Bounds on charged higgs boson in the 2HDM type III from Tevatron

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Abstract

We consider the Two Higgs Doublet Model (2HDM) which leads to Flavour Changing Neutral Currents (FCNC) at tree level. In the framework of this model we can use an appropiate form of the Yukawa Lagrangian that makes the type II model limit of the general type III couplings apparent. This way is useful in order to compare with the experimental data which is model dependent. The analytical expressions of the partial width $\Gamma(t \to H^+ b)$ are derived and we compare with the data available at this energy range. We examine the limits on the new parameters λ_{ij} fron the validness of perturbation theory.

The Standard Model (SM) of particle physics based on the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ accommodates 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. Furthermore, the scalar sector suffers from two serious problems, known as: the gauge hierarchy problem and the triviality problem [1]. The scalars involved in electroweak symmetry breaking should therefore be a party to new physics at some finite energy scale. Thus the SM would be merely a low-energy effective field theory, and the dynamics responsible for generating mass might lie in physics beyond the SM. There is the option of a model like the SM but including a richer scalar sector, which includes one more Higgs doublet, it is called generically 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 quarks and the other one to the down quarks. These two models have a discrete symmetry to avoid FCNC at tree level [2]. However, the discrete symmetry is not necessary in whose case both doublets generate the masses of the quarks of up-type and down-type, simultaneously. In the literature, the latter model is known as the model type III [3]. It has been used to look for physics beyond the SM and specifically for FCNC at tree level [4–6]. In general, both doublets could acquire a vacuum expectation value (VEV), but one of them can be absorbed redefining the Higgs fields properly. Nevertheless, we have showed that from the case in which both doublets get the VEV is possible to study the models type I and II in an specific limit [6]. Therefore we consider the model type III in two basis. In the first base, the two Higgs doublets acquire VEV (case (a) in ref.[6]). In the second one, only one Higgs doublet acquire VEV (case (b) in ref[6])[4]. In the latter case the free parameter $\tan\beta \equiv v_2/v_1$ is removed from the theory making its phenomenological analysis simpler. But in the former one is possible to get bounds for the model type III using the experimental bounds which have been gotten in the framework of the model type II.

In these kind of models (2HDM) additional degrees of freedom appear, providing a total of five observable Higgs fields: two neutral CP-even scalars h^0 and H^0, a neutral CP-odd scalar A^0, and two charged scalars $H^\pm$. Direct searches have carried out by LEP experiments, and report a combined lower limit on $M_{H^\pm}$ of 78.6 GeV [7]. The CDF collaboration has also reported a direct search for charged Higgs boson, setting an upper limit on $B(t \to H^+ b)$ around 0.6 at 95 % C.L. for masses in the range 60-160 GeV [8]. On the other hand, indirect and direct searches have been carried out by D0 looking for a decrease in the $t\bar{t} \to W^+W^-b\bar{b}$ signal expected from the SM and the direct search for the decay mode $H^\pm \to \tau^\pm\nu$. They exclude most regions of the plane $M_{H^\pm} - \tan\beta$ where the $B(t \to bH^+) > 0.36$ [9]. We should note that all the 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 $\Gamma(t \to bH^+)$ including all sources for large Yukawa couplings were presented in references [10, 11]. In what follows we concentrate on the charged sector, with the relevant parameters being its mass $M_{H^\pm}$ and the ratio of the VEV's of the doublets, $\tan\beta$ and the coupling intensities λ_{tt} and λ_{bb}.

In the present work, we study the process $t \to bH^+$ in the 2HDM type III. If $m_{H^\pm} < m_t - m_b$ then the charged Higgs boson $H^\pm$ can be produced in the decay of the top quark via $t \to bH^+$. This decay can be competitive with the dominant SM decay mode, $t \to bW^+$. The Higgs boson production in top decays has been studied in the framework of the 2HDM type II and under considerations that also apply to the MSSM [11]. We are going to work in the Higgs mass range 60-160 GeV, assuming that $B(t \to bW^+) + B(t \to bH^+) = 1$ and the masses of the neutral scalars are assumed to be large enough to be suppressed in $H^\pm$ decays. In this way the only available decays of $H^\pm$ are fermionic.

