Rare decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$ in a top quark two-Higgs-doublet model*

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Abstract In the framework of T2HDM, we calculated the new physics contributions involving neutral Higgs bosons to the branching ratios of $\bar{B}^0_{s,d} \to \ell^+\ell^-$ ($\ell = e, \mu$) decays. Comparing the theoretical predictions with the experimental upper-limits, we found that (a) The data of $Br(\bar{B}^0_d \to \ell^+\ell^-)$ give the upper bound on $\tan\beta$: $\tan\beta \leqslant 22$, while $Br(\bar{B}^0_s \to \ell^+\ell^-)$ give $\tan\beta \leqslant 12$ for fixed $\delta = 0°$, $m_{H^+} = 350$ GeV, $m_{H^0} = 160$ GeV, $m_{h^0} = 115$ GeV and $m_{A^0} = 120$ GeV; (b) A light neutral Higgs boson mass m_{h^0} (m_{A^0}) less than 50 GeV (120 GeV) is excluded by the data of branching ratios for $\bar{B}^0_{s,d} \to \ell^+\ell^-$ ($\ell = \mu$) decays with $\tan\beta = 10$; (c) The bounds on m_{h^0} and $\tan\beta$, or m_{A^0} and $\tan\beta$ are strongly correlated: a smaller (larger) $\tan\beta$ means a lighter (heavier) neutral Higgs boson.

Key words rare decay, branching ratio, two-Higgs-doublet model, new physics

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1 Introduction

The decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$ are flavor-changing neutral current (FCNC) processes. In the context of the Standard Model (SM), they proceed via Z^0 penguin and box-type diagrams, and their decays rates are expected to be greatly suppressed. However, in models with an extended Higgs sector these observables may receive sizable contributions from new particles, and thus provide a promising means to search for new physics.

The leptonic decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$ are particularly interesting among rare decays, since the prediction of the decay rate in the SM can be obtained with a relatively small error, due to the limited impact of long-distance hadronic corrections [1]. The SM expectations for the branching ratios with muons in the final state are [2]:

$$Br(\bar{B}^0_s \to \mu^+\mu^-) = (3.42 \pm 0.54) \times 10^{-9},$$
$$Br(\bar{B}^0_d \to \mu^+\mu^-) = (1.00 \pm 0.14) \times 10^{-10}, \quad (1)$$

which are smaller by one order of magnitude than the current experimental sensitivity. The corresponding branching ratios of the e^+e^- modes can be obtained from (1) by scaling with (m_e^2/m_μ^2).

Experimentally, the upper bounds on the $\bar{B}^0_{s,d} \to \ell^+\ell^-$ branching fractions at 90% confidence level have been given as [3, 4]

$$Br(\bar{B}^0_s \to e^+e^-) < 2.8 \times 10^{-7},$$
$$Br(\bar{B}^0_s \to \mu^+\mu^-) < 5.8 \times 10^{-8}, \quad (2)$$

$$Br(\bar{B}^0_d \to e^+e^-) < 11.3 \times 10^{-8},$$
$$Br(\bar{B}^0_d \to \mu^+\mu^-) < 5.2 \times 10^{-8}. \quad (3)$$

Up to now, the leptonic decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$ have been extensively studied in new physics models, for example, in the two-Higgs-doublet model (2HDM) [5, 6] and suspersymmetry (SUSY) models [7–10]. In SUSY, contributions from diagrams involving the new particles can provide several orders of magnitude to these branching ratios at large $\tan\beta$. The main purpose of this paper is to study these decays in the top quark two-Higgs-doublet model (T2HDM) [11]. We will calculate the new particles contributions to the branching ratios of these decays and compare them

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with the SM predictions and the experimental data.

