THE CP ASYMMETRY IN $b \to s l^+ l^-$ DECAY

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Abstract

Using the experimental upper bound on the neutron EDM and experimental result on $b \to s\gamma$ branching ratio we have calculated CP asymmetry and $\Gamma^{2HDM}(b \to s l^+ l^-)/\Gamma^{SM}(b \to s l^+ l^-)$. It is shown that in the invariant dilepton mass q^2 region $(m_{\psi'}^2 + 0.2\ GeV^2) < q^2 < m_b^2$ the CP asymmetry is maximal and quite detectable.

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1 Introduction

The experimental discovery of the inclusive and exclusive decays $B \to X_s\gamma$ and $B \to K^*\gamma$ by the CLEO collaboration [1,2] has triggered a lot of theoretical and the experimental activity in the field of rare decays of B- mesons. These decays are interesting for checking the predictions of SM at one-loop level, for determining the CKM matrix elements, and for looking for the "new physics" beyond the SM. From the experimental point of view another promising decay in this direction is the semileptonic decay $b \to X_s l^+ l^-$, because this decay is easier to measure provided that we are given a good electromagnetic detector and a large number of B hadrons. Theoretically this decay has been the subject of many works in the framework of the SM [3,4,5,6] and its extensions, particulary in Two Higgs Doublet Model (2HDM).

$b \to s l^+ l^-$ decay is an FCNC process which appears only at the one-loop level of pertubation theory. The basic thing about this decay is that the penguin diagrams provide the two key ingredients needed for partial rate asymmetries. Being a loop diagram, it involves all three generations, each generation contributing with different elements of the CKM matrix. At the same time the loop effects that involve on- shell particle rescatterings provide the necessary absorbtive parts.

It is well known that in 2HDM, $b \to s l^+ l^-$ decay receives significant contributions from the charged Higgs ($H^{\pm}$) exchange [7]. Another interesting pecularity of 2HDM is the appearence of new sources of CP violation [8] in addition to the one in SM. An interesting version of 2HDM, so called the most general 2HDM, which was proposed in [9], has a new source of CP violation, arising from the relative phase between the vacuum expectation values of two Higgs scalars.

In this work we shall work out $b \to s l^+ l^-$ decay. In particular we shall determine the CP asymmetry A and the ratio $r = \Gamma^{2HDM}(b \to s l^+ l^-)/\Gamma^{SM}(b \to s l^+ l^-)$ as functions of the charged Higgs mass.

In the calculation of the CP asymmetry we shall consider both the SM and 2HDM contributions simultaneously. In determining r and A we shall make use of the experimental results on $BR(b \to s\gamma)$ [1,2], and the neutron electric dilpole moment (EDM).

Section 2 is devoted to the derivation of basic theoretical results and Section 3 contains the numerical analysis of them.

2 Formalism

In the most general 2HDM [8,9] the couplings of $H^\pm$ with t_R and b_R is characterised by the coefficients ξ_f defined by

$$\xi_f = \frac{sin\delta_f}{sin\beta cos\beta sin\delta} e^{i\sigma_f(\delta-\delta_f)} - cot\beta \tag{1}$$

where f= t or b, $\sigma_f = +$ for b and - for t, and $\delta_f = h_2/h_1$ where h_2 and h_1 are the diagonal elements of the matrices Γ_2^u and Γ_1^u respectively. Here Γ^u are the matrices in the flavour space, and determine the Yukawa couplings (for more detail see [9]), and δ is the relative phase between the vacuum expectations of the two Higgs scalars:

$$\begin{aligned} < \phi_1^0 > &= \frac{v}{\sqrt{2}} cos\beta e^{i\delta} \\ < \phi_2^0 > &= \frac{v}{\sqrt{2}} sin\beta \end{aligned} \tag{2}$$

The most general 2HDM reduces to the well-known 2HDM's in the current literature, in certain limiting cases [9]. Namely, if $\delta_t = \delta_b = 0$, then $\xi_t = \xi_b =$

$-cot\beta$ (Model I) and, if $\delta_b = \delta, \delta_t = 0$, then $\xi_t = -cot\beta, \xi_b = tan\beta$ (Model II).

