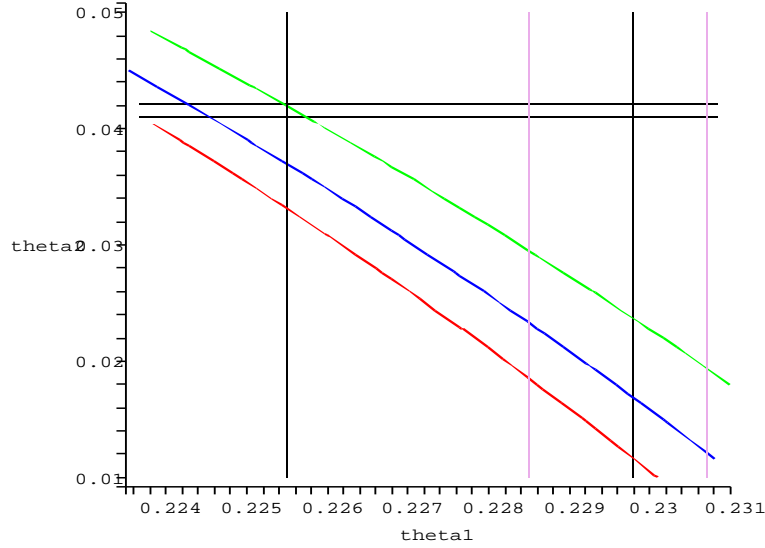


Fig. 3:  $\theta_{23}$  for quarks as a function of  $\theta_{12}$ , general case.  $\theta_{13} = 0$  (red),  $0.004$  (blue) and  $0.01$  (green)

We conclude that:

- \* The discrepancy between our predictions and experiments is smaller than the experimental uncertainty;
- \* a slightly larger value of  $\theta_{13}$  and/or slightly smaller values of  $\theta_{23}$  and/or  $\theta_{12}$  still increase the agreement between our predictions and experimental measurements;
- \* the determination of  $\theta_{12}$  from  $V_{us}$  seems preferred to that from  $V_{ud}$ .

Another confirmation of the relevance of our criterion is given in the next section concerning neutrino mixing angles.

## 4 A neutrino-like pattern; quark-lepton complementarity

In the “quark case”, we dealt with three “Cabibbo-like” angles. The configuration that we investigate here is the one in which  $\theta_{23}$  is, as observed experimentally [7], (close to) maximal, and  $\theta_{12}$  and  $\theta_{13}$  are Cabibbo-like (see subsection 2.5).

### 4.1 The case $\theta_{13} = 0 = \tilde{\theta}_{13}$

We explore the vicinity of this solution, slightly departing from the corresponding unitary mixing matrix, by considering that  $\tilde{\theta}_{12}$  now slightly differs from  $\theta_{12}$ , and  $\tilde{\theta}_{23}$  from its maximal value

$$\tilde{\theta}_{12} = \theta_{12} + \epsilon,$$

$$\theta_{23} = \pi/4 \quad , \quad \tilde{\theta}_{23} = \theta_{23} + \eta. \quad (27)$$

The l.h.s.'s of eqs. (7) (8) (9) (10) and (11) no longer vanish, and become respectively

$$-\frac{1}{2}\eta^2(s_{12} + \epsilon c_{12}), \quad (28a)$$

$$\frac{1}{2}\eta^2(c_{12} - \epsilon s_{12}), \quad (28b)$$

$$* - \eta s_{12} c_{12} + \epsilon(s_{12}^2 - c_{12}^2)(1 + \eta), \quad (28c)$$

$$* - \eta(c_{12}^2 - s_{12}^2) + 4\epsilon s_{12} c_{12}(1 + \eta), \quad (28d)$$

$$\eta(1 + c_{12}^2) - 2\epsilon s_{12} c_{12}(1 + \eta), \quad (28e)$$

showing by which amount the five conditions under scrutiny are now violated. Some care has to be taken concerning the accurateness of equations (28). Indeed, we imposed a value of  $\theta_{13}$  which is probably not the physical one (even if close to). It is then reasonable to consider that channel (1, 2) is the less sensitive to this approximation and that, accordingly, of the five equations above, (28c) and (28d), marked with an “\*”, are the most accurate<sup>9</sup>.