This paper is arranged as follows. In Section 2, we give a brief review for the top quark two-Higgs-doublet model. In Section 3, we present the effective hamiltonian describing the quark transition $b \to s(d)\ell^+\ell^-$ in the presence of nonstandard Higgs bosons. We also show the hadronic matrix elements and the branching ratio required for the decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$. In Section 4, we give the numerical results of the branching ratios, compare the theoretical predictions with the experimental data and obtain the bounds on the T2HDM free parameters. The conclusions are included in the final section.

2 Outline of the top quark two-Higgs-doublet model

The model to be considered here is a top quark two-Higgs-doublet model, which is proposed in Ref. [11] and studied for example in Refs. [12–17], which is also a special case of the 2HDM of type III [18]. In this model, the top quark is much heavier than the other quarks and leptons because it couples to the second Higgs doublet with a much larger vacuum expectation value (VEV), while all the other fermions couple to the first Higgs doublet.

The Lagrangian density of Yukawa interactions of the T2HDM can be of the following form [11]:

$$\mathcal{L}_{\rm Y} = -\overline{L}_{\rm L}\phi_1 E l_{\rm R} - \overline{Q}_{\rm L}\phi_1 F d_{\rm R} - \overline{Q}_{\rm L}\widetilde{\phi}_1 G 1^{(1)} u_{\rm R} - \overline{Q}_{\rm L}\widetilde{\phi}_2 G 1^{(2)} u_{\rm R} + \text{h.c}\,, \tag{4}$$

where ϕ_i $(i=1,2)$ are the two Higgs doublets with $\widetilde{\phi}_i = \mathrm{i}\tau_2\phi_i^*$; and E, F, G are the generation space 3×3 matrices; $Q_{\rm L}$ and $L_{\rm L}$ are 3-vector of the left-handed quark and lepton doublets; $1^{(1)} \equiv {\rm diag}(1,1,0)$; $1^{(2)} \equiv {\rm diag}(0,0,1)$ are the two orthogonal projection operators onto the first two and the third families respectively.

In T2HDM, there are five physical Higgs bosons: the charged scalar $H^\pm$, the neutral CP-even scalars (H^0, h^0) and the CP-odd pseudoscalar A^0. After the rotation that diagonalizes the mass matrix of the quark fields, the Yukawa couplings for quarks are of the form [11]

$$\begin{aligned}\mathcal{L}_{\rm Y} = & -\sum_{\rm D=d,s,b} m_{\rm D}\bar{D}D - \sum_{\rm U=u,c,t} m_{\rm U}\bar{U}U - \sum_{\rm D=d,s,b}\frac{m_{\rm D}}{v}\bar{D}D[H^0-\tan\beta h^0] - \mathrm{i}\sum_{\rm D=d,s,b}\frac{m_{\rm D}}{v}\bar{D}\gamma_5 D[G^0-\tan\beta A^0] - \\ & \frac{m_{\rm u}}{v}\bar{u}u[H^0-\tan\beta h^0] - \frac{m_{\rm c}}{v}\bar{c}c[H^0-\tan\beta h^0] - \frac{m_{\rm t}}{v}\bar{t}t[H^0+\cot\beta h^0] + \mathrm{i}\frac{m_{\rm u}}{v}\bar{u}\gamma_5 u[G^0-\tan\beta A^0] + \\ & \mathrm{i}\frac{m_{\rm c}}{v}\bar{c}\gamma_5 c[G^0-\tan\beta A^0] + \mathrm{i}\frac{m_{\rm t}}{v}\bar{t}\gamma_5 t[G^0+\cot\beta A^0] + \frac{g}{\sqrt{2}M_W}\{-\overline{U}_{\rm L}Vm_{\rm D}D_{\rm R}[G^+-\tan\beta H^+] + \\ & \overline{U}_{\rm R}m_{\rm U}VD_{\rm L}[G^+-\tan\beta H^+] + \overline{U}_{\rm R}\Sigma^\dagger VD_{\rm L}[\tan\beta+\cot\beta]H^+ + \text{h.c.}\},\end{aligned} \tag{5}$$

with $G^\pm$ and G^0 Goldstone bosons. Here $M_{\rm U}$ and $M_{\rm D}$ are the diagonal up- and down-type mass matrices, V is the usual Cabibbo-Kobayashi-Maskawa (CKM) matrix and $\Sigma \equiv M_{\rm U}U_{\rm R}^\dagger 1^{(2)} U_{\rm R}$. $U_{\rm R}^\dagger$ is the unitary matrix which diagonalizes the right-handed up-type quarks as defined in Ref. [12].