As mentioned above the penguin diagrams provide the necessary absorbtive parts for the calculation of the CP asymmetry. In this decay the dilepton invariant mass q^2 ranges from $4m_l^2$ to m_b^2; therefore, u and c loops give rise to nonzero absorbtive parts which are described, at the point $\mu = m_b$, by

$$F = i4\sqrt{2}G_F\lambda_u\frac{\alpha}{4\pi}A_9\bar{s}_L\gamma_\mu b_L l^+\gamma_\mu l^- \tag{3}$$

where $\lambda_i = V_{is}V_{ib}^*$ and the function A_9 is given by

$$A_9 = w_u[Q(m_c^2/q^2) - Q(m_u^2/q^2)] \tag{4}$$

where

$$Q(x) = \frac{2\pi}{9}(2+4x)\sqrt{1-4x}\theta(1-4x) \tag{5}$$

and w_u, having the numerical value of 0.3864, comes from the RGE movement of the Wilson coefficients from $\mu = M_W$ to $\mu = m_b$ point.

It is well- known that in the range $(4m_l^2, m_b^2)$ one can create real low lying charmonium states [10,11]. In this work we shall discard that portion of total dilepton mass range including J/ψ and ψ' poles and the region between them to avoid the addition of new hadronic uncertainities to the decay amplitude. Thus we restrict ourselves to the following kinematical regions [6]:

$$\begin{aligned} Region\ I &: \ 4m_l^2 \leq q^2 \leq (m_\psi^2 - \tau) \\ Region\ II &: \ (m_{\psi'}^2 + \tau) \leq q^2 \leq m_b^2 \end{aligned} \tag{6}$$

where $\tau = 0.2GeV^2$ is the cut- off parameter.

Taking into account the 2HDM contributions and absorbtive part described by F in (3), the amplitude for $b \to s l^+ l^-$ can be written as

$$\begin{aligned} M_{b\to sl^+l^-} &= 4\sqrt{2}G_F \frac{\alpha}{4\pi} \times \\ & \{C_9^{eff}(\mu)\bar{s}_L\gamma_\mu b_L l^+\gamma_\mu l^- + \\ & C_{10}(\mu)\bar{s}_L\gamma_\mu b_L l^+\gamma_\mu\gamma_5 l^- + \\ & \frac{q^\nu}{q^2}\times C_7(\mu)\bar{s}\sigma_{\mu\nu}(m_b R + m_s L) b l^+\gamma_\mu l^-\} \end{aligned} \tag{7}$$

The Wilson coefficients appearing in (7) are given by

$$\begin{aligned} C_7(\mu) &= \lambda_t[C_7^{SM}(\mu) + C_7^{2HDM}(\mu)] \\ C_9^{eff}(\mu) &= \lambda_t[C_9^{SM}(\mu) + C_9^{2HDM}(\mu)] + i\lambda_u A_9 \\ C_{10}(\mu) &= \lambda_t[C_{10}^{SM}(\mu) + C_{10}^{2HDM}(\mu)] \end{aligned} \tag{8}$$

The explicit forms of $C_i^{SM}(\mu)$, (i=7,9,10) including leading and next-to-leading order QCD corrections can be found in [3,12,13,14]. The 2HDM contributions, $C_i^{2HDM}(\mu)$, in the framework of the most general 2HDM [9] are given by

$$\begin{aligned} C_7^{2HDM}(\mu) &= \mid \xi_t \mid^2 K_7^{tt} + (R_{tb} + iI_{tb})K_7^{tb} \\ C_9^{2HDM}(\mu) &= \mid \xi_t \mid^2 K_9^{tt} \\ C_{10}^{2HDM}(\mu) &= \mid \xi_t \mid^2 K_{10}^{tt} \end{aligned} \tag{9}$$

