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Introduction to flavour and CP violation in the Standard Model and beyond

Irréversibilité microscopique, violation de CP et saveurs de quarks : Approche dans le cadre du Modèle Standard et au-delà

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ABSTRACT

We define the Cabibbo–Kobayashi–Maskawa matrix and briefly discuss its present determination. We recall the formalism for heavy-meson mixing and CP violation in meson decays. Finally we extend our discussion of flavour and CP violation beyond the Standard Model.

RÉSUMÉ

Nous définissons la matrice de Cabibbo–Kobayashi–Maskawa et nous présentons brièvement la connaissance actuelle de ces paramètres. Nous rappelons ensuite le formalisme du mélange des mésons neutres ainsi que celui de la violation de CP dans leur désintégration. Finalement nous discutons la violation de la saveur et de la symétrie de CP au delà du Modèle Standard.

1. The Cabibbo–Kobayashi–Maskawa matrix

In the Standard Model (SM), flavour non-diagonal and CP-violating couplings only appear in the Cabibbo–Kobayashi–Maskawa (CKM) matrix V [1] entering the charged current weak interactions

$$\mathcal{L}^{cc} = \frac{g_2}{\sqrt{2}} \bar{u}^i_L \gamma_\mu V_{ij} d^j_L W^\mu + \text{H.c.} \tag{1}$$

where u^i_L, (d^i_L) are the left-handed components of the up-type (down-type) quark fields $u^i = (u, c, t)$ $(d^j = (d, s, b))$, W is the W boson field and g_2 is the $SU(2)$ coupling constant.

V is related as $V = U_L D^\dagger_L$ to the matrices $U_{L,R}$, $D_{L,R}$ relating weak interaction eigenstates to mass eigenstates through a unitary transformation in flavour space. A generic 3×3 unitary matrix can be represented using 3 Euler angles and 6 phases, but 5 of the latter can be removed by re-phasing the quark fields. Therefore V can be parameterized by 3

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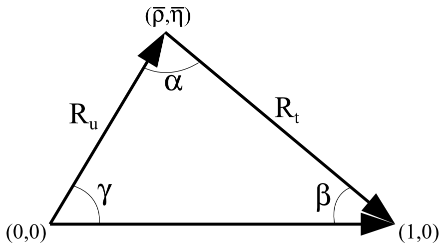

Fig. 1. The unitarity triangle.

angles $\theta_{12}, \theta_{23}, \theta_{13} \in [0, \pi/2]$ and 1 phase $\delta \in [-\pi, \pi)$. Notice that the choice of the parameters is not unique. In the parameterization used by the PDG [2], the CKM matrix reads

$$V \equiv \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}\,e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{13}s_{23}\,e^{i\delta} & c_{12}c_{23} - s_{12}s_{13}s_{23}\,e^{i\delta} & c_{13}s_{23} \\ s_{12}s_{23} - c_{12}s_{13}c_{23}\,e^{i\delta} & -c_{12}s_{23} - s_{12}s_{13}c_{23}\,e^{i\delta} & c_{13}c_{23} \end{pmatrix} \tag{2}$$

where $s_{ij} = \sin\theta_{ij}$ and $c_{ij} = \cos\theta_{ij}$.

The CP-violating phase δ is not a physical parameter as it depends on the phase conventions of the quark fields. An invariant condition for CP violation in the SM is given by [3]

$$J\,(m_u^2 - m_c^2)(m_c^2 - m_t^2)(m_t^2 - m_u^2)(m_d^2 - m_s^2)(m_s^2 - m_b^2)(m_b^2 - m_d^2) \neq 0 \tag{3}$$

where the Jarlskog invariant $J = |\,\mathrm{Im}(V_{ij}V_{kl}V_{il}^*V_{kj}^*)|$ (for all $i \neq k$ and $j \neq l$) is independent of the convention chosen for the CKM matrix. In terms of the parameters used in Eq. (2), $J = s_{12}s_{13}s_{23}c_{12}c_{13}^2c_{23}\sin\delta$. Therefore CP violation in the SM requires $\delta \neq 0, \pi$, non-degenerate masses in the up and down sectors and non-trivial mixing angles $\theta_{ij} \neq 0, \pi/2$.

