Higgs hunting with B decays

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B physics is sensitive to the effects of Higgs bosons in the Minimal Supersymmetric Standard Model, if the parameter $\tan \beta$ is large. I briefly summarise the role of $B \rightarrow \mu^+\mu^-$ and $B^+ \rightarrow \tau^+\nu_\tau$ in the hunt for new Higgs effects and present new results on the decay $B \rightarrow D\tau\nu_\tau$: Using the analyticity properties of form factors one can predict the ratio $R \equiv B(B \rightarrow D\tau\nu_\tau)/B(B \rightarrow D\ell\nu_\ell)$, $\ell = e, \mu$, with small hadronic uncertainties. In the Standard Model one finds $R = 0.31 \pm 0.02$, $B(B^- \rightarrow D^0\tau^-\bar{\nu}_\tau) = (0.71 \pm 0.09)\%$ and $B(B^0 \rightarrow D^+\tau^+\bar{\nu}_\tau) = (0.66 \pm 0.08)\%$, if the vector form factor of the Heavy Flavor Averaging Group is used. $B \rightarrow D\tau\nu_\tau$ is competitive with $B^+ \rightarrow \tau^+\nu_\tau$ in the search for effects of charged Higgs bosons. Especially sensitive to the latter is the differential distribution in the decay chain $B \rightarrow D\bar{\nu}_\tau\tau^- \rightarrow \pi^-\nu_\tau$.

1. Higgs effects in B physics

Weakly-coupled extensions of the Standard Model (SM) typically possess a richer Higgs sector than the latter. The easiest extension of the SM Higgs sector involves one additional Higgs doublet and is realised in the Minimal Supersymmetric Standard Model (MSSM). At tree-level the MSSM Higgs sector coincides with a Two-Higgs-doublet model (2HDM) of type II, in which down-type fermions receive their masses solely from one doublet, while up-type fermion masses exclusively stem from Yukawa interactions with the other Higgs doublet. An important parameter is the ratio $\tan \beta$ of the two vacuum expectation values. In the type-II 2HDM the bottom and top Yukawa couplings $y_b$ and $y_t$ satisfy the relation

$$\frac{y_b}{y_t} = \frac{m_b}{m_t} \tan \beta.$$  

Values around $\tan \beta = O(60)$ correspond to $y_b$-$y_t$ unification, which occurs in grand-unified theories (GUTs) with a minimal Yukawa sector. The idea of grand unification seems to call for low-energy supersymmetry, which stabilises the electroweak scale against radiative corrections from heavy GUT particles, improves the unification of the gauge couplings and reconciles the prediction of the proton lifetime with its experimental bounds. Probing the large-$\tan \beta$ region of the MSSM is therefore of great interest, since the question of Yukawa unification sheds light on the Yukawa sector of the underlying GUT theory. Yet large values of $\tan \beta$ are also interesting from purely phenomenological considerations: The tension between the measured anomalous magnetic moment of the muon, $a_\mu$ [1], and the Standard Model prediction [2] invites supersymmetry with $\tan \beta \gtrsim 10$, and larger values of $\tan \beta$ allow to saturate $a_\mu$ with heavier superpartners. Recent global fits of electroweak and B-physics observables to the constrained MSSM and the model with minimal gauge-mediated supersymmetry breaking gave best fits for values of $\tan \beta = 54$ and $\tan \beta = 55$, respectively [3].

B physics is excellently suited to study large-$\tan \beta$ scenarios, because down-type Yukawa couplings grow with $\tan \beta$ and $\tan \beta = O(50)$ corresponds to $y_b \sim 1$ [4]. Most dramatic effects can be expected in the leptonic decays $B_q \rightarrow \ell^+\ell^-$ (with $q = d$ or $s$ and $\ell = e, \mu$ or $\tau$), which are not only loop-suppressed in the Standard Model but also suffer from an additional helicity suppression. In particular the 95% CL limit

$$B(B_s \rightarrow \mu^+\mu^-) \leq 5.8 \cdot 10^{-8} \approx 18 \cdot B_{SM}(B_s \rightarrow \mu^+\mu^-).$$