where $R_{tb} = Re[\xi_t \xi_b^*]$, $I_{tb} = Im[\xi_t \xi_b^*]$ and

$$\begin{aligned} K_7^{tb} &= \eta^{16/23}[G(y) - \frac{8}{3}(1-\eta^{-2/23})E(y)] \\ K_7^{tt} &= \frac{1}{6}\eta^{16/23}[A(y) + \frac{8}{3}(1-\eta^{-2/23})D(y)] \\ K_9^{tt} &= -\frac{-1+4s_W^2}{s_W^2}\frac{x}{2}B(y) + yF(y) \\ K_{10}^{tt} &= -\frac{1}{s_W^2}\frac{x}{2}B(y) \end{aligned} \tag{10}$$

with $x = m_t^2/M_W^2$, $y = m_t^2/M_H^2$, $s_W^2 = 0.2315$, $\eta = \alpha_s(M_W)/\alpha_s(m_b)$ and the explicit expressions for functions A, B, D, E, F, G can be found in [12].

As noted in [9], ξ_t is expected to be of order of unity or less, if the Yukawa couplings of the top quark is reasonable. We have shown that this happens to hold also for the decay process under consideration. Thus, without loosing generality, in what follows we set $\mid \xi_t \mid^2 = 0$ (all the conclusions remain in force for the case of $\mid \xi_t \mid^2 =1$ as well).

Using (7), the differential decay rate for $b \to s l^+ l^-$ is obtained as

$$\begin{aligned} \frac{d\Gamma^{2HDM}}{ds} &= \lambda_0 (1-s)^2 \{ 4(\frac{2}{s}+1) \mid C_7(\mu) \mid^2 + (1+2s)(\mid C_9^{eff}(\mu) \mid^2 + \mid C_{10}(\mu) \mid^2) \\ &+ 12 Re[C_7(\mu) C_9^{eff}(\mu)] \} \end{aligned} \tag{11}$$

where $s = q^2/m_b^2$, and $\lambda_0 = \frac{\alpha^2 G_F^2}{768\pi^5}$.

After integrating (11) over s we get

$$\begin{aligned} \gamma &= \gamma_0 + 4\rho I^2 + 2I(6I_9 + 6a_9^{(1)} R_{tu}) \\ &+ 4\rho R^2 + 2R(6R_9 + 6a_9^{(1)} I_{tu} + 4\rho C_7^{SM}) \\ &+ 12 a_9^{(1)} C_7^{SM} I_{tu} + a_9^{(2)} f_{tu} + 2(a_{r9} I_{tu} + a_{i9} R_{tu}) \end{aligned} \tag{12}$$

where

$$\begin{aligned} \gamma &= \frac{\Gamma^{2HDM}}{\lambda_0 \mid \lambda_t \mid^2} \\ \gamma_0 &= (\frac{\Gamma^{SM}}{\lambda_0 \mid \lambda_t \mid^2}) \mid_{A_9=0} \\ I &= I_{tb} K_{tb}^7 \\ R &= R_{tb} K_{tb}^7 \\ I_{tu} &= \frac{Im[\lambda_t \lambda_u^*]}{\mid \lambda_t \mid^2} \\ R_{tu} &= \frac{Re[\lambda_t \lambda_u^*]}{\mid \lambda_t \mid^2} \end{aligned} \tag{13}$$

$$f_{tu} = \frac{|\lambda_u|^2}{|\lambda_t|^2}$$

and the other parameters in (12) are defined by the following integrals:

$$\begin{aligned}
\rho &= \int ds(1-s)^2(\frac{2}{s}+1) \\
R_9 &= \int ds(1-s)^2 Re(C_9^{SM}) \\
I_9 &= \int ds(1-s)^2 Im(C_9^{SM}) \\
a_9^{(1)} &= \int ds(1-s)^2 A_9 \\
a_9^{(2)} &= \int ds(1-s)^2(1+2s)A_9^2 \\
a_{r9} &= \int ds(1-s)^2(1+2s)Re(C_9^{SM})A_9 \\
a_{i9} &= \int ds(1-s)^2(1+2s)Im(C_9^{SM})A_9
\end{aligned} \tag{14}$$

