Flavour and CP Violation

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Abstract. In this talk I review the status of $\mathrm{B}-\overline{\mathrm{B}}$ mixing in the Standard Model and the room for new physics in $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ and $\mathrm{B_d}-\overline{\mathrm{B}}_\mathrm{d}$ mixing in the light of recent LHCb data.

1. Introduction

The only source of flavour-changing transition in the Standard Model (SM) is the Yukawa interaction of the Higgs doublet H with the fermion fields. The Yukawa lagrangian of the quark fields reads

$$-L_Y = Y^d_{jk}\,\overline{Q}^j_L\,H\,d^k_R \;+\; Y^u_{jk}\,\overline{Q}^j_L\,\widetilde{H}\,u^k_R \;+\; \text{h.c.} \tag{1}$$

Here Q^j_L denotes the SU(2) doublet of the left-handed quark fields of the j-th generation and d^j_R, u^j_R are the corresponding right-handed singlet fields. The 3×3 Yukawa matrices $Y^{d,u}$ and the Higgs field vacuum expectation value $v = 174$ GeV combine to the quark mass matrices $M^{d,u} = Y^{d,u}v$. The gauge interactions of the quark fields do not change under unitary rotations of any of Q^j_L, d^j_R, and u^j_R in flavour space. Using these unphysical transformations one can bring Y^u and Y^d to the form

$$Y^u = \widehat{Y}^u = \begin{pmatrix} y_u & 0 & 0 \\ 0 & y_c & 0 \\ 0 & 0 & y_t \end{pmatrix} \qquad \text{and} \qquad Y^d = V^\dagger \widehat{Y}^d \qquad \text{with} \quad \widehat{Y}^d = \begin{pmatrix} y_d & 0 & 0 \\ 0 & y_s & 0 \\ 0 & 0 & y_b \end{pmatrix} \tag{2}$$

and $y_i \geq 0$. With the choice adopted in Eq. (2) the mass matrix M^u of up, charm, and top quark is diagonal. The diagonalisation of M^d requires the additional rotation $d^j_L = V_{jk} d^{k\prime}_L$, which puts the unitary Cabbibbo-Kobayashi-Maskawa (CKM) matrix V into the W boson vertices: $W_\mu \overline{u}^j_L \gamma^\mu d^j_L = W_\mu V_{jk} \overline{u}^j_L \gamma^\mu d^{k\prime}_L$. Eq. (2) defines a basis of weak eigenstates which is particularly suited to display the smallness of flavour violation in the SM. All flavour-changing transitions originate from Y^d which reads

$$Y^d = V^\dagger \widehat{Y}^d \;=\; \begin{pmatrix} 10^{-5} & -7\cdot 10^{-5} & (12+6i)\cdot 10^{-5} \\ 4\cdot 10^{-6} & 3\cdot 10^{-4} & -6\cdot 10^{-4} \\ (2+6i)\cdot 10^{-8} & 10^{-5} & 2\cdot 10^{-2} \end{pmatrix}, \tag{3}$$

Figure 1. Left: SM box diagram describing $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing, a $|\Delta B| = 2$ process. Right: $b \to s$ penguin amplitude, a $|\Delta B| = 1$ transition.

if all Yukawa couplings are evaluated at the scale $\mu = m_t$. The off-diagonal element largest in magnitude is $V_{ts}^* y_b \equiv V_{32}^* y_b = -6 \cdot 10^{-4}$. Moreover, in the SM flavour-changing neutral current (FCNC) transitions involve an additional loop suppression, making FCNC processes an extremely sensitive probe of physics beyond the SM. This feature puts flavour physics into a win-win position: If CMS and ATLAS find new particles, FCNC observables will be used to probe the flavour pattern of the BSM theory to which these particles belong. If instead CMS and ATLAS do not find any BSM particles, flavour physics will indirectly probe new physics to scales exceeding 100 TeV, well beyond the center-of-mass energy of the LHC.

The weak interaction makes the neutral $K \sim \overline{s}d$, $D \sim c\overline{u}$, $B_d \sim \overline{b}d$ and $B_s \sim \overline{b}s$ mesons mix with their antiparticles, $\overline{K}$, $\overline{D}$, $\overline{B}_d$, and $\overline{B}_s$, respectively. This means that, say, a B_d meson evolves in time into a quantum-mechanical superposition of a B_d and a $\overline{B}_d$. This feature is a gold-mine for new-physics searches, as it permits to probe CP phases in any decay to a final state which is accessible from both the B_d and the $\overline{B}_d$ component of the decaying state. In a tree-level decay like $\overset{(-)}{B}_d \to J/\psi K_S$ the time-dependent CP asymmetry probes new physics in the mixing amplitude itself. If instead a rare decay (with FCNC decay amplitude) is studied, the time-dependent CP asymmetry may reveal a new physics contribution to the decay amplitude. Meson-antimeson mixing is a $|\Delta F| = 2$ process, meaning that the relevant flavour quantum numbers change by two units. In the case of $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing depicted in Fig. 1 these flavour quantum numbers are beauty B and strangeness S. Weak decays are $|\Delta F| = 1$ processes, Fig. 1 shows a $|\Delta B| = |\Delta S| = 1$ FCNC decay amplitude. It is instructive to compare the reach of $|\Delta F| = 2$ and $|\Delta F| = 1$ amplitudes to new physics: The flavour violation in the SM the amplitudes $A_{\rm SM}^{|\Delta B|=2}$ and $A_{\rm SM}^{|\Delta B|=1}$ in Fig. 1 is governed by the small CKM element V_{ts} and both diagrams scale as $1/M_W^2$. A new-physics (NP) contribution may instead involve some new flavour-violating parameter $\delta_{\rm FCNC}$. If the scale of NP is Λ and the NP contribution enters at the one-loop level, one finds

