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Constraining the charged Higgs mass in the left–right symmetric model from $b \to s\gamma$

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ABSTRACT

In the context of the left–right symmetric model, the decay $b \to s\gamma$ receives contributions from the gauge interactions mediated mainly by the W_L, through W_L–W_R mixing and also from the Yukawa interactions of the charged and the neutral (flavour-changing) scalars (the latter type of Yukawa interaction has been overlooked in the previous literature). Following the recent CLEO measurement of the inclusive $b \to s\gamma$ process and the measurement of the top-quark mass by the CDF and D0 collaborations, the parameter space of the left–right symmetric model is more squeezed than before.

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In the $SU(2)_L \otimes SU(2)_R \otimes U(1)$ left–right symmetric model (LRM) [1] there is a new scale at which the gauge group breaks to the $SU(2)_L \otimes U(1)$ Standard Model (SM). The sensitivity to this scale of low-energy phenomena such as K–$\bar{K}$ mixing and neutrino masses [2] has been a subject of wide interest over the past few years. Of late, a particularly interesting channel to examine various species of new physics, including this LRM scenario, has been provided by the inclusive B-decay measurement by the CLEO collaboration, $B(b \to s\gamma) = (2.32 \pm 0.57 \pm 0.35) \times 10^{-4} \to (1.0 - 4.2) \times 10^{-4}$ (at 95% C.L.) [3]. It has already been pointed out ([4]–[8]) that this rare decay has a strong influence on restricting the parameter space of the LRM. Recently, Cho and Misiak [6] investigated the effects of W_L–W_R mixing on $b \to s\gamma$ with an extensive analysis of QCD corrections which are very important for this process; however, they have not considered the contributions from the scalar sector. Babu et al. [7] included the charged scalars in the analysis, but their treatment of QCD corrections is incomplete. In this paper we attempt to improve upon the previous analyses by

- including all the above contributions coherently in a single analysis,
- incorporating the contribution of the flavour-violating neutral scalars, which has so far been *overlooked*, and
- reexamining the parameter space in the light of the new CLEO measurement of the $b \to s\gamma$ inclusive branching ratio [3] and the recent CDF and D0 measurements of the top-quark mass as $m_t = 180 \pm 12$ GeV [9] [1].

In the $SU(2)_L \otimes SU(2)_R \otimes U(1)$ gauge model the quarks (q) and the leptons (l) transform as $q_L(2,1,1/3)$, $q_R(1,2,1/3)$, $l_L(2,1,-1)$, and $l_R(1,2,-1)$. The scalar sector consists of the following Higgs fields: $\Delta_L(3,1,2)$, $\Delta_R(1,3,2)$ and $\Phi(2,2,0)$, of which only the latter participates in the Yukawa interaction and is explicitly shown as:

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix}. \tag{1}$$

Δ_R is used to break $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ to $SU(2)_L \otimes U(1)$, while Δ_L is introduced to maintain a discrete parity invariance. The vev v_R of Δ_R sets the LRM breaking scale. The vevs of Φ are given by $\langle \phi_1^0 \rangle = k$ and $\langle \phi_2^0 \rangle = k'$. One requires $v_R \gg k, k'$ from the excellent agreement between the $(V - A)$ theory and the experimental data, while the hierarchy $v_L \ll k, k'$ is set from the ρ parameter constraint. The SM gauge group is reproduced in the limit $v_R \to \infty, v_L \to 0$. We neglect any small phase difference between k and k'.

Introducing $\tan\beta = k/k'$, we obtain the physical charged scalars and the second set of neutral scalar and pseudoscalar [2] as

$$H^\pm = \cos\beta\ \phi_1^\pm + \sin\beta\ \phi_2^\pm,$$

[1] The value is the weighted average of $m_t = 176 \pm 13$ GeV (CDF) and $m_t = 199 \pm 30$ GeV (D0).

[2] The neutral scalar and the pseudoscalar in the first set are identified with the Higgs and the longitudinal component of the lighter Z, respectively, whose Yukawa couplings are flavour-diagonal.

$$
\begin{aligned}
H_2^0 &= \sqrt{2}\left(-\sin\beta \ \mathrm{Re}\phi_2^0 + \cos\beta \ \mathrm{Re}\phi_1^0\right), \\
G_2^0 &= \sqrt{2}\left(\cos\beta \ \mathrm{Im}\phi_1^0 + \sin\beta \ \mathrm{Im}\phi_2^0\right).
\end{aligned} \tag{2}
$$

The Yukawa interaction in the quark sector is given by the Lagrangian [7] [3]:

$$
L_Y = \bar{q}_L h \Phi q_R + \bar{q}_L \tilde{h} \tilde{\Phi} q_R + \text{h.c.} \tag{3}
$$

where $\tilde{\Phi} \equiv \tau_2 \Phi^* \tau_2$ and $h, \tilde{h}$ are 3×3 Hermitian matrices in flavour space. The up- and down-type quark mass matrices are given by:

$$
\begin{aligned}
M_u &= hk + \tilde{h}k', \\
M_d &= hk' + \tilde{h}k.
\end{aligned} \tag{4}
$$