3 Effective hamiltonian and branching ratio

The effective weak hamiltonian describing the flavor changing processes $b \to q\ell^+\ell^-$, with $q = s, d$ and $\ell = e, \mu, \tau$ in the presence of non-standard interactions can be written as [6, 19]:

$$\mathcal{H} = -\frac{4G_{\rm F}}{\sqrt{2}}V_{\rm tb}V_{\rm tq}^*[C_{10}(\mu)\mathcal{O}_{10}(\mu) + C_{\rm S}(\mu)\mathcal{O}_{\rm S}(\mu) + C_{\rm P}(\mu)\mathcal{O}_{\rm P}(\mu)], \tag{6}$$

where $V_{\rm tb}V_{\rm tq}^*$ is the CKM factor, and $G_{\rm F}$ is the Fermi coupling constant. C_i and $\mathcal{O}_i$ are the Wilson coefficients and local operators, respectively.

The explicit expressions of the operators are given by [6, 19–21]

$$\begin{aligned}\mathcal{O}_{10} &= \frac{e^2}{16\pi^2}(\bar{q}\gamma^\mu P_{\rm L}b)(\bar{\ell}\gamma_\mu\gamma_5\ell)\,, \\ \mathcal{O}_{\rm S} &= \frac{e^2}{16\pi^2}m_{\rm b}(\bar{q}P_{\rm R}b)(\bar{\ell}\ell)\,, \\ \mathcal{O}_{\rm P} &= \frac{e^2}{16\pi^2}m_{\rm b}(\bar{q}P_{\rm R}b)(\bar{\ell}\gamma_5\ell)\,,\end{aligned} \tag{7}$$

with $P_{\rm L,R} \equiv (1\mp\gamma_5)/2$. For the decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$, the matrix element is to be taken between vacuum and $|B^0_q\rangle$ state.

In addition to the operators in Eq. (7), there are additional operators such as $(\bar{q}\sigma_{\mu\nu}P_{\rm R}b)(\bar{\ell}\sigma^{\mu\nu}P_{\rm L,R}\ell)$ and $(\bar{q}\gamma_\mu P_{\rm L}b)(\bar{\ell}\gamma^\mu\ell)$. However, they give no contributions to the $\bar{B}_q \to \ell^+\ell^-$ decays and unmix with those appearing in Eq. (6). The matrix element that involves the antisymmetric tensor $\sigma^{\mu\nu}$ must vanish since $p^\mu\equiv p^\mu_{B_q}$ is the only four-momentum vector available. Similarly, the matrix element $\langle\ell^+\ell^-|\bar{\ell}\gamma^\mu\ell|0\rangle$ vanishes when contracted with $\langle 0|\bar{q}\gamma_\mu P_L b|\bar{B}_q(p)\rangle\propto p^\mu$ by the equations of motion. Therefore, the decays $\bar{B}_q\to\ell^+\ell^-$ are governed by the operators defined Eq. (7).

The evolution of the Wilson coefficients valuated at the matching scale $\mu=M_W$ down to the low-energy scale at $\mu=m_b$ can be performed by means of the renormalization group methods (see, e.g., [19, 21]).