$$\frac{|A_{\rm NP}^{|\Delta B|=2}|}{|A_{\rm SM}^{|\Delta B|=2}|} = \frac{|\delta_{\rm FCNC}|^2}{|V_{ts}|^2}\frac{M_W^2}{\Lambda^2} \qquad \text{and} \qquad \frac{|A_{\rm NP}^{|\Delta B|=1}|}{|A_{\rm SM}^{|\Delta B|=1}|} = \frac{|\delta_{\rm FCNC}|}{|V_{ts}|}\frac{M_W^2}{\Lambda^2}. \tag{4}$$

One sees that $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing is more sensitive to generic NP than an FCNC $b \to s$ decay amplitude, if $|\delta_{\rm FCNC}| > |V_{ts}| \approx 0.04$. The estimate in Eq. (4) applies, for example, to the (N)MSSM with generic flavour structure in the bilinear SUSY-breaking terms, which induce new $b_R \to s_R$ or $b_L \to s_L$ transitions through squark-gluino loops [1, 2]. But $|\Delta B| = 1$ transitions can be more sensitive probes of NP in models in which $b_R \to s_L$ or $b_L \to s_R$ transitions are parametrically enhanced over their SM counterparts. This situation occurs in the (N)MSSM with a large value of the parameter $\tan\beta$ or large trilinear SUSY-breaking terms. Theories with such chirally enhanced FCNC transitions are efficiently probed with the radiative decays $b \to s\gamma, d\gamma$ [3–10], leptonic decays like the recently observed [11] decay $B_s \to \mu^+\mu^-$ [7, 12–21] and the chromomagnetic contribution to non-leptonic decays like $B_d \to \phi K_S$ [18].

2. $\mathrm{B}-\overline{\mathrm{B}}$ mixing: formalism and Standard-Model prediction

$\mathrm{B_q}-\overline{\mathrm{B}}_\mathrm{q}$ mixing with $q = d$ or $q = s$ is governed by the 2×2 matrix $M^q - i\Gamma^q/2$, with the hermitian mass and decay matrices M^q and Γ^q. Due to non-vanishing off-diagonal elements M_{12}^q and Γ_{12}^q a meson tagged as a B_q at time $t = 0$ will for $t > 0$ evolve into a superposition of B_q and $\overline{B}_q$. In the SM M_{12}^s is calculated from the dispersive part of the box diagram in Fig. 1, which amounts to discarding the imaginary part of the loop integral. Γ_{12}^s is instead obtained from the absorptive part, meaning that the real part of the loop integral is dropped. M_{12}^d and Γ_{12}^d are calculated in the same way from the box diagram with the external s quarks replaced by d quarks. M_{12}^q is dominated by the top contribution, while Γ_{12}^q stems from the diagrams with internal up and charm. Γ_{12}^q is made up of all final states into which both B_q and $\overline{B}_q$ can decay. There are three physical quantities in $\mathrm{B_q}-\overline{\mathrm{B}}_\mathrm{q}$ mixing:

$$|M_{12}^q|, \quad |\Gamma_{12}^q|, \quad \phi_q \equiv \arg\left(-\frac{M_{12}^q}{\Gamma_{12}^q}\right).$$

By diagonalising $M - i\,\Gamma/2$ one finds the two mass eigenstates B_q^H and B_q^L, where the superscripts stand for "heavy" and "light". These eigenstates differ in their mass and width with

$$\Delta m_q \;=\; M_H^q - M_L^q \;\simeq\; 2|M_{12}^q|, \qquad\qquad \Delta\Gamma_q \;=\; \Gamma_L^q - \Gamma_H^q \;\simeq\; 2|\Gamma_{12}^q|\cos\phi_q\,.$$