After a straightforward calculation, the charged and neutral current (which are relevant to $b \to s\gamma$) Yukawa Lagrangian, in a more transparent form, are given by:

$$
\begin{aligned}
L_Y^C &= -\frac{N}{\cos 2\beta} \bar{u}_i \left[\sin 2\beta \left(\hat{M}_u V P_L - V \hat{M}_d P_R\right)\right. \\
&+ \left.\left(\hat{M}_u V P_R - V \hat{M}_d P_L\right)\right] H^+ d_i + \text{h.c.}
\end{aligned} \tag{5}
$$

and

$$
L_Y^N(d) = \frac{N}{\sqrt{2}\cos 2\beta} \bar{d}_i \left[(V^\dagger \hat{M}_u V)_{ij} - \sin 2\beta \hat{M}_d \delta_{ij}\right] \left(H_2^0 - i\gamma_5 G_2^0\right) d_j, \tag{6}
$$

where $\hat{M}_u$ and $\hat{M}_d$ are diagonal up- and down-type quark mass matrices, V is the standard Cabibbo-Kobayashi-Maskawa mixing matrix and $N = 1/\sqrt{k^2 + k'^2}$.

The contribution of the flavour-violating neutral scalars has been overlooked in the previous literature, which we find to be significant for some region of the parameter space. It may be noted that H_2^0 and G_2^0 mediate flavour violation even in the limit $\beta \to 0$. This originates from the fact that the Yukawa interaction mediated by the bidoublet scalar in LRM does not reproduce the SM scenario in the above limit.

In the charged gauge boson sector, $W_L^\pm$ and $W_R^\pm$ mix to give the lighter and heavier mass eigenstates as [4]

$$
\begin{aligned}
W_1^\pm &= \cos\xi \ W_L^\pm + \sin\xi \ W_R^\pm, \\
W_2^\pm &= -\sin\xi \ W_L^\pm + \cos\xi \ W_R^\pm.
\end{aligned} \tag{7}
$$

where the mixing angle ξ is given by

$$
\xi \simeq \sin 2\beta \frac{m_{W_1}^2}{m_{W_2}^2} \equiv \sin 2\beta \frac{m_W^2}{m_{W_2}^2}. \tag{8}
$$

[3]We use the same notation as in [7] as much as possible.

[4]Indeed, $W_1^\pm$are identified with the $W^\pm$ of the SM.

The effective Lagrangian relevant for the process $b \to s\gamma$ can be written as [5]:

$$
\begin{aligned}
\mathcal{L}_{eff} &= \sqrt{\frac{G_F^2}{8\pi^3}} V_{tb} V_{ts}^* \, m_b \Big[\sqrt{\alpha} \left\{ A_{\gamma L} \, \overline{s}_L \sigma^{\mu\nu} b_R + A_{\gamma R} \, \overline{s}_R \sigma^{\mu\nu} b_L \right\} F_{\mu\nu} \\
&+ \sqrt{\alpha_S} \left\{ A_{gL} \, \overline{s}_L T_a \sigma^{\mu\nu} b_R + A_{gR} \, \overline{s}_R T_a \sigma^{\mu\nu} b_L \right\} G^a_{\mu\nu} \Big] + \text{h.c.},
\end{aligned} \tag{9}
$$

where the contributions to $A_{\gamma L}, A_{\gamma R}, A_{gL}$ and A_{gR} from the different sectors [6] are given by ($x = m_t^2/m_W^2, y = m_t^2/m_{H^+}^2, z_1 = m_b^2/m_{H_2^0}^2, z_2 = m_b^2/m_{G_2^0}^2$) [7]:

$$
\begin{aligned}
A_{\gamma L} &= A_\gamma^{\rm SM}(x) + \xi \frac{m_t}{m_b} A_\gamma^{\rm mix}(x) + \tan^2 2\beta \; F_\gamma(y) + \frac{m_t \sin 2\beta}{m_b \cos^2 2\beta} \; G_\gamma(y) \\
&+ \frac{Q_b}{4} \frac{m_t}{m_{H_2^0}^2} \left(\frac{m_t}{\cos^2 2\beta} - \frac{m_b \sin 2\beta}{\cos^2 2\beta} \right) \left\{ H(z_1) + G(z_2) \right\}, \\
A_{\gamma R} &= \xi \frac{m_t}{m_b} A_\gamma^{\rm mix}(x) + \sec^2 2\beta \; F_\gamma(y) + \frac{m_t \sin 2\beta}{m_b \cos^2 2\beta} \; G_\gamma(y) \\
&+ \frac{Q_b}{4} \frac{m_t}{m_{H_2^0}^2} \left(\frac{m_t}{\cos^2 2\beta} - \frac{m_b \sin 2\beta}{\cos^2 2\beta} \right) \left\{ H(z_1) + G(z_2) \right\}, \\
A_{gL} &= A_g^{\rm SM}(x) + \xi \frac{m_t}{m_b} A_g^{\rm mix}(x) + \tan^2 2\beta \; F_g(y) + \frac{m_t \sin 2\beta}{m_b \cos^2 2\beta} \; G_g(y) \\
&+ \frac{1}{4} \frac{m_t}{m_{H_2^0}^2} \left(\frac{m_t}{\cos^2 2\beta} - \frac{m_b \sin 2\beta}{\cos^2 2\beta} \right) \left\{ H(z_1) + G(z_2) \right\}, \\
A_{gR} &= \xi \frac{m_t}{m_b} A_g^{\rm mix}(x) + \sec^2 2\beta \; F_g(y) + \frac{m_t \sin 2\beta}{m_b \cos^2 2\beta} \; G_g(y) \\
&+ \frac{1}{4} \frac{m_t}{m_{H_2^0}^2} \left(\frac{m_t}{\cos^2 2\beta} - \frac{m_b \sin 2\beta}{\cos^2 2\beta} \right) \left\{ H(z_1) + G(z_2) \right\}.
\end{aligned} \tag{10}
$$

The functions in the above expressions are given by [8]

$$
\begin{aligned}
A_\gamma^{\rm SM} &= \frac{x(8x^2 + 5x - 7)}{24(1-x)^3} + \frac{x^2(3x-2)}{4(1-x)^4} \ln x, \\
A_g^{\rm SM} &= \frac{x(x^2 - 5x - 2)}{8(1-x)^3} - \frac{3x^2}{4(1-x)^4} \ln x, \\
A_\gamma^{\rm mix} &= \frac{-20 + 31x - 5x^2}{12(1-x)^2} + \frac{x(3x-2)}{2(1-x)^3} \ln x,
\end{aligned}
$$

[5] We neglect the contributions proportional to m_s.

[6] We do not show, in the expressions for $A_{\gamma R}$ and A_{gR}, the W_2-induced contributions which are sufficiently damped since m_{W_2} is set to 1.6 TeV [10] all along our analysis. However, we have included it in our numerical code.

[7] Whenever we encounter a b-quark inside a triangle loop, we take the running mass $m_b(m_W) \sim 3$ GeV, as in [7].

[8] $A_\gamma^{\rm SM}$ and $A_g^{\rm SM}$ were first calculated by Inami and Lim [11].

$$\begin{aligned}
A_g^{\rm mix} &= -\frac{4+x+x^2}{4(1-x)^2} - \frac{3x}{2(1-x)^3}\ln x, \\
F_\gamma(x) &= \frac{x(-25+53x-22x^2)}{72(1-x)^3} + \frac{x^2(2-x)}{12(1-x)^4}\ln x, \\
G_\gamma(x) &= \frac{x(5x-3)}{12(1-x)^2} - \frac{x(2-3x)}{6(1-x)^3}\ln x, \\
F_g(x) &= \frac{-20x+19x^2-5x^3}{24(1-x)^3} - \frac{x}{4(1-x)^4}\ln x, \\
G_g(x) &= \frac{-3x+x^2}{4(1-x)^2} - \frac{x}{2(1-x)^3}\ln x, \\
H(x) &= \frac{16-29x+7x^2}{12(1-x)^3} + \frac{2-3x}{2(1-x)^4}\ln x, \\
G(x) &= \frac{20-19x+5x^2}{12(1-x)^3} + \frac{2-x}{2(1-x)^4}\ln x.
\end{aligned} \tag{11}$$

It should be noted that there is a *disagreement of sign* between Babu et al.'s [7] and Asatryan et al.'s [5] calculation of the charged scalar-induced contribution. Our calculation *agrees with the latter*.

The effective Wilson coefficients can now be written as:

$$C_{7L}^{\rm eff} = \eta^{16/23}A_{\gamma L} + \frac{8}{3}(\eta^{14/23}-\eta^{16/23})A_{gL} + C + C', \tag{12}$$

$$C_{7R}^{\rm eff} = \eta^{16/23}A_{\gamma R} + \frac{8}{3}(\eta^{14/23}-\eta^{16/23})A_{gR} + C'. \tag{13}$$

