

Leptonic Decays of the B Charged Meson and $B \rightarrow X_s \gamma$ in the Two Higgs Doublet Model Type III

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We consider the Two Higgs Doublet Model (2HDM) of type III which leads to Flavour Changing Neutral Currents (FCNC) at tree level. In the framework of this model we calculate the NLO contribution for $b \rightarrow s \gamma$ and the branchings for the meson decays $B^+ \rightarrow l^+ \nu$. We examine the limits on the new parameters λ_{bb} and M_{H^\pm} . We take into account the relationship between λ_{tt} and λ_{bb} coming from the validness of perturbation theory.

Keywords: B Meson decay; 2HDM type III; Charged Higgs mass

The Standard Model (SM) of particle physics based on the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ makes fit the symmetry breaking by including a fundamental weak doublet of scalar Higgs bosons ϕ with a scalar potential $V(\phi) = \lambda(\phi^\dagger \phi - \frac{1}{2}v^2)^2$. However, the SM does not explain the dynamics responsible for the generation of masses. Between the models beyond the SM, many of them include more than one scalar Higgs doublet; for instance, the case of the minimal supersymmetric standard model (MSSM)). We consider one of these extensions which has a richer scalar sector, called generically as the Two Higgs Doublet Model (2HDM). There are several kinds of such 2HDM models. In the model called type I, one Higgs doublet provides masses to the up and down quarks, simultaneously. In the model type II, one Higgs doublet gives masses to the up type quarks and the other one to the down type quarks. These two models include a discrete symmetry to avoid flavour changing neutral currents (FCNC) at tree level [1]. However, the addition of these discrete symmetries is not compulsory and in this case both doublets are contributing to generate the masses for up-type and down-type quarks. In the literature, such a model is known as the 2HDM type III [2]. It has been used to search for physics beyond the SM and specifically for FCNC at tree level [3–5]. In general, both doublets can acquire a vacuum expectation value (VEV), but one of them can be absorbed redefining the Higgs boson fields properly. Nevertheless, other studies on 2HDM-III using different basis have been done and there is a case where both doublets get VEVs that allows to study the models type I and II in a specific limit [5, 6].

In the 2HDM models, the two complex Higgs doublets correspond to eight scalar states. Spontaneous Symmetry breaking procedure leads to five Higgs fields: two neutral CP-even scalars h^0 and H^0 , a neutral CP-odd scalar A^0 , and two charged scalars H^\pm . While the neutral Higgs bosons may be difficult to distinguish from the one of the SM, the

charged Higgs bosons would have a distinctive signal for physics beyond the SM. Therefore the direct or indirect evidence of a charged Higgs boson would play an important role in the discovery of an extended Higgs sector. Direct searches have carried out by LEP collaborations and they reported a combined lower limit on M_{H^\pm} of 78.6 GeV [7] assuming $H^+ \rightarrow \tau^+ \nu_\tau (c\bar{s})$. At the Tevatron, the direct searches for charged Higgs boson are based on $p\bar{p} \rightarrow t\bar{t}$ where at least one top quark decay can be used in order to look for the channel $t \rightarrow H^+ b$. The CDF collaboration has reported a direct search for charged Higgs boson, setting an upper limit on $B(t \rightarrow H^+ b)$ around 0.36 at 95 % C.L. for masses in the range of 60-160 GeV [8]. On the other hand, indirect and direct searches have been done by D0 looking for a decrease in the $t\bar{t} \rightarrow W^+ W^- b\bar{b}$ signal expected from the SM and also the direct search for the decay mode $H^\pm \rightarrow \tau^\pm \nu$. We should note that all these bounds have been gotten in the framework of the 2HDM type II. And, in the framework of the 2HDM type II and MSSM a full one loop calculation of $\Gamma(t \rightarrow bH^+)$ including all sources for large Yukawa couplings were presented in references [10, 11]. Other indirect bounds on the charged Higgs boson mass come from processes where the charged Higgs boson is a virtual particle which is the case of the process $b \rightarrow s \gamma$. However, the indirect limits which have been obtained from the branching ratio $B \rightarrow X_s \gamma$ are strongly model dependent [12]. Finally, the search for the charged Higgs boson mass above the top quark mass will continue at LHC. The main production mechanisms would be the processes $gg \rightarrow tbH^+$ and $gb \rightarrow tH^+$ which have been studied using simulations of the LHC detectors [14].