The mass difference Δm_q simply equals the frequency at which B_q and $\overline{B}_q$ oscillate into each other. The width difference $\Delta\Gamma_s$ is sizable, so that in general untagged B_s decays are governed by the sum of two exponentials. There is no useful data on the tiny width difference $\Delta\Gamma_d$ in the B_d system yet, which is predicted to be around 0.5% of $\Gamma_{L,H}^d$ [22–25]. The CP-violating phase ϕ_q can be determined by measuring the CP asymmetry in flavour-specific decays:

$$a_{\mathrm{fs}}^q \;=\; \frac{|\Gamma_{12}^q|}{|M_{12}^q|}\sin\phi_q\,.$$

A decay $B_q \to f$ is called flavour-specific, if the decay $\overline{B}_q \to f$ is forbidden. The standard method to determine a_{fs}^q uses semileptonic decays, so that a_{fs}^q is often called semileptonic CP asymmetry. While this measurement simply requires to count the numbers of positive and negative leptons in B_q decays, it is still very difficult, because $|\Gamma_{12}^q| \ll |M_{12}^q|$ renders $|a_{\mathrm{fs}}^q|$ small, even if NP contributions to M_{12}^q enhance ϕ_q over its small SM value.

2.1. Mass difference Δm_s

The theoretical prediction of Δm_q requires the separation of short-distance and long-distance QCD effects. To this end one employs an operator product expansion, which factorises M_{12}^q as

$$M_{12}^q = \left(V_{tq}^* V_{tb}\right)^2\, C\langle B_q|Q|\overline{B}_q\rangle. \tag{5}$$

The Wilson coefficient C comprises the short-distance physics, with all dependence on the heavy particle masses. QCD corrections to C have been calculated reliably in perturbation theory [26]. Since the CKM elements are factored out in Eq. (5), $C = C(m_t, M_W, \alpha_s)$ is real in the SM.

$$Q \;=\; \bar{q}_L\gamma_\nu b_L\;\bar{q}_L\gamma^\nu b_L \tag{6}$$

is a local four-quark operator describing a point-like $|\Delta B| = 2$ transition, see Fig. 2. The formalism is the same for the B_d and B_s complex; for definiteness I only discuss the case $q = s$ in the following. The average of the CDF and LHCb measurements of $\Delta m_s \simeq 2|M_{12}^s|$ is [27]

$$\Delta m_s^{\mathrm{exp}} = (17.719 \pm 0.043)\,\mathrm{ps}^{-1}, \tag{7}$$

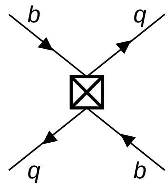

Figure 2. The $|\Delta B| = 2$ operator Q of Eq. (6). It is pictorially obtained by shrinking the $\mathrm{B_q}-\overline{\mathrm{B}}_\mathrm{q}$ mixing box diagram to a point.

meaning that $|M_{12}^s|$ is known very precisely. However, we need the hadronic matrix element $\langle B_s|Q|\overline{B}_s\rangle$ to confront Eq. (7) with the SM. The latter is usually parametrised as

$$\langle B_s|Q|\overline{B}_s\rangle \;\; = \;\; \frac{2}{3} M_{B_s}^2 \, f_{B_s}^2 \, B_{B_s}.$$

Here M_{B_s} and f_{B_s} are mass and decay constant of the B_s meson, respectively. The hadronic parameter B_{B_s} depends on the renormalisation scheme and scale used to define Q, in this talk $B_{B_s} \approx 0.85$ is understood to be evaluated at the scale $\mu = m_b$ in the $\overline{\mathrm{MS}}$ scheme. Noting that $|V_{ts}|$ is fixed by CKM unitarity from the well-measured element $|V_{cb}|$, we can write

$$\Delta m_s = \left(\, 18.8 \pm 0.6_{V_{cb}} \pm 0.3_{m_t} \pm 0.1_{\alpha_s} \right) \, \mathrm{ps}^{-1} \, \frac{f_{B_s}^2 \, B_{B_s}}{(220\,\mathrm{MeV})^2}, \tag{8}$$

where the uncertainties from the experimental errors of the input parameters indicated. Averaging various calculations from lattice gauge theory one finds [28]

$$f_{B_s}^2 B_{B_s} = \left[(211 \pm 9)\,\mathrm{MeV}\right]^2 . \tag{9}$$

Inserting this result into Eq. (8) implies

$$\Delta m_s = (17.3 \pm 1.5)\,\mathrm{ps}^{-1} \tag{10}$$

complying excellently with the experimental result in Eq. (7). However, the average in Eq. (9) mainly involves different calculations of f_{B_s}, which are combined with two fairly old results for B_{B_s}. With the recent preliminary Fermilab/MILC result [29]

$$f_{B_s}^2 \, B_{B_s} = 0.0559(68)\,\mathrm{GeV}^2 \simeq \left[(237 \pm 14)\,\mathrm{MeV}\right]^2$$

one finds instead

$$\Delta m_s = (21.7 \pm 2.6)\,\mathrm{ps}^{-1}. \tag{11}$$

Therefore more effort on lattice-QCD calculations of $f_{B_s}^2 \, B_{B_s}$ is highly desirable. As long as modern calculations of this quantity are unavailable, I recommend to inflate the error in Eq. (10) to $2.5\,\mathrm{ps}^{-1}$. One concludes that the precise measurement in Eq. (7) still permits a NP contribution of 15% in Δm_s.