The unitarity of V implies 9 conditions involving products of rows (or columns) of the form $\sum_k V_{ik}V_{jk}^* = \delta_{ij}$. The off-diagonal relations are called triangular as they define triangles in the complex plane. Remarkably, the area of all these triangles is a constant equal to $J/2$ and thus a measure of CP violation in the SM. One of them, of particular phenomenological interest, is referred to as the Unitarity Triangle (UT):

$$V_{ud}V_{ub}^* + V_{cd}V_{cb}^* + V_{td}V_{tb}^* = 0 \tag{4}$$

The UT can be rewritten in the normalized form

$$R_t e^{-i\beta} + R_u e^{i\gamma} = 1 \tag{5}$$

with

$$R_t = \left|\frac{V_{td}V_{tb}^*}{V_{cd}V_{cb}^*}\right|, \qquad R_u = \left|\frac{V_{ud}V_{ub}^*}{V_{cd}V_{cb}^*}\right|, \qquad \beta = \arg\left(-\frac{V_{cd}V_{cb}^*}{V_{td}V_{tb}^*}\right), \qquad \gamma = \arg\left(-\frac{V_{ud}V_{ub}^*}{V_{cd}V_{cb}^*}\right) \tag{6}$$

being two sides and two angles as sketched in Fig. 1. The third side is the unity vector, while the third angle is $\alpha = \pi - \beta - \gamma = \arg(-V_{td}V_{tb}^*/(V_{ud}V_{ub}^*))$.[1] Similarly, it is useful to define the angle $\beta_s = \arg(-V_{ts}V_{tb}^*/(V_{cs}V_{cb}^*))$ from the triangle $V_{us}V_{ub}^* + V_{cs}V_{cb}^* + V_{ts}V_{tb}^* = 0$, as the phase of the B_s–$\bar{B}_s$ mixing amplitude is $-2\beta_s$ in the phase convention of Eq. (2). The UT sides and angles are observables and we discuss below their present determination from K- and B-meson physics.

Given the definition in Eq. (5), all the information related to the UT is encoded in one complex number

$$\bar{\rho} + i\bar{\eta} = R_u e^{i\gamma} \tag{7}$$

corresponding to the coordinates $(\bar{\rho}, \bar{\eta})$ in the complex plane of the only non-trivial apex of the UT.

It is worth mentioning another popular CKM parameterization introduced by Wolfenstein [4] which allows one to write an expansion of the CKM matrix in terms of a small parameter λ:

$$V = \begin{pmatrix} 1 - \frac{\lambda^2}{2} & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1 - \frac{\lambda^2}{2} & A\lambda^2 \\ A\lambda^3(1 - \rho - i\eta) & -A\lambda^2 & 1 \end{pmatrix} + \mathcal{O}(\lambda^4) \tag{8}$$

This parameterization makes explicit the hierarchy of CKM matrix elements, showing that quark flavour-changing transitions are suppressed in the SM. The exact relations between the PDG and the Wolfenstein parameterizations are

[1] The UT angles are also denoted as $\phi_1 = \beta$, $\phi_2 = \alpha$ and $\phi_3 = \gamma$ in the literature.

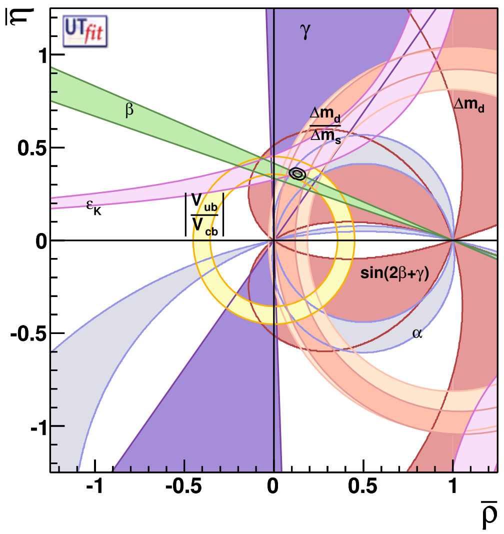

Fig. 2. Determination of $\bar{\rho}$ and $\bar{\eta}$ selected from constraints on $|V_{ub}/V_{cb}|$, Δm_d, Δm_s, ε_K, α, β, γ, $2\beta+\gamma$ and $\mathrm{BR}(B\to\tau\nu)$. 68% and 95% total probability contours are shown, together with 95% probability regions from the individual constraints.

