A brief review of theoretical aspects of new-physics searches in flavor physics is given. Special attention is thereby devoted to $B_s$ mixing and the interplay of high- and low-$p_T$ observations.

1 Introduction

The combined 2.1 fb$^{-1}$ of 2011 and 2012 LHCb data represent both a great success and a big disappointment: a success because analyses of the data sets led to textbook measurements of $B$-meson observables such as $B_s$-$\bar{B}_s$ mixing and $B \rightarrow K^* \mu^+ \mu^-$; a disappointment because the wealth of data agrees well with the standard model (SM) predictions in essentially all channels, casting doubt on some of the (2-4)$\sigma$ anomalies reported by the $B$ factories and the Tevatron experiments. The success story of LHCb culminated recently in finding the first evidence for the decay $B_s \rightarrow \mu^+ \mu^-$. While the observation of $B_s \rightarrow \mu^+ \mu^-$ with a branching fraction close to the SM expectation excludes the possibility of spectacular new-physics (NP) effects in the $B$-meson sector, deviations of $O(50\%)$, i.e., NP effects of “natural” size, are only started to being probed by LHCb as well as ATLAS and CMS. The measurements of $B \rightarrow K^* \mu^+ \mu^-$ and the evidence for $B_s \rightarrow \mu^+ \mu^-$ hence mark the beginning of the flavor precision era at the LHC.

2 NP in $B_s$-$\bar{B}_s$ mixing

The phenomenon of $B_s$-$\bar{B}_s$ oscillations is encoded in the elements $M_{12}^s$ and $\Gamma_{12}^s$ of the hermitian mass and decay rate matrices. In the SM, $M_{12}^s$ is calculated from the dispersive part of electroweak box diagrams with “off-shell” top quarks, while $\Gamma_{12}^s$ is due to the absorptive part related to “on-shell” light up-type quarks. Sufficiently heavy NP in the off-diagonal element $M_{12}^s$ ($\Gamma_{12}^s$) can be described via effective $\Delta B = 2$ ($\Delta B = 1$) interactions. Schematically, one has

\[
(M_{12}^s)_{\text{NP}} \propto C_2^i \begin{bmatrix} b \\ s \\ b \end{bmatrix} \begin{bmatrix} Q^i_2 \\ Q^i_2 \\ Q^i_2 \end{bmatrix} \sim \frac{1}{\Lambda_{\text{NP}}^2},
\]

\[
(\Gamma_{12}^s)_{\text{NP}} \propto C_1^i C_1^j \text{Im} \begin{bmatrix} b \\ s \\ f' \\ f \\ b \\ s \\ b \\ s \end{bmatrix} \begin{bmatrix} Q^i_1 \\ Q^i_1 \\ f' \\ f \\ Q^i_1 \\ Q^i_1 \end{bmatrix} \sim \frac{1}{(4\pi)^2} \frac{1}{\Lambda_{\text{NP}}^4},
\]

where $\Lambda_{\text{NP}}$ denotes the suppression scale of the higher-dimensional operators.
The off-diagonal elements $M_{12}^s$ and $\Gamma_{12}^s$ can be related to the physical observables in the $B_s$-meson sector by introducing the following model-independent parameterization:

$$M_{12}^s = (M_{12}^s)_{\text{SM}} + (M_{12}^s)_{\text{NP}} = (M_{12}^s)_{\text{SM}} R_M e^{i \phi_M},$$
$$\Gamma_{12}^s = (\Gamma_{12}^s)_{\text{SM}} + (\Gamma_{12}^s)_{\text{NP}} = (\Gamma_{12}^s)_{\text{SM}} R_R e^{i \phi_R}. \quad (2)$$