In the above equations, $\eta = \alpha_S(M_Z)/\alpha_S(\mu)$, where μ is the QCD renormalization scale; C corresponds to the leading log QCD corrections in SM [12], and C' refers to the extra contribution from mixing of additional operators in LRM, which has been computed in ref. [6]. These are given by

$$C = \sum_{i=1}^{8} h_i \eta^{a_i}, \tag{14}$$

where

$$a_i = \frac{14}{23}, \frac{16}{23}, \frac{6}{23}, -\frac{12}{23}, 0.4086, -0.4230, -0.8994, 0.1456 \tag{15}$$

$$h_i = \frac{626126}{272277}, -\frac{56281}{51730}, -\frac{3}{7}, -\frac{1}{14}, -0.6494, -0.0380, -0.0186, -0.0057 \tag{16}$$

and

$$C' = \xi\frac{m_c}{m_b}\sum_{i=1}^{4} h_i' \eta^{a_i'}, \tag{17}$$

where

$$a_i' = (0.6957, 0.6087, -1.0435, 0.1304) \tag{18}$$

$$h'_i = (-0.6615, 1.3142, 0.0070, 1.0070). \tag{19}$$

Finally, the branching ratio of $b \to s\gamma$ is given, in units of the semileptonic b-decay branching ratio, by

$$\frac{B(b \to s\gamma)}{B(b \to ce\overline{\nu})} = \frac{6\alpha}{\pi\rho\lambda} \left|\frac{V_{tb}V_{ts}^*}{V_{bc}}\right|^2 \left[|C_{7L}^{\rm eff}|^2 + |C_{7R}^{\rm eff}|^2\right], \tag{20}$$

where $\rho = (1 - 8r^2 + 8r^6 - r^8 - 24r^4 \ln r)$ with $r = m_c/m_b$ and $\lambda = 1 - 1.61\ \alpha_S(m_b)/\pi$. It may be noted that the ${m_b}^5$ dependence in the partial decay widths of the b quark cancels out in eq. (20). An $\mathcal{O}(m_s^2/m_b^2)$ part in the branching ratio is neglected. We take $B(b \to ce\overline{\nu}) = 0.107$.

In Figs. 1–3 we have plotted the branching ratio $B(b \to s\gamma)$ as a function of m_{H^+}. We have fixed $m_{W_2} = 1.6$ TeV. To demonstrate the magnitude of the H_2^0-induced effect, we work with two representative values: $m_{H_2^0} = 300$ GeV and 800 GeV and have displayed their effects in Figs. 1 and 2, respectively, for $m_t = 180$ GeV. On the other hand, the chirality-flip of the top quark inside the loop due to the presence of the right-handed current is responsible for a strong m_t-dependence of our prediction. To illustrate this point, we demonstrate in Fig. 3 how a line curve in Fig. 1 corresponding to $\beta = -10°$, as an example, becomes a thick band due to the variation of m_t in the range 168–192 GeV. In Fig. 4 we fix the charged Higgs mass to 800 GeV and plot the branching ratio $B(b \to s\gamma)$ as a function of $m_{H_2^0}$. The salient features of the relative contributions of the different sectors of the LRM that emerge from the above figures are listed below:

1. The contributions induced by the W_L–W_R mixing and the charged scalar are very sensitive to the choice of β. The reason is that the chirality-flipped (m_t/m_b)-enhancement factor, which constitutes the potentially largest contribution, multiplies $\sin 2\beta$. Even a choice of $|\beta| \sim 5°$ can rule out a charged Higgs up to a mass of several hundred GeV. Evidently, choosing a larger $|\beta|$ pushes it up even further.

2. Contrary to the Yukawa couplings of H_1^0 or G_1^0, those of H_2^0 and G_2^0 are *not* totally flavour-diagonal and hence they contribute to $b \to s\gamma$. For the sake of simplicity we have assumed H_2^0 and G_2^0 to be mass-degenerate. With increasing $m_{H_2^0}$ the contribution from the neutral scalar sector decouples fast, which is the origin of the relative shifts between the curves in Fig. 2 ($m_{H_2^0} = 800$ GeV) and in Fig. 1 ($m_{H_2^0} = 300$ GeV). It may be noted, though, that the individual contributions from H_2^0 and G_2^0 are in the same direction and roughly of the same order of magnitude. So if one relaxes the condition of their mass degeneracy and their joint contribution is thought to be of the same order of magnitude as their individual contributions, the curves should lie somewhere between their corresponding positions in Fig. 1 and Fig. 2.

3. In the limit of $\beta = 0°$, the dominant contribution in the LRM comes from the new operator-mixing effect, represented by C', and also from the flavour-violating neutral scalar sector (when those neutral scalar masses are not too heavy).

4. The m_t-dependent (dominant) contributions multiply $1/(\cos^2 2\beta)$ in the neutral scalar sector and $\sin 2\beta/(\cos^2 2\beta)$ in the charged scalar sector. A close look at the relative signs of the associated factors (see eqs. (10) and (11)) in those contributions reveals that for a positive (negaive) β, there is a constructive (destructive) interference between the effects induced by the neutral and the charged scalars. The W_L–W_R mixing contribution also depends on the sign of β through ξ. A comparison between the curves for $\beta = -5^\circ$ and 5° in Fig. 4, as an example, amply demonstrates this analytic interplay involving the sign of β.