The question: is there a special value of  $\theta_{12} = \tilde{\theta}_{12}$  Cabibbo-like for which small deviations ( $\epsilon, \eta$ ) from unitarity entail equal strength violations of

- \* the absence of  $\{12\}, \{21\}$  non-diagonal neutral currents;
- \* the universality of  $\{11\}$  and  $\{22\}$  neutral currents ?

gets then a simple answer

$$s_{12}c_{12} = c_{12}^2 - s_{12}^2 \Rightarrow \tan(2\theta_{12}) = 2. \quad (29)$$

We did not take into account the terms proportional to  $\epsilon$  because we assumed that the mass splittings between the first and second generations (from which the lack of unitarity originates) are much smaller than the ones between the second and the third generation<sup>10</sup>.

In the case of two generations, only  $\epsilon$  appears, and one immediately recovers from (28c) and (28d) the condition fixing  $\tan(2\theta_c) = 1/2$  for the Cabibbo angle.

Accordingly, the same type of requirement that led to a value of the Cabibbo angle for two generations very close to the observed value leads, for three generations, to a value of the first mixing angle satisfying the quark-lepton complementarity relation (3) [6].

The values of  $\theta_{12}$  and  $\theta_{23}$  determined through this procedure are very close to the observed neutrino mixing angles [7].

Though we only considered the two equations that are *a priori* the least sensitive to our choice of a vanishing third mixing angle (which is not yet confirmed experimentally), it is instructive to investigate the sensitivity of our solution to a small non-vanishing value of  $\theta_{13}$ . This is done in Appendix C in which, for this purpose, we made the simplification  $\tilde{\theta}_{13} \approx \theta_{13}$ . It turns out that the terms proportional to

<sup>9</sup>The limitation of this approximation also appears in the fact that (28b), of second order in  $\eta$ , is not compatible with (28e), which is of first order.

<sup>10</sup>Since the three angles play *a priori* symmetric roles, the simultaneous vanishing of  $\theta$  and  $\tilde{\theta}$ , which we demonstrated for  $\theta_{13}$  and  $\tilde{\theta}_{13}$  (see Appendix A), should also occur for the other angles. Two competing effects accordingly contribute to the magnitude of the parameters  $\epsilon, \eta \dots$ : on one hand, they should be proportional to (some power of) the corresponding  $\theta$ , and, on the other hand, one reasonably expects them to increase with the mass splitting between the fermions mixed by this  $\theta$ . So, in the quark sector, that the violation of unitarity should be maximal for  $\theta_{13}$  is not guaranteed since the corresponding mixing angle is also very small (as expected from hierarchical mixing matrices [10]). A detailed investigation of this phenomenon is postponed to a further work. In the neutrino sector, however, since  $\theta_{23}$  is maximal (large), the assumption that the mass splitting between the second and third generation is larger than between the first and second is enough to guarantee  $\epsilon \ll \eta$ .

$s_{13}$  in the two equations  $[12] = 0 = [21]$  and  $|[11]| = |[22]|$  are also proportional to  $(c_{23}^2 - s_{23}^2)$ , such that our solution with  $\theta_{23}$  maximal is very stable with respect to a variation of  $\theta_{13}$  around zero. This may of course not be the case for the other three equations, which are expected to be more sensitive to the value of  $\theta_{13}$ .

## 4.2 Prediction for $\theta_{13}$

We now consider, like we did for quarks, the general case  $\theta_{13} \neq 0 \neq \tilde{\theta}_{13}(\rho \neq 0)$ ,  $\tilde{\theta}_{12} \neq \theta_{12}(\epsilon \neq 0)$ ,  $\tilde{\theta}_{23} \neq \theta_{23}(\eta \neq 0)$ , while assigning to  $\theta_{12}$  and  $\theta_{23}$  their values predicted in subsection 4.1.

We investigate the eight different relations between  $\theta_{12}$ ,  $\theta_{23}$  and  $\theta_{13}$  which originate from the  $2 \times 2 \times 2$  possible sign combinations in the conditions (25) (the r.h.s. is now replaced by a condition  $F(\theta_{12}, \theta_{23}, \theta_{13}) = 1$  involving the three mixing angles), where each modulus can be alternatively replaced by “+” or “-”.

Among the solutions found for  $\theta_{13}$ , only two (up to a sign) satisfy the very loose experimental bound

$$\sin^2(\theta_{13}) \leq 0.1. \quad (30)$$

They correspond respectively to the sign combinations  $(+/-/-)$ ,  $(+/+/+)$ ,  $(-/+/+)$  and  $(-/-/-)$

$$\begin{aligned} \theta_{13} &= \pm 0.2717 \quad , \quad \sin^2(\theta_{13}) = 0.072, \\ \theta_{13} &= \pm 5.7 \cdot 10^{-3} \quad , \quad \sin^2(\theta_{13}) = 3.3 \cdot 10^{-5}. \end{aligned} \quad (31)$$

The most recent experimental bounds can be found in [11]. They read

$$\sin^2(\theta_{13}) \leq 0.05, \quad (32)$$

which only leaves the smallest solution in (31).