To leading power in $|\Gamma_{12}^s|/|M_{12}^s|$ the mass difference $\Delta M_s$, the CP-violating phase $\phi_{J/\psi\phi}^s$, the decay-width difference $\Delta \Gamma_s$, and the flavor-specific CP asymmetry $a_{f_s}^s$ are then given by

$$\Delta M_s = (\Delta M_s)_{\text{SM}} R_M, \quad \phi_{J/\psi\phi}^s = (\phi_{J/\psi\phi}^s)_{\text{SM}} + \phi_M,$$
$$\Delta \Gamma_s = (\Delta \Gamma_s)_{\text{SM}} R_R \sin(\phi_M - \phi_R),$$
$$a_{f_s}^s = (a_{f_s}^s)_{\text{SM}} \frac{R_R}{R_M} \frac{\sin(\phi_M - \phi_R)}{\phi_{\text{SM}}^s}, \quad (3)$$

with $\phi_{\text{SM}}^s = (0.22 \pm 0.06) \circ$. \textsuperscript{[4]}

The $\Lambda_{\text{NP}}$-scaling \textsuperscript{[1]} combined with \textsuperscript{[2]} and \textsuperscript{[3]} suggests that new particles are more likely to leave a visible imprint in $\Delta M_s$ and $\phi_{J/\psi\phi}^s$ than in $\Delta \Gamma_s$ and $a_{f_s}^s$. Experimentally, however, the mass difference and the CP phase in $B_s \to B_s^*$ mixing are spot on the SM predictions,\textsuperscript{[41]} while $a_{f_s}^s$, if extracted from the latest measurement\textsuperscript{[5]} of the like-sign dimuon charge asymmetry $A_{0\,\text{SL}}^s$, assuming the absence of NP in the $B_d$-meson system, shows an anomaly of almost 4$\sigma$. In view of these results, it seems worthwhile to ask: how big can NP in $\Gamma_{12}^s$ be afterall?

While any operator $(\bar{s}b) f$ with $f$ leading to a flavor-neutral final state of two or more fields and mass below the bottom-quark threshold can alter $\Gamma_{12}^s$, the possible final states are in practice limited, because essentially all $B_s \to f$ and $B_d \to X_s f$ decay modes involving light states in the final state are strongly constrained experimentally. An exception are $B$ decays to tau pairs.\textsuperscript{[42]}

The possibility of large $b \to s \tau^+\tau^-$ contributions to $\Gamma_{12}^s$ can be analyzed in a model-independent fashion\textsuperscript{[5]} by adding the complete set of dimension-six operators $(A, B = L, R)$

$$Q_{S,AB} = (\bar{s} P_A b) (\bar{\tau} P_B \tau) , \quad Q_{V,AB} = (\bar{s} \gamma\mu P_A b) (\bar{\tau} \gamma\mu P_B \tau) , \quad Q_{T,AB} = (\bar{s} \sigma^{\mu\nu} P_A b) (\bar{\tau} \sigma^{\mu\nu} P_B \tau) , \quad (4)$$

to the SM Lagrangian. Here $P_{L,R} = (1 \mp \gamma_5)/2$ project onto left- and right-handed chiral fields and $\sigma^{\mu\nu} = i [\gamma^\mu, \gamma^\nu]/2$.

The ten operators entering \textsuperscript{[4]} govern the purely leptonic $B_s \to \tau^+\tau^-$ decay, the inclusive semileptonic $B \to X_s \tau^+\tau^-$ decay, and its exclusive counterpart $B^+ \to K^+ \tau^+\tau^-$, making these channels potentially powerful probes of $(\bar{s}b)(\bar{\tau}\tau)$ operators. In practice, these direct constraints however turn out to be rather loose at present. Explicitly, one obtains

$$\text{BR}(B_s \to \tau^+\tau^-) < 3\%,$$
$$\text{BR}(B \to X_s \tau^+\tau^-) < 2.5\%,$$
$$\text{BR}(B^+ \to K^+ \tau^+\tau^-) < 3.3 \cdot 10^{-3}. \quad (5)$$

Here the first limit derives\textsuperscript{[5]} from comparing the SM prediction $\tau_{B_s}/\tau_{B_d} = 1 - [-0.4, 0.0]\%$ with the corresponding experimental result $\tau_{B_s}/\tau_{B_d} = (0.4 \pm 1.9)\%$, while the second (crude) bound follows from estimating\textsuperscript{[5]} the possible contamination of the exclusive and inclusive semileptonic decay samples by $B \to X_s \tau^+\tau^-$ events. The final number corresponds to the 90% confidence level (CL) upper limit on the branching ratio of $B^+ \to K^+ \tau^+\tau^-$ as measured by BaBar\textsuperscript{[43]}.