At this point, a few words about the theoretical uncertainties in this process are in order [13]. The evolution of $\mathcal{L}_{\rm eff}$ from m_W to a lower momentum scale ($\mu \sim m_b$) by the QCD renormalization group analysis involves a significant theoretical uncertainty regarding a precise choice of μ at which α_S is to be determined. The SM branching ratio for a leading log calculation is quoted as $B^{\rm SM}(b \to s\gamma) = (2.8 \pm 0.8) \times 10^{-4}$ [14] where the error comes mainly from the uncertainty of μ in the range $m_b/2 < \mu < 2m_b$. We take $\mu = m_b$ to obtain $B^{\rm SM}(b \to s\gamma) = 2.9 \times 10^{-4}$ for $m_t = 180$ GeV. Recently a part of the next-to-leading order QCD corrections has been estimated with a consequence of reducing the QCD enhancement in the SM yielding $B^{\rm SM}(b \to s\gamma) = (1.9 \pm 0.5) \times 10^{-4}$ [15]. It may be noted that all our estimates of the LRM contributions stand above the base value of the SM approximated at the leading order QCD corrections. If, for instance, the full calculation of the next-to-leading order QCD corrections pulls the SM estimate down, the total effect including the LRM contributions will also come down by the same absolute amount. However, for a non-negligible β, the total LRM contribution or even the contribution from each individual sector as well is shown to be numerically significant and hence their impact are not likely to be masked by the uncertainties from the next-to-leading order QCD corrections.

In conclusion: we have investigated the parameter space of the LRM in the context of $b \to s\gamma$. Different sectors of the LRM contributing to this process have been added coherently. We have included the effect of the flavour-violating neutral scalars, which was missing in the previous analyses. Upon imposition of the CLEO measurement of the inclusive $b \to s\gamma$ rate, the lower limit of the charged Higgs mass is pushed up to several hundred GeV even for a small value of $|\beta|$. Although there are lots of parameters in the LRM which can conspire, leading to cancellations and reducing the power of prediction in this model, still a reduction of errors in the inclusive $b \to s\gamma$ measurement at CLEO, a more precise determination of the top-quark mass and, indeed, a better understanding of the next-to-leading order QCD corrections will all serve to constrain the parameter space even more strongly.

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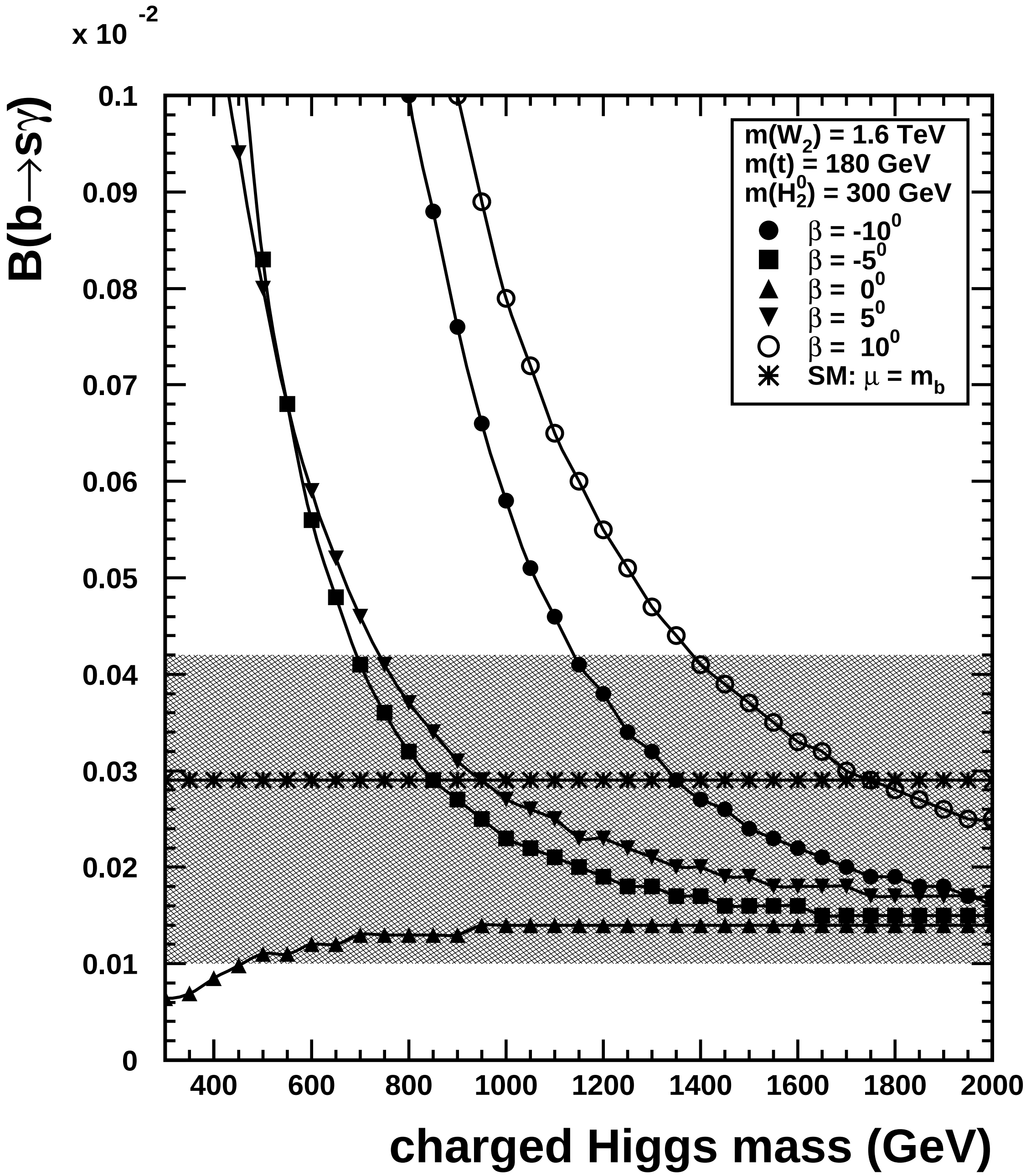