The 2HDM type III is an extension of the SM plus 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 depends on two matrices or two Yukawa couplings. The rotation of the quarks and lepton gauge eigenstates allow us to diagonalize one of the matrices but not both simultaneously, so one of the Yukawa couplings remains non-diagonal, generating the FCNC at tree level.

The Higgs couplings to fermions are model dependent. The most general structure for the Higgs-fermion Yukawa couplings, 2HDM type-III [3], is as follow:

$$\begin{aligned} -\pounds_Y &= \eta_{ij}^{U,0}\overline{Q}_{iL}^0\widetilde{\Phi}_1 U_{jR}^0 + \eta_{ij}^{D,0}\overline{Q}_{iL}^0\Phi_1 D_{jR}^0 + \eta_{ij}^{E,0}\overline{l}_{iL}^0\Phi_1 E_{jR}^0 \\ &+ \xi_{ij}^{U,0}\overline{Q}_{iL}^0\widetilde{\Phi}_2 U_{jR}^0 + \xi_{ij}^{D,0}\overline{Q}_{iL}^0\Phi_2 D_{jR}^0 + \xi_{ij}^{E,0}\overline{l}_{iL}^0\Phi_2 E_{jR}^0 \\ &+ h.c. \end{aligned} \quad (1)$$

where $\Phi_{1,2}$ are the Higgs doublets, $\widetilde{\Phi}_i \equiv i\sigma_2\Phi_i^*$, 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 non-dimensional matrices 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 $\xi_{ij}^{D,0}$, $\eta_{ij}^{U,0}$, $\xi_{ij}^{E,0}$. We next shift the scalar fields according to their VEV's, as

$$\langle\Phi_1\rangle_0 = \begin{pmatrix} 0 \\ v_1/\sqrt{2} \end{pmatrix} \;,\; \langle\Phi_2\rangle_0 = \begin{pmatrix} 0 \\ v_2/\sqrt{2} \end{pmatrix} \quad (2)$$

and we take the complex phase of v_2 equal to zero since we are not interested in CP violation. Then re-express the scalars in terms of the physical Higgs states and would-be Goldstone bosons,

$$\begin{aligned} \begin{pmatrix} G_W^\pm \\ H^\pm \end{pmatrix} &= \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} \phi_1^\pm \\ \phi_2^\pm \end{pmatrix}, \\ \begin{pmatrix} G_Z^0 \\ A^0 \end{pmatrix} &= \begin{pmatrix} \cos\beta & \sin\beta \\ -\sin\beta & \cos\beta \end{pmatrix} \begin{pmatrix} \sqrt{2}Im\phi_1^0 \\ \sqrt{2}Im\phi_2^0 \end{pmatrix}, \\ \begin{pmatrix} H^0 \\ h^0 \end{pmatrix} &= \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} \sqrt{2}Re\phi_1^0 - v_1 \\ \sqrt{2}Re\phi_2^0 - v_2 \end{pmatrix} \end{aligned} \quad (3)$$

where $\tan\beta \equiv t_\beta = v_2/v_1$ and α is the mixing angle of the h^0, H^0 CP-even neutral Higgs sector. $G_{Z(W)}^{0(\pm)}$ are the would-be Goldstone bosons for Z^0 ($W^\pm$), respectively. And A^0 is the CP-odd neutral Higgs. $H^\pm$ are the charged physical Higgses.

In addition, we diagonalize the quark mass matrices and define the quark mass eigenstates. The resulting Higgs-fermion Lagrangian can be written in several ways [6]. We choose to display the form that makes the type-II model limit of the general type-III couplings apparent. In the model type-II (where $\eta_{ij}^{U,0} = \xi_{ij}^{D,0} = 0$) tree-level Higgs mediated flavor-changing neutral currents are automatically absent, whereas these are generally present for type-III couplings. The fermion mass eigenstates are related to the interaction eigenstates by biunitary transformations:

$$\begin{aligned} U_L &= V_L^U U_L^0 , & U_R &= V_R^U U_R^0 , \\ D_L &= V_L^D D_L^0 , & D_R &= V_R^D D_R^0 , \end{aligned} \quad (4)$$

and the Cabibbo-Kobayashi-Maskawa matrix is defined as $K \equiv V_L^U V_L^{D\dagger}$. It is also convenient to define "rotated" coupling matrices:

$$\begin{aligned} \eta^U(\xi^U) &\equiv V_L^U \eta^{U,0}(\xi^{U,0}) V_R^{U\dagger} , \\ \eta^D(\xi^D) &\equiv V_L^D \eta^{D,0}(\xi^{D,0}) V_R^{D\dagger} . \end{aligned} \quad (5)$$