Table 1
Summary of different measurements and corresponding SM predictions. In the last column the pull is explicitly indicated.

Observable	Prediction	Measurement	Pull (σ)
γ [°]	69.6 ± 3.1	74 ± 11	-0.4
α [°]	85.4 ± 3.7	91.4 ± 6.1	-0.8
$\sin 2\beta$	0.771 ± 0.036	0.654 ± 0.026	$+2.6$
$\|V_{ub}\|$ $[10^{-3}]$	3.55 ± 0.14	3.76 ± 0.20	-0.9
$\|V_{cb}\|$ $[10^{-3}]$	42.69 ± 0.99	40.83 ± 0.45	$+1.6$
ε_K $[10^{-3}]$	1.92 ± 0.18	2.23 ± 0.010	-1.7
$\mathrm{BR}(B\to\tau\nu)$ $[10^{-4}]$	0.805 ± 0.071	1.72 ± 0.28	-3.2

$$\lambda = \sin\theta_{12}, \qquad A = \frac{\sin\theta_{23}}{\sin^2\theta_{12}}, \qquad \rho = \frac{\sin\theta_{13}\cos\delta}{\sin\theta_{12}\sin\theta_{23}}, \qquad \eta = \frac{\sin\theta_{13}\sin\delta}{\sin\theta_{12}\sin\theta_{23}} \tag{9}$$

Notice that $\lambda \sim 0.22$ (the sine of the Cabibbo angle θ_{12}) is indeed a good expansion parameter. The relation between the UT apex coordinates and the Wolfenstein parameters is

$$\rho + i\eta = \sqrt{\frac{1-A^2\lambda^4}{1-\lambda^2}}\,\frac{\bar{\rho}+i\bar{\eta}}{1-A^2\lambda^4(\bar{\rho}+i\bar{\eta})} \simeq \left(1+\frac{\lambda^2}{2}\right)(\bar{\rho}+i\bar{\eta}) + \mathcal{O}(\lambda^4) \tag{10}$$

showing that $\bar{\rho}=\rho$ and $\bar{\eta}=\eta$ at the lowest order in λ.

The CKM matrix elements $|V_{ud}|$ and $|V_{us}|$ can be measured from super-allowed β decays [5] and semileptonic/leptonic kaon decays [6], respectively, determining accurately the sine of the Cabibbo angle. The other CKM parameters are determined through a fit to the UT in Eq. (5), as shown in Fig. 2, using the latest determinations of the theoretical and experimental parameters [7]. The basic constraints are $|V_{ub}/V_{cb}|$ from semileptonic B decays, Δm_d and Δm_s from $B_{d,s}$ oscillations, ε_K from K^0–$\bar{K}^0$ mixing, α from charmless hadronic B decays, γ and $2\beta+\gamma$ from charm hadronic B decays, $\sin 2\beta$ from $B^0 \to J/\psi K^0$ decays and the $\mathrm{BR}(B\to\tau\nu)$ [8].

The consistency between the different constraints clearly establishes the CKM matrix as the dominant source of flavour mixing and CP violation, described by a single parameter $\bar{\eta}$. The CKM picture can be quantitatively tested comparing the direct measurement of a flavour observable entering the UT fit with the SM prediction obtained from the UT fit without using the constraint being tested.

Table 1 contains a set of flavour observables: for each of them, the SM prediction, the measurement, and the pull (the difference of prediction and measurement in unit of σ) are shown. At present, the main tensions in the UT fit come from $\mathrm{BR}(B\to\tau\nu)$, which is found to deviate from the measurement by $\sim 3.2\sigma$, and $\sin 2\beta$, which is larger than the experimental value by $\sim 2.6\sigma$. It is interesting to note that the former prefers large values of $|V_{ub}/V_{cb}|$, while the latter wants $|V_{ub}/V_{cb}|$ to be small. Any change of the measured value of $|V_{ub}|$ and $|V_{cb}|$ will not improve the situation.