The hadronic matrix elements responsible for the decays $\bar{B}_q\to\ell^+\ell^-$ can be characterized by the decay constant of the $\bar{B}_q$ meson, which is defined by the axial vector current matrix element [22].

$$\langle 0|\bar{q}\gamma_\mu\gamma_5 b|\bar{B}_q(p)\rangle=\mathrm{i}p_\mu f_{B_q}, \tag{8}$$

and, by using the equations of motion for quark fields, we have

$$\langle 0|\bar{q}\gamma_5 b|\bar{B}_q(p)\rangle=-\mathrm{i}f_{B_q}\frac{M^2_{B_q}}{m_b+m_q}. \tag{9}$$

Here f_{B_q} is the B_q meson decay constant, which can be obtained from lattice QCD computations and given such as in Ref. [23].

Using the effective hamiltonian in Eq. (6) together with Eqs. (8) and (9), we can write the decay amplitude for $\bar{B}_q\to\ell^+\ell^-$ in the following.

$$A=\frac{G_F\alpha}{2\sqrt{2}\pi}M_{B_q}f_{B_q}V_{tb}V^*_{tq}[F_S\bar{\ell}\ell+(F_P+2\hat{m}_\ell F_A)\bar{\ell}\gamma_5\ell], \tag{10}$$

here $\hat{m}_\ell\equiv m_\ell/M_{B_q}$, M_{B_q} is the mass of the B_q meson, and the F_i's are a function of Lorentz-invariant quantities.

The corresponding branching ratio then takes the form

$$Br(\bar{B}_q\to\ell^+\ell^-)=\frac{G_F^2\alpha^2M^3_{B_q}f^2_{B_q}\tau_{B_q}}{64\pi^3}|V_{tb}V^*_{tq}|^2\sqrt{1-4\hat{m}^2_\ell}\times \{(1-4\hat{m}^2_\ell)|F_S|^2+|F_P+2\hat{m}_\ell F_A|^2\}. \tag{11}$$

here τ_{B_q} refers to the B_q lifetime.

The scalar, pseudoscalar, and axial vector form factors are given by $(i=S,P)$

$$F_i=M_{B_q}\frac{m_bC_i}{m_b+m_q},\quad F_A=C_{10}, \tag{12}$$

where C_i, C_{10} are the Wilson coefficients. In the SM, the Wilson coefficient C_{10} that appears in Eq. (6) can be found easily in Refs. [1, 2]. At the one-loop level, the neutral Higgs bosons A^0, h^0, and H^0 also contribute to the rare decays $\bar{B}_q\to\ell^+\ell^-$ in the T2HDM. The T2HDM corrections induced by penguin diagrams involving neutral Higgs bosons to the Wilson coefficients have been given in Ref. [16].

4 Numerical results

The input parameters that we use in our numerical calculations are listed in Table 1.

In the T2HDM, the free parameters are charged Higgs mass m_{H^+}, $\tan\beta$, $|\xi|$, a new CP-violating phase δ and neutral Higgs masses (m_{H^0}, m_{h^0}, m_{A^0}). We fix $|\xi|=1$ throughout the paper and consider the other three as variable parameters. In previous papers [14–16], we have found strong constraints on the parameter space of the T2HDM from the well-measured radiative decay $B\to X_s\gamma$ [25], the measurements of $B\to X_s\ell^+\ell^-$ decays [26, 27] and the $B^0_{s(d)}$-$\bar{B}^0_{s(d)}$ mixing [25, 28]. Here we will consider these constraints in our choice for the free parameters of the T2HDM.

