Neckarzimmern *B* physics workshop 2017

Effective theories for flavour within & beyond the SM

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Outline

1. EFT for weak decays
   - Introduction
   - Example: BR($B_s \to \mu^+\mu^-$) in the SM
2. Applications to BSM phenomenology
   - Example: $B_s \to \mu^+\mu^-$ beyond the SM
   - Standard Model as EFT
   - Model-independent analysis of $B_s \to \mu^+\mu^-$
   - $B_s \to \mu^+\mu^-$ in specific NP models

1. EFT for weak decays
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   - Example: BR($B_s \to \mu^+\mu^-$) in the SM
2. Applications to BSM phenomenology

Multiple scales

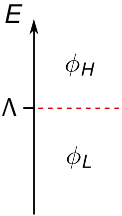
Weak decays always involve physics at vastly disparate energy scales, e.g.

- non-perturbative QCD interactions describing the hadrons, $\Lambda_{\text{QCD}} \sim 0.2$ GeV
- b quark mass ~ 4 GeV
- mass of the W mediating FCNCs ~ 80 GeV
- top quark mass ~ 170 GeV
- new heavy particles in loops $\sim$ TeV?

Such a multitude of scales can only be tackled with a powerful tool:
Effective field theory

Effective field theory

We want to study physics at energies much lower than some scale Λ in a theory where particles lighter and heavier than Λ are present.



To this end, we can replace the complicated Lagrangian of the “full” theory by an **effective Lagrangian** containing only the light fields and a series of local operators built out of the light fields

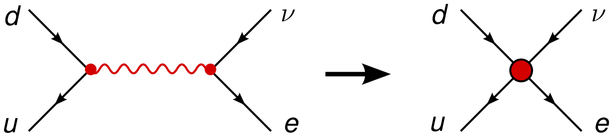
$$\mathcal{L}(\varphi_L, \varphi_H) \to \mathcal{L}(\varphi_L) + \mathcal{L}_{\text{eff}} = \mathcal{L}(\varphi_L) + \sum_i C_i\, Q_i(\varphi_L)$$

This expansion is called the **operator product expansion**

Example: modern view of Fermi theory

In Fermi's model of β decay, the full weak Lagrangian (that he didn't know of course) is effectively replaced by the low-energy (QED) Lagrangian plus a single operator

$$\mathcal{L}_{\text{ew}} \to \mathcal{L}_{\text{QED}} + \frac{G_F}{\sqrt{2}}(\bar{u}d)(e\bar{\nu})$$



Local operator $\equiv$ effective vertex!

More about the OPE

$$\mathcal{L}_{\text{eff}} = \sum_i C_i\, Q_i(\varphi_L)$$

- the local operators have mass dimension > 4, i.e. they are **non-renormalizable**
- operators with dimension $4 + n$ contribute with strength $(E/\Lambda)^n$ to a process with energy E; thus, the OPE can be **truncated** at some dimension d and typically a small number of operators is important
- C_i are called **Wilson coefficients** $\equiv$ effective coupling constants
- $\mathcal{H}_{\text{eff}} \equiv -\mathcal{L}_{\text{eff}}$

Weak effective Hamiltonian

$$\mathcal{H}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} \sum_i \xi^i_{\text{CKM}}\, C_i\, Q_i$$

- we only have to consider operators up to **dimension 6**
- since flavour-change is always mediated by the W boson, one can factor out the Fermi constant $\frac{4G_F}{\sqrt{2}} = \frac{g^2}{2m_W^2}$ $\Rightarrow$ the WC of dimension-6 operators are **dimensionless**
- factoring out the CKM elements, the WC are **real** in the SM
- the amplitude of a weak decay takes the generic form

$$A(i \to f) = \langle f|\mathcal{H}_{\text{eff}}|i\rangle = \frac{4G_F}{\sqrt{2}} \sum_i \xi^i_{\text{CKM}}\, C_i(\mu)\, \langle f|Q_i(\mu)|i\rangle$$

Calculating Wilson coefficients: matching

- The values of the effective coupling constant should be such that amplitudes in the effective theory reproduce the ones in the full theory.

$$\mathcal{L}_{\text{eff}} = \frac{4G_F}{\sqrt{2}} V_{ud}\, C\, Q$$
$$Q = (\bar{u}_L \gamma^\mu d_L)(\bar{e}_L \gamma_\mu \nu_L)$$

- Requiring the amplitudes to coincide, one finds

$$\mathcal{A}_{\text{full}} = \frac{g^2}{2m_W^2} V_{ud} \langle Q \rangle \stackrel{!}{=} \frac{4G_F}{\sqrt{2}} V_{ud} C \langle Q \rangle = \mathcal{A}_{\text{eff}}$$
$$\Rightarrow C = 1$$