Future experiments will confirm, or infirm, for neutrinos, the properties that we have shown to be satisfied with an impressive accuracy by quark mixing angles.

## 5 Comments, open questions and problems

### 5.1 How close are mixing matrices to unitarity?

An important characteristic of the conditions that fix the mixing angles is that they do not depend on the strength of the violation of the two properties under concern, namely, the absence of non-diagonal neutral currents and the universality of the diagonal ones in the space of mass eigenstates. Since only their ratio is concerned, each violation can be infinitesimally small.

This is fortunate since we have not yet been able to calculate the magnitude of the violation of the unitarity of the mixing matrices from, for example, mass ratios. The issue, for fundamental particles, turns indeed to be much more difficult conceptually than it was for composite particles like neutral kaons, for which standard Feynman diagrams provided the estimate  $\epsilon_L - \epsilon_S \approx 10^{-17}$  for the difference of the  $CP$  violating parameters of  $K_L$  and  $K_S$  mesons [1]. This problem is under investigation.

### 5.2 The measured mixing angles are those of charged currents

The results that have been exposed are valid for fermions of both electric charges. They concern the mixing angles which parametrize

- \* for quarks, the mixing matrix  $K_u$  of  $u$ -type quarks as well as  $K_d$  of  $d$ -type quarks;
- \* for leptons, the mixing matrix  $K_\nu$  of neutrinos as well as that of charged leptons  $K_\ell$ ,

and we have shown that our approach allows to obtain on purely theoretical grounds the values of the mixing angles which are experimentally determined.

However, there arises a non negligible problem : indeed, the matrix elements which are measured correspond to charged currents, that is, to the product of the two corresponding mixing matrices,  $K_u^\dagger K_d$  for quarks and  $K_\ell^\dagger K_\nu$  for leptons. Thus, they are *a priori* related to an entanglement of the mixing angles of quarks (or leptons) of different charges.

Nevertheless, what happens in the common approach of the SM is the following : in the expression of the charged current, the product  $K_u^\dagger K_d$  (or  $K_\ell^\dagger K_\nu$ ) is applied as a whole "to the right" and thus practically redefined as the mixing matrix of the type-d fermions (e.g. neutrinos) ; it takes *de facto* the place of  $K_d$  (or  $K_\nu$ ), such that the angles which parametrize it are defined as the mixing angles for the type-d fermions. Such a procedure is equivalent to assuming that only one of the two types of fermions undergoes a mixing, while the other has its mass and flavour eigenstates aligned. Though it is difficult to agree with this opportunistic statement (since the two species should play *a priori* similar roles), our results tend to confirm it (see also Appendix D).

### 5.3 Why are quarks different from leptons?

The generality of our procedure and, in particular, its being independent of the type of fermions (quarks or leptons) raises another well-known but still unanswered question : why is the mixing pattern of leptons so different from that of quarks ?

A sketch of solution could be provided by considering configurations of the two quark mixing matrices  $K_u$  and  $K_d$  which, on one hand, reproduce the angles of the CKM matrix  $K = K_u^\dagger K_d$  as they are generally accounted for in data books [7]<sup>11</sup>, and, on the other hand, are both leptonic-like. A simple example of such a configuration is given by the symmetrical pattern:  $K_u = (\pi/4, \theta_\nu, \theta)$ ,  $K_d = (\theta_\nu, \pi/4, \varphi)$ , with  $\tan(2\theta_\nu) = 2$ . Since, according to (3),  $\theta_c = \pi/4 - \theta_\nu$ , the Cabibbo angle appears in  $K_u^\dagger K_d$ <sup>12</sup>.