The charged Higgs boson can also be revealed through contributions to low energy processes such as $B^0 - \bar{B}^0$, $D^0 - \bar{D}^0$ and $K^0 - \bar{K}^0$ where bounds on the charged Higgs sector have been found [2]. Moreover, there are other options through leptonic decays of the charged B mesons. They occur via

the annihilation process $B^\pm \rightarrow W^*(H^*) \rightarrow l^\pm \nu_l$. Then, it is possible to use the upper limits on these branching ratios obtained at CLEO [13], BELLE [15] and BABAR [16] in order to get bounds on the charged Higgs boson mass. Recent experimental results on $B(B_u \rightarrow \tau \nu)$ were reported by BELLE [17] and it is the first evidence of this kind of decays. The decays $B^\pm \rightarrow l^\pm \nu_l$ are sensitive at tree level to charged Higgs bosons and can be enhanced up to the current experimental limits by multi-Higgs models [18, 19]. On the other hand, the rare decay $B \rightarrow X_s \gamma$ is sensitive to charged higgs bosons at one loop level through electromagnetic and chromomagnetic penguin diagrams, and therefore the decay $B \rightarrow X_s \gamma$ can put strong constraints on the parameters of any charged Higgs sector because its high precision measurement done by CLEO [20].

In the present work, we study the processes $B \rightarrow X_s \gamma$ and $B^+ \rightarrow l^+ \nu$ in the framework of the 2HDM type III. And we concentrate on the charged Higgs boson sector of this model, with the relevant parameters being its mass M_{H^\pm} and the Yukawa coupling.

The 2HDM type III is an extension of the SM which adds a new Higgs doublet and three new Yukawa couplings in the quark and leptonic sectors. The mass terms for the up-type or down-type sector depend on two Yukawa coupling matrices. The rotation of quarks and leptons to mass eigenstates allow us to diagonalize one of the matrices but not both simultaneously, then one of the Yukawa coupling matrix remains non-diagonal, generating FCNC at tree level. The most general structure for the Higgs-fermion Yukawa couplings in the so called 2HDM type-III is as follow [2]:

$$\begin{aligned}
-\mathcal{L}_Y &= \eta_{ij}^{U,0} \bar{Q}_{iL}^0 \tilde{\Phi}_1 U_{jR}^0 + \eta_{ij}^{D,0} \bar{Q}_{iL}^0 \Phi_1 D_{jR}^0 + \eta_{ij}^{E,0} \bar{l}_{iL}^0 \Phi_1 E_{jR}^0 \\
&+ \xi_{ij}^{U,0} \bar{Q}_{iL}^0 \tilde{\Phi}_2 U_{jR}^0 + \xi_{ij}^{D,0} \bar{Q}_{iL}^0 \Phi_2 D_{jR}^0 + \xi_{ij}^{E,0} \bar{l}_{iL}^0 \Phi_2 E_{jR}^0 \\
&+ h.c.
\end{aligned} \tag{1}$$

$$\begin{aligned}
\Gamma(B^\pm \rightarrow l^\pm \nu_l)^{III} &= \frac{G_F^2 m_B m_l^2 f_B^2}{8\pi} |V_{ub}|^2 \left(1 - \frac{m_l^2}{m_B^2}\right) \left[1 - \frac{|d||b|M_B}{2\sqrt{2}G_F m_l m_H^2}\right]^2 \\
&= \Gamma(B^\pm \rightarrow l^\pm \nu_l)^{SM} \times R
\end{aligned} \tag{3}$$