For the CP conjugate process, the analog of (12) can be obtained by the following replacements:

$$\bar{\gamma} = \gamma(I \to -I; \;\; I_{tu} \to -I_{tu}) \tag{15}$$

Now we introduce the parameter r that measures the relative strength of 2HDM and SM rates

$$r = \frac{\gamma}{\gamma_{SM}} \tag{16}$$

where γ_{SM} is obtained by setting $I = R = 0$ in (12).

Next we define the CP asymmetry by

$$A = \frac{\bar{\gamma} - \gamma}{\bar{\gamma} + \gamma} \tag{17}$$

Substituting the expressions for γ and γ_{SM} into (16) we obtain a circle for fixed values of r:

$$(R + R_0)^2 + (I + I_0)^2 = t(r-1) + R_0^2 + I_0^2 \tag{18}$$

where the parameters R_0 and I_0 are given by

$$\begin{aligned} R_0 &= \frac{3}{2\rho}(R_9 + \frac{2}{3}\rho C^7_{SM}) + r_0 \\ I_0 &= \frac{3}{2\rho}(I_9 + a_9^{(1)} R_{tu}) \end{aligned} \tag{19}$$

and the quantity $r_0 = \frac{3}{2\rho} a_9^{(1)} I_{tu}$ is introduced for later use.

On the other hand, insertion of (12) and (15) into (17) yields another circle

$$(R + R_0')^2 + (I + I_0')^2 = -t + \epsilon(1 - \frac{1}{A}) + R_0'^2 + I_0'^2 \tag{20}$$

where

$$\begin{aligned} I_0' &= \frac{I_0}{A} \\ R_0' &= \frac{3}{2\rho}(R_9 + \frac{2}{3}\rho C_7^{SM}) + \frac{r_0}{A} \end{aligned} \tag{21}$$

The parameters ϵ and t in (19) and (20) are given by

$$\begin{aligned} \epsilon &= \frac{I_{tu}}{4\rho}(12 a_9^{(1)} C_7^{SM} + 2 a_{r9}) \\ t &= -\frac{(1 - A_s)}{A_s}\epsilon \end{aligned} \tag{22}$$

where A_s is the CP asymmetry in SM which is obtained from (17) by:

$$A_s = A \mid_{I=R=0} \tag{23}$$

Up to this point, our analysis of $b \to s l^+ l^-$ decay parallels that of $b \to s\gamma$ in [9] except for the definition of A. We shall, however, analyze the circles in (18) and (20) in a different context by exploiting the relation between I and neutron EDM, and experimental results on $b \to s\gamma$ branching ratio [1,2].

First we obtain the expression for the CP asymmetry in (17) by subtracting (20) from (18) and solving for A:

$$A = \frac{1}{1-a} \tag{24}$$

where

$$a = \frac{tr}{\epsilon + 2II_0 + 2Rr_0} \tag{25}$$

Now we turn to the determination of I with the use of the experimental upper bound on neutron EDM. Weinberg has proposed a CP violating 6 dimensional gluonic operator [15]

$$O_6 \sim f_{abc} G_a^{\mu\rho} G_{b\rho}^{\nu} \tilde{G}_{c\mu\nu} \tag{26}$$

which has been shown to give very large contribution to neutron EDM, d_n by the neutral [15] or charged [16] Higgs exchange. Weinberg, after relating the hadronic matrix elements of O_6 to d_n, predicts the value of d_n on the basis of a Naive Dimensional Analysis (NDA). However a detailed analysis by Bigi and Uraltsev [17] reports a different value for d_n which equals $\frac{1}{30}$ of that of Weinberg's. The big difference between the results of these analyses is an indication of the existence of hadronic uncertainities which are mainly introduced by the matrix elements of O_6 between the nucleon states. In addition to these theoretical uncertainities, we have also problems with experimental data (in that experiment yields only an upper bound on neutron EDM). These can be summarized as