<sup>11</sup>in which the CKM angles are generally attributed to the sole mixing of d-type quarks

<sup>12</sup>Consider, for example, the simplified case  $\theta = 0 = \varphi$

$$K_u = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_2 & s_2 \\ 0 & -s_2 & c_2 \end{pmatrix} \times \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix} \times \begin{pmatrix} c_1 & s_1 & 0 \\ -s_1 & c_1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (33)$$

and

$$K_d = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_1 & s_1 \\ 0 & -s_1 & c_1 \end{pmatrix} \times \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix} \times \begin{pmatrix} c_2 & s_2 & 0 \\ -s_2 & c_2 & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (34)$$

One has

$$K_u^\dagger K_d = \begin{pmatrix} c_1 c_2 + s_1 s_2 \cos(\theta_1 - \theta_2) & c_1 s_2 - s_1 c_2 \cos(\theta_1 - \theta_2) & -s_1 \sin(\theta_1 - \theta_2) \\ s_1 c_2 - c_1 s_2 \cos(\theta_1 - \theta_2) & s_1 s_2 + c_1 c_2 \cos(\theta_1 - \theta_2) & c_1 \sin(\theta_1 - \theta_2) \\ s_2 \sin(\theta_1 - \theta_2) & -c_2 \sin(\theta_1 - \theta_2) & \cos(\theta_1 - \theta_2) \end{pmatrix} \stackrel{(\theta_1 - \theta_2) \text{ small}}{\approx} \begin{pmatrix} \cos(\theta_1 - \theta_2) & -\sin(\theta_1 - \theta_2) & -s_1 \sin(\theta_1 - \theta_2) \\ \sin(\theta_1 - \theta_2) & \cos(\theta_1 - \theta_2) & c_1 \sin(\theta_1 - \theta_2) \\ s_2 \sin(\theta_1 - \theta_2) & -c_2 \sin(\theta_1 - \theta_2) & \cos(\theta_1 - \theta_2) \end{pmatrix}. \quad (35)$$

For  $\theta_1 = \pi/4$  and  $\theta_2 = \theta_\nu$ ,  $\tan(2\theta_\nu) = 2$ , the Cabibbo angle  $(\theta_1 - \theta_2)$  naturally appears. This simplified case of course needs improvement since  $V_{ub}, V_{td}, V_{cb}, V_{ts}$  are far from being suitably predicted.

## 5.4 A multiscale problem

Recovery of the present results by perturbative techniques (Feynman diagrams) stays, as already mentioned, an open issue.

All the subtlety of the problem lies in the inadequacy of using a single constant mass matrix; because non-degenerate coupled systems are multiscale systems, as many mass matrix should be introduced as there are poles in the (matricial) propagator [13]<sup>13</sup>.

The existence of different scales makes the use of an “on-shell” renormalized Lagrangian [14] hazardous, because each possible renormalization scale optimizes the calculation of parameters at this scale, while, for other scales, one has to rely on renormalization group equations.

Unfortunately, these equations have only been approximately solved with the simplifying assumption that the renormalized mass matrices are hermitian<sup>14</sup> and that the renormalized mixing matrices are unitary [14]. Performing the same job dropping these hypotheses looks rather formidable and beyond the scope of the present work. It also unfortunately turns out that, as far as the Yukawa couplings are concerned, the expressions that have been obtained at two loops for their  $\beta$  functions (which start the evolution only up from the top quark mass) [15] have poles in  $(m_i - m_j)$ , which makes them inadequate for the study of subsystems with masses below the top quark mass.

## 5.5 Using a $q^2$ -dependent renormalized mass matrix

Departure from the inappropriate Wigner-Weisskopf approximation can also be done by working with an effective renormalized  $q^2$ -dependent mass matrix  $M(q^2)$ . It however leads to similar conclusions as the present approach.

Its eigenvalues are now  $q^2$ -dependent, and are determined by the equation  $\det[M(q^2) - \lambda(q^2)] = 0$ <sup>15</sup>. Let them be  $\lambda_1(q^2) \dots \lambda_n(q^2)$ . The physical masses satisfy the  $n$  self-consistent equations  $q^2 = \lambda_{1\dots n}(q^2)$ , such that  $m_1^2 = \lambda_1(m_1^2) \dots m_n^2 = \lambda_n(m_n^2)$ . At each  $m_i^2$ ,  $M(m_i^2)$  has  $n$  eigenvectors, but only one corresponds to the physical mass eigenstate; the others are “spurious” states [1]. Even if the renormalized mass matrix is hermitian at any given  $q^2$ , the physical mass eigenstates corresponding to different  $q^2$  belong to as many different orthonormal sets of eigenstates and thus, in general, do not form an orthonormal set. The discussion proceeds like in the core of the paper.

Determining the exact form of the renormalized mass matrix could accordingly be a suitable way to recover our predictions via perturbative techniques (like was done in [1] for the quantitative prediction of the ratio  $\epsilon_S/\epsilon_L$ ). As already mentioned, the difficulty is that hermiticity assumptions should be dropped, which open the possibility of departing from the unitarity of the mixing matrix. This is currently under investigation.