where in the framework of the 2HDM-III, we have the factors

$$d = \xi_{ll} \tag{4}$$

$$b = \frac{g}{2m_W} m_b \lambda_{bb}. \tag{5}$$

And R can be written as

$$R = \left[1 - \frac{|\lambda_{bb}||\lambda_{ll}|M_B^2}{m_H^2}\right] \tag{6}$$

where we have used the Cheng and Sher parametrization [4]. When $R = 1$, i.e., $|\lambda_{bb}||\lambda_{ll}|M_B^2 = 2m_H^2$ the new physics predicts the same value for the width $\Gamma(B^\pm \rightarrow l^\pm \nu_l)^{III}$ that the SM, and

where $\Phi_{1,2}$ are the Higgs doublets, $\tilde{\Phi}_i \equiv i\sigma_2 \Phi_i^*$, Q_L^0 is the weak isospin quark doublet, and U_R^0 , D_R^0 are weak isospin quark singlets, whereas η_{ij}^0 and ξ_{ij}^0 are non-diagonal 3×3 and i, j are family indices. The superscript 0 indicates that the fields are not mass eigenstates yet. In the so-called model type I, the discrete symmetry forbids the terms proportional to η_{ij}^0 , meanwhile in the model type II the same symmetry forbids terms proportional to $\xi_{ij}^{D,0}$, $\eta_{ij}^{U,0}$, $\xi_{ij}^{E,0}$. On the other hand, we are considering the 2HDM-III in a basis where only one Higgs doublet acquire VEV and then it does not have the usual parameter $\tan\beta = v_2/v_1$ of the 2HDM type II. In this way we have the usual 2HDM type III [4], where the Lagrangian of the charged sector is given by

$$-L_{H^\pm ud}^{III} = H^+ \bar{U} [K \xi^D P_R - \xi^U K P_L] D + h.c. \tag{2}$$

where K is the Cabbibo-Kobayashi-Maskawa (CKM) matrix and $\xi^{U,D}$ the flavour changing matrices. In the framework of the 2HDM type III is useful to use the parameterization proposed by Cheng and Sher for the Yukawa couplings given by $\xi_{ii} = \lambda_{ii} g m_i / (2m_W)$ [4].

The leptonic decays of the B^\pm mesons are possible via annihilation processes into W^\pm gauge bosons or H^\pm scalar bosons. The first one is the usual SM contribution and the second one in our case is own to the 2HDM type III. We should mention that the branching fractions for $e^- \bar{\nu}_e$ and $\mu^- \bar{\nu}_\mu$ in the framework of the SM are helicity suppressed by factors of $\sim 10^{-8}$ and $\sim 10^{-3}$, respectively. But physics beyond the SM can enhance these branching fractions through the introduction of a charged Higgs boson, as we will notice. The decay width can be written as

there is an allow region in the plane $m_H - \lambda_{bb}$ with the same SM values. In this form the decay width depends only on the free parameters λ_{ll} , m_{H^\pm} and λ_{bb} .

About the experimental data for the B meson decays $B^- \rightarrow l^- \bar{\nu}_l$, they are experimentally challenging because there are at least two undetectable neutrinos in the final state. These kind of decays has been searched at BELLE, BABAR and CLEO-b. Bounds on the braching fraction $B(B \rightarrow \mu \nu)$ have been reported. The stringent bounds come from BABAR measurements and they are $B(B \rightarrow \mu \nu_\mu) \leq 6.8 \times 10^{-6}$ [16] and the SM prediction is $B(B \rightarrow \mu \nu_\mu) = 3.9 \times 10^{-7}$. For the decay $B_u \rightarrow \tau \nu$, the first evidence has been reported by BELLE [17], they report an experimental result of $B(B \rightarrow$