Further constraints on the Wilson coefficients of the $(\bar{s}b)(\bar{\tau}\tau)$ operators arise indirectly from the experimentally available information on the $b \to s\gamma$, $b \to s\ell^+\ell^-$ ($\ell = e, \mu$), and $b \to s\gamma\gamma$ transitions, because some of the effective operators introduced in \textsuperscript{[4]} mix into the electromagnetic
dipole operators $Q_{7,A}$ and the vector-like semileptonic operators $Q_{9,A}$. An explicit calculation\cite{ref1} shows that the operators $Q_{S,AB}$ mix neither into $Q_{7,A}$ nor $Q_{9,A}$, while $Q_{V,AB}$ ($Q_{T,A}$) mixes only into $Q_{9,A}$ ($Q_{7,A}$). As a result of the particular mixing pattern, the stringent constraints from the radiative decay $B \rightarrow X_s \gamma$ rule out large contributions to $\Gamma_{12}^s$ only if they arise from the tensor operators $Q_{T,A}$. Similarly, the rare decays $B \rightarrow X_s \ell^+\ell^-$ and $B \rightarrow K^{(*)}(\bar{\ell}^+\ell^-)$ primarily limit contributions stemming from the vector operators $Q_{V,AB}$. In contrast to $B \rightarrow X_s \gamma$, all $(\bar{s}b)(\bar{\tau}\tau)$ operators contribute to the double-radiative $B_s \rightarrow \gamma\gamma$ decay at the one-loop level. A detailed study\cite{ref2} shows however that the limits following from $b \rightarrow s\gamma\gamma$ are in practice not competitive with the bounds obtained from the other tree- and loop-level mediated $B_s,d$-meson decays.

The off-diagonal element $\Gamma_{12}^s$ is related via the optical theorem to the absorptive part of the forward-scattering amplitude which converts a $B_s$ into a $B_s$ meson. Working toward leading order in the strong coupling constant and $\Lambda_{QCD}/m_b$, the contributions from the operators (4) to $\Gamma_{12}^s$ is found by computing the matrix elements of the $(Q_i, Q_j)$ double insertions between quark states. Such a calculation\cite{ref3} leads to the following 90\% CL bounds

$$
(R_{\Gamma})_{S,AB} < 1.15, \quad (R_{\Gamma})_{V,AB} < 1.35, \quad (R_{\Gamma})_{T,L} < 1.02, \quad (R_{\Gamma})_{T,R} < 1.04.
$$

These numbers imply that $(s\bar{b})(\bar{\tau}\tau)$ operators of scalar (vector) type can lead to enhancements of $|\Gamma_{12}^s|$ over its SM value by 15\% (35\%) without violating any existing constraint. In contrast, contributions from tensor operators can alter $|\Gamma_{12}^s|$ by at most 4\%. Since an explanation of the experimentally observed large negative values of $\alpha_{s,t}^s$ (or equivalent $A_{t,s}^s$) calls for NP in $\Gamma_{12}^s$ that changes the SM value by a factor of 3 or more, absorptive NP in the form of $(\bar{s}b)(\bar{\tau}\tau)$ operators obviously does not provide a satisfactory description of the large dimuon charge asymmetry observed by the DØ collaboration. This is a model-independent conclusion that can be shown to hold in explicit models of NP with modification of the $(Q_i, Q_j)$ double insertions between quark states or $Z'$ models.\cite{ref4}

### 3 Interplay between high- and low-$p_T$ measurements

The discovery of the Higgs boson with mass of around 125 GeV\cite{ref5,ref6} combined with the direct limits on new particles rule out most of the simple and natural implementations of NP. At this stage, it is hence very important to examine the plethora of available experimental data and to look for hints that can guide us towards special regions where NP may still hide. Indirect probes of NP as provided by rare $B$-meson decays can play a key role in this endeavor, as will be described in the following for the case of the minimal supersymmetric SM (MSSM).