The diagonal quark mass matrices are obtained by replacing the scalar fields with their VEV's:

$$M_D = \frac{1}{\sqrt{2}}(v_1\eta^D + v_2\xi^D) , \quad M_U = \frac{1}{\sqrt{2}}(v_1\eta^U + v_2\xi^U) . \quad (6)$$

After eliminating η^D, ξ^U, the resulting Yukawa couplings are [6]:

$$\begin{aligned} \pounds_Y = &\frac{1}{v}\overline{D}M_D D\left(\frac{s_\alpha}{c_\beta}h^0 - \frac{c_\alpha}{c_\beta}H^0\right) + \frac{i}{v}\overline{D}M_D\gamma_5 D(t_\beta A^0 - G_Z^0) \\ &-\frac{1}{\sqrt{2}c_\beta}\overline{D}(\xi^D P_R + \xi^{D\dagger}P_L)D(c_{\beta-\alpha}h^0 - s_{\beta-\alpha}H^0) - \frac{i}{\sqrt{2}c_\beta}\overline{D}(\xi^D P_R - \xi^{D\dagger}P_L)D\,A^0 \\ &-\frac{1}{v}\overline{U}M_U U\left(\frac{c_\alpha}{s_\beta}h^0 + \frac{s_\alpha}{s_\beta}H^0\right) + \frac{i}{v}\overline{U}M_U\gamma_5 U(t_\beta^{-1}A^0 + G_Z^0) \\ &+\frac{1}{\sqrt{2}s_\beta}\overline{U}(\eta^U P_R + \eta^{U\dagger}P_L)U(c_{\beta-\alpha}h^0 - s_{\beta-\alpha}H^0) - \frac{i}{\sqrt{2}s_\beta}\overline{U}(\eta^U P_R - \eta^{U\dagger}P_L)U\,A^0 \\ &+\frac{\sqrt{2}}{v}\left[\overline{U}KM_D P_R D(t_\beta H^+ - G_W^+) + \overline{U}M_U K P_L D(t_\beta^{-1}H^+ + G_W^+) + \text{h.c.}\right] \\ &-\left[\frac{1}{s_\beta}\overline{U}\eta^{U\dagger}KP_L D\,H^+ + \frac{1}{c_\beta}\overline{U}K\xi^D P_R D\,H^+ + \text{h.c.}\right] . \end{aligned} \quad (7)$$

where we have used the notation $s(c)_\alpha = \sin(\cos)\alpha$ and $\sin(\beta - \alpha) = s_{\beta-\alpha}$ and so on.

In the 2HDM type III after using the parameterization proposed by Cheng and Sher [5] for the couplings $\xi(\eta)_{ii} = \lambda_{ii} g m_i/(2m_W)$, we get the following expression for the decay width $t \to bH^+$,

$$\Gamma(t \to bH^+) = \frac{G_F K_{tb}^2}{4\pi\sqrt{2}} \quad \left[a^2(m_t^2 + m_b^2 - m_{H^\pm}^2) + 4abm_t m_b + b^2(m_t^2 + m_b^2 - m_{H^\pm}^2)\right] |\vec{p}_H| \quad (8)$$

where

$$\begin{aligned} |\vec{p}_H| &= \left[(m_t^2 - (m_b + m_{H^\pm})^2)(m_t^2 - (m_{H^\pm} - m_b)^2)\right]^{1/2}/(2m_t) , \quad (9) \\ a &= \cot\beta - \frac{\lambda_{tt}}{\sqrt{2}}\csc\beta , \\ b &= \frac{m_b}{m_t}\left(\tan\beta - \frac{\lambda_{bb}}{\sqrt{2}\cos\beta}\right) . \quad (10) \end{aligned}$$

Further we have taken the products $(\eta K)_{33} \sim \eta_{tt}K_{tb}$ and $(K\xi)_{33} \sim \xi_{bb}K_{tb}$, neglecting the off-diagonal terms because they are suppressed by the CKM entries. From the expression (8) is possible to get the decay width in the framework of the 2HDM-II just replacing $\lambda_{ii} = 0$.