$\tau\nu_\tau) = 1.06_{-0.28}^{+0.34}(\text{stat})_{-0.16}^{+0.18}(\text{sys}) \times 10^{-4}$. In addition, the values predicted by the SM is $B(B \rightarrow \tau\nu_\tau) = 1.59 \times 10^{-4}$, which is consistent with the experiment value. This new measurement could guide to a deeper understanding of flavour and electroweak dynamics, and it could provide evidence of a non-standard Higgs sector. As we already mentioned in the B meson decays is possible to reduce the number of parameters to λ_{bb} and the charged Higgs boson mass m_H where we have used the flavour changing couplings for the leptonic sector from the literature[5]. These couplings are bounded by $-0.12 \leq \xi_{22} \leq 0.12$ and $-1.8 \times 10^{-2} \leq \xi_{33} \leq 2.2 \times 10^{-2}$. These bounds can also be written as $|\lambda_{\mu\mu}| \leq 139.82$ and $|\lambda_{\tau\tau}| \leq 1.52$. For our numerical analysis we always take the upper bounds $\lambda_{\mu\mu} = 139.82$ and $\lambda_{\tau\tau} = 1.52$.

The Fig. 1 shows the allowed values (fullfilled) for m_H vs λ_{bb} according to the experimental result from BELLE [17] for the $B \rightarrow \mu\nu_\mu$ (left) and $B \rightarrow \tau\nu_\tau$ (right) decays where the allow regions correspond to the fullfilled regions. For the leptonic parameter we take the values $\lambda_{\mu\mu} = 139.82$ and $\lambda_{\tau\tau} = 1.52$, respectively. For $B \rightarrow \mu\nu_\mu$ (left), the lower solid line corresponds to the experimental limit and the upper solid line corresponds to the solutions given by $R = 1$ which predicts for the new physics the same SM width. For $B \rightarrow \tau\nu_\tau$ (right), the upper curve corresponds to experimental limit and the lower solid line corresponds to $R = 1$ solution which is excluded.

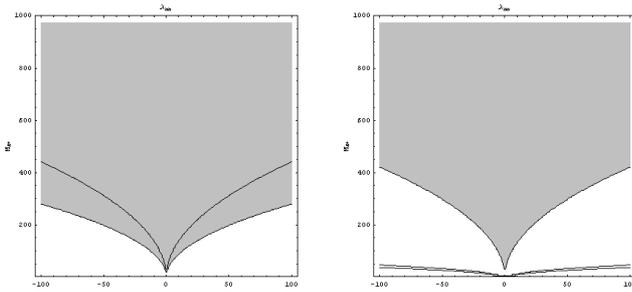


FIG. 1: The plane $m_H - \lambda_{bb}$ for the $B \rightarrow \mu\nu_\mu$ (left) and $B \rightarrow \tau\nu_\tau$ (right) decays in the 2HDM-III with $\lambda_{\mu\mu} = 139.82$ and $\lambda_{\tau\tau} = 1.52$. The allow values are the fullfilled region.

On the other hand, for the radiative decay $B \rightarrow X_s\gamma$ we follow references [23, 24]. The $B \rightarrow X_s\gamma$ process as any FCNC process does not arise at the tree level in the SM. In the framework of the SM it is generated by the one-lopp W-exchange diagrams but these contributions are small enough to be comparable to nonstandard contributions, in our case the exchange of a charged scalar Higgs boson. The branching ratio of the inclusive radiative decay $B \rightarrow X_s\gamma$ is

$$B(B \rightarrow X_s\gamma)_{LO} = B_{SL} \left| \frac{V_{ts}^* V_{tb}}{V_{cb}} \right|^2 \frac{6\alpha_{em}}{\pi} |C_7^{0,eff}(\mu_b)|^2 \quad (7)$$

at the leading order level, where $C_7^{0,eff}(\mu_b)$ is the effective coefficient at the scale μ_b ,

$$\begin{aligned}
 C_7^{0,eff}(\mu_b) &= \eta_{23}^{16} C_7^{0,eff}(\mu_W) + \frac{8}{3} \left(\eta_{23}^{14} - \eta_{23}^{16} \right) C_8^{0,eff}(\mu_W) \\
 &+ \sum_{i=1}^8 h_i \eta^i C_2^{0,eff}(\mu_W), \quad (8)
 \end{aligned}$$