- This process is called **matching**

Detour: renormalization

In a QFT, infinities in calculations have to be removed by **renormalizing** bare parameters in the Lagrangian, e.g. in QCD

$$G^a_{0\mu} = \sqrt{Z_3} G^a_\mu \quad q_0 = \sqrt{Z_q} q \quad g_{0s} = Z_g g_s \mu^\varepsilon \quad m_0 = Z_m m$$

- 0: unrenornormalized = bare fields/parameters
- Z_i: renormalization constants
- μ: renormalization scale

Dimensional regularization + minimal subtraction: only poles in $\varepsilon = 2 - d/2$ subtracted

$$Z_i = \frac{\alpha_s}{4\pi} \frac{a_{1i}}{\varepsilon} + O(\alpha_s^2)$$

Renormalization scale

- g_s and m (in fact, all couplings in a QFT) become μ-dependent
- μ dependent terms are renormalization scheme dependent
- physical **observables** have to be μ**-independent**
- values of the parameters at different scales are connected by **renormalization group equations (RGE)**, e.g.

$$\frac{dg_s(\mu)}{d\ln\mu} = \beta(g_s(\mu))\,, \quad \frac{dm(\mu)}{d\ln\mu} = -\gamma_m(g_s(\mu))\,m(\mu)$$

Operator renormalization

- Also Q_i have to be renormalized:

$$Q_i^0 = Z_{ij} Q_j$$

- C_i become scale-dependent
- The scale dependence is cancelled by the scale dependence of the matrix element

$$A(i \to f) = \frac{G_F}{\sqrt{2}} \sum_i \xi^i_{\text{CKM}} \, C_i(\mu) \, \langle f | Q_i(\mu) | i \rangle$$

Renormalizing non-renormalizable op.s?

- In a **renormalizable** theory, all infinities can be removed to all orders in perturbation theory by a finite number of counterterms
- In a **non-renormalizable** theory, infinities have to be removed at any order and the number of subtraction terms is infinite
- In the OPE, once we renormalize the (finite number of) operators of a given dimension, the higher-dimensional ones are still divergent, but we don't care since they are **not relevant for low-energy physics**

Modern view of renormalization: the low-energy limit of every EFT is a renormalizable theory

RGE for the Wilson coefficients

- Recall our multi-scale problem: which renormalization scale μ to choose for $C_i(\mu)$?
- C_i obey a RGE

$$\frac{d}{d\ln\mu}C_j(\mu) = \sum_i C_i(\mu)\,\gamma_{ij}(\mu)$$
$$\Rightarrow C_j(\mu_1) = U_{ji}(\mu_1,\mu_2)\,C_i(\mu_2)$$

- we calculate (match) C_i at a high scale where QCD is perturbative and use the RGE to evolve it down to the appropriate scale
- by "running" the RGEs, we are in effect running through a series of EFTs where μ playing the role of the scale Λ

Generic weak decay amplitude

$$A(i \to f) = \frac{4G_F}{\sqrt{2}} \sum_i \xi^i_{\rm CKM} \quad C_i(m_W) \quad U(\mu_l, m_W) \quad \langle f|Q_i(\mu_l)|i\rangle$$

CKM factors	short-distance	QCD corrections	hadronic matrix element
	— perturbative —		non-perturbative
	— indep. of external states —		specific for ext. state
	sensitive to NP	— independent of NP —	

The OPE has achieved a **separation of scales**

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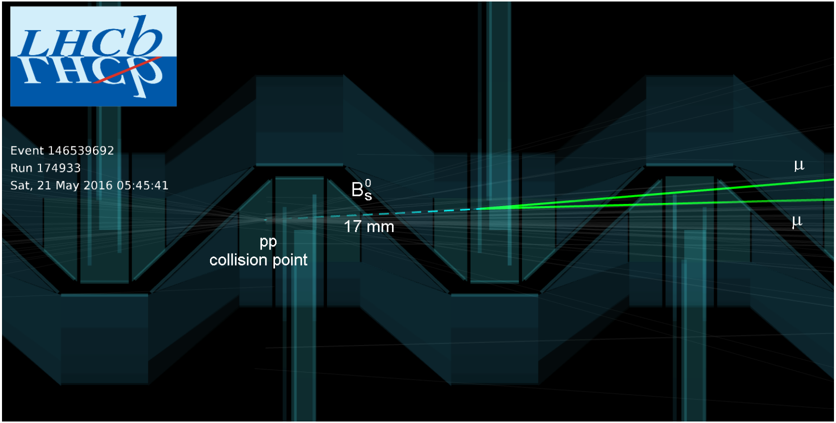
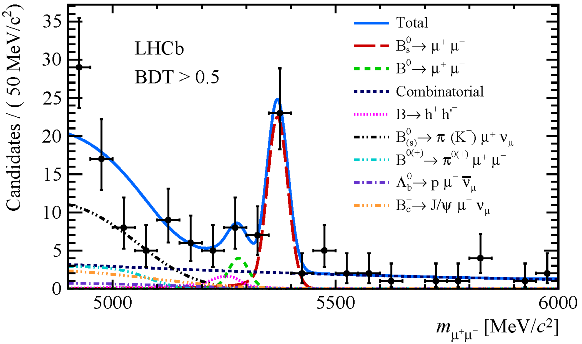