$$\begin{aligned} d_n^{theor} &= c_{theor} \times I_{tb} K(y) 10^{-25}\, e\, cm & (27) \\ d_n^{actual} &= c_{exp} \times d_n^{max} & (28) \end{aligned}$$

where c_{theor} and c_{exp} are constants and $\mid c_{exp} \mid$ is known to be less than unity. Let us note that c_{theor} is related to the theoretical uncertainities and c_{exp}

to the experimental uncertainities. Experiment yields $d_n^{max} = 1.1\ 10^{-25} e\, cm$ [18]. The function $K(y)$ in (27) is given by [16,17]:

$$K(y) = \frac{y}{(y-1)^3}[3/2 - 2y + y^2/2 + ln(y)] \tag{29}$$

The common point for the analyses in [16] and [17] is the presence of the function $K(y)$ which is equal to $\frac{1}{3}$ as $y \to 1$.

Equating (27) to (28) and defining $\beta = 1.1 \frac{c_{exp}}{c_{theor}}$, we obtain

$$I = \beta f(y) \tag{30}$$

where

$$f(y) = \frac{K_{tb}^7(y)}{K(y)} \tag{31}$$

Note that the constant β in (30) includes now both theoretical and experimental undeterminicies. We shall not make any assumption concerning the value of β; instead we are going to fix it through the use of the experimental results on $b \to s\gamma$ branching ratio.

The $b \to s\gamma$ decay amplitude is given by

$$M = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{4\pi} C_7(\mu) \bar{s}(p') \sigma_{\mu\nu} (m_b R + m_s L) b(p) F^{\mu\nu} \tag{32}$$

where $C_7(\mu)$ is defined in (8). Using the experimental result on the braching ratio of $b \to s\gamma$ decay [1,2] we get the following circle

$$(C_7^{SM} + R)^2 + I^2 = (C_7^{ex})^2 \tag{33}$$

where C_7^{ex} is the experimental value of $C_7(\mu)$

$$0.22 \leq \mid C_7^{ex} \mid \leq 0.30 \tag{34}$$

We shall determine the central values of β, r and A which are defined in equations (20), (16) and (17) respectively. In doing this, we will make use of circles in equations (18), (20) and (33) together with equation (30). Let us note that (30) is obtained by the use of the experimental upper bound on neutron EDM [18], and (33) is constructed with the use of the experimental data on $b \to s\gamma$ branching ratio [1].

Let us first determine β. For this purpose we consider the circle in (33) in the limit of infinitely large M_H or equivalently $y \to 0$. As $y \to 0$, $R \to 0$ and through (30), $I \to \beta f_0$, where numerically $f_0 = 0.2706$. Then equation (33), which is valid for any value of M_H, yields

$$\beta = \pm\{\frac{(C_7^{ex})^2 - (C_7^{SM})^2}{f_0^2}\}^{1/2} \tag{35}$$

With (35), I in (30) has now become a completely known function of M_H. Now we solve (33) for R, yielding

$$R = -C_7^{SM} + \sqrt{(C_7^{ex})^2 - I^2} \tag{36}$$

where the choice of plus sign is necessary to satisfy asymptotic condition on R.

Using (36) for R, and (30) for I we can solve equation (18) for r

$$r = 1 + \frac{(R + R_0)^2 + (I + I_0)^2 - R_0^2 - I_0^2}{t} \tag{37}$$

whose M_H dependence shall be discussed in the next section.

Finally, taking r from (37), R from (36) and I from (30) we determine the CP asymmetry A in (24) whose dependence on M_H shall also be studied in the next section.

3 Numerical Analysis

In the numerical analysis we shall use $m_u = 10MeV, \;\; m_c = 1.5GeV, \;\; m_b = 4.6GeV$. For the top quark mass we rely on the CDF data [19] and for the W mass we use $M_W = 80.22GeV$ [18].