For massless neutrinos, the SM lepton sector is flavour diagonal and CP conserving. Once the SM is trivially extended to include right-handed neutrinos and account for neutrino masses, the formalism for lepton flavour and CP violation becomes similar to the one used for quarks. The lepton mixing matrix, called the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix [9], is parameterized much as the CKM matrix. The only difference is that the possibility of having Majorana mass terms for neutrinos allows for two additional CP-violating phases. In spite of these formal similarities, the flavour phenomenology of the lepton sector, in particular of neutrinos, is rather different from the quark one and its discussion goes beyond the scope of this primer. We just recall that, given the smallness of neutrino masses, charged lepton flavour violation is negligible in the SM.

2. Meson mixing and CP violation

CP violation in B decays appears as a difference between the rates into a given final state and its CP-conjugate, accounted for by the direct CP asymmetry

$$A_{\mathrm{CP}} = \frac{\Gamma(\bar{B} \to \bar{f}) - \Gamma(B \to f)}{\Gamma(\bar{B} \to \bar{f}) + \Gamma(B \to f)} = \frac{|\bar{A}|^2 - |A|^2}{|\bar{A}|^2 + |A|^2} \tag{11}$$

where f is the final state, $\bar{f}$ is its CP-conjugate and A and $\bar{A}$ are the two decay amplitudes. Charged mesons can only violate CP in the decay. On the other hand, neutral B mesons are subject to the mixing phenomenon, i.e. the mass eigenstates are a superposition of the flavour ones. In this case, CP violation can also occur in the mixing itself and in the interference between mixing and decay, giving additional opportunities to observe it. We briefly summarize in the following the main formulae related to mixing and CP violation in neutral B decays.

The time evolution of a system of unstable neutral meson–antimeson states such as B–$\bar{B}$ can be described by a 2×2 non-Hermitean matrix Hamiltonian $\hat{H}$

$$i\frac{\mathrm{d}}{\mathrm{d}t}\begin{pmatrix} |B(t)\rangle \\ |\bar{B}(t)\rangle \end{pmatrix} = \hat{H}\begin{pmatrix} |B(t)\rangle \\ |\bar{B}(t)\rangle \end{pmatrix} = \left(\hat{m} - \frac{i}{2}\hat{\Gamma}\right)\begin{pmatrix} |B(t)\rangle \\ |\bar{B}(t)\rangle \end{pmatrix} \tag{12}$$

$\hat{H}$ can be decomposed using the two Hermitean matrices $\hat{m}$ and $\hat{\Gamma}$ representing its dispersive and absorptive part respectively. In particular, assuming CPT invariance, one can write

$$\hat{m} = \begin{pmatrix} m & m_{12} \\ m_{12}^* & m \end{pmatrix}, \qquad \hat{\Gamma} = \begin{pmatrix} \Gamma & \Gamma_{12} \\ \Gamma_{12}^* & \Gamma \end{pmatrix} \tag{13}$$

The eigenstates of $\hat{H}$, denoted as $|B_{L,H}\rangle$, can be written as

$$|B_{L,H}\rangle = \frac{1}{\sqrt{1+|q/p|^2}}\big(|B\rangle \pm q/p\,|\bar{B}\rangle\big), \qquad q/p = -\sqrt{\frac{m_{12}^* - \frac{i}{2}\Gamma_{12}^*}{m_{12} - \frac{i}{2}\Gamma_{12}}} \tag{14}$$

where L (H) corresponds to the upper (lower) sign. Solving the eigenvalue problem and taking into account that $\Gamma_{12} \ll m_{12}$ for neutral B mesons, one finds the following mass and width differences:

$$\Delta m = m_H - m_L \simeq 2|m_{12}|, \qquad \Delta\Gamma = \Gamma_H - \Gamma_L \simeq \Delta m\,\mathrm{Re}\left(\frac{\Gamma_{12}}{m_{12}}\right) \tag{15}$$