In order to proceed with the numerical evaluations, we wonder about the perturbation regime. First of all, we are calculating a decay width at tree level, therefore we should take into account the possible values of λ which should be consistent with perturbation theory. Looking at the coupling $\bar{t}bH^+$ from (7), we get

$$\frac{m_b^2}{m_t^2}\left|\sqrt{2}\frac{t_\beta}{\sqrt{1+t_\beta^2}} - \lambda_{bb}\right|^2 + t_\beta^{-2}\left|\frac{\sqrt{2}}{\sqrt{1+t_\beta^2}} - \lambda_{tt}\right|^2 < \frac{8}{1+t_\beta^2} . \quad (11)$$

The allowed region in the plane $\lambda_{tt} - \lambda_{bb}$ depends on t_β, and it is shown in figure 1. This constraints will be taken into account in the following numerical evaluations.

On the other hand, we can consider the 2HDM-III in a basis where only one Higgs doublet acquire VEV and then it does not have the parameter $\tan\beta$. It is the usual 2HDM type III [5], where the Lagrangian of the charged sector is given by

$$-L_{H^\pm ud}^{III} = H^+\overline{U}[K\xi^D P_R - \xi^U K P_L]D + h.c. \quad (12)$$

With the above Lagrangian the model is simpler due to the absence of the $\tan\beta$ parameter, and it can be obtained from a Lagrangian written in a basis where both doublets acquire VEV different from zero [6]. For the decay width (8), it is reduced because now we have $a = \frac{\lambda_{tt}}{\sqrt{2}}$ and $b = \frac{\lambda_{bb}m_b}{\sqrt{2}m_t}$. And from the perturbation theory consideration we get

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

which is an ellipse with $|\lambda_{bb}| \le 100$ and $|\lambda_{tt}| \le \sqrt{8}$; this constraint might be satisfied. In this context, taking into account the bound from D0 experiment, $B(t \to H^+b) \le 0.36$ [9] we present in figure 2 the plane $m_{H^\pm}$ vs λ_{bb} for different values of λ_{tt}. We are using the constraint from perturbation in order to get the lower limit allowed in this plane. The allowed region is above the curve shown. For the allowed values of λ_{bb} the charged Higgs boson mass has a lower bound of ~ 140 GeV for $\lambda_{tt} = \pm 2.8$, but in general the bound is depending on the parameters λ_{ii}. Further in this plot when $a = b = 0$ the charged Higgs mass is not bounded as it is expected.

But notice that the bounds coming from the experiment (LEP and Tevatron) are really gotten in a model dependent way, specifically they have been gotten in the framework of the 2HDM-II. Then we should use the parameterization given by the Lagrangian (7) which is the 2HDM-II plus new changing flavor interactions. This reason makes useful the Lagrangian (7) in order to do comparisons of 2HDM-III with the experimental values obtained using the 2HDM-II. In figures 3 and 4, we show the fraction $B(t \to bH^+)$ vs $\tan\beta$ for different values of λ_{bb} and λ_{tt}. The $B(t \to bH^+)$ branching fraction for different values of $M_{H^\pm}$ is significant for very small or very large values of $\tan\beta$, while it is suppressed for intermediate values of $\tan\beta$. This fact is because the fraction $B(t \to bH^+)$ is proportional to $(m_b^2 + m_t^2 - m_{H^\pm}^2)(m_t^2\cot^2\beta + m_b^2\tan^2\beta) + 4m_t^2m_b^2$ which has a minimum around $\tan\beta = \sqrt{m_t/m_b}$ and is symmetric in $\log(\tan\beta)$ about this point. We note that for a charged Higgs 2HDM-II like, lighter than the top quark, it would be detected if $\tan\beta$ is substantially different from $\sqrt{m_t/m_b}$, it is because the branching fraction is very suppressed around this point, by around 10^{-4}. In what follows we are going to show that model type III can modify this scenario. For $\lambda_{ii} = 0$ we obtain the prediction of the 2HDM type II and it is symmetric around $\tan\beta = 6$, but it is not the case for $\lambda_{ii} \neq 0$ where the minimum is shifted. In this analysis we are considering two different set of values for the λ_{ii} parameters. In figure 3, λ_{tt} is the order of λ_{bb} and in figure 4 λ_{tt} is one order of magnitude smaller than λ_{bb}, and in both cases the charged Higss boson mass is fixed to 140 GeV. In figure 3 and 4, it is drawn the most stringent bound coming from D0 collaboration on $B(t \to bH^+)$ in the range of $0.3 \le \tan\beta \le 150$ which should be less than 0.36 (horizontal line)[9] and in the framework of the 2HDM-III there are some cases for the λ_{ii} where the lines are amputated due to the perturbative constraints coming from equation (11).