$f(z) = 1 - 8z^2 + 8z^6 - z^8 - 24z^4 \log z$ is the phase space factor in the semileptonic b -decay parameterized in terms of $z = m_c^{pole}/m_b^{pole}$ and α_{em} is the fine-structure constant. The coefficients $C_{7,8}^{0,eff}(\mu_b)$ have an important property and it is that they are quite similar in many interesting extensions of the SM, such as 2HDM or the MSSM [21–23] and therefore it is possible to parametrize the new contributions using new functions $C_{i,j}^{0,1}(\mu_W)$ with $i = 7, 8$ and $j = YY, XY$. These functions depend on the unknown parameter m_H^\pm and also on the size and sign fo the couplings X and Y that in the case of the model III under study they are $X = -\lambda_{bb}$, and $Y = \lambda_{tt}$. To get these couplings we assume that the flavour changing parameters for the light quarks are negligible and $\lambda_{bb} > 1$, $\lambda_{tt} < 1$ which is the case discussed by Atwood, Reina and Soni as their third case [3]. Then the LO Wilson coefficients at the matching energy scale m_W are [23, 24],

$$\begin{aligned}
 C_2^{0,eff}(\mu_W) &= 1, \\
 C_i^{0,eff}(\mu_W) &= 0, \quad (i = 1, 3, 4, 5, 6), \\
 C_7^{0,eff}(\mu_W) &= C_{7,SM}^0(m_W) + |Y|^2 C_{7,YY}^0(m_W) + (XY^*) C_{7,XY}^0(m_W), \\
 C_8^{0,eff}(\mu_W) &= C_{8,SM}^0(m_W) + |Y|^2 C_{8,YY}^0(m_W) + (XY^*) C_{8,XY}^0(m_W), \quad (9)
 \end{aligned}$$

with

$$\begin{aligned}
 C_{7,SM}^0 &= \frac{3x_t^3 - 2x_t^2}{4(x_t - 1)^4} \ln x_t + \frac{-8x_t^3 - 5x_t^2 + 7x_t}{24(x_t - 1)^3}, \\
 C_{8,SM}^0 &= \frac{-3x_t^2}{4(x_t - 1)^4} \ln x_t + \frac{-x_t^3 + 5x_t^2 + 2x_t}{8(x_t - 1)^3}, \\
 C_{7,YY}^0 &= \frac{3y_t^3 - 2y_t^2}{12(y_t - 1)^4} \ln y_t + \frac{-8y_t^3 - 5y_t^2 + 7y_t}{72(y_t - 1)^3}, \\
 C_{7,XY}^0 &= \frac{y_t}{12} \left[\frac{-5y_t^2 + 8y_t - 3 + (6y_t - 4) \ln y_t}{(y_t - 1)^3} \right], \quad (10) \\
 C_{8,YY}^0 &= \frac{-3y_t^2}{12(y_t - 1)^4} \ln y_t + \frac{-y_t^3 + 5y_t^2 + 2y_t}{24(y_t - 1)^3}, \\
 C_{8,XY}^0 &= \frac{y_t}{4} \left[\frac{-y_t^2 + 4y_t - 3 - 2 \ln y_t}{(y_t - 1)^3} \right], \quad (11)
 \end{aligned}$$

where $x_t = m_t^2/m_W^2$, $y_t = m_t^2/m_h^2$, and these leading order functions have no explicit μ_W dependence.