Figure 1: The branching ratio for the process $b \to s\gamma$ as a function of m_{H^+} for different values of β in the LRM. The value of m_t has been fixed to 180 GeV and $m_{H_2^0} = m_{G_2^0}$ has been set to 300 GeV. The SM line relies on leading log QCD calculation at $\mu = m_b$. The shaded area is the experimentally allowed region at 95% C.L. [3].

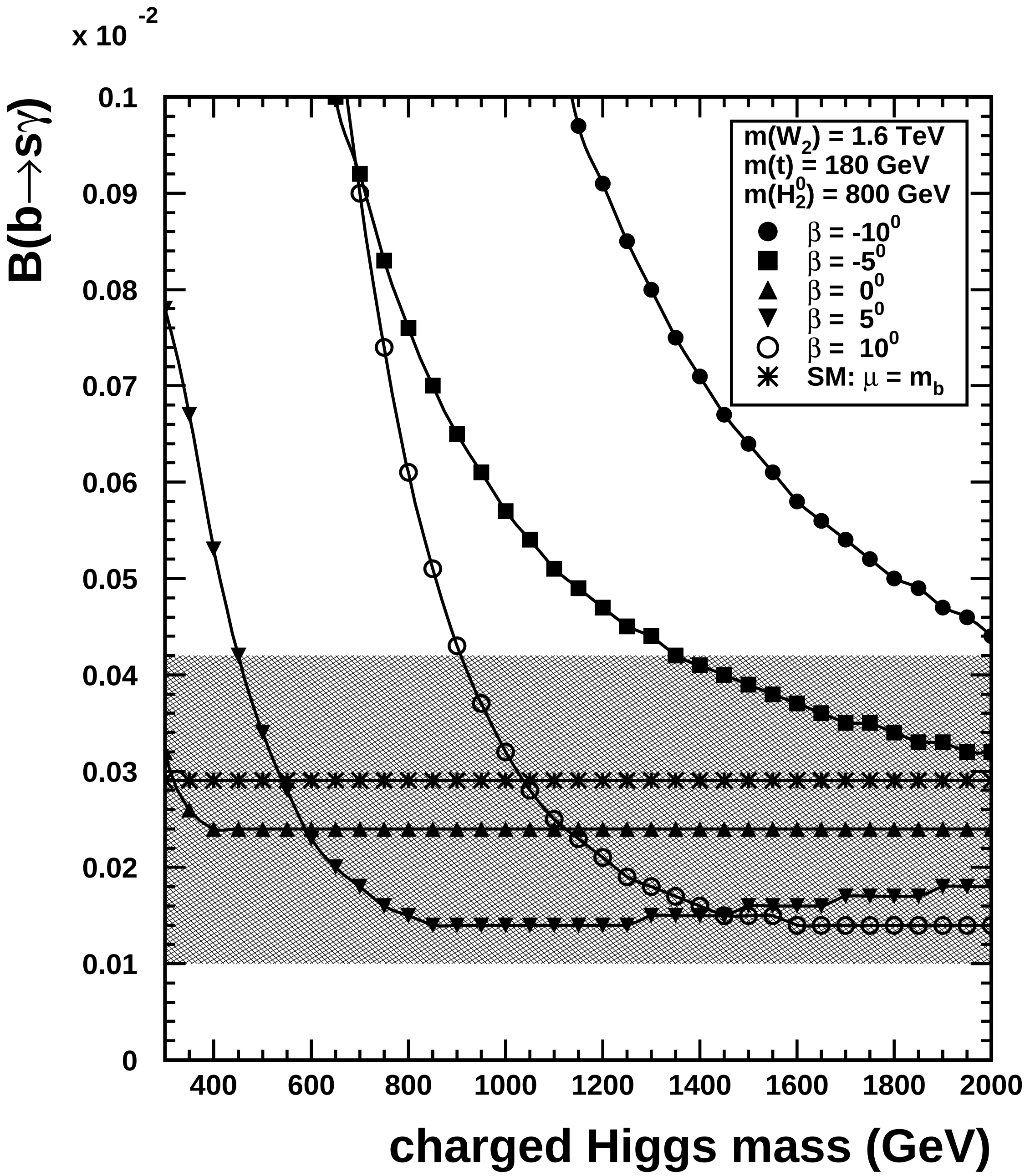