In the decoupling limit of the MSSM, $i.e., M_H^2 \gg M_Z^2$, the lightest CP-even Higgs boson acquires the loop-corrected squared mass\cite{ref7}

$$
M_h^2 \approx M_Z^2 c_{2\beta}^2 + \frac{3 G_F}{\sqrt{2} \pi^2} m_t^4 \left[ - \ln \left( \frac{m_t^2}{m_t^2} \right) + \frac{X_t^2}{m_t^2} \left( 1 - \frac{X_t^2}{12 m_t^2} \right) \right].
$$

Here $c_{2\beta} = \cos (2\beta)$, $m_t^2 = m_{t1}^2 m_{t2}^2$, and $X_t = A_t - \mu/t_\beta$ denotes the stop-mixing parameter, which depends on the trilinear stop-Higgs boson coupling $A_t$ and the higgsino mass parameter $\mu$. From the above expression, one concludes that to raise $M_h$ from the $Z$-boson mass $M_Z$ to 125 GeV requires sizable values of $t_\beta = \tan \beta$, a heavy stop spectrum ($m_t \gtrsim 1$ TeV), and/or large stop mixing ($|A_t| \gtrsim 2$ TeV).

The mentioned MSSM parameters also play an important role in the production and the decay of the Higgs. Applying Higgs low-energy theorems, it is easy to derive that the modification of the $gg \rightarrow h$ production cross section due to stops, can be written as\cite{ref8}

$$
R_h = \frac{\sigma(gg \rightarrow h)_{\text{MSSM}}}{\sigma(gg \rightarrow h)_{\text{SM}}} \approx \left[ 1 + \frac{m_t^2}{4} \left( \frac{1}{m_{t1}^2} + \frac{1}{m_{t2}^2} - \frac{X_t^2}{m_{t1}^2 m_{t2}^2} \right) \right]^2.
$$


the SM expectation. For the choices of (10) implies that for higgsino-like charginos, leading to $\Delta m_{\mu}$ (lower center), and $\Omega h^2$ (lower right). The dotted black lines indicate the parameter regions with $R_\tau > 1$, while the dashed black lines correspond to the 95% CL regions favored by $B \to X_s \gamma$. See text for further explanations.

Figure 1: Predictions for $M_h$ in GeV (upper left), $R_\tau$ (upper center), $R_{X_s}$ (upper right), $R_{\mu^+\mu^-}$ (lower left), $\Delta m_{\mu}$ (lower center), and $\Omega h^2$ (lower right). The dotted black lines indicate the parameter regions with $R_\tau > 1$, while the dashed black lines correspond to the 95% CL regions favored by $B \to X_s \gamma$. See text for further explanations.

One infers that the amount of mixing in the stop sector determines whether $\tilde{\tau}_1$ is smaller or larger than $1$. For no mixing ($X_t = 0$), one has $R_h > 1$, while if $X_t$ is parametrically larger than the mass eigenvalues $m_{\tilde{t}_1,2}$ (with $m_{\tilde{t}_1} < m_{\tilde{t}_2}$) then one has $R_h < 1$. The fact that in the MSSM, to make the Higgs sufficiently heavy, one needs large/maximal mixing, then implies that for a random MSSM parameter point with $M_h \approx 125$ GeV one should find a suppression of $\sigma(gg \to h)$. In fact, this is precisely what happens in large parts of the parameter space.

For $M_A > M_Z$ and $t_\beta > 1$, charged Higgs as well as chargino effects are strongly suppressed in $h \to \gamma\gamma$, but stau loops can have a notable impact on the diphoton signal strength

$$R_\gamma = \frac{\sigma(pp \to h) BR(h \to \gamma\gamma)}{\sigma(pp \to h) BR(h \to \gamma\gamma)}_{\text{SM}} \approx 1 + 0.10 \frac{m_{\tilde{\tau}_1}^2 X_t^2}{m_{\tilde{\tau}_1}^2 m_{\tilde{\tau}_2}^2},$$

if mixing in the stau sector is large, i.e., $|X_{\tau}| = |A_\tau - \mu t_\beta| \gg m_{\tilde{\tau}_1,2}$ and the stau mass eigenstate $\tilde{\tau}_1$ is sufficiently light. In (9) we have assumed that $M_h \approx 125$ GeV and explicitly included only stau effects.