Now, at the next leading order level that is necessary in order to use the experimental data, the branching ratio is

$$B(B \rightarrow X_s\gamma)_{NLO} = B_{SL} \left| \frac{V_{ts}^* V_{tb}}{V_{cb}} \right|^2 \frac{6\alpha_{em}}{\pi} |f(z)\kappa(z)|^2 [\bar{D}]^2 + A + \Delta \quad (12)$$

where B_{SL} is the measured semileptonic branching ratio of B mesons, and $\kappa(z)$ is the QCD correction for the semileptonic

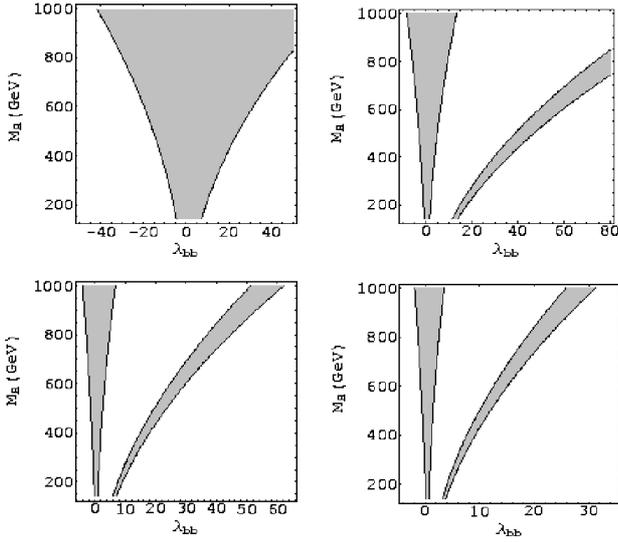


FIG. 2: The allowed values (fullfilled) in the plane λ_{bb} - M_H taking into account the experimental values for $B(B_s \rightarrow X_s \gamma)$ and $B(B \rightarrow \tau \nu)$. Different values of $\lambda_{tt} = 0.1, 0.5, 1, 2$ are considered for the figures from left-upper to right-lower, respectively.

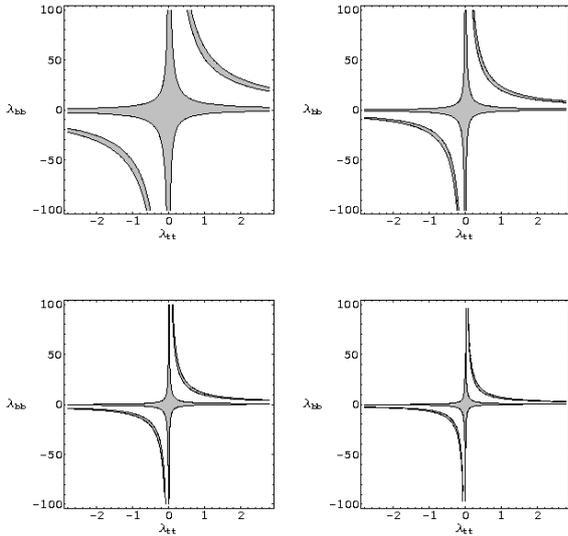


FIG. 3: The plane λ_{bb} - λ_{tt} taking into account the experimental value for $B(B_s \rightarrow X_s \gamma)$ and $B(B \rightarrow \tau \nu)$. The allowed regions are full-filled. Different values of the charged Higgs boson mass $M_H = 120, 250, 500, 1000$ GeV are considered for the figures from left-upper to right-lower, respectively.

B decay. The term \bar{D} corresponds to the subprocesses $b \rightarrow s\gamma$ which involves the NLO Wilson coefficient $C_7^{eff}(\mu_b)$, the virtual correction functions r_i and $\gamma_{i7}^{0,eff}$ the elements of the anomalous dimension matrix which govern the evolution of the Wilson coefficients from the matching scale μ_W to lower scale μ_b . The term A in equation (12) is the correction coming

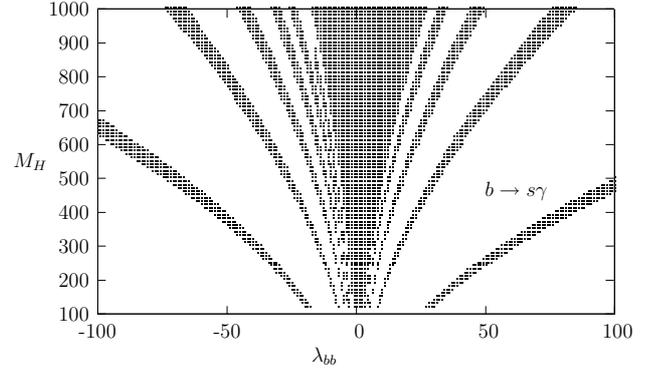


FIG. 4: The allowed values in the plane λ_{bb} - M_H taking into account the experimental value for $B(B_s \rightarrow X_s \gamma)$ and the condition coming from validness of perturbation theory.