Event 146539692
Run 174933
Sat, 21 May 2016 05:45:41
B_s^0
17 mm
pp
collision point
μ
μ

$B_s \to \mu^+\mu^-$ LHCb Run 2



$$\mathrm{BR}(B_s \to \mu^+\mu^-) = (3.0 \pm 0.6^{+0.3}_{-0.2}) \times 10^{-9}$$ Aaij et al. 1703.05747

What does this imply for the SM; for NP models?

$B_s \to \mu^+\mu^-$ branching ratio in the SM

$$
\begin{aligned}
\text{BR}(B_s \to \mu^+\mu^-) &= \Gamma(B_s \to \mu^+\mu^-)/\Gamma(B_s \to \text{anything}) \\
&= \tau_{B_s}\,\Gamma(B_s \to \mu^+\mu^-) \\
&= \tau_{B_s}\,\Phi(m_{B_s}, m_\mu)\,|\langle \mu\mu|A|B_s\rangle|^2
\end{aligned}
$$

- $\tau_{B_s} = 1/\Gamma_s$ – lifetime
- Φ – phase space
- A – amplitude

$$
\Phi(m_{B_s}, m_\mu) = \frac{1}{16\pi}\frac{1}{m_{B_s}}\sqrt{1 - \frac{4m_\mu^2}{m_{B_s}^2}}
$$

$B_s \to \mu^+\mu^-$ amplitude in the EFT

$$\langle \mu\mu | A | B_s \rangle = i \langle \mu\mu | C_{10} O_{10} | B_s \rangle + O(m_b^2/m_W^2)$$

- $O_{10} = (\bar{s}_L \gamma^\mu b_L)(\bar{\mu}\gamma_\mu \gamma_5 \mu)$ – semi-leptonic axial vector operator
- C_{10} – Wilson coefficient

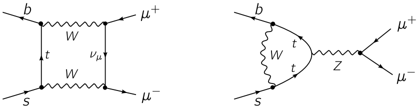
$$\langle \mu\mu | C_{10} O_{10} | B_s \rangle = C_{10} \langle 0 | \bar{s}_L \gamma^\mu b_L | B_s \rangle (\bar{\mu}\gamma_\mu \gamma_5 \mu)$$

- $\langle 0 | \bar{s}_L \gamma^\mu b_L | B_s \rangle$ – hadronic matrix element

$$\langle 0 | \bar{s}_L \gamma^\mu b_L | B_s \rangle = \frac{1}{2} \langle 0 | \bar{s} \gamma^\mu b | B_s \rangle - \frac{1}{2} \langle 0 | \bar{s} \gamma^\mu \gamma_5 b | B_s \rangle = 0 - \frac{1}{2} i f_{B_s} p^\mu$$

- f_{B_s} decay constant

$B_s \to \mu^+\mu^-$: matching



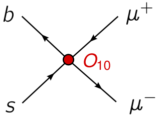
“Partonic” amplitude in the full theory

$$\mathcal{A}_{\text{full}} = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} \frac{1}{s_W^2} V_{tb} V_{ts}^* \, Y(x_t) (\bar{s}_L \gamma^\mu b_L)(\bar{\mu}\gamma_\mu\gamma_5\mu)$$

- G_F – Fermi constant
- V_{tq} – CKM elements
- $x_t = m_t^2/m_W^2$
- Y – Inami-Lim function

$$Y(x_t) = Y_0(x_t)\left[1 + O(\alpha_s) + O(\alpha_s^2) + O(\alpha_{em}) + \ldots\right]$$

$B_s \to \mu^+\mu^-$: matching



"Partonic" amplitude in the effective theory

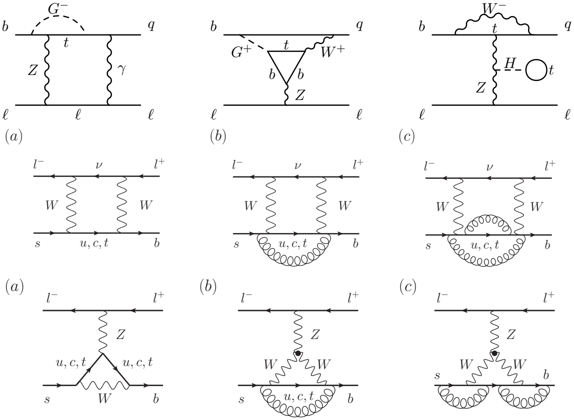
$$\mathcal{H}_{\text{eff}} = -C_{10} O_{10}$$

$$\mathcal{A}_{\text{eff}} = -C_{10} (\bar{s}_L \gamma^\mu b_L)(\bar{\mu}\gamma_\mu\gamma_5\mu)$$