2.2. CP phase from $B_s \to J/\psi\phi$

While Δm_s fixes the magnitude of M_{12}^s, the phase of M_{12}^s can be probed through the mixing-induced CP asymmetry in $B_s \to J/\psi\phi$. The final state contains two vector mesons, so that the orbital angular momentum L can assume the values 0,1, and 2. The final states $(J/\psi\,\phi)_{L=0,2}$

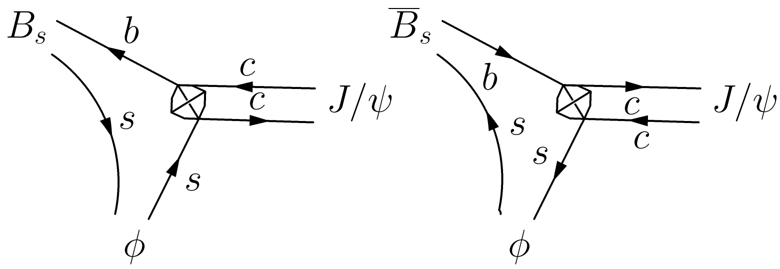

Figure 3. Amplitudes of $B_s \to J/\psi\,\phi$ and $\overline{B}_s \to J/\psi\,\phi$. The cross denotes the W-mediated $b \to c\bar{c}s$ decay.

have CP quantum number $\eta_{CP} = 1$, while $(J/\psi\,\phi)_{L=1}$ is CP-odd. Denoting a meson which was a B_s at time $t = 0$ by $B_s(t)$ (with an analogous definition of $\overline{B}_s(t)$), one can define the time-dependent CP asymmetry

$$a_{CP}(t) = \frac{\Gamma(B_s(t) \to (J/\psi\,\phi)_L) - \Gamma(\overline{B}_s(t) \to (J/\psi\,\phi)_L)}{\Gamma(B_s(t) \to (J/\psi\,\phi)_L) + \Gamma(\overline{B}_s(t) \to (J/\psi\,\phi)_L)}. \tag{12}$$

The two interfering amplitudes giving rise to $a_{CP}(t)$ are shown in Fig. 3. The analytical result reads

$$a_{CP}(t) = \eta_{CP} \frac{\sin(2\beta_s)\sin(\Delta m_s t)}{\cosh(\Delta\Gamma_s t/2) - \cos(2\beta_s)\sinh(\Delta\Gamma_s t/2)}. \tag{13}$$

The CP phase entering $a_{CP}(t)$ is $\beta_s = \arg(-V_{tb}^* V_{ts}/(V_{cb}^* V_{cs}))$, in the standard phase convention of the CKM matrix β_s is essentially just the phase of $-V_{ts}$. It should be mentioned that Eq. (12) is a theoretical definition, in practice the experimental separation of the $L = 0, 1, 2$ amplitudes involves a complicated angular analysis of the J/ψ and ϕ decays.

The CKM matrix can be parametrised in terms of the four Wolfenstein parameters λ,A,$\overline{\rho}$, and $\overline{\eta}$ [30,31]. Expanding to leading order in $\lambda = 0.225$ one finds $\sin(2\beta_s) \simeq 2\lambda^2\overline{\eta}$. The parameter $\overline{\eta}$ is the height of the CKM unitarity triangle defined by $\overline{\rho} + i\overline{\eta} = -V_{ub}^* V_{ud}/(V_{cb}^* V_{cd})$. The suppression by λ^2 renders $\sin(2\beta_s)$ small; determining the CKM elements from a global fit to the data gives $\overline{\eta} = 0.343 \pm 0.015$ and leads to the SM prediction [28]

$$2\beta_s = 2.1^\circ \pm 0.1^\circ. \tag{14}$$

The combination of CDF, D0, ATLAS, and LHCb data on $B_s \to J/\psi\phi$ and of LHCb data on $B_s \to J/\psi\pi^+\pi^-$ gives [27]

$$2\beta_s^{\rm exp} = 0.7^\circ{}^{+5.2^\circ}_{-4.8^\circ}. \tag{15}$$

Of course, this average is dominated by the LHCb data. At this conference Jeroen van Leerdam has reported $2\beta_s = 0.1^\circ \pm 4.8^\circ_{\rm stat} \pm 1.5^\circ_{\rm syst}$ from the combined LHCb analysis of $B_s \to J/\psi K^+K^-$ and $B_s \to J/\psi\pi^+\pi^-$ [32]. In Eq. (13) an additional small contribution, the "penguin pollution" has been neglected. This –presently uncalculable– contribution inflicts an additional error of order 1° on $2\beta_s^{\rm exp}$ in Eq. (15).