Table 1. The input parameters entering the numerical calculations. A, λ, $\bar{\rho}$ and $\bar{\eta}$ are the Wolfenstein parameters of the CKM mixing matrix [24].

m_e	0.511×10^{-3} GeV	α	1/137
m_μ	0.106 GeV	A	0.818
m_W	80.425 GeV	λ	0.2272
m_b	(4.8 ± 0.2) GeV	$\bar{\rho}$	0.221 ± 0.070
m_t	(174.2 ± 3.3) GeV	$\bar{\eta}$	0.340 ± 0.048
m_d	0.0076 GeV	f_{B_d}	0.216 ± 0.022 GeV
m_s	0.122 GeV	f_{B_s}	0.259 ± 0.032 GeV
m_{B_d}	5.279 GeV	τ_{B_d}	1.530 ps
m_{B_s}	5.367 GeV	τ_{B_s}	1.466 ps

Based on the studies about the constraints on the parameter space of the T2HDM and the formulae presented in previous sections, we are ready to perform our numerical analysis.

In Table 2, we give the SM predictions and the theoretical values by assuming $\delta=0^\circ$, $\tan\beta=30$, $m_{H^+}=350$ GeV, $m_{h^0}=115$ GeV, $m_{H^0}=160$ GeV, and $m_{A^0}=120$ GeV for the above leptonic decay branching ratios. The relevant experimental upper limits [29] are also listed for comparison. The dominant source of uncertainty comes from the decay constant dependence. After including the neutral Higgs bosons contributions, the theoretical predictions are about 3 orders of magnitude above the SM expectations. For decay modes $\bar{B}^0_d\to e^+e^-$ and $\bar{B}^0_s\to e^+e^-$, the theoretical predictions are still about 4 orders of magnitude away from the current experimental upper-limits. The theoretical values of branching ratio $Br(\bar{B}^0_d \to \mu^+\mu^-)$ and $Br(\bar{B}^0_s \to \mu^+\mu^-)$ are larger than the experimental upper limits because these decay modes are very sensitive to the parameter $\tan\beta$ and can put strong constraints on it.

Table 2. The SM predictions, the T2HDM theoretical values and experimental average [29] of branching ratios for leptonic decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$.

	SM prediction	T2HDM value	experimental average
$Br(\bar{B}^0_d \to e^+e^-)$	$(2.57^{+0.55}_{-0.49})\times10^{-15}$	$(1.44^{+0.31}_{-0.27})\times10^{-12}$	$<8.3\times10^{-8}$
$Br(\bar{B}^0_d \to \mu^+\mu^-)$	$(1.10^{+0.23}_{-0.21})\times10^{-10}$	$(6.14^{+1.31}_{-1.19})\times10^{-8}$	$<1.5\times10^{-8}$
$Br(\bar{B}^0_s \to e^+e^-)$	$(9.39^{+2.47}_{-2.17})\times10^{-14}$	$(5.35^{+1.40}_{-1.24})\times10^{-11}$	$<2.8\times10^{-7}$
$Br(\bar{B}^0_s \to \mu^+\mu^-)$	$(4.01^{+1.05}_{-0.93})\times10^{-9}$	$(2.28^{+0.60}_{-0.52})\times10^{-6}$	$<4.7\times10^{-8}$

When the new contributions are included, the size of the corresponding branching ratios depends mainly on the parameters $\tan\beta$, m_{h^0} and m_{A^0} and is not very sensitive to other parameters. As an illustration, assuming $\delta=0°$, $m_{H^0}=160$ GeV, $m_{H^+}=350$ GeV and using the input parameters given in Table 1, we find

$$\begin{cases} 1.13\times10^{-12} \leqslant Br(\bar{B}^0_s \to e^+e^-) \leqslant 3.18\times10^{-10}, \\ 4.83\times10^{-8} \leqslant Br(\bar{B}^0_s \to \mu^+\mu^-) \leqslant 1.36\times10^{-5}, \\ 3.05\times10^{-14} \leqslant Br(\bar{B}^0_d \to e^+e^-) \leqslant 8.54\times10^{-12}, \\ 1.30\times10^{-9} \leqslant Br(\bar{B}^0_d \to \mu^+\mu^-) \leqslant 3.64\times10^{-7}, \end{cases} \tag{13}$$

if $m_{h^0}=115$ GeV, $m_{A^0}=120$ GeV, and $10\leqslant\tan\beta\leqslant 50$.