Figure 2: Same as in Fig. 1, but with $m_{H_2^0} = m_{G_2^0} = 800$ GeV.

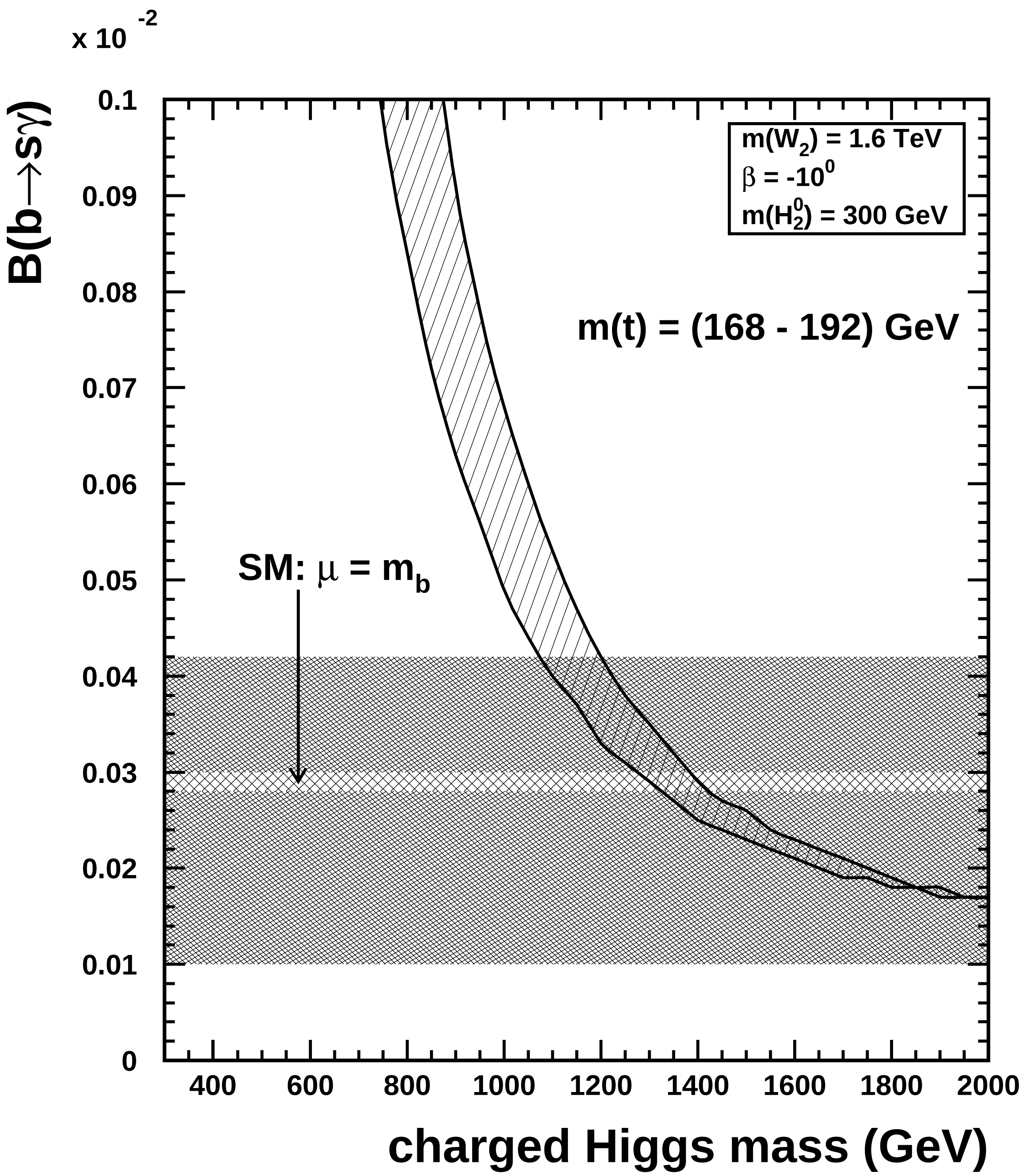

Figure 3: The m_t dependence of the result is exhibited by showing how the curve corresponding to $\beta = -10°$ (as an example) in Fig. 1 becomes a band (hatched). The SM line for $\mu = m_b$ also becomes a band (horizontal hatched strip within the shaded area) due to the same effect. The shaded area is the experimentally allowed region at 95% C.L. [3].