2.3. Decay matrix Γ_{12}^q

The calculation of Γ_{12}^q, $q = d, s$, is needed for the width difference $\Delta\Gamma_q \simeq 2|\Gamma_{12}^q| \cos\phi_q$ and the semileptonic CP asymmetry $a_{\rm fs}^q = \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin\phi_q$. The theoretical results of Refs. [22, 23, 33] have recently been updated to [25]

$$\begin{array}{rclcrcl} \phi_s &=& 0.22^\circ \pm 0.06^\circ, & & \phi_d &=& -4.3^\circ \pm 1.4^\circ, \\ a_{\rm fs}^s &=& (1.9 \pm 0.3) \cdot 10^{-5}, & \text{and} & a_{\rm fs}^d &=& -(4.1 \pm 0.6) \cdot 10^{-4}. \end{array} \tag{16}$$

The prediction of Γ_{12}^q involves new operators in addition to Q in Eq. (6). However, the matrix element of Q comes with the largest coefficient, so that the ratio $\Delta\Gamma_q/\Delta m_q = |\Gamma_{12}^q|/|M_{12}^q|$ suffers from smaller hadronic uncertainties than $\Delta\Gamma_q$. Predicting $\Delta\Gamma_q$ with the help of the experimental values $\Delta m_d^{\rm exp} = 0.507\,{\rm ps}^{-1}$ and $\Delta m_s^{\rm exp}$ in Eq. (7) one infers from Ref. [25]:

$$\begin{aligned} \Delta\Gamma_d &= \frac{\Delta\Gamma_d}{\Delta m_d} \Delta m_d^{\rm exp} = (27 \pm 5) \cdot 10^{-4}\,{\rm ps}^{-1} \\ \Delta\Gamma_s &= \frac{\Delta\Gamma_s}{\Delta m_s} \Delta m_s^{\rm exp} = (0.090 \pm 0.018)\,{\rm ps}^{-1} \end{aligned} \tag{17}$$

The more recent update in Ref. [34] differs from Ref. [25] in two respects: The quoted results are obtained in a different renormalisation scheme and use the preliminary lattice results of [29] instead of the averages from [28]. Ref. [34] finds for $\Delta\Gamma_s$:

$$\Delta\Gamma_s = \frac{\Delta\Gamma_s}{\Delta m_s} \Delta m_s^{\rm exp} = (0.078 \pm 0.018)\,{\rm ps}^{-1}, \tag{18}$$

which is consistent with Eq. (17). In Eqs. (17) and (18) all uncertainties are added in quadrature, which may not be a conservative estimate of the overall error. The ranges comply with the LHCb measurement [32]

$$\Delta\Gamma_s^{\rm LHCb} = [0.116 \pm 0.018_{\rm stat} \pm 0.006_{\rm syst}]\,{\rm ps}^{-1} \tag{19}$$

and the world average [27]:

$$\Delta\Gamma_s^{\rm exp} = [0.089 \pm 0.012]\,{\rm ps}^{-1}. \tag{20}$$

3. New physics in $\mathrm{B}-\overline{\mathrm{B}}$ mixing

3.1. A model-independent analysis

If some NP amplitude adds to the $\mathrm{B_q}-\overline{\mathrm{B}}_\mathrm{q}$ mixing box diagram, both $|M_{12}^q|$ and ϕ_q may deviate from their SM predictions. The DØ experiment has measured [35–37]

$$\begin{aligned} A_{\rm SL}^{\rm D0} &= (0.532 \pm 0.039) a_{\rm fs}^d + (0.468 \pm 0.039) a_{\rm fs}^s \\ &= (-7.87 \pm 1.72 \pm 0.93) \cdot 10^{-3}, \end{aligned} \tag{21}$$

which is 3.9σ off the SM prediction inferred from Eq. (16),

$$A_{\rm SL} = (-0.20 \pm 0.03) \cdot 10^{-3}. \tag{22}$$

Adding a NP contribution $\phi_{d,s}^\Delta$ to either ϕ_d or ϕ_s in Eq. (16) may enhance $|A_{\rm SL}^{\rm D0}|$, but the very same contribution will also affect the value of β_s in Eq. (15) and the value of $\beta = \arg = \arg(-V_{tb}V_{td}^*/(V_{cb}V_{cd}^*))$ found from $a_{CP}(t)(B_d \to J/\psi K_S)$. One can parametrise NP in $\mathrm{B_d}-\overline{\mathrm{B}}_\mathrm{d}$ mixing and $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing by two complex parameters Δ_d and Δ_s:

$$M_{12}^q \equiv M_{12}^{\rm SM,q} \cdot \Delta_q\,, \qquad \Delta_q \equiv |\Delta_q| e^{i\phi_q^\Delta}.$$