$$\begin{cases} 6.29\times10^{-12} \leqslant Br(\bar{B}^0_s \to e^+e^-) \leqslant 1.93\times10^{-10}, \\ 2.69\times10^{-7} \leqslant Br(\bar{B}^0_s \to \mu^+\mu^-) \leqslant 8.22\times10^{-6}, \\ 1.69\times10^{-13} \leqslant Br(\bar{B}^0_d \to e^+e^-) \leqslant 5.18\times10^{-12}, \\ 7.24\times10^{-9} \leqslant Br(\bar{B}^0_d \to \mu^+\mu^-) \leqslant 2.21\times10^{-7}, \end{cases} \tag{14}$$

if $\tan\beta = 20$, $m_{A^0} = 120$ GeV, and $50 \leqslant m_{h^0} \leqslant 600$ GeV.

$$\begin{cases} 5.87\times10^{-12} \leqslant Br(\bar{B}^0_s \to e^+e^-) \leqslant 2.01\times10^{-10}, \\ 2.50\times10^{-7} \leqslant Br(\bar{B}^0_s \to \mu^+\mu^-) \leqslant 8.57\times10^{-6}, \\ 1.57\times10^{-13} \leqslant Br(\bar{B}^0_d \to e^+e^-) \leqslant 5.39\times10^{-12}, \\ 6.72\times10^{-9} \leqslant Br(\bar{B}^0_d \to \mu^+\mu^-) \leqslant 2.31\times10^{-7}. \end{cases} \tag{15}$$

if $\tan\beta=20$ GeV, $m_{h^0}=115$ GeV, and $50\leqslant m_{A^0}\leqslant 600$ GeV.

From the numerical results, we can see that although the present limit of $Br(\bar{B}^0_{s,d}\to\ell^+\ell^-)(\ell=\mu)$ is 1–2 orders of magnitude higher than the SM predictions, the upper-limit of the decay mode $\bar{B}^0_d\to\mu^+\mu^-$ will be very helpful to test the SM and to probe the effects of the new physics beyond the SM, or at least to put some limits on the parameters of new physics since these decays are sensitive to new pseudo-scalar interactions.

Figure 1 shows the $\tan\beta$ dependence of the branching ratios for $\bar{B}^0_{s,d} \to \ell^+\ell^-$ ($\ell = \mu$) decays. The dots and dash-dot curves correspond to the central values of the theoretical values in the SM and T2HDM for $\delta=0°$, $m_{H^+}=350$ GeV, $m_{H^0}=160$ GeV, $m_{h^0}=115$ GeV and $m_{A^0}=120$ GeV. The three dash-dot lines refer to the new physics contributions when uncertainties of decay constants f_{B_q} (q = s,d) are taken into account. The solid lines represent the experimental upper-limits as given in Table 2.

From Fig. 1 and the numerical results as given in Eqs. (13)–(15), one can see that

(i) The T2HDM contributions can provide large enhancements to the corresponding branching ratios at large $\tan\beta$.

(ii) The theoretical predictions depend strongly on the value of the parameter $\tan\beta$ and are not very sensitive to the parameters δ, m_{H^+}, and m_{H^0}.

(iii) Comparing the theoretical predictions with the experimental upper-limit of branching ratio $Br(\bar{B}^0_d \to \ell^+\ell^-)$, we can read off the upper bound on $\tan\beta$, $\tan\beta\leqslant22$, after the inclusion of the uncertainties of decay constant.

(iv) The experimental upper-limit of branching ratio $Br(\bar{B}^0_d \to \ell^+\ell^-)$ provides much stronger constraint on $\tan\beta$. For fixed $\delta=0°$, $m_{H^+}=350$ GeV, $m_{H^0}=160$ GeV, $m_{h^0}=115$ GeV and $m_{A^0}=120$ GeV, the upper bound on $\tan\beta$ is $\tan\beta\leqslant12$.