Matching

$$\mathcal{A}_{\text{full}} \stackrel{!}{=} \mathcal{A}_{\text{eff}} \Rightarrow C_{10} = -\frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} V_{tb} V_{ts}^* \frac{1}{s_w^2} Y(x_t)$$

Higher order corrections: full theory



Bobeth et al. 1311.1348, Hermann et al. 1311.1347

Recipe: how to predict $\text{BR}(B_s \to \mu^+\mu^-)_{\text{SM}}$

$$\text{BR}(B_s \to \mu^+\mu^-)_{\text{SM}} = \tau_{B_s} \frac{G_F^2}{\pi} \left(\frac{\alpha}{4\pi s_w^2}\right)^2 m_\mu^2 \sqrt{1 - \frac{4m_\mu^2}{m_{B_s}^2}}\, m_{B_s} f_{B_s}^2 \,|V_{tb}V_{ts}^*|^2\, Y(x_t)^2$$

- Liftetime τ_{B_s}: take from experiment

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- $G_F, \alpha, s_w, m_{B_s}^2, m_\mu$: take from PDG

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- $Y(x_t)$: include NNLO QCD and NLO EW corrections and RG evolution

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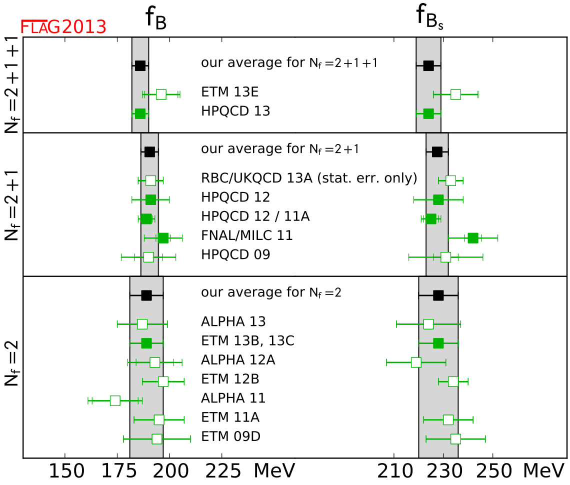
- $f_{B_s}^2$: from lattice QCD
- $|V_{tb}V_{ts}^*|^2$: from experiment

Lattice determinations of f_{B_s}

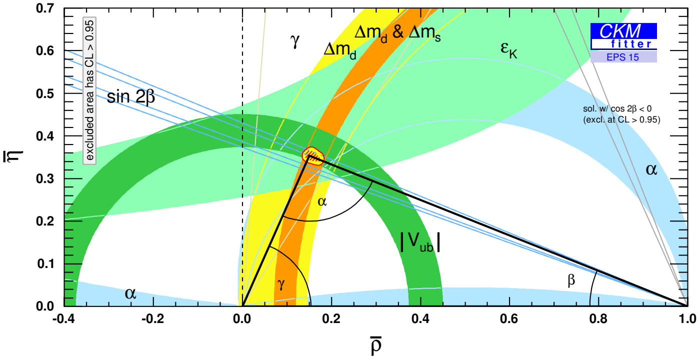


Determining $|V_{tb}V_{ts}^*|$

- There is no direct measurement of V_{ts}
- But CKM elements can be extracted from a global fit of the CKM matrix

$$|V_{tb}V_{ts}^*| = A\lambda^2 \left[1 + \lambda^2 \left(\bar{\rho} - \frac{1}{2}\right)\right] + O(\lambda^6)$$

Global CKM fits



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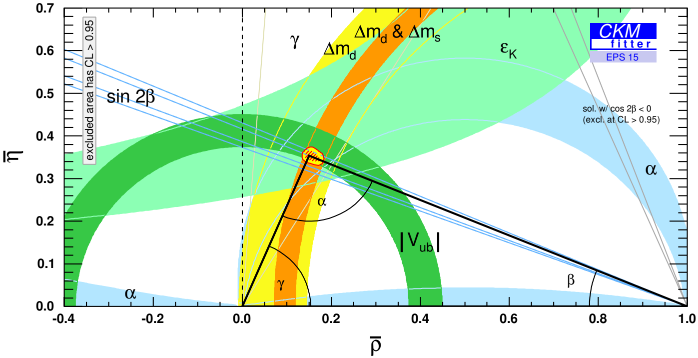


Using the global fit result assumes that neutral meson mixing is free from physics BSM