Finally, in figure 5 we plot $M_{H^\pm}$ vs $\tan\beta$ for different λ_{ii} using the bound from Tevatron $B(t \to bH^+) \le 0.36$. We show the excluded regions at 95 % C.L. by Tevatron from Run I and the limits that will be reached on this plane in Run II using 2 and 10 fb^{-1} for the integrated luminosity at $\sqrt{s} = 2$ TeV. We should clarify that the exclusion regions taken from D0 at Tevatron are a combination of two searches. An indirect search, looking for a decrease in $t\bar{t} \to W^+W^-b\bar{b}$ signal expected from the SM, this search excludes simultaneously both large and small $\tan\beta$. And a direct search, that look for the $H^\pm \to \tau^\pm\nu$ in the region $0.3 < \tan\beta < 150$. This is because the fraction rate of the leptonic decay of $H^\pm$ is around 0.95 for large $\tan\beta$, and the other option $H^+ \to c\bar{s}$ is important for $\tan\beta < 0.4$ and low Higgs boson mass. In our case we could enhanced the influence of $H^+ \to c\bar{s}$ channel due to the appearance of new couplings, it does not matter the value of $\tan\beta$, but that possibility corresponds to a very unusual set of values of the new couplings λ_{cc}, λ_{ss} and ξ^E, instead of that we conserve the hierarchy of the decay channels according to the value of $\tan\beta$ in order to consider a more general and conservative scenario where the experimental limits can be used. The solid line inside the future explored region corresponds to $\lambda_{ii} = 0$ which is the prediction of the 2HDM-II with a fraction rate of 0.36. A charged Higgs 2HDM-III like could have different scenarios, as we can see from figure 5 for λ_{ii} different from zero. The excluded region is below the curve. Again we are considering two cases: λ_{tt} of the order of λ_{bb} and λ_{tt} one order of magnitude smaller than λ_{bb}. Some lines are amputated because of equation (11) which is a contraint for the values of the parameters in the model taking into account the validness of perturbation theory. Cases like $\lambda_{tt} = 1, 2$ which are excluded for small $\tan\beta$. The case of $\lambda_{tt} = 4$ is valid in the whole range of $\tan\beta$ but it has a minimum value for the charged Higgs boson mass around 105 GeV in contrast the case of $\lambda_{tt} = 1$, $\lambda_{bb} = 10$ is only valid for $\tan\beta \le 10$ but it is not presenting any lower bound for the chargued Higgs boson mass. On the other hand, the case of $\lambda_{tt} = 3$ is a weak constraint for small and intermediate values of $\tan\beta$ because the experimental limit from LEP on the chargued Higgs mass (horizontal line) is stronger in these regions but it is stronger than the LEP limit in the region of large $\tan\beta$.