The above discussion should have made clear that the part of the MSSM parameter space with $M_A, t_\beta \to \infty$ represents a phenomenologically interesting region for Higgs physics. How do the dominant MSSM corrections to $B \to X_s \gamma$ and $B_s \to \mu^+\mu^-$ look like in this limiting case? For $B \to X_s \gamma$, one obtains

$$R_{X_s} = \frac{\text{BR}(B \to X_s \gamma)_{\text{MSSM}}}{\text{BR}(B \to X_s \gamma)_{\text{SM}}} \approx 1 - 2.61 \Delta C_7 + 1.66 (\Delta C_7)^2,$$

where the dominant one-loop contributions are provided by diagrams with top squarks and higgsino-like charginos, leading to $\Delta C_7 \propto -\mu A_t t_\beta m_{\tilde{\tau}_1}^2 / m_{\tilde{\tau}_1}^2$. The sign of $\Delta C_7$ combined with that of (10) implies that for $\mu A_t > 0$ ($\mu A_t < 0$) the MSSM branching ratio is larger (smaller) than the SM expectation. For the choices $t_\beta = 50$, $m_{\tilde{t}} = 1.5$ TeV, $|\mu| = 1$ TeV, and $|A_t| = 3$ TeV,
one finds numerically an enhancement/suppression of $\text{BR}(B \to X_s\gamma)_{\text{MSSM}}$ by $\mathcal{O}(30\%)$. Shifts of this size are detectable given the present theoretical calculations and experimental extractions.

In the case of the purely leptonic $B_s$ decay, one can write

$$R_{\mu^+\mu^-} = \frac{\text{BR}(B_s \to \mu^+\mu^-)_{\text{MSSM}}}{\text{BR}(B_s \to \mu^+\mu^-)_{\text{SM}}} \approx 1 - 13.2 C_P + 43.6 \left(C_S^2 + C_P^2\right), \quad (11)$$

where $C_{P,S}$ denote the dimensionless Wilson coefficients of the semileptonic pseudo-scalar and scalar operators. The term linear in $C_P$ arises from the interference with the SM contribution to the semileptonic axial-vector operator. For $C_P > 0$ it interferes destructively with the term proportional to $(C_S^2 + C_P^2)$, which implies that a pseudo-scalar contribution of the correct sign and size will lead to a suppression of $\text{BR}(B_s \to \mu^+\mu^-)_{\text{MSSM}}$ below its SM value.

Within the MSSM the contributions to $C_{P,S}$ with the strongest $t_\beta$ dependence arise from neutral Higgs double penguins. In the decoupling limit, one has

$$C_P \approx -C_S \propto \mu A_t \left(\frac{t_\beta^2}{1 + \epsilon_b t_\beta}\right) \frac{m_{\tilde{b}}^2}{m_t^2} \frac{mgm_\mu}{M_W^2 M_A^2}, \quad \epsilon_b \propto \frac{\alpha_s \mu M_3}{\pi m_b^2}. \quad (12)$$

Here $\epsilon_b$ encodes loop-induced non-holomorphic terms due to gluino exchange which introduce a dependence of $C_{P,S}$ on $\text{sgn}(\mu M_3)$. From (12) one deduces that the sign of $C_S (C_P)$ is opposite to (follows) that of $\mu A_t$ and that both coefficients are suppressed (enhanced) for $\mu M_3 > 0$ ($\mu M_3 < 0$). The recent LHCb measurement $^9$ of $\text{BR}(B_s \to \mu^+\mu^-) = (3.2^{+1.3}_{-1.2}) \cdot 10^{-9}$ hence clearly favors $\mu A_t > 0$ and $\mu M_3 > 0$, if Higgs exchange gives the dominant contribution to (11).