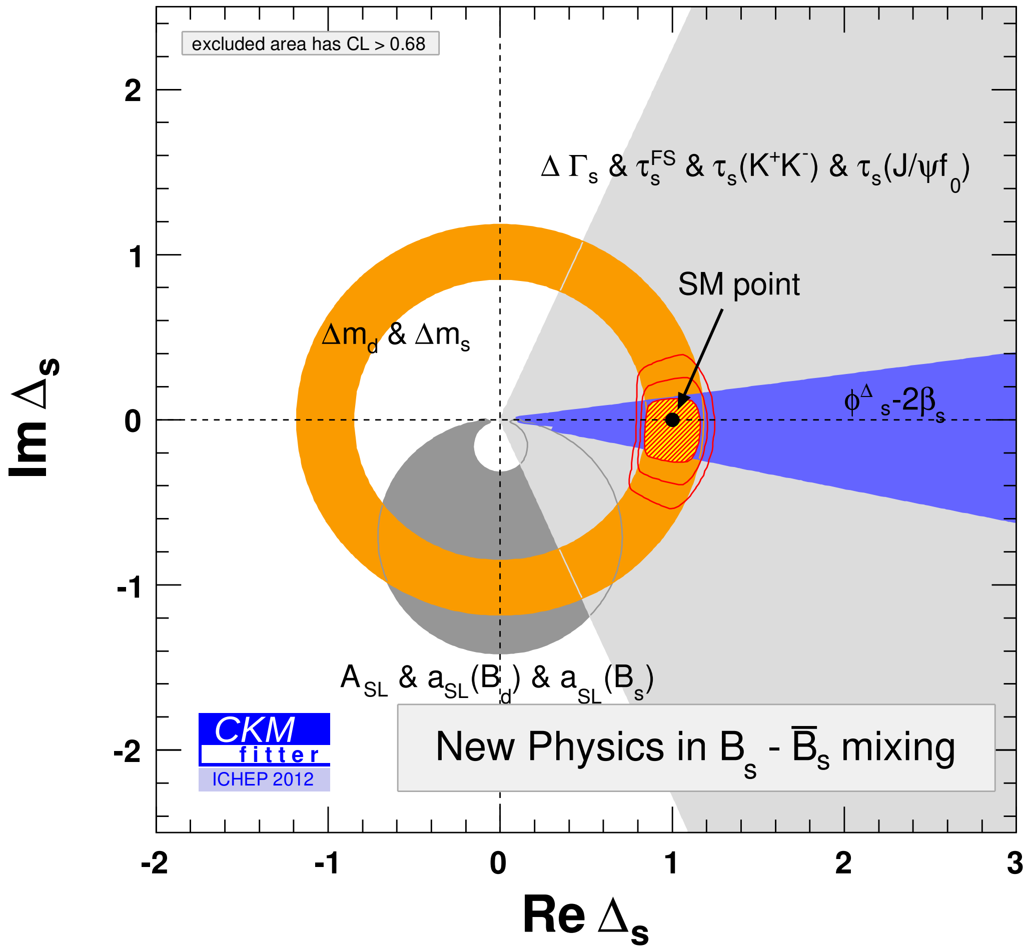

Figure 4. Allowed range for Δ_s, see [28, 39] for details.

In summer 2010, before the advent of precision data from LHCb, a global fit of all relevant flavour data to the CKM elements and Δ_d and Δ_s has resulted in a 3.6σ evidence of NP, with a large negative NP phase ϕ_s^Δ [38]. In spring 2012 this fit has been repeated [39]; Figs. 4 and 5 show an update of the results in Ref. [39] with the data of 2012 summer conferences [28]. The fit tries to accomodate $A_{\rm SL}$ in Eq. (21) and a slightly high value of the world average for $B(B \to \tau\nu)$ through $\phi_d^\Delta < 0$. Since $a_{CP}(t)(B_d \to J/\psi K_S)$ precisely fixes $2\beta + \phi_d^\Delta = 42.8^\circ \pm 1.6^\circ$, $\phi_d^\Delta < 0$ comes with a higher value of β than in the SM. ($\beta > \beta^{\rm SM}$ entails a larger value of $|V_{ub}|$ which governs $B(B \to \tau\nu)$.) Contrary to the situation in 2010, the Standard Model point $\Delta_s = \Delta_d = 1$ is merely disfavoured by 1 standard deviation, consistent with natural statistical fluctuations. At the best-fit point the problem with $A_{\rm SL}$ is only marginally alleviated.

One could relax the problem with $A_{\rm SL}$ without affecting $a_{CP}(t)(B_d \to J/\psi K_S)$ and $a_{CP}(t)(B_s \to J/\psi\phi)$ by postulating new physics in Γ_{12}^d or Γ_{12}^s [39–41]. Since Γ_{12}^s originates from Cabibbo-favoured tree-level decays, it can hardly be changed in a significant way without spoiling the value of the average decay width $\Gamma_s = (\Gamma_H^s + \Gamma_L^s)/2$ [38, 39]. The LHCb measurement