Analogous to Fig. 1, Figs. 2 and 3 show the m_{h^0} and m_{A^0} dependence of the branching ratios $Br(\bar{B}^0_{s,d}\to\ell^+\ell^-)$ ($\ell=\mu$), respectively. For $\tan\beta=10$ in T2HDM, the light neutral Higgs boson mass m_{h^0} (m_{A^0}) less than 50 GeV (120 GeV) is excluded by the current experimental data of these decays. The bounds on m_{h^0} and $\tan\beta$, or m_{A^0} and $\tan\beta$ are indeed strongly correlated: a smaller (larger) $\tan\beta$ means a lighter (heavier) neutral Higgs boson.

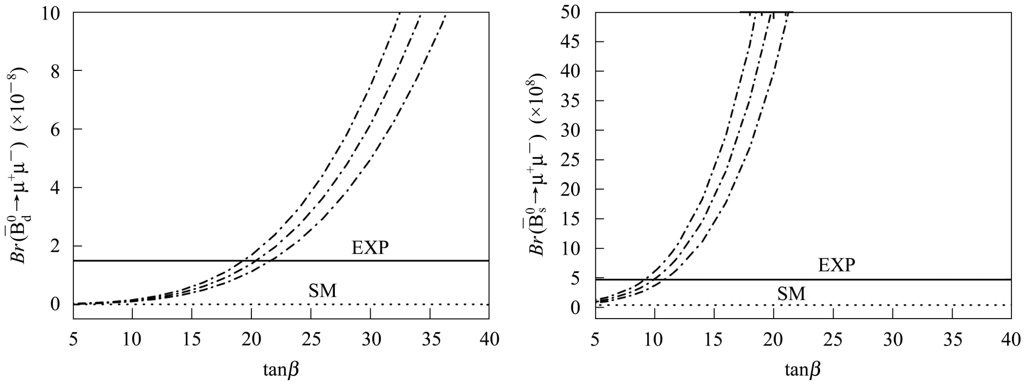

Fig. 1. The $\tan\beta$ dependence of the branching ratios of decays $\bar{B}^0_{s,d} \to \ell^+\ell^-$ $(\ell = \mu)$ in the SM and T2HDM, respectively. The dots and dash-dot curves show the central values of the SM and T2HDM predictions. The solid lines represent the experimental upper-limits as given in Table 2.

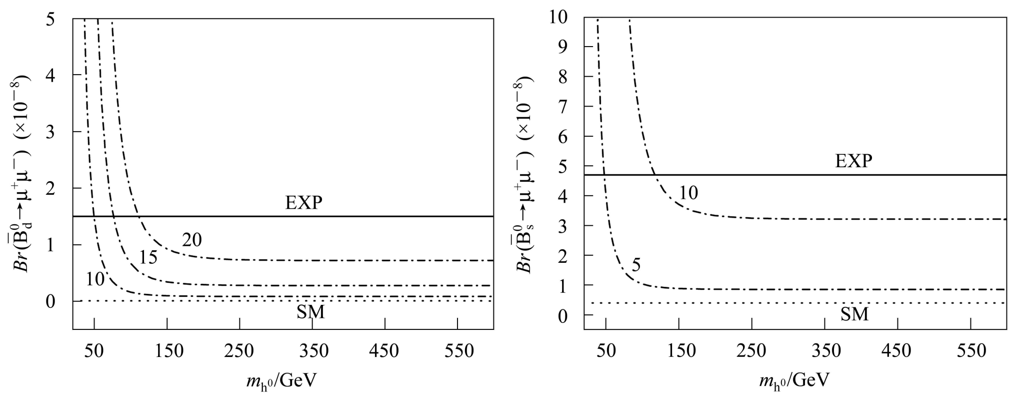

Fig. 2. The same as Fig. 1, but it shows the m_{h^0} dependence of the branching ratio of $\bar{B}^0_{s,d} \to \ell^+\ell^-$ $(\ell = \mu)$ in the SM and T2HDM. The dash-dot lines correspond to the predictions for different values of $\tan\beta$.