In Figure 1 we show the results of numerical scans in the $A_t-\mu$ plane. We restrict ourselves to the quadrant with $A_t > 0$ and $\mu > 0$ since it shows the most interesting effects and correlations. The predictions correspond to $t_\beta = 60$, $M_A = 1$ TeV, $M_1 = 50$ GeV, $M_2 = 300$ GeV, $M_3 = 1.2$ TeV, $m_{\tilde{Q}_3} = m_{\tilde{u}_3} = 1.5$ TeV, $m_{\tilde{L}_3} = m_{\tilde{t}_3} = 350$ GeV, while we take common soft masses of $1.5$ TeV and $2$ TeV ($1$ TeV) for the other “left-” and “right-handed” squark (sleptons). We furthermore employ $A_b = 2.5$ TeV and $A_t = 500$ GeV, while the first and second generation trilinear couplings take the same values as those of the third generation. We see that for the above choice of parameters, $A_t$ has to lie in the range of $[2, 5]$ TeV to accommodate $M_h \in [123, 129]$ GeV and that $R_\gamma > 1$ can only be obtained in a narrow sliver around $\mu = 900$ GeV. For our choice of soft masses $m_{\tilde{L}_3}$ and $m_{\tilde{t}_3}$ such large $\mu$ values lead to stau masses $m_{\tilde{\tau}_1} \in [80, 120]$ GeV. This strong correlation between notable enhancements in $R_\gamma$ and the presence of a very light stau is a smoking gun signal of the discussed scenario. From the panel showing $R_{X_s\gamma}$, one infers that $\text{BR}(B \to X_s\gamma)$ is always enhanced (by about 20% to 60%) with respect to the SM. This is an interesting and potentially important finding, since the $B \to X_s\gamma$ constraint starts cutting into the already narrow region in the $A_t-\mu$ plane with $M_h \approx 125$ GeV and $R_\gamma > 1$. We also observe that solutions that satisfy $B \to X_s\gamma$ typically feature a suppression of $B_s \to \mu^+\mu^-$. In fact, asking for an agreement with $R_{X_s\gamma}$ at the 95% CL as well as $R_\gamma > 1$, implies $R_{\mu^+\mu^-} \in [0.6, 1.0]$. The figure furthermore shows that assuming a light slepton spectrum, the long-standing discrepancy of the anomalous magnetic moment of the muon $a_\mu$ is reduced in the parameter region selected by $M_h$ and the enhanced diphoton signal. Finally, notice that in this very region of parameter space also the correct thermal dark matter relic density $\Omega h^2$ can be achieved, but only if one assumes the hierarchy $|M_1| \ll |M_2| \ll |\mu|$. A typical MSSM spectrum leading to a significantly enhanced $h \to \gamma\gamma$ rate as well as the correct value of $\Omega h^2$, hence contains a light bino as the dark matter candidate, a light and maximally mixed stau, and a heavy higgsino. The aforementioned deviations and found correlations could be tested in the near future and hence may become very valuable as guidelines and consistency checks, in particular if the high-$p_T$ LHC experiments start to see supersymmetric partners.
4 Conclusions

While the 7 TeV and 8 TeV LHC runs took the hope to find spectacular effects in heavy flavor physics, the low-\(p_T\) data is at present not precise enough to exclude NP contaminations of \(\mathcal{O}(50\%)\). This still leaves room for visible and interesting effects in rare \(B\)-meson decays, given the theoretical cleanness of these observables. Since there is no direct sign of NP at the LHC and also the Higgs signal strengths look very much SM-like, such indirect NP probes are more important than ever and only the synergy between high- and low-\(p_T\) observations may give us the key to solving the puzzles of fundamental physics. The expected LHC precision measurements of the \(B\)-mixing observables, \(B_s \rightarrow \mu^+\mu^-\), \(B \rightarrow K^\ast \ell^+\ell^-\), the angle \(\gamma\), etc. may play a crucial role in this endeavor.

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