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from the CDF experiment [5] already cuts into the large-tan $\beta$ region of the MSSM parameter space [4]. This is even true for the popular scenario of Minimal Flavour Violation (MFV) [6], in which the supersymmetric contribution to the $B_s \to \mu^+\mu^-$ amplitude involves the same elements of the Cabibbo-Kobayashi-Maskawa (CKM) matrix as the SM amplitude. For large tan $\beta$ one finds

$$B(B_s \to \mu^+\mu^-) \propto \epsilon_Y^2 \tan^6 \beta/M_A^6,$$  \hspace{1cm} (1)

where $M_A$ is the mass of the CP-odd Higgs boson and $\epsilon_Y$ is a loop function which depends on several MSSM parameters [4]. The pattern of Eq. (1) is in sharp contrast with the case of the naive type-II 2HDM, in which $B(B_s \to \mu^+\mu^-)$ is proportional to only four powers of tan $\beta$ [7]. In the LHCb experiment it will be possible to measure $B(B_s \to \mu^+\mu^-)$ for any value of tan $\beta$. One can then test the MFV hypothesis by checking whether $B(B_d \to \mu^+\mu^-)/B(B_s \to \mu^+\mu^-)$ agrees with $|V_{td}/V_{tb}|^2 f_{B_d}^2/f_{B_s}^2$, where $f_{B_s}$ is the decay constant of the $B_s$ meson.

Due to the overall loop factor of $\epsilon_Y$ the bound on tan $\beta/M_A^6$ derived from Eq. (1) depends on a plethora of other MSSM parameters. By contrast, effects of the charged Higgs boson $H^+$ enter $B$ decays at the tree-level. The information gained from charged-Higgs-mediated processes is therefore more directly related to the parameters of the MSSM Higgs sector, with smaller dependences on e.g. superpartner masses. $H^+$ effects are best studied in leptonic and semi-leptonic $B$ decays, in which hadronic uncertainties are under sufficient control. We specify our discussion to the case of a $\tau$ lepton in the final state, because the Yukawa couplings of the third fermion generation are largest. The B factories have observed the decay $B \to \tau\nu$ with $[8]$.

$$B(B \to \tau\nu) = (1.41 \pm 0.43) \times 10^{-4},$$ \hspace{1cm} (2)

which allows to place first useful constraints on tan $\beta/M_{H^+}$ [9]. In the following sections I will elaborate on another promising charged-Higgs hunting ground, the decay $B \to D\tau\nu$, and compare this mode with $B \to \tau\nu$. The presented work has been performed in collaboration with Stéphanie Trine and Susanne Westhoff [10].

2. Charged-Higgs effects at large tan $\beta$

The $q_b H^+$ coupling (with $q = u$ or $c$) is given by

$$L_{q_b H^+} = -\frac{g}{2\sqrt{2} M_W} \frac{\tan \beta}{1 + \epsilon_b \tan \beta} V_{qb} \phi \frac{(1 + \gamma_5) b H^+}{\phi},$$ \hspace{1cm} (3)

where the small Yukawa coupling $\epsilon_b$ is set to zero. The bottom quark coupling $y_b$ is defined in the same QCD renormalisation scheme as the current $\phi(1 + \gamma_5) b$. In the 2HDM of type-II the parameter $\epsilon_b$ vanishes. In the MSSM with MFV $\epsilon_b$ is a loop factor: the typically dominant squark-gluino contribution to $\epsilon_b$ is proportional to $\mu^*/M_{\tilde{g}}$, where $\mu$ is the Higgsino mass parameter and $M_{\tilde{g}}$ is the gluino mass [4]. A priori $\epsilon_b$ could be complex, but experimental constraints from electric dipole moments severely constrain the phase of $\mu^*/M_{\tilde{g}}$ and thereby of $\epsilon_b$. Since $|\epsilon_b|\tan \beta$ can be of order 1, the charged-Higgs phenomenology does involve genuine supersymmetric parameters, yet with much less impact than in $B_s \to \mu^+\mu^-$. In the MSSM with a generic flavour structure $\epsilon_b$ is different for $q = u$ and $q = c$ and may obtain a sizable phase. In the generic 2HDM $\epsilon_b$ is generated at tree-level and is typically complex. For an early extensive study of $B \to D\tau\nu$ and $B \to D\tau\nu$ in the MFV-MSSM see Ref. [11].

The effective hamiltonian describing $b \to q\tau\nu$ transitions mediated by $W^+$ or $H^+$ can be written as

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} V_{qb} \left\{ \phi \gamma^\mu (1 - \gamma_5) b \phi \gamma^\mu (1 - \gamma_5) \nu_\tau - \frac{m_b m_\tau}{m_{H^+}^2} \phi [g_s + g_P \gamma_5] b \phi (1 - \gamma_5) \nu_\tau \right\} + \text{h.c.}$$ \hspace{1cm} (4)

In the MSSM the effective couplings $g_s$ and $g_P$ read

$$g_s = g_P = \frac{m_b^2}{M_{H^+}^2 (1 + \epsilon_b \tan \beta)(1 + \epsilon_\tau \tan \beta),}$$ \hspace{1cm} (5)

Here $\epsilon_\tau$ is the analogue of $\epsilon_b$ for the $\tau_\tau H^+$ coupling. $B^+ \to \tau^+\nu$ probes the coupling $g_P$ [9].

$$B(B^+ \to \tau^+\nu) \propto |V_{ub}|^2 f_{B^+}^2 [1 - g_P]^2$$ \hspace{1cm} (6)
The $B$ meson decay constant $f_B = 216 \pm 38$ MeV \cite{13} is the dominant source of theoretical uncertainty in the extraction of $|1 - g_P|$ from $\mathcal{B}(B^+ \to \tau^+\nu)$. Eq. \eqref{eq:6} is confronted with the experimental result of Eq. \eqref{eq:2} in Fig. \ref{fig:1}.