These observables, or the corresponding dimensionless variables $x = \Delta m/\Gamma$ and $y = \Delta\Gamma/2\Gamma$ (Γ is the average lifetime), characterize the mixing. In addition, the absolute value $|q/p|$ is also observable. The deviation of $|q/p|$ from one is a measure of CP violation in the mixing, as shown by the expression of the semileptonic CP asymmetry

$$A_{\mathrm{SL}} = \frac{\Gamma(\bar{B} \to X\ell^+\nu_\ell) - \Gamma(B \to X\ell^-\bar{\nu}_\ell)}{\Gamma(\bar{B} \to X\ell^+\nu_\ell) - \Gamma(B \to X\ell^-\bar{\nu}_\ell)} = \frac{1-|q/p|^4}{1+|q/p|^4} \tag{16}$$

Expanding $|q/p|$ to first order in Γ_{12}/m_{12}, one finds

$$|q/p| \simeq 1 - \frac{1}{2}\,\mathrm{Im}\,\frac{\Gamma_{12}}{m_{12}} \tag{17}$$

Finally, the third manifestation of CP violation in B decays is through the interference between mixing and decay [10]. The key observable in this case is the time-dependent CP asymmetry of a meson state $B(t)$ decaying into a final state f. Taking a CP-eigenstate final state with the approximation $|q/p| \simeq 1$ and $\Delta\Gamma \simeq 0$, the asymmetry is given by

$$A_{\mathrm{CP}}^{B\to f}(t) = \frac{\Gamma(B(t) \to f) - \Gamma(\bar{B}(t) \to f)}{\Gamma(B(t) \to f) + \Gamma(\bar{B}(t) \to f)} \simeq C_{B\to f}\cos(\Delta m t) - S_{B\to f}\sin(\Delta m\, t) \tag{18}$$

where

$$C_{B\to f} = \frac{1-|\lambda_{B\to f}|^2}{1+|\lambda_{B\to f}|^2}, \qquad S_{M\to f} = \frac{2\,\mathrm{Im}(\lambda_{B\to f})}{1+|\lambda_{B\to f}|^2}, \quad \text{with } \lambda_{B\to f} = q/p\,\frac{\bar{A}}{A} \tag{19}$$

q/p is the B mixing parameter and A ($\bar{A}$) is the amplitude for $B \to f$ ($\bar{B} \to f$). The coefficient $C_{B\to f} \simeq -A_{\rm CP}$, the direct CP asymmetry defined in Eq. (11). The coefficient $S_{B\to f}$, instead, is a new measurement of CP violation generated by the interference of mixing and decay amplitudes.

Time-dependent CP asymmetries for B decays give access to the UT angles. For example, the amplitude for $B_d \to J/\psi K_S$ is dominated by a single term with a definite weak phase so that one finds $C_{B_d\to J/\psi K_S} = 0$ and

$$A_{\rm CP}^{B_d\to J/\psi K_S}(t) = -S_{B_d\to J/\psi K_S}\sin(\Delta m t), \qquad S_{B_d\to J/\psi K_S} \simeq \sin 2\beta \tag{20}$$

up to doubly Cabibbo-suppressed corrections. In other cases, the decay amplitude is not dominated by a single term and the extraction of the UT angles is less straightforward, as hadronic amplitudes no longer cancel in $\lambda_{B\to f}$. Previous formulae can also be generalized to non-CP-eigenstate final states [11].

3. Flavour and CP violation beyond the Standard Model

The Standard Model (SM) of electroweak and strong interactions is extremely successful in describing all presently available experimental data. However, its validity can extend at most to energies of the order of the Planck scale, where gravity comes into play. Let us therefore consider the SM as an effective theory valid up to a scale Λ. We can then write the SM Lagrangian as

$$\mathcal{L} = C_2\Lambda^2 H^\dagger H + \lambda\big(H^\dagger H\big)^2 + \mathcal{L}_{\rm gauge} + \mathcal{L}_{\rm Yukawa} + \sum_{d=5}^{\infty}\sum_{i=1}^{n_d}\frac{C_d^i}{\Lambda^{(d-4)}}O_d^i \tag{21}$$