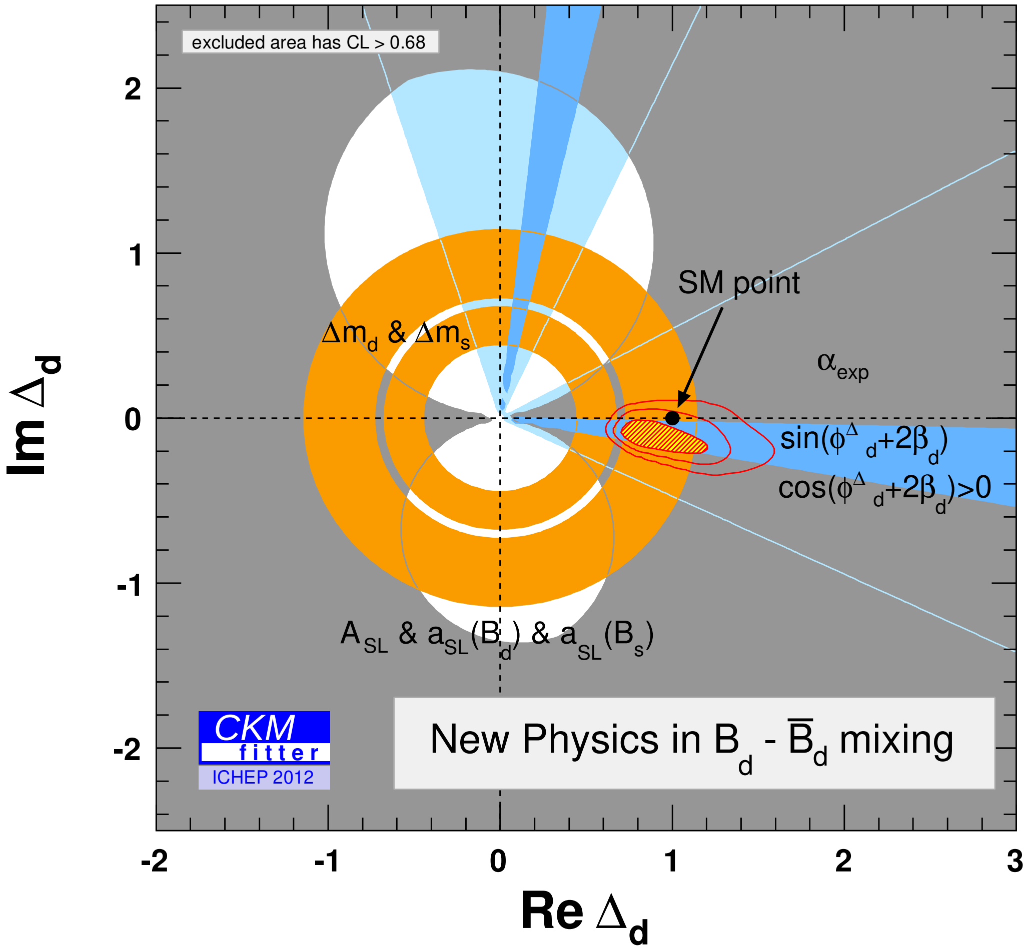

Figure 5. Allowed range for Δ_d, see [28, 39] for details.

$\Gamma_s^{\rm LHCb} = (0.6580 \pm 0.0054 \pm 0.0066)\,{\rm ps}^{-1}$ [32] implies

$$\frac{\Gamma_d}{\Gamma_s} = \frac{\tau_{B_s}}{\tau_{B_d}} = 0.997 \pm 0.013 \tag{23}$$

in excellent agreement with the SM prediction $\tau_{B_s}/\tau_{B_d} = 0.998 \pm 0.003$ [25, 42]. NP in the doubly Cabibbo-suppressed quantity Γ_{12}^d is phenomenologically only poorly constrained, but requires a somewhat contrived model of NP.

3.2. A supersymmetric SO(10) model

The Minimal Supersymmetric Standard Model (MSSM) has many new sources of flavour violation, which all reside in the supersymmetry-breaking sector. It is no problem to get a big effect in a chosen FCNC process, but rather to suppress big effects elsewhere. This *supersymmetric flavour problem* is substantially alleviated with the lower bounds on the squark massed placed by ATLAS and CMS. Grand Unified Theories (GUT) offer the possibility to have "controlled" deviations from the CKM pattern of flavour violation in the quark sector: In GUTs quarks and leptons are unified in common symmetry multiplets, opening the possibility to observe the large lepton-flavour mixing encoded in the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix $U_{\rm PMNS}$ in quark flavour physics [43, 44]. Consider SU(5) multiplets:

$$\mathbf{\overline{5}_1} = \begin{pmatrix} d_R^c \\ d_R^c \\ d_R^c \\ e_L \\ -\nu_e \end{pmatrix}, \qquad \mathbf{\overline{5}_2} = \begin{pmatrix} s_R^c \\ s_R^c \\ s_R^c \\ \mu_L \\ -\nu_\mu \end{pmatrix}, \qquad \mathbf{\overline{5}_3} = \begin{pmatrix} b_R^c \\ b_R^c \\ b_R^c \\ \tau_L \\ -\nu_\tau \end{pmatrix}.$$