To summarize, in the present work we have examined a 2HDM type III which produces FCNC at tree level, in general these new interactions are governed by parameters λ_{ij}. The experimental analysis has been carried out using the 2HDM type II as a framework, so they are model dependent. We have already presented a form of the Yukawa Lagrangian of the 2HDM-III (7) which can be reduced to the 2HDM-II as a limit [6], in this way it can be used to compare the model type III with the experimental analysis based on 2HDM type II. We have shown that 2HDM type III can modify the situation for the branching fraction $B(t \to bH^+)$ in the case of a charged Higgs boson in the allowed region from kinematic considerations. For $\lambda_{ii} = 0$, we obtain the prediction of the 2HDM type II and it is symmetric around $\tan\beta = \sqrt{m_t/m_b} \sim 6$, but it is not the case for $\lambda_{ii} \neq 0$. Finally, we have presented the parameter plane $\tan\beta - m_{H^+}$ showing the experimental limits for Run I and II and, in the same plot we show the solutions obtained in the cases of 2HDM-II ($\lambda_{ii} = 0$) and 2HDM-III ($\lambda_{tt} \neq 0$) taking into account the criterium of validness of perturbation theory.

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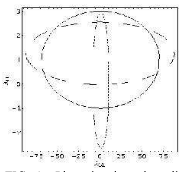

FIG 1. Plot showing the allowed region for λ_{tt} - λ_{bb} according to equation (11), using different values of $\tan\beta$ (1 solid line, 0.5 long dashed line and 10 short dashed line).

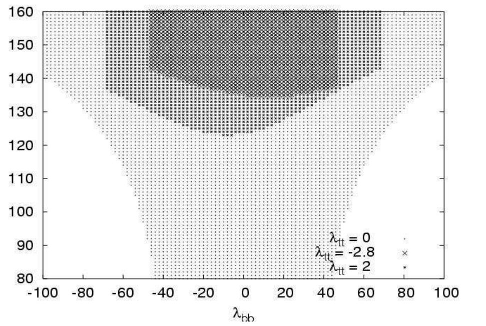

FIG 2. Plot for the charged Higgs boson mass $M_{H\pm}$ versus the parameter λ_{bb} . In the 2HDM-III when $\tan\beta$ is not present. The allowed region is shown when λ_{tt}=0 (points), λ_{tt}=-2.8 (crosses) and λ_{tt}=2 (stars)

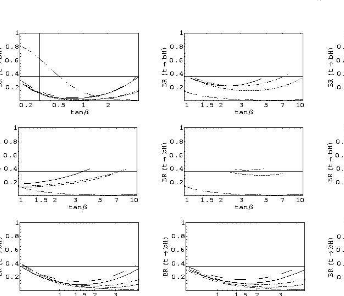

FIG 3. Plot for $B(t \to bH^+)$ versus $\tan\beta$ parameter for M_{H+} =140 GeV and different values of the parameter λ_{ii}. Each plane corresponds to a fixed λ_{tt}, downward= ±1, ±2, ± 0.1 right side positive values, left side negatives values. And λ_{bb}=-2 dashed-line, = -1 dot dashed line =1 short-dashed line, and = 2 solid line. We also show the 2HDM-II case, dot-dot-dashed line, showing the minimum around $\tan\beta \sim 6$. The excluded region by Tevatron (horizontal line) corresponds to the right-up corner and the region for $\tan\beta<0.3$ (vertical line) is an unexplored region by experiments

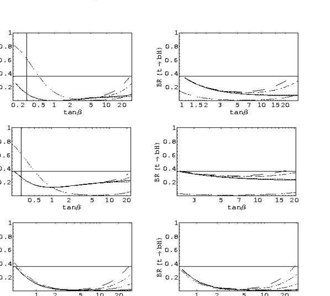

FIG. 4. Like figure 3 but λ_{bb}=-20 dashed line = 10 short dashed line, = 20 solid line.

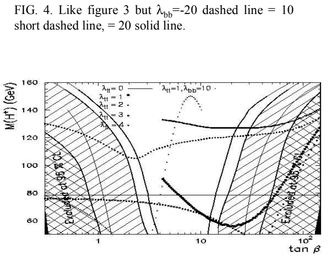

FIG. 5. Plot fon $M_{H\pm}$ versus the tan when $B(t \to bH^+) \le$ 0.36 for different set of values of λ_{ii}. Inside the figure are the labels for λ_{tt} and λ_{bb}=1 except for the last one wich λ_{bb}=10. We overlap the expected limits from D0 for mt=175 GeV and several values of the integrated luminosity and $\sqrt{s}$ (0.1 fb^{-1} , 1.8 TeV), (2 fb^{-1} , 2 TeV) , assuming $\sigma(tt)$=7pb. The limits were taken from reference [11]

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