Using tree-level CKM determinations

- $|V_{cb}|$ from inclusive & exclusive $b \to c\ell\nu$
- $|V_{ub}|$ from inclusive & exclusive $b \to u\ell\nu$
- $|V_{us}|$ from $K \to \pi\ell\nu$
- γ from $B \to DK$

$$|V_{tb}V_{ts}^*| = |V_{cb}|\left(1 - \frac{|V_{us}|^2}{2} + \frac{|V_{ub}|}{|V_{cb}|}|V_{us}|\cos\gamma\right) \approx |V_{cb}|(1 - 0.025 + 0.007)$$

Subtle issue: B_s lifetime difference

Due to B_s-$\bar{B}_s$ mixing, there is a sizable lifetime difference between the two B_s mass eigenstates:

$$\tau_{B_s^L} = \Gamma_{B_s^L}^{-1} = 1.42\,\text{ps} \qquad \tau_{B_s^H} = \Gamma_{B_s^H}^{-1} = 1.61\,\text{ps}$$

$$\tau_{B_s} = \Gamma_{B_s}^{-1} = \left[\frac{1}{2}\left(\Gamma_{B_s^L} + \Gamma_{B_s^H}\right)\right]^{-1}$$

Which lifetime should we use?

Time-dependent untagged decay rate

$$\Gamma(B_s(t) \to \mu^+\mu^-) = R_H\, e^{-t/\tau_{B_s^H}} + R_L\, e^{-t/\tau_{B_s^L}}$$

So far, we have computed

$$\text{BR}(B_s \to \mu^+\mu^-) = \frac{\tau_{B_s}}{2}\, \Gamma(B_s(t=0) \to \mu^+\mu^-)$$

But experiments actually measure

$$\overline{\text{BR}}(B_s \to \mu^+\mu^-) = \frac{1}{2}\int_0^\infty \Gamma(B_s(t) \to \mu^+\mu^-)\, dt$$

It turns out that De Bruyn et al. 1204.1737

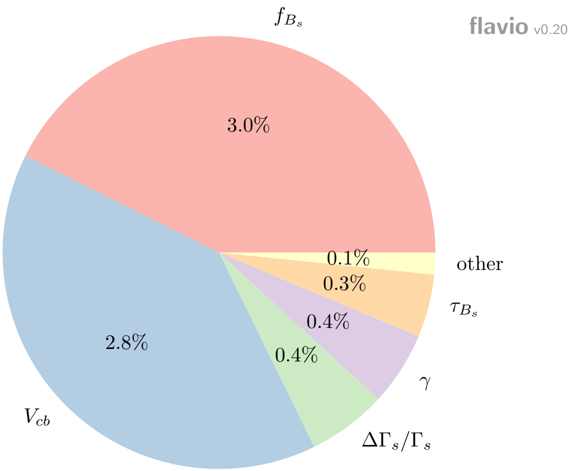
$$\frac{\overline{\text{BR}}(B_s \to \mu^+\mu^-)}{\text{BR}(B_s \to \mu^+\mu^-)} = \frac{\tau_{B_s^H}}{\tau_{B_s}}$$

Result: $B_s \to \mu^+\mu^-$ SM vs. experiment

$$\overline{\text{BR}}(B_s \to \mu^+\mu^-)_{\text{LHCb}}(3.0 \pm 0.6^{+0.3}_{-0.2}) \times 10^{-9}$$

$$\overline{\text{BR}}(B_s \to \mu^+\mu^-)_{\text{SM}} = (3.59 \pm 0.18) \times 10^{-9}$$

BR($B_s \to \mu^+\mu^-$) error budget



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2 Applications to BSM phenomenology

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Physics beyond the SM in $B_s \to \mu^+\mu^-$

Assuming no new particles below 5 GeV, new physics does **not** affect

- Matrix element (f_{B_s})
- CKM extraction based on tree-level decays *
- QCD corrections
- Phase space

All “short-distance” physics enters through modified **Wilson coefficients**

* see however Brod et al. 1412.1446

All possible contributing operators ($D = 6$)

$$O_{10} = (\bar{s}_L \gamma^\mu b_L)(\bar{\mu}\gamma_\mu \gamma_5 \mu) \qquad O'_{10} = (\bar{s}_R \gamma^\mu b_R)(\bar{\mu}\gamma_\mu \gamma_5 \mu)$$
$$O_S = m_b(\bar{s}_R b_L)(\bar{\mu}\mu) \qquad O'_S = m_b(\bar{s}_L b_R)(\bar{\mu}\mu)$$
$$O_P = m_b(\bar{s}_R b_L)(\bar{\mu}\gamma_5 \mu) \qquad O'_P = m_b(\bar{s}_L b_R)(\bar{\mu}\gamma_5 \mu)$$