The semi-tauonic decay $B \to D\tau\nu$ instead probes the coupling $g_S$, because $B$ and $D$ have the same parity. In the MFV-MSSM $B \to D\tau\nu$ and $B^+ \to \tau^+\nu$ probe the same parameters, because $\epsilon_b$ in Eqs. \eqref{eq:3} and \eqref{eq:4} is the same for $b \to u$ and $b \to c$ transitions. Thus within the MFV-MSSM we can combine the information from both decay modes to constrain the combination of $\tan\beta$, $M_{H^+}$, $\epsilon_b$ and $\epsilon_\tau$ in Eq. \eqref{eq:6}. If new physics is found, a comparison of $g_S$ extracted from $B \to D\tau\nu$ with $g_P$ obtained from $B^+ \to \tau^+\nu$ will probe physics beyond the MFV-MSSM. $B \to D\tau\nu$ compares to $B^+ \to \tau^+\nu$ as follows:

i) $\mathcal{B}(B \to D\tau\nu) \approx 50 \mathcal{B}(B^+ \to \tau^+\nu)$. 

ii) $|V_{cb}|$ entering $B \to D\tau\nu$ is much better known than $|V_{ub}|$.

iii) The uncertainty in lattice calculations of $f_B^2$ needed for $B \to \tau\nu$ is 30%. Instead $B \to D\tau\nu$ involves hadronic form factors. We will see below in Sect. \ref{sect:3} that the associated theoretical uncertainty is smaller, if experimental data on $B \to D\ell\nu$ with $\ell = e,\mu$ are exploited.

iv) The three–body decay $B \to D\tau\nu$, has decay distributions which discriminate between $W^+$ and $H^+$ exchange.

v) $B \to \tau\nu$ is mildly helicity–suppressed. In $B \to D\tau\nu$, the (transverse) $W^+$ contribution is helicity suppressed in the kinematic region with slow $D$ meson (P–wave suppression) \cite{14}:

\begin{equation}
\langle D(p_D)|\bar{c}\gamma^\mu b|\bar{B}(p_B)\rangle = V_1(w)\frac{m_B + m_D}{2\sqrt{m_Bm_D}}(p'_D + p'_D - \frac{m_B - m_D}{q^2}q^\mu) + S_1(w)(1 + w)\sqrt{m_Bm_D}\frac{m_B - m_D}{q^2}q^\mu.
\end{equation}

\begin{equation}
\langle D(p_D)|\bar{c}\bar{b}(\mu)|\bar{B}(p_B)\rangle = S_1(w)(1 + w)\sqrt{m_Bm_D}\frac{m_B - m_D}{m_b(\mu) - m_c(\mu)}, \quad (7)
\end{equation}

where $m_D$ and $m_B$ are the meson masses, $q = p_B - p_D$ is the momentum transfer and the running quark masses $m_b$ and $m_c$ are evaluated at the renormalisation scale $\mu$ at which the scalar current $\bar{c}\bar{b}$ is defined. The kinematic variable

\begin{equation}
w = \frac{m_B^2 + m_D^2 - q^2}{2m_Dm_B}, \quad (8)
\end{equation}

is defined in such a way that the kinematic endpoint $q^2 = (m_B - m_D)^2$, corresponds to $w = 1$. Heavy quark symmetry implies \cite{15}

\begin{equation}
S_1(1) = V_1(1) = 1 + \mathcal{O}(1/m_c, \alpha_s). \quad (9)
\end{equation}

For $m_\ell = 0$ the other kinematic endpoint is at $q^2 = 0$ corresponding to $w_{\text{max}} = (m_B^2 + m_D^2)/(2m_Dm_B) = 1.59$. The absence of a pole at $q^2 = 0$ in Eq. \eqref{eq:7} implies

\begin{equation}
S_1(w_{\text{max}}) = V_1(w_{\text{max}}). \quad (10)
\end{equation}

The next step towards a precision analysis of $B \to D\tau\nu$ exploits the analyticity properties of form factors \cite{16}. The location of poles and

\footnote{There is a typo in the definition of $S_1$ in Ref. \cite{16}.}
branch points in $V_1$ and $S_1$ can be inferred from the particle spectrum. The conformal mapping

$$z = \frac{\sqrt{(m_B + m_D)^2 - q^2} - \sqrt{(m_B + m_D)^2 - t_0}}{\sqrt{(m_B + m_D)^2 - q^2} + \sqrt{(m_B + m_D)^2 - t_0}}$$

and the elimination of subthreshold poles renders the form factors analytic in the new variable $z$ in the entire kinematic region (see Ref. [10] for details). That is, we can parametrize $V_1$ and $S_1$ in terms of power series in $z$. With a proper choice of the free parameter $t_0$ the kinematic range $1 \leq w \leq 1.59$ is mapped onto $0 \leq |z| \leq 0.032$.