## 6 Conclusion and perspective

This work does not, obviously, belong to what is nowadays referred to as “Beyond the Standard Model”, since it does not incorporate any “new physics” such as supersymmetry, “grand unified theories (GUT)”

<sup>13</sup>In QFT, as opposed to a Quantum Mechanical treatment (in which a single constant mass matrix is introduced – this is the Wigner-Weisskopf approximation–), a constant mass matrix can only be introduced in a linear approximation to the inverse propagator in the vicinity of each of its poles. When several coupled states are concerned, the (matricial) propagator having several poles, as many (constant) mass matrices should be introduced; only one of the eigenstates of each of these mass matrices corresponds to a physical (mass) eigenstate.

<sup>14</sup>One can go to hermitian mass matrices by rotating right-handed fermions *as far as they are not coupled*; however, at 3 loops, the charged weak currents also involve right-handed fermions, which cannot be anymore freely rotated.

<sup>15</sup>This is the simple case of a normal mass matrix, which can be diagonalized by a single ( $q^2$ -dependent) unitary matrix. When it is non-normal, the standard procedure uses a bi-unitary diagonalization, in which case the so-called “mass eigenstates” are non longer the eigenstates of the mass matrix.

or extra-dimensions. However it does not strictly lie within the SM either, even if it is very close to. Of course, it shares with the latter its general framework (mathematical background and physical content), and also borrows from it the two physical conditions of universality for diagonal neutral currents and absence of FCNC's, which play a crucial role in the process. But, on the basis of the most general arguments of QFT, we make a decisive use of the essential non-unitarity of the mixing matrices, whereas only unitary matrices are present in the SM. This property may be considered, in the SM, as an "accidental" characteristic of objects which are intrinsically non-unitary.

The mixing angles experimentally observed get constrained in the vicinity of this "standard" situation, a slight departure from which being due to mass splittings. Hence our approach can be considered to explore the "Neighborhood of the Standard Model", which is likely to exhibit low-energy manifestations of physics "Beyond the Standard Model".

While common approaches limit themselves to guessing symmetries for the mass matrix (see for example [16] and references therein), we showed that special patterns are instead likely to reveal themselves in the violation of some (wrongly) intuitive properties<sup>16</sup>. In each given  $(i, j)$  channel of *mass eigenstates*, the characteristic pattern that emerges is that two *a priori* different properties are violated with the same strength, which can even be arbitrarily small: the absence of  $\{ij\}$  and  $\{ji\}$  non-diagonal neutral currents and the universality of diagonal neutral currents  $\{ii\} = \{jj\}$ .

The way of proceeding exposed here is reminiscent of Gell-Mann's approach to  $SU(3)$  flavour symmetry [4], in which the interesting structures were to be looked for in its violation. The equivalent here would be that "symmetries" relevant for flavour physics should not be looked for, or implemented, at the level of the mass matrices and Yukawa couplings, but at the level of *deviations* from properties which are usually taken for granted.

To conclude, the present work demonstrates that flavor physics satisfies very simple criteria which had been, up to now, unnoticed. Strong arguments have been presented in both the quark and leptonic sectors, which will be further tested when the third mixing angle of neutrinos is accurately determined. These features nature offers to our perspicacity and to our quest for still hidden symmetries.

*Acknowledgments:* Discussions with A. Djouadi, J. Orloff and M.I. Vysotsky are gratefully acknowledged.

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<sup>16</sup>For a (constant unique) mass matrix, unitarity of the mixing matrix has indeed always been linked with the unitarity of the theory. In the case of coupled systems, this fundamental feature is instead linked to the property that, at any given  $q^2$ , the renormalized  $q^2$ -dependent mixing matrix linking flavor states to the ( $q^2$ -dependent) eigenvectors of the mass matrix at this given  $q^2$  is unitary. This set of eigenvectors however never contains more than one physical mass eigenstate.

## Appendix

### A $\tilde{\theta}_{13} = 0 \Rightarrow \theta_{13} = 0$

Using the notations of section 2, we start with the following system of equations:

$$\frac{[11] + [22]}{2} = [33] \Leftrightarrow s_{13}^2 + s_{23}^2 + \tilde{c}_{23}^2 = 1; \quad (36a)$$

$$[11] = [22] \Leftrightarrow c_{13}^2 \cos(2\theta_{12}) = (c_{23}^2 + \tilde{s}_{23}^2) \cos(2\tilde{\theta}_{12}); \quad (36b)$$

$$[12] = 0 = [21] \Leftrightarrow c_{13}^2 \sin(2\theta_{12}) = (c_{23}^2 + \tilde{s}_{23}^2) \sin(2\tilde{\theta}_{12}); \quad (36c)$$

$$[13] = 0 = [31] \Leftrightarrow \tilde{s}_{12} \left( \sin(2\theta_{23}) - \sin(2\tilde{\theta}_{23}) \right) = c_{12} \sin(2\theta_{13}); \quad (36d)$$

$$[23] = 0 = [32] \Leftrightarrow \tilde{c}_{12} \left( \sin(2\tilde{\theta}_{23}) - \sin(2\theta_{23}) \right) = s_{12} \sin(2\theta_{13}). \quad (36e)$$