In calculating I_{tu} and R_{tu} we use the parametrisation in [18], and in doing this we take the mid values of the quantities. For the phase δ_{13} of CKM matrix in [18] we shall use the the mid value of $cos\delta_{13} = 0.47 \pm 0.32$ given in [20] which icludes a large uncertainity. A straightforward calculation shows that corresponding to the uncertainity in $cos\delta_{13}$, R_{tu} and I_{tu} are uncertain by 3.87% and 23.75% respectively. Thus, the standard model asymmtery A_s in (23) is uncertain by 23.75%, and we shall use its central value in our calculations. This choice is justified by the closeness of I_{tu} and R_{tu} calculated in this way to that obtained by the use of Wolfenstein parametrisation [21].

Fig. 1 shows the variation of $f(y)$ in (29) with M_H for the lowest, central and the highest values of m_t permitted by the CDF data [19]. As we see from Fig. 1 dependence of $f(y)$ on m_t is very weak; thus, insensitivity of results to the variation of y with m_t is guaranteed. In what follows we shall use therefore the central value of CDF data $m_t = 176GeV$.

For $m_t = 176GeV$ we obtain $C_7^{SM} = -0.2686$. The $b \to s\gamma$ branching ratio has approximately 50% error [1] which is tranferred into a range of values that C_7^{ex} may take, as described by (34).

With the use of above-mentioned data we calculate SM CP asymmetry in (23) to be $A_s = 0.0714\%$ in Reg. I, and $A_s = 0.0223\%$ in Reg. II.

In the second column of Table 1 we give the values of β as $\mid C_7^{ex} \mid$ moves from its maximum value 0.30 towards $\mid C_7^{SM} \mid = 0.2686$. We see that $\mid \beta \mid$ decreases gradually with decreasing $\mid C_7^{ex} \mid$. Moreover, it is seen that the maximum value that $\mid \beta \mid \approx 0.5$.

Regarding the present calculations in [16] and [17] as the possible candidates for c_{theor} in (27), we can make certain predictions for c_{exp} in (28). A simple calculation yields $c_{theor} = 9.9$ and $c_{theor} = 0.33$ for Weinberg's NDA and Bigi-Uraltsev calculations respectively. In the case of NDA, a solution for c_{exp} exist only for $\mid \beta \mid <\sim 0.27$ at which d_n^{actual} turns out to be very close to its experimental upper bound. On the other hand, for Bigi-Uraltsev calculation, being a more detailed analysis, for all values of $\mid C_7^{ex} \mid$ ranging from $\mid C_7^{SM} \mid$ to 0.30 there exists a solution for c_{ex} with the help of which, through (28), one determine the value d_n^{actual}. In the third column of Table 1 we give the values of d_n^{actual} as $\mid C_7^{ex} \mid$ moves from its maximum value 0.30 towards $\mid C_7^{SM} \mid = 0.2686$. We observe that for $\mid C_7^{ex} \mid = 0.3$ $\mid d_n^{actual} \mid$ reaches its maximum value of $1.63\ 10^{-26}$ which is one order of magnitude less than the present experimental upper bound.

In our numerical analysis we use the range of values of M_H from $44 GeV$ [18] to $10 m_t$ [15]. In Fig. 2 and Fig.3 we show the variation of r in (37) with M_H in Regions I and II respectively. We observe that in both figures r is fairly high at low M_H and lands rapidly to a lower value after $M_H \sim 500 GeV$.

As we see from Fig.2, dependence of r on the sign of β in Region I is very weak. Moreover, for $M_H >\sim 1 TeV$, r attains the values ~ 1.056, ~ 1.0050, ~ 1.020, and ~ 1.016 for $\beta = +0.4938$, -0.4938, 0.2922, and -0.2922 respectively.