from the bremsstrahlung process $b \rightarrow s\gamma$ and in the Δ have been included the nonperturbative corrections obtained with the method of the heavy-quark effective theory relating the actual hadronic process to the quark decay rate. The whole set of functions already mentioned have been given in references [23, 24]. With the set of above equations we can estimate the ratio $B(B \rightarrow X_s \gamma)$ and use the experimental world average $B(B \rightarrow X_s \gamma)^{exp} = (3.52 \pm 0.30) \times 10^{-4}$ [27].

From the perturbation theory considerations in the eq. (2) and using the Cheng Sher parametrization, we have already gotten the inequality [25]

$$\frac{m_b^2}{m_t^2} |\lambda_{bb}|^2 + |\lambda_{tt}|^2 < 8, \quad (13)$$

which will be taken into account in the numerical evaluations for the $B(B \rightarrow X_s \gamma)$ in order to reduce the number of variables. It defines an ellipse with $|\lambda_{bb}| \leq 100$ and $|\lambda_{tt}| \leq \sqrt{8}$ bounds. In this case we consider the inequality from perturbation theory validness in order to reduce the space of parameters. This link between the parameter λ_{tt} and λ_{bb} allows to get the plane λ_{bb} versus m_H using the experimental measurement for the branching ratio $B(B_s \rightarrow X_s \gamma)$.

In Fig. 2, we present the allowed regions (fullfilled) for the $B(B \rightarrow X_s \gamma)$ in the plane M_H versus λ_{bb} for different values of λ_{tt} . And in Fig. 3, we present the allowed regions (fullfilled) for the $B(B \rightarrow X_s \gamma)$ in the plane λ_{bb} - λ_{tt} for different values of the charged Higgs boson mass.

Finally, in Fig. 4 we show the case of the induced decay $B \rightarrow X_s \gamma$ decay. The fullfilled regions are the allowed regions, it means these are the regions satisfying the experimental value and the perturbation theory constraint eq. (12). We notice that these regions in Fig. 4 correspond to a different choice of λ_{tt} as it was presented in Fig. 2.

In conclusion, We have studied in the framework of the 2HDM type III, the allowed region for the parameters λ_{bb} and m_H using the processes $B \rightarrow \tau \nu \tau$, $B \rightarrow \mu \nu \mu$ and $B \rightarrow X_s \gamma$. We have used the condition on the parameter space coming from the fact that the Yukawa couplings should be perturba-

tive, eq. (13), in to order to reduce the number of free parameters. Finally, we have compared the plots looking for the stringest regions in the plane $\lambda_{bb}-m_H$ and we have noticed that the $B \rightarrow X_s \gamma$ decay is the most restrictive process constraining the parameters of the charged Higgs sector in the 2HDM-III. But however there are small regions for small values of λ_{bb} and light m_H that leptonic decays can exclude. We also have found that in case of the leptonic decays $B \rightarrow lv$, there are values of the parameters λ_{bb} and m_H given a 2HDM prediction

which cannot be distinguishable from the SM prediction. It is because the factor $(1 - dbM_B/(2\sqrt{2}G_F m_l m_H^2))^2$ in eq. (3) could get the value equal one for some values of λ_{bb} and m_H and then reach out the SM prediction. For the process $B \rightarrow \mu\nu_\mu$, these values are in the allowed experimental region of the plane $\lambda_{bb} - m_H$ as it was showed in Fig. 1. But for the process $B \rightarrow \tau\nu_\tau$ these ones are excluded.

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