From equation (36a), we have  $c_{23}^2 + \tilde{s}_{23}^2 \neq 0$ , which entails  $c_{13}^2 \neq 0$ <sup>17</sup>. Let us study the consequence on the two equations (36b) and (36c).

• the two sides of (36b) vanish for  $\cos(2\theta_{12}) = 0 = \cos(2\tilde{\theta}_{12})$ , i.e.  $\theta_{12} = \frac{\pi}{4}[\frac{\pi}{2}] = \tilde{\theta}_{12}$ . (36c) then gives  $c_{13}^2 = c_{23}^2 + \tilde{s}_{23}^2$ , which, associated with (36a), yields the following solution<sup>18</sup>:  $\theta_{13} = 0[\pi]$  and  $\tilde{\theta}_{23} = \pm\theta_{23}[\pi]$ .

• the two sides of (36c) vanish for  $\sin(2\theta_{12}) = 0 = \sin(2\tilde{\theta}_{12}) = 0$ , i.e.  $\theta_{12} = 0[\frac{\pi}{2}] = \tilde{\theta}_{12}$ . (36b) gives then  $c_{13}^2 = c_{23}^2 + \tilde{s}_{23}^2$ , hence, like previously,  $\theta_{13} = 0[\pi]$  and  $\tilde{\theta}_{23} = \pm\theta_{23}[\pi]$ .

• in the other cases we can calculate the ratio (36b) / (36c), which gives  $\tan(2\theta_{12}) = \tan(2\tilde{\theta}_{12})$ , hence  $\theta_{12} = \tilde{\theta}_{12}[\pi]$  or  $\theta_{12} = \frac{\pi}{2} + \tilde{\theta}_{12}[\pi]$ :

\*  $\theta_{12} = \frac{\pi}{2} + \tilde{\theta}_{12}[\pi]$  implies for (36b)(36c)  $c_{13}^2 = -c_{23}^2 - \tilde{s}_{23}^2$ , which, together with (36a) ( $c_{13}^2 = s_{23}^2 + \tilde{c}_{23}^2$ ), gives a contradiction :  $2 = 0$ :

\*  $\theta_{12} = \tilde{\theta}_{12}(\neq 0)[\pi]$  implies, like previously,  $c_{13}^2 = c_{23}^2 + \tilde{s}_{23}^2$ , which gives, when combined with (36a):  $\theta_{13} = 0[\pi]$  and  $\tilde{\theta}_{23} = \pm\theta_{23}[\pi]$ .

Hence, it appears that whatever the case, the solution gives rise to  $\theta_{13} = 0[\pi]$ .

Let us now look at (36d) and (36e). Since  $\theta_{13} = 0$ , the two r.h.s.'s vanish, and we obtain the twin equations  $\tilde{s}_{12}(\sin(2\theta_{23}) - \sin(2\tilde{\theta}_{23})) = 0$  and  $\tilde{c}_{12}(\sin(2\theta_{23}) - \sin(2\tilde{\theta}_{23})) = 0$ , which, together, imply  $\sin(2\theta_{23}) = \sin(2\tilde{\theta}_{23})$ . It follows that, either  $\theta_{23} = \tilde{\theta}_{23}[\pi]$  or  $\theta_{23} = \frac{\pi}{2} - \tilde{\theta}_{23}[\pi]$ ;

\*  $\theta_{23} = \tilde{\theta}_{23}[\pi]$  matches the result of the previous discussion in the “+” case, whereas, in the “-” case, the matching leads to  $\theta_{23} = \tilde{\theta}_{23} = 0$ , which is to be absorbed as a particular case in the “+” configuration;

\*  $\theta_{23} = \frac{\pi}{2} - \tilde{\theta}_{23}[\pi]$  matches the result of the previous discussion in the “+” configuration, in which case it leads to  $\theta_{23} = \tilde{\theta}_{23} = \frac{\pi}{4}[\frac{\pi}{2}]$ , i.e. maximal mixing between the fermions of the second and third generations.