where O_d^i is a generic gauge-invariant operator of dimension d. Now, it turns out that the Lagrangian truncated at $d \leqslant 4$ has some very important "accidental" symmetries that are violated by $O^i_{d>4}$. Most notable examples of such symmetries are given by baryon and lepton number conservation. The agreement of the SM with experimental data would suggest a very high value of Λ, so that the breaking of SM accidental symmetries gets strongly suppressed by the inverse powers of Λ in front of the higher-dimensional operators. However, we see from the first term in Eq. (21) that $C_2\Lambda$ controls the scale of electroweak symmetry breaking. Thus, unless we are willing to accept an extremely small value of C_2 (which means an extremely large amount of fine-tuning, since radiative corrections within the effective theory naturally generate $C_2 \sim \mathcal{O}(1)$), we are forced to consider values of the New Physics (NP) scale Λ not too far above the electroweak scale. But then the SM accidental symmetries require that NP has a peculiar structure, so that the coefficients of symmetry-breaking higher-dimensional operators are strongly suppressed and the phenomenological success of the SM remains unscathed. Turning the argument around, the coefficients of those higher-dimensional operators that break SM accidental symmetries provide the most stringent constraints on the NP scale and couplings (or better, on a combination thereof).

Let us now concentrate on two accidental symmetries of the SM: (i) the absence of tree-level Flavour Changing Neutral Currents (FCNC), and the GIM suppression of loop-mediated FCNC; (ii) The absence of tree-level CP violation in weak interactions. These accidental symmetries ensure that flavour physics is extremely sensitive to NP. Generically, NP will generate contributions to higher-dimensional operators mediating FCNC and CP-violating processes at a level that may well exceed by several orders of magnitude the present experimental bounds (see Ref. [12] for bounds on the coefficients of these operators). This implies that either NP is well above the TeV scale, or its flavour structure is non-trivial so that the coefficients of the relevant operators are generated with tiny couplings. In particular, one may define a class of models denoted by Minimal Flavour Violation (MFV) [13], in which the CKM matrix and quark masses remain the only source of flavour violation.

As an explicit example consider minimal supersymmetric extensions of the Standard Model (MSSM). Squark and slepton masses contain new sources of flavour and CP violation, that can be parameterized by the off-diagonal mass terms in the so-called super-CKM basis, in which quark masses and neutral current couplings are diagonal. Off-diagonal terms in general connect squarks of different chiralities, while in the SM flavour change only occurs in left-handed currents. This has dramatic implications in all processes where chirality-flipping amplitudes are enhanced, such as Kaon mixing (or where a chirality flip is required, such as $B \to X_s\gamma$). As a result, for $\mathcal{O}(1)$ off-diagonal terms, lower bounds on squark masses from ε_K are of the order of 10^4 TeV, while for squark masses at the TeV scale off-diagonal terms must be of $\mathcal{O}(10^{-4})$.

In Table 2 we report upper bounds on the parameters $\delta^d_{ij} \equiv \Delta^d_{ij}/m^2_{\tilde{d}}$, where Δ^d_{ij} is the flavour-violating off-diagonal entry appearing in the down-type squark mass matrices and $m^2_{\tilde{d}}$ is the average squark mass. The mass insertions include the LL/LR/RL/RR types, according to the chirality of the corresponding SM quarks. Detailed bounds on the individual δs have been derived by considering limits from various FCNC processes as described in Ref. [14].

Viable models can be constructed in several ways: either by allowing squarks of the first two generations to be very heavy (in order to evade the stringent constraints from Kaon physics) [15], or by requiring squark masses to have definite quantum numbers under (a subset of) the flavour symmetry group of SM fermion gauge interactions [16].

Table 2
95% probability bounds on $|(\delta^d_{ij})_{AB}|$ taken from Ref. [14].

$ij\backslash AB$	*LL*	*LR*	*RL*	*RR*
12	1.4×10^{-2}	9.0×10^{-5}	9.0×10^{-5}	9.0×10^{-3}
13	9.0×10^{-2}	1.7×10^{-2}	1.7×10^{-2}	7.0×10^{-2}
23	1.6×10^{-1}	4.5×10^{-3}	6.0×10^{-3}	2.2×10^{-1}

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