If the observed large atmospheric neutrino mixing angle stems from a rotation of $\mathbf{\overline{5}_2}$ and $\mathbf{\overline{5}_3}$ in flavour space, it will induce a large $\tilde{b}_R - \tilde{s}_R$-mixing. Contrary to the situation with right-handed quark fields, rotations of right-handed squark fields in flavour space are physical because of the soft SUSY-breaking terms. The Chang–Masiero–Murayama (CMM) model has implemented this idea in a GUT based on the symmetry breaking chain SO(10)$\to$ SU(5) $\to$ SU(3)$\times$SU(2)$_L\times$U(1)$_Y$ [45–47]. In some weak basis the Yukawa matrix of down (s)quarks is diagonalised as

$$\mathsf{Y}_d = V_{\rm CKM}^* \begin{pmatrix} y_d & 0 & 0 \\ 0 & y_s & 0 \\ 0 & 0 & y_b \end{pmatrix} U_{\rm PMNS}$$

and the right-handed down squark mass matrix has the following diagonal form:

$$\mathsf{m}^2_{\tilde{d}}(M_Z) = \mathrm{diag}\left(m^2_{\tilde{d}},\, m^2_{\tilde{d}},\, m^2_{\tilde{d}} - \Delta_{\tilde{d}} \right).$$

The real parameter $\Delta_{\tilde{d}}$ is calculated from renormalisation-group effects driven by the top-Yukawa coupling.

Rotating Y_d to diagonal form puts the large atmospheric neutrino mixing angle into $\mathsf{m}^2_{\tilde{d}}$:

$$U^\dagger_{\rm PMNS}\, \mathsf{m}^2_{\tilde{d}}\, U_{\rm PMNS} = \begin{pmatrix} m^2_{\tilde{d}} & 0 & 0 \\ 0 & m^2_{\tilde{d}} - \frac{1}{2}\Delta_{\tilde{d}} & -\frac{1}{2}\Delta_{\tilde{d}}\, e^{i\xi} \\ 0 & -\frac{1}{2}\Delta_{\tilde{d}}\, e^{-i\xi} & m^2_{\tilde{d}} - \frac{1}{2}\Delta_{\tilde{d}} \end{pmatrix} \qquad (24)$$

The CP phase ξ affects CP violation in $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing!

In Ref. [47] we have confronted the CMM model with flavour data. The analysis involves seven parameters, which we have fixed by choosing values for the squark masses $M_{\tilde{u}}$, $M_{\tilde{d}}$ of right-handed up and down squarks, the trilinear term a_1^d of the first generation, the gluino mass $m_{\tilde{g}_3}$, $\tan\beta$, and the sign of the higgsino mass parameter μ. We have considered $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing, $b \to s\gamma$, $\tau \to \mu\gamma$, vacuum stability bounds, lower bounds on sparticle masses and the lower bound on the lightest Higgs boson. From these inputs first universal SUSY-breaking terms defined at a fundamental scale near the Planck scale have been determined through the renormalisation group equations (RGE). Subsequently we have used the RGE to determine all low-energy parameters, with the MSSM as the low-energy theory.

Two experimental results of the year 2012 put the CMM model under pressure: First, the sizable neutrino mixing angle θ_{13} leads to an unduly large effect in $B(\mu \to e\gamma)$. In Eq. (24) I have tacitly assumed a $U_{\rm PMNS}$ with tri-bimaximal mixing, corresponding to $\theta_{13} = 0$. For the actual value $\theta_{13} \approx 8^\circ$ the $(1, 2)$ element of the charged slepton mass matrix gets large, and one has to resort to much larger sfermion masses than those considered in Ref. [47]. Second, the Higgs mass of 126 GeV challenges the model and we are unable to find parameters which simultaneously satisfy the Higgs mass constraint and the experimental upper bound on $B(\mu \to e\gamma)$ [48]. The problem with the Higgs mass could be circumvented by considering the NMSSM as the low-energy theory.

4. Conclusions

LHCb has provided us with a significantly better insight into the $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ mixing complex. Δm_s and $\Delta\Gamma_s$ comply with the SM, but we need better lattice data for the hadronic matrix elements involved. Theoretical uncertainties still permit an $\mathcal{O}(20\%)$ NP contribution to the $\mathrm{B_s}-\overline{\mathrm{B}}_\mathrm{s}$ amplitude M_{12}^s. While in 2010 the DØ result for A_{SL} could be explained in scenarios with NP only in $M_{12}^{d,s}$, the LHCb data on $B_s \to J/\psi\phi$ now prohibit this solution. An alternative explanation invoking new physics in Γ_{12}^s is not viable, because this will spoil the ratio Γ_s/Γ_d which agrees well with the SM prediction. Maybe it is worthwile to look at NP in Γ_{12}^d, although this possibility leads to somewhat contrived models. Models of GUT flavour physics with $\widetilde{b} \to \widetilde{s}$ transition driven by the atmospheric neutrino mixing angle are under pressure from the large value of θ_{13}, which induces a too large $B(\mu \to e\gamma)$. Moreover, in the studied CMM model it seems impossible to accomodate the measured value of the lightest Higgs mass, if one insists on the MSSM as the low-energy theory.

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