<sup>17</sup>Indeed, let us suppose that  $c_{13}$  vanishes. Then  $\cos(2\tilde{\theta}_{12})$  and  $\sin(2\tilde{\theta}_{12})$  must vanish simultaneously, which is impossible.

<sup>18</sup>
$$\begin{cases} c_{13}^2 = c_{23}^2 + \tilde{s}_{23}^2 \\ s_{13}^2 + s_{23}^2 + \tilde{c}_{23}^2 = 1 \end{cases} \implies \begin{cases} s_{23}^2 + \tilde{c}_{23}^2 = 1 \\ s_{13}^2 = 0 \end{cases}$$



## B $(\theta_{12}, \theta_{23})$ solutions of eqs. (7) (8) (9) (10) (11) for $\theta_{13} = 0 = \tilde{\theta}_{13}$

Excluding  $\tilde{\theta}_{12} = 0$ , (13a) and (13b) require  $\sin(2\theta_{23}) = \sin(2\tilde{\theta}_{23}) \Rightarrow \tilde{\theta}_{23} = \theta_{23} + k\pi$  or  $\tilde{\theta}_{23} = \pi/2 - \theta_{23} + k\pi$ .

• for  $\tilde{\theta}_{23} = \theta_{23} + k\pi$  Cabibbo-like,

(13c) requires  $\sin(2\theta_{12}) = \sin(2\tilde{\theta}_{12}) \Rightarrow \tilde{\theta}_{12} = \theta_{12} + n\pi$  or  $\tilde{\theta}_{12} = \pi/2 - \theta_{12} + n\pi$ ;

(13d) requires  $\cos(2\theta_{12}) = \cos(2\tilde{\theta}_{12}) \Rightarrow \tilde{\theta}_{12} = \pm\theta_{12} + p\pi$ ;

(13e) requires  $s_{12}^2 + \tilde{c}_{12}^2 - 1 = 0 \Rightarrow \tilde{\theta}_{12} = \pm\theta_{12} + r\pi$ .

The solutions of these three equations are  $\theta_{12} = \tilde{\theta}_{12} + k\pi$  Cabibbo-like or  $\theta_{12} = \pi/4 + q\pi/2$  maximal.

• for  $\tilde{\theta}_{23} = \pi/2 - \theta_{23} + k\pi$ ,

(13c) requires  $s_{12}c_{12} = 2c_{23}^2\tilde{s}_{12}\tilde{c}_{12}$ ;

(13d) requires  $c_{12}^2 - s_{12}^2 = 2c_{23}^2(\tilde{c}_{12}^2 - \tilde{s}_{12}^2)$ ;

(13e) requires  $s_{12}^2 + 2c_{23}^2\tilde{c}_{12}^2 - 2s_{23}^2 = 0$ .

The first two conditions yield  $\tan(2\theta_{12}) = \tan(2\tilde{\theta}_{12}) = 2c_{23}^2 \Rightarrow \tilde{\theta}_{12} = \theta_{12} + k\pi/2 + n\pi$ , which entails  $2c_{23}^2 = 1 \Rightarrow \theta_{23} = \pm\pi/4 + p\pi/2$  maximal;  $\tilde{\theta}_{23}$  is then maximal, too, and the third condition is automatically satisfied.

$\tilde{\theta}_{12} = \theta_{12} + n\pi$  is Cabibbo-like, while, for  $\tilde{\theta}_{12} = \theta_{12} + (2k+1)\pi/2$ , the second condition becomes  $(c_{12}^2 - s_{12}^2) = 0$ , which means that  $\theta_{12}$  must be maximal.

## C Sensitivity of the neutrino solution to a small variation of $\theta_{13}$

If one allows for a small  $\theta_{13} \approx \tilde{\theta}_{13}$ , (9) and (10) become

$$\begin{aligned} & -2\eta s_{12}c_{12}s_{23}c_{23} + \epsilon(s_{12}^2 - c_{12}^2) + \eta s_{13}(c_{23}^2 - s_{23}^2)(c_{12}^2 - s_{12}^2), \\ & -2\eta s_{23}c_{23}(c_{12}^2 - s_{12}^2) + 4\epsilon s_{12}c_{12} - 2\eta s_{13}(c_{23}^2 - s_{23}^2)(2s_{12}c_{12} + \epsilon(c_{12}^2 - s_{12}^2)). \end{aligned} \quad (37)$$

For  $\theta_{23}, \tilde{\theta}_{23}$  maximal, the dependence on  $\theta_{13}$  drops out.