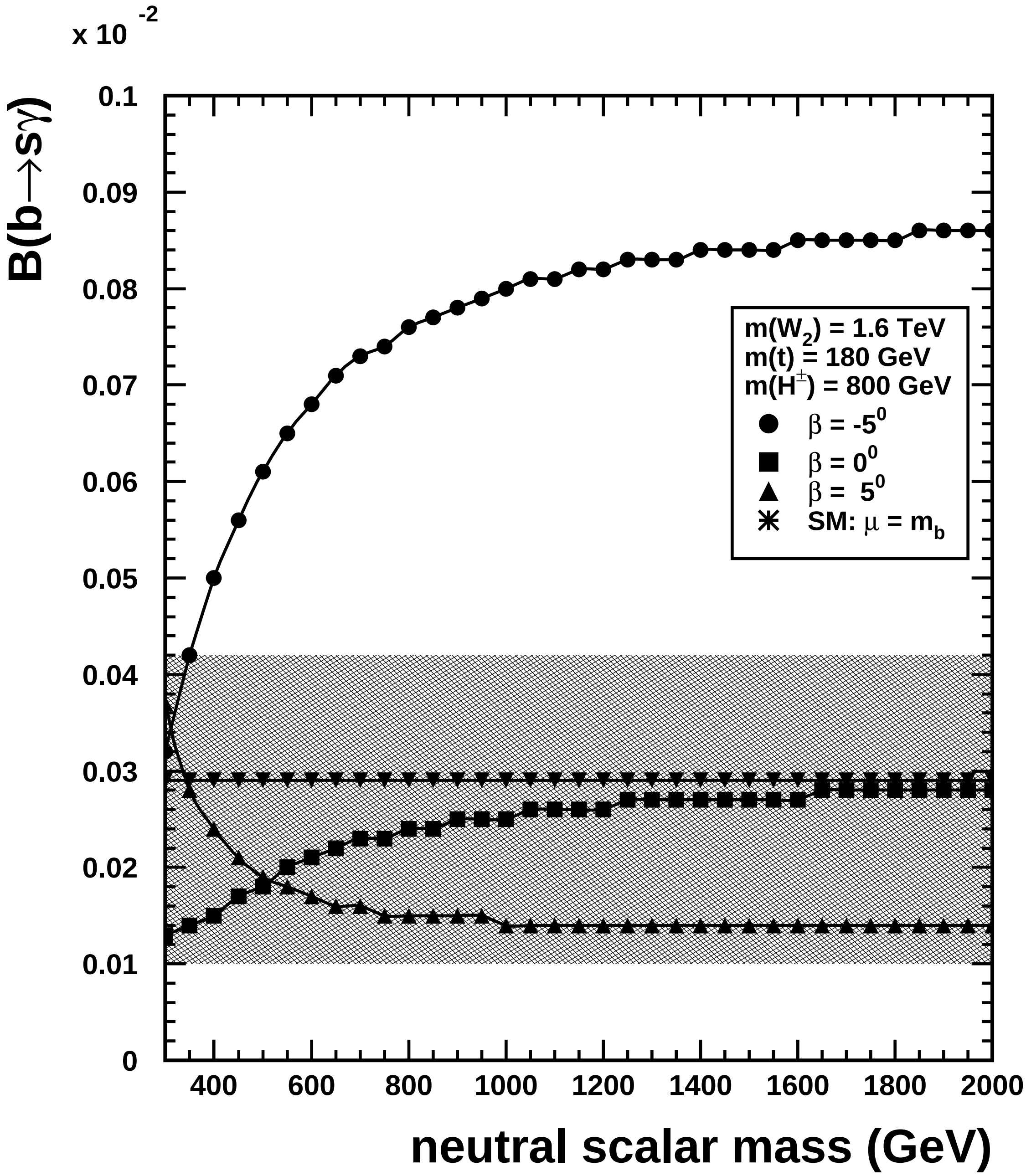

Figure 4: The branching ratio for the process $b \to s\gamma$ as a function of $m_{H_2^0}$ for different values of β in the LRM. H_2^0 and G_2^0 have been assumed to be mass-degenerate. The value of m_t has been fixed to 180 GeV and $m_{H^\pm}$ has been set to 800 GeV. The SM line relies on leading log QCD calculation at $\mu = m_b$. The shaded area is the experimentally allowed region at 95% C.L. [3].