$B \rightarrow D\ell\nu_{\ell}$ with $\ell = e, \mu$ only involves $V_1$. We can use experimental data to verify that the first two coefficients $a^S_0$ and $a^V_1$ of the power series are sufficient to describe the normalisation and shape of $V_1$. Moreover, the dependence on $a^V_1$ is moderate. The result is shown in Fig. 2. Finally we need the scalar form factor $S(w)$. The corresponding parameters $a^S_0$ and $a^S_1$ are fixed through Eqs. (9) and (10), using the $1/m_c$ and $\alpha_s$ corrections from Ref. [19]. An alternative approach uses lattice data to fix the form factors near $w = 1$ [20]. Note that our analysis faces much smaller hadronic uncertainties than the extraction of $|V_{cb}|$ from $B \rightarrow D\ell\nu_{\ell}$: We first fix $|V_{cb}|V_1(w)$ from experiment. Then Eq. (10) determines the normalisation of $|V_{cb}|S_1(w)$ in terms of measured quantities. Keeping only $a^S_0$ already reproduces $S(w)$ to 90%, and neither hadronic uncertainties nor the error in $|V_{cb}|$ have entered the prediction of $B \rightarrow D\tau\nu_\tau$ yet. Hadronic uncertainties only enter when we increase the accuracy further and fix the slope parameter $a^S_1$ by including Eq. (9).

These uncertainties are suppressed by $1/m_c$. In this step one also needs the experimental value of $|V_{cb}|$, whose error is around 2%. A remaining source of hadronic uncertainty is the next term $a^S_2$ of the series in $z$. In the data for $V_1$ in Fig. 2 we found the influence of $a^V_1$ negligible and expect the same behaviour for $a^S_1$ and $S_1$.

The first application of our analysis in Ref. [10] is a new prediction of branching fractions in the SM:

$$B(B^- \rightarrow D^0\tau^-\bar{\nu}_\tau) = (0.71 \pm 0.09)\%$$

$$B(B^0 \rightarrow D^{\pm}\tau^{\mp}\bar{\nu}_\tau) = (0.66 \pm 0.08)\%$$

$$R \equiv \frac{B(B \rightarrow D\tau\nu_{\tau})}{B(B \rightarrow D\ell\nu_{\ell})} = 0.31 \pm 0.02$$

using the HFAG form factor $V_1$ [18]. Note the small uncertainty in $R$ which is the relevant quantity for charged-Higgs hunting. This has to be compared with the $O(40%)$ error of $|V_{ub}|f_B^2$ entering the SM prediction of $B \rightarrow \tau\nu_\tau$. The uncertainties in Eq. (11) will decrease further with better data on $B \rightarrow D\ell\nu_{\ell}$. The dependence of $R$ on $g_S$ is shown in Fig. 3. A home-use formula for $R$ as a function of $g_S$ can be found in Eq. (7) of Ref. [10]. We notice that the dependence of $R$ on $g_S$ is weaker than the dependence of $B(B^+ \rightarrow \tau^+\nu)$ on $g_D$, but the present $1\sigma$ upper bounds on the effective coupling constants are similar. However, if nature has opted for a large charged-Higgs contribution suppressing $R$ to values below 0.2, it will not be easy to determine $g_S$ because the curve in Fig. 3 is quite flat. At present there is also an open experimental issue in $B \rightarrow D\tau\nu_{\tau}$: The Monte Carlo simulations for a lattice study of $a^V_1$ in the $B \rightarrow \pi$ form factor, which involves the much larger range $0 \leq |z| \leq 0.28$ than our $B \rightarrow D$ transition, see Paul Mackenzie’s talk at this conference.
However, in formation on the $\tau$ differential decay rate $\Gamma(D/\nu)$ derived from $R$ with a displaced vertex, so that the $\tau$ allows to constrain the phase of charged-Higgs effects in an excellent way, it also is not accessible to experiment. To my knowledge, the only theory paper addressing this problem is Ref. [23], which proposes to study the differential efficiencies for them for proofreading. Stimulating discussions with Christoph Schwanda on experimental issues are gratefully acknowledged.

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