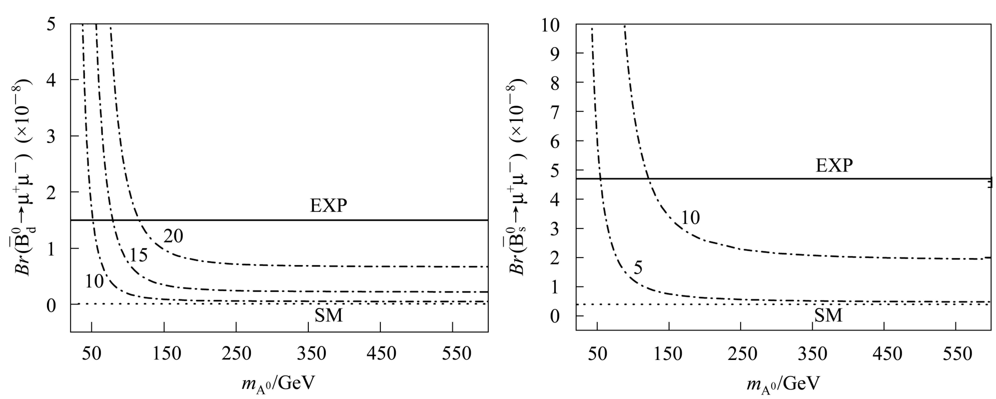

Fig. 3. The same as Fig. 2, but it shows the m_{A^0} dependence of the branching ratio of $\bar{B}^0_{s,d} \to \ell^+\ell^-$ $(\ell = \mu)$ in the SM and T2HDM.

5 Summary

In this paper, we calculated the new physics contributions to the branching ratios $Br(\bar{B}^0_{s,d} \to \ell^+\ell^-)$ $(\ell = e, \mu)$ induced by the one-loop box diagrams involving neutral Higgs boson propagators in the framework of T2HDM.

We firstly presented a brief review about the basic structure of the top-quark two-Higgs-doublet model, then gave the effective Hamiltonian, hadronic matrix elements and the branching ratios responsible for the decays $\bar{B}_q \to \ell^+\ell^-$ in the T2HDM. The new physics contributions are incorporated mainly through the Wilson coefficients C_S and C_P. Finally, we computed the branching ratios in the SM and the T2HDM, and made phenomenological analysis for these decays.

From the numerical results and the figures, we found that the T2HDM contributions to the branching ratios $Br(\bar{B}^0_{s,d} \to \ell^+\ell^-)$ ($\ell = e, \mu$) are very significant at large $\tan\beta$. The experimental upper-limits can put strong constraints on the free parameters of T2HDM. For fixed $\delta = 0°$, $m_{H^+} = 350$ GeV, $m_{H^0} = 160$ GeV, $m_{h^0} = 115$ GeV and $m_{A^0} = 120$ GeV, the data of $Br(\bar{B}^0_d \to \ell^+\ell^-)$ give the upper bound on $\tan\beta$: $\tan\beta \leqslant 22$, while $Br(\bar{B}^0_s \to \ell^+\ell^-)$ give $\tan\beta \leqslant 12$. A light neutral Higgs boson mass m_{h^0} (m_{A^0}) less than 50 GeV (120 GeV) is excluded by the data of branching ratios for $\bar{B}^0_{s,d} \to \ell^+\ell^-$ ($\ell = \mu$) decays with $\tan\beta = 10$. The bounds on m_{h^0} and $\tan\beta$, or m_{A^0} and $\tan\beta$ are indeed strongly correlated: a smaller (larger) $\tan\beta$ means a lighter (heavier) neutral Higgs boson.

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