From Fig.3 we observe that in Region II dependence of r on the sign of β is large. Specificially, we see that, for large M_H, r becomes practically independent of M_H and attains the values ~ 1.021, ~ 0.998, ~ 1.01, and ~ 0.9996 corresponding to $\beta = +0.4938$, -0.4938, 0.2922 and -0.2922 respectively.

In Fig.4 and Fig.5 we show the variation of A in (24) with M_H in Regions I and II respectively. What we observe to be common between them is the saturation of CP asymmetry A to a certain value after $M_H \sim 500GeV$.

From Fig.4 we observe that the $2HDM$ CP asymmetry A, practically for all M_H, is of the same order as the SM CP asymmetry A_s. Indeed, especially for large M_H, corresponding to the values of β, $\beta = +0.4938,\ -0.4938,\ 0.2922$ and -0.2922, A attains the percentage values of $\sim -0.27,\ \sim 0.40,\ \sim -0.14,$ and ~ 0.28.

In Fig. 5 we observe that asymmetry A, as compared to the previous figure, is completey different in that it is positive and takes higher values for all values of M_H. Actually, we see that for small M_H, $2HDM$ CP asymmetry is larger than the SM CP asymmetry by approximately three orders of magnitude. For large M_H, however, A gets values which are larger than SM asymmetry by two orders of magnitue. Indeed, for large M_H, corresponding to the values of β, $\beta = +0.4938,\ -0.4938,\ 0.2922$ and -0.2922, A gets the following percentage values $\sim 1.1,\ \sim 3.25,\ \sim 0.2,$ and ~ 1.5

The last point to be noted about the Figs. 2-5 is that negative β gives rise to larger r and A than positive β does.

To decern a CP asymmetry A at the σ significance level with only statistical errors, the number of B hadrons N_B needed to demonstrate the asymmetry is given by[22]

$$N_B \approx \frac{\sigma^2}{BR \times A^2} \tag{38}$$

Now denoting the number of B hadrons to observe A_s, A in I and A in II by N_B^s, N_B^I and N_B^{II} respectively, we get, using the values of r and A we have obtained already, the following ratios

$$\begin{aligned} \frac{N_B^I}{N_B^s} &\approx 1 \\ \frac{N_B^{II}}{N_B^s} &\approx 10^{-4} \end{aligned} \qquad (39)$$

which clearly prove that Region II is more suitable for experimental investigations on A.

In conclusion we have determined the 2HDM CP asymmetry A, ratio of 2HDM decay rate to SM decay rate r and actual value of neutron EDM. In doing these we have utilized the experimental results on $b \to s\gamma$ branching ratio, and on the upper bound of neutron EDM. Both r and A relax to constant values after $M_H \sim 500 GeV$. This saturation property of quantities shows that if charged Higgs mass happens to be large ($\sim 1TeV$) then the most general $2HDM$ merely shifts the SM values of r and A to some other value which may be important for establishing 2HDM. Boldly speaking, in the high dilepton mass region (Region II) r is closer to unity and asymmetry is very large as compared to those in low dilepton mass region (Region I). Thus on the basis of the order of magnitude analysis carried out for N_B, we conclude that the high dilepton mass region is important and appropriate for experimental check of the quantities under concern. Region II [6] is accessible to the B experiments which will be carried out with hadron beams in CDF, HERA and LHC.

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Figure Captions

Figure 1: The M_H dependence of $f(y)$ for $m_t = 194 GeV$ (with circles), $m_t = 176 GeV$ (bare solid curve) and $m_t = 158 GeV$ (with squares).

Figure 2: The M_H dependence of r in Region I. Here labes 1, 2, 3 and 4 correspond to $\beta = 0.4938,\ -0.4938,\ 0.2922$ and -0.2922 respectively.

Figure 3: The same as in Fig. 2 but for Region II.

Figure 4: The M_H dependence of A in Region I. Labels have the same meaning as in Fig.1. Here the unlabled solid line shows the SM asymmetry.

Figure 5: The same as in Fig. 4 but for Region II.