- In the SM, $C'_{10} = C_S = C'_S = C_P = C'_P = 0$
- f_{B_s} remains the only required matrix element because

$$\langle 0|\bar{s}\gamma_\mu\gamma_5 b|\bar{B}_s\rangle = ip^\mu f_{B_s}, \qquad \langle 0|\bar{s}\gamma_5 b|\bar{B}_s\rangle = -\frac{if_{B_s} m_{B_s}^2}{m_b + m_s},$$

- Other operators (tensor, dipole) have vanishing matrix elements

Branching ratio beyond the SM

$$\mathrm{BR}(B_s \to \mu^+\mu^-) = \mathrm{BR}(B_s \to \mu^+\mu^-)_{\mathrm{SM}} \left[|A|^2 + |B|^2 \left(1 - \frac{4m_\mu^2}{m_{B_s}^2} \right) \right]$$

$$A = \frac{1}{C_{10}^{\mathrm{SM}}} \left[(C_{10} - C'_{10}) + \frac{m_{B_s}^2}{2m_\mu} (C_P - C'_P) \right]$$

$$B = \frac{1}{C_{10}^{\mathrm{SM}}} \left[\frac{m_{B_s}^2}{2m_\mu} (C_S - C'_S) \right]$$

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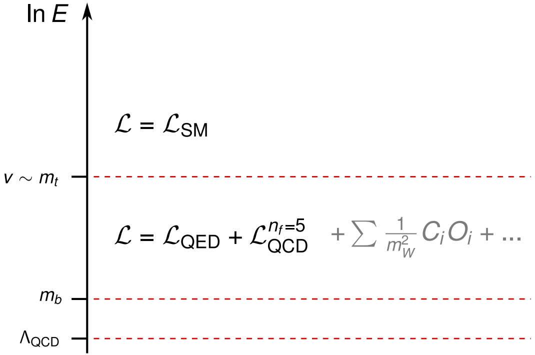
Hierarchy of effective theories



- $m_b \ll v$: weak effective Hamiltonian

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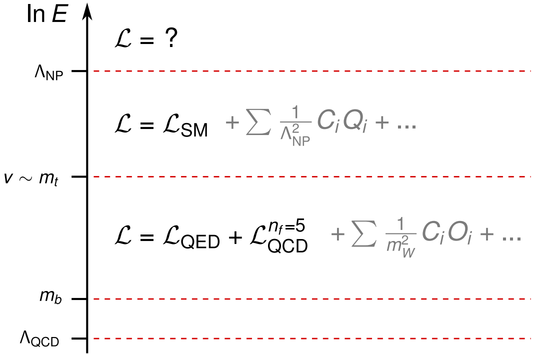

- $m_b \ll v$: weak effective Hamiltonian
- $v \ll \Lambda_{\rm NP}$: “SMEFT”

SMEFT operators matching onto $O_{10}^{(\prime)}$

$$
\begin{aligned}
Q_{Hq}^{(1)} &= \left(H^\dagger\, iD_\mu H\right)\left(\bar{q}_s \gamma^\mu q_b\right)\\
Q_{Hq}^{(3)} &= H^\dagger\, iD_\mu^I H(\bar{q}_s \tau^I \gamma^\mu q_b)\\
Q_{Hd} &= \left(H^\dagger\, iD_\mu H\right)\left(\bar{s}_R \gamma^\mu b_R\right)
\end{aligned}
$$

$$
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Q_{\ell q}^{(1)} &= (\bar{\ell}\gamma_\mu \ell)(\bar{q}_s \gamma^\mu q_b), & Q_{\ell q}^{(3)} &= (\bar{\ell}\gamma_\mu \tau^I \ell)(\bar{q}_s \gamma^\mu \tau^I q_b),\\
Q_{ed} &= (\bar{l}_R \gamma_\mu l_R)(\bar{s}\gamma^\mu b_R), & Q_{\ell d} &= (\bar{\ell}\gamma_\mu \ell)(\bar{s}\gamma^\mu b_R),\\
Q_{qe} &= (\bar{q}_s \gamma_\mu q_b)(\bar{l}_R \gamma^\mu l_R)
\end{aligned}
$$

$$
\begin{aligned}
C_{10}^{\text{NP}} &= C_{qe} - C_{\ell q}^{(1)} - C_{\ell q}^{(3)} + (C_{Hq}^{(1)} + C_{Hq}^{(3)})\\
C_{10}^{\prime\text{NP}} &= C_{ed} - C_{\ell d} + C_{Hd}
\end{aligned}
$$

SMEFT operators matching onto $O_{10}^{(\prime)}$

$$Q_{Hq}^{(1)} = \left(H^\dagger\, iD_\mu H\right)\left(\bar q_s \gamma^\mu q_b\right)$$
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$$Q_{Hd} = \left(H^\dagger\, iD_\mu H\right)\left(\bar s_R \gamma^\mu b_R\right)$$

