FTUV/97-61

# CP Asymmetries in $B^0$ Decays in the Left-Right Model

G. Barenboim[1], J. Bernabéu[1] and M. Raidal[1,2]

[1] *Department of Theoretical Physics, University of Valencia, E-46100 Burjassot, Valencia, Spain*

[2] *KBFI, Academy of Estonia, Rävala 10, EE-0001 Tallinn, Estonia*

(December 23, 2013)



## Abstract

We study the time dependent CP asymmetries in $B^0_{d,s}$ decays in the left-right model with spontaneous breakdown of CP. Due to the new contributions to $B^0$-$\bar{B}^0$ mixing, the CP asymmetries can be substantially modified. Moreover, there can be significant new contributions to the $B$-meson decay amplitudes from the magnetic penguins. Most promising for detection of the new physics in the planned $B$ factories is that the CP asymmetries in the decays $B \to J/\psi K_S$ and $B \to \phi K_S$ which are supposed to be equal in the standard model can differ significantly in this class of models independently of the results in the measurements of $B \to X_s\gamma$.

11.30.Er, 12.60.-i, 13.20.Hw

CP violation, currently observed only in the neutral kaon system, is one of the least tested aspects of Nature. The standard model (SM) has specific predictions on the size as well as on the patterns of CP violation in $B$ meson decays [1] which, if disproved in the future $B$ factories, would signal the existence of new physics [2]. In $B^0$ decays new physics can possibly contribute to the $B^0_q$-$\bar{B}^0_q$ ($q = d, s$) mixing as well as to the decay amplitudes. The effect of the new physics in the mixing is universal, i.e., the time dependent rate asymmetries between $B^0_q$ and $\bar{B}^0_q$ in all their decays to the common CP eigenstates receive the same contribution. On the other hand, the effects of new physics in the decay amplitudes are non-universal and can show up in the comparison of the CP asymmetries in different decay modes [3,4].

In this Letter we analyze the CP asymmetries in $B^0$ decays in the $\mathrm{SU}(2)_R \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_{B-L}$ left-right symmetric model (LRSM) [5] with spontaneous breakdown of CP [6–8]. Indeed, in such a model with spontaneous parity violation it is natural to consider also CP as a spontaneously broken symmetry. We show that with the present constraints on the parameters of the right-handed sector the new contribution to $B$ meson mixing can be large and time dependent CP asymmetries can vary from -1 to 1 in both $B^0_d$ and $B^0_s$ systems. In addition, due to the new penguins contributing to the flavor changing decay $b \to s\bar{s}s$ the CP asymmetries in $B \to J/\psi K_S$ and $B \to K_S\phi$ which with high accuracy measure the same unitary triangle angle, $\beta$, in the SM may differ from each other almost by unity in the LRSM even in the case in which the measurements in $B \to X_s\gamma$ correspond exactly to the SM predictions. These two effects are complementary, while the former one is dominated by the new heavy particle exchange, the latter one is due to the left-right mixing.

The Higgs sector of the LRSM contains a bidoublet $\Phi(\frac{1}{2}, \frac{1}{2}, 0)$ and two triplets, $\Delta_L(1, 0, 2)$ and $\Delta_R(0, 1, 2)$. In order to have parity as a spontaneously broken symmetry, a discrete left-right symmetry, $\Psi_{iL} \leftrightarrow \Psi_{iR}$, $\Delta_L \leftrightarrow \Delta_R$, $\Phi \leftrightarrow \Phi^\dagger$, should be imposed. After spontaneous symmetry breaking, the vacuum expectation values (vev) of the neutral components of $\Phi$, $k_1$ and $k_2 e^{i\omega}$ give masses to the quarks and left-handed gauge bosons. The phase $\omega$ which is the relative phase between the vev's is the only source of CP-violation in our model. The left- and right-handed Cabbibo-Kobayashi-Maskawa (CKM) matrices $V_L$ and $V_R$, respectively, are related as $|V_L| = |V_R|$, due to the discrete left-right symmetry. They contain all together six CP phases which are related to $\omega$. In the following it would be convenient to think about $V_L$ as the SM CKM matrix and to shift all the phases but one to $V_R$. The charged current Lagrangian in the LRSM is given by $\mathcal{L}_{cc} = g/\sqrt{2}\overline{u}(\cos\xi V_L \gamma^\mu P_L - e^{i\omega}\sin\xi V_R\gamma^\mu P_R)d\, W_{1\mu} + g/\sqrt{2}\overline{u}(e^{-i\omega}\sin\xi V_L\gamma^\mu P_L + \cos\xi V_R\gamma^\mu P_R)d\, W_{2\mu} + \text{H.c.}$, where $P_{L,R} \equiv (1 \mp \gamma_5)/2$, $W_1$, $W_2$ are the charged vector boson fields with the masses $M_1$, $M_2$, respectively, and $\xi$ denotes their mixing. The appearance of $\omega$ in the charged current Lagrangian is pure convention since it can be removed to $V_R$. The most stringent lower bound on $W_2$ mass, $M_2 \gtrsim 1.6$ TeV, is derived from the $K_L$-$K_S$ mass difference [9]. The experimental upper bound on the mixing angle $\xi$ depends on the phase $\omega$. For small phases it is $\xi \lesssim 0.0025$ while for large phases $\xi \lesssim 0.013$ [10]. All these results are subject of large hadronic uncertainties. The best limit on $\xi$, free of these uncertainties, arises from the muon decay data and is $\xi \lesssim 0.033$ [11]. However, for our numerical evaluations we use the appropriate stringent bounds from Ref. [10]. There are two neutral flavor changing Higgs bosons in the model whose masses are constrained as $M_H \gtrsim 12$ TeV [8,12].