## D Charged weak currents

Charged weak currents can be written in six different forms that are all strictly equivalent, but nonetheless refer to different physical pictures. As an example, for two generations of leptons :

$$\begin{aligned} & \overline{\begin{pmatrix} \nu_{ef} \\ \nu_{\mu f} \end{pmatrix}} W_{\mu}^{+} \gamma_L^{\mu} \begin{pmatrix} e_f^- \\ \mu_f^- \end{pmatrix} = \overline{\begin{pmatrix} \nu_{em} \\ \nu_{\mu m} \end{pmatrix}} W_{\mu}^{+} \gamma_L^{\mu} K_{\nu}^{\dagger} K_{\ell} \begin{pmatrix} e_m^- \\ \mu_m^- \end{pmatrix} \\ & = \overline{\begin{pmatrix} \nu_{ef} \\ \nu_{\mu f} \end{pmatrix}} W_{\mu}^{+} \gamma_L^{\mu} \left[ K_{\ell} \begin{pmatrix} e_m^- \\ \mu_m^- \end{pmatrix} \right] = \overline{\begin{pmatrix} \nu_{em} \\ \nu_{\mu m} \end{pmatrix}} \left[ K_{\nu} \begin{pmatrix} \nu_{ef} \\ \nu_{\mu f} \end{pmatrix} \right] W_{\mu}^{+} \gamma_L^{\mu} \begin{pmatrix} e_f^- \\ \mu_f^- \end{pmatrix} \\ & = \overline{\begin{pmatrix} \nu_{em} \\ \nu_{\mu m} \end{pmatrix}} W_{\mu}^{+} \gamma_L^{\mu} \left[ K_{\nu}^{\dagger} \begin{pmatrix} e_f^- \\ \mu_f^- \end{pmatrix} \right] = \overline{\begin{pmatrix} \nu_{ef} \\ \nu_{\mu f} \end{pmatrix}} \left[ K_{\ell}^{\dagger} \begin{pmatrix} \nu_{em} \\ \nu_{\mu m} \end{pmatrix} \right] W_{\mu}^{+} \gamma_L^{\mu} \begin{pmatrix} e_m^- \\ \mu_m^- \end{pmatrix}. \end{aligned} \quad (38)$$

In the case where one of the  $SU(2)$  partners, for example the charged lepton, is undoubtedly a mass eigenstate<sup>19</sup>, the last expression of (38) shows that it is coupled to the so-called *electronic and muonic*

<sup>19</sup>This is the case inside the sun [5] where, because of the limited available energy, only massive electrons can be produced, and also in the detection process of neutrinos on earth, which always proceeds via charged currents and the detection of produced physical (mass eigenstates) charged leptons.

*neutrinos*

$$\begin{aligned}\nu_e &= (K_\ell^\dagger K_\nu)_{11}\nu_{em} + (K_\ell^\dagger K_\nu)_{12}\nu_{\mu m} = K_{\ell,11}^\dagger\nu_{ef} + K_{\ell,12}^\dagger\nu_{\mu f}, \\ \nu_\mu &= (K_\ell^\dagger K_\nu)_{21}\nu_{em} + (K_\ell^\dagger K_\nu)_{22}\nu_{\mu m} = K_{\ell,21}^\dagger\nu_{ef} + K_{\ell,22}^\dagger\nu_{\mu f}.\end{aligned}\quad (39)$$

The latter are neither flavour eigenstates, nor mass eigenstates, but a third kind of neutrinos, precisely defined as the ones which couple to electron and muon mass eigenstates in the weak charged currents

$$\overline{\begin{pmatrix} \nu_e \\ \nu_\mu \end{pmatrix}} W_\mu^+ \gamma_L^\mu \begin{pmatrix} e_m^- \\ \mu_m^- \end{pmatrix}.\quad (40)$$

It then occurs that neutral currents of both charged leptons mass eigenstates and  $\nu_e, \nu_\mu, \nu_\tau$  make appear the product  $K_\ell^\dagger K_\ell$ , which only contains the mixing matrix of the former. It may accordingly happen that, if the properties that we have implemented in this work for both types of mass eigenstates are implemented now for  $(\nu_e, \nu_\mu, \nu_\tau)$  instead of  $(\nu_{em}, \nu_{\mu m}, \nu_{\tau m})$  (and still for charged lepton mass eigenstates) the only mixing angles that get constrained are the ones of *charged leptons* and not the ones of neutrinos.

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