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$$Q_{ed} = (\bar l_R\gamma_\mu l_R)(\bar s \gamma^\mu b_R), \qquad Q_{\ell d} = (\bar\ell\gamma_\mu \ell)(\bar s \gamma^\mu b_R),$$
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For C_{10}, we have not gained anything

SMEFT operators matching onto $O_{S,P}^{(\prime)}$

$$Q_{\ell edq}^{ijkl} = (\bar{l}_a^i e^j)(\bar{d}^k q_a^l)$$

$$C_S = -C_P = C_{\ell edq}^{3323}$$

$$C_S' = C_P' = C_{\ell edq}^{3332}$$

SMEFT operators matching onto $O_{S,P}^{(\prime)}$

$$Q_{\ell edq}^{ijkl} = (\bar{l}_a^i e^j)(\bar{d}^k q_a^l)$$

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- At dimension 6 in the SMEFT, there are **only 2** independent scalar/pseudoscalar operators (as opposed to 4 in the low-energy EFT).
- The SM gauge symmetries **restrict** the form of scalar NP contributions (valid if $\Lambda_{\text{NP}} \gg v$)

Alonso et al. 1407.7044

1 EFT for weak decays
- Introduction
- Example: BR($B_s \to \mu^+\mu^-$) in the SM

2 Applications to BSM phenomenology
- Example: $B_s \to \mu^+\mu^-$ beyond the SM
- Standard Model as EFT
- Model-independent analysis of $B_s \to \mu^+\mu^-$
- $B_s \to \mu^+\mu^-$ in specific NP models

Fitting the Wilson coefficients

- We can obtain model-independent constraints on new physics by considering the x^2 function

$$x^2(C_i) = \frac{(x(C_i) - x_{\text{exp}})^2}{\sigma_{exp}^2 + \sigma_{\text{th}}^2}$$

where $x = \overline{\text{BR}}(B_s \to \mu^+\mu^-)$ and C_i are the Wilson coefficients.

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$$x^2(C) - x^2(C^*) < 1 \; (< 4)$$

where C^* is the value that minimizes x^2.

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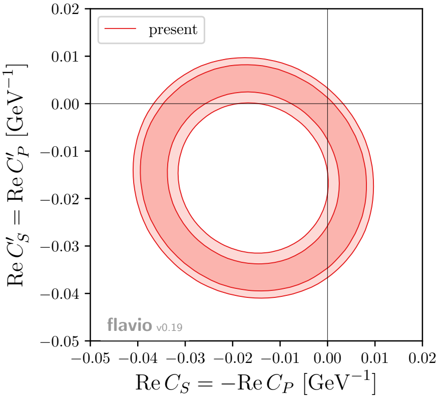
  where C^* is the value that minimizes x^2.
- For two coefficients, the 1σ (2σ) regions are given by

$$x^2(\vec{C}) - x^2(\vec{C}^*) < 2.3 \; (< 6)$$

Fit results



Constraint on real part of scalar operators imposing SMEFT relation

Altmannshofer et al. 1702.05498

Advertisement break: `flavio`

- A Python package for flavour phenomenology in the SM & beyond
  - repository: `http://github.com/flav-io/flavio`
  - documentation: `http://flav-io.github.io`
- Features
  - SM predictions with uncertainties
  - NP predictions for arbitrary Wilson coefficients
  - Fitting SM parameters and Wilson coefficients to data
  - Plotting library to visualize fit results

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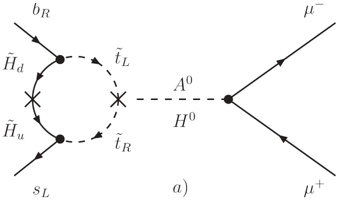
Minimal Supersymmetric SM (MSSM)

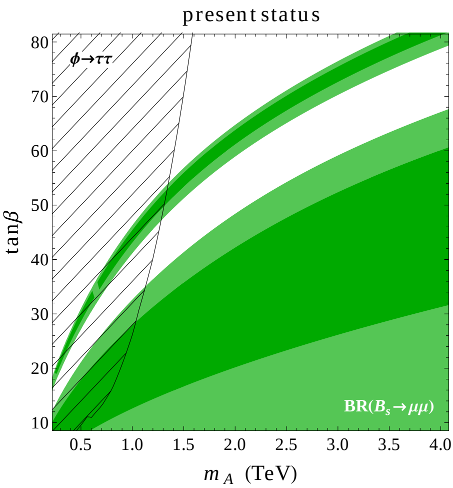
Even for a degenerate spectrum: Higgsino contribution

$$C_S \approx -C_P \approx \frac{G_F^2 m_t^2}{8\pi^2} \frac{m_\mu}{m_A^2} \frac{A_t \mu \tan^3\beta}{m_{\tilde{t}}^2} f\left(\frac{\mu^2}{m_{\tilde{t}}^2}\right)$$



MSSM with Minimal Flavour Violation: $C'_{S,P} = 0$, $C_S \approx -C_P \in \mathbb{R}$

Constraints on parameter space



- $B_s \to \mu^+\mu^-$ is complementary to Higgs physics ($H/A \to \tau^+\tau^-$)
- Two disjoint solutions corresponding to different overall signs of the amplitude. How to disentangle?

The $B_s \to \mu^+\mu^-$ time-dependent rate

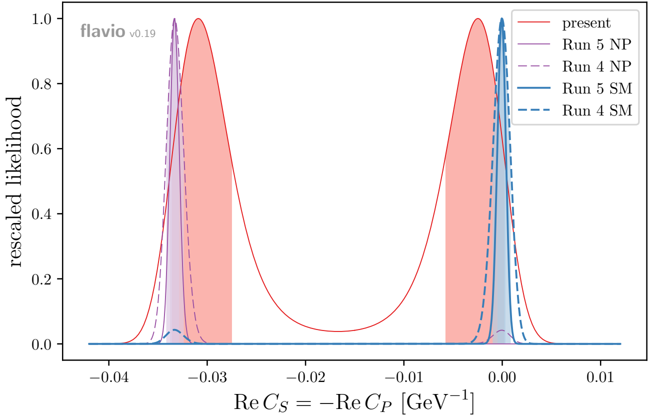
The non-zero B_s^H-B_s^L lifetime difference, $y_s = \frac{\Delta\Gamma_s}{2\Gamma_s} \approx 7\%$ gives access to an additional observable

$$\Gamma(B_s(t) \to \mu^+\mu^-) + \Gamma(\bar{B}_s(t) \to \mu^+\mu^-) \propto \left[\cosh\left(\frac{y_s t}{\tau_{B_s}}\right) + A_{\Delta\Gamma} \sinh\left(\frac{y_s t}{\tau_{B_s}}\right)\right] \times e^{-t/\tau_{B_s}}$$

Mass-eigenstate rate asymmetry $A_{\Delta\Gamma}$

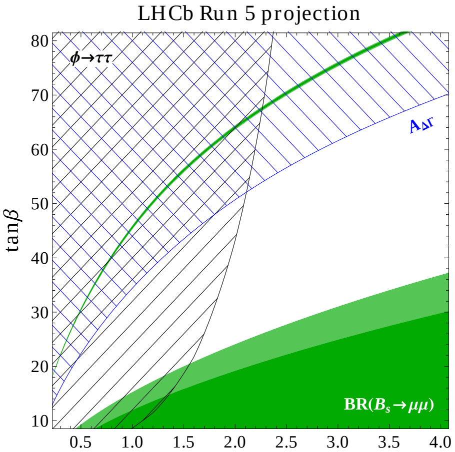
cf. Fleischer 0802.2882

Impact of future measurement of $A_{\Delta\Gamma}$



Constraint on real part of C_S imposing SMEFT relation Altmannshofer et al. 1702.05498

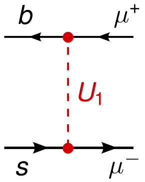
Future constraints on MSSM parameter space



- $A_{\Delta\Gamma}$ can exclude (or confirm??) second solution

Example 2: Leptoquarks

$$\mathcal{L}_{U_1} = \hat{\lambda}_{\rm L}^{ij} \left(\bar{Q}_{\rm L}^i \gamma_\mu L_{\rm L}^j\right) U_1^\mu + \hat{\lambda}_{\rm R}^{ij} \left(\bar{d}_{\rm R}^i \gamma_\mu e_{\rm R}^j\right) U_1^\mu + \text{h.c.}$$



$C_9^{\rm NP}$	$-\frac{1}{2}\mathcal{N}\lambda_L^{s\mu}\lambda_L^{b\mu\,*}$
C_9'	$-\frac{1}{2}\mathcal{N}\lambda_R^{s\mu}\lambda_R^{b\mu\,*}$
$C_{10}^{\rm NP}$	$\frac{1}{2}\mathcal{N}\lambda_L^{s\mu}\lambda_L^{b\mu\,*}$
C_{10}'	$-\frac{1}{2}\mathcal{N}\lambda_R^{s\mu}\lambda_R^{b\mu\,*}$
$C_S = -C_P$	$\mathcal{N}\lambda_L^{s\mu}\lambda_R^{b\mu\,*}\,m_b^{-1}$
$C_S' = C_P'$	$\mathcal{N}\lambda_R^{s\mu}\lambda_L^{b\mu\,*}\,m_b^{-1}$

Constraint on LQ parameter space

Scenario with $C'_S = C'_P = 0$