CP violation in $B^0$ decays takes place due to the interference between mixing and decay. The corresponding CP asymmetry depends on the parameter $\lambda$ defined as [2]

$$\lambda = \left(\sqrt{\frac{M_{12}^* - \frac{i}{2}\Gamma_{12}^*}{M_{12} - \frac{i}{2}\Gamma_{12}}}\right)\frac{\bar{A}}{A} = e^{-2i\phi_M}\frac{\bar{A}}{A}\,, \tag{1}$$

where $A$ and $\bar{A}$ are the amplitudes of $B^0$ and $\bar{B}^0$ decay to a common CP eigenstate, respectively, and we have used $\Gamma_{12} \ll M_{12}$ to introduce the $B$-$\bar{B}$ mixing phase $\phi_M$. If $|\lambda| = 1$ also $\bar{A}/A = e^{-2i\phi_D}$ is a pure phase and the time dependent CP asymmetry takes a particularly simple form

$$a_{CP}(t) = -Im\lambda\sin(\Delta M t) = \sin 2(\phi_M + \phi_D)\sin(\Delta M t), \tag{2}$$

where $\Delta M$ is the mass difference between the two physical states. From Eq. (1) and Eq. (2) it is clear that any new physics effect in the mixing will translate into $\phi_M \to \phi_M + \delta_M$ and will be universal to all decays while the effect in the decay, $\phi_D \to \phi_D + \delta_D$, will depend on the process. In the SM the mixing is already one-loop effect and therefore new physics contribution to it may be sizeable. Without rigourous arguments some of the recent reviews [2] claim that the LRSM contributions to the $B^0$ mixing are negligible. We show the opposite by performing an explicit calculation.

Let us assume that the off-diagonal element $M_{12}$ of the $B_q$-$\bar{B}_q$ mixing is changed by a factor of $\Delta_q$ as a result of the new contribution from the LRSM, $M_{12} = M_{12}^{LL} + M_{12}^{LR} = M_{12}^{LL}\Delta_q$. Here $LL$ denotes the contribution from the left-handed sector which in our convention is equal to the SM result, and $LR$ denotes the dominant new contribution from the box diagrams with one $W_1$ and one $W_2$ and from the tree level flavor changing Higgs exchange. $M_{12}^{LR}$ including the LO QCD corrections has been calculated in Ref. [8] using the vacuum insertion approximation. With $m_b(m_b) = 4.4$ GeV, $m_t(M_1) = 170$ GeV, $M_B = 5.3$ GeV, $\Lambda_{\overline{MS}}^{(5)} = 225$ MeV, $\sqrt{B_B} f_B = 200$ MeV and the SM input as in Ref. [1] the LO QCD improved result reads [8]

$$\kappa = F(M_2)\left(\frac{1.6\mathrm{TeV}}{M_2}\right)^2 + \left(\frac{12\mathrm{TeV}}{M_H}\right)^2 , \tag{3}$$

where $\kappa = |M_{12}^{LR}|/|M_{12}^{LL}|$ and the function $F(M_2)$ is a complicated function of $W_2$. Numerically $F(1.6\,\mathrm{TeV}) = 0.2$ and $F(10.\,\mathrm{TeV}) = 0.5$. Note that this estimate holds for both $B_d^0$ and $B_s^0$ systems. One can write $\Delta_q = 1 + \kappa e^{i\sigma_q}$, where $\sigma_q = \mathrm{Arg}(M_{12}^{LR}/M_{12}^{LL})$. Consequently the phase $\phi_M^q$ in the mixing in the LRSM becomes $\phi_M^q = \phi_M^{SM,q} + \delta_M^q$ where

$$\delta_M^q = \arctan\left(\frac{\kappa\sin\sigma_q}{1+\kappa\cos\sigma_q}\right) . \tag{4}$$

The phase $e^{i\sigma_q} \simeq -(V_{R,tq}V_{R,tb}^*)/(V_{L,tq}V_{L,tb}^*)$ in our model has been calculated in terms of the quark masses and phase $\omega$ and reads [7] $\sin\sigma_d \simeq \pm k_2/k_1 \sin\omega[2\mu_c/\mu_s(1+s_1^2\mu_s/(2\mu_d))+\mu_t/\mu_b]$, $\sin\sigma_s \simeq \pm k_2/k_1 \sin\omega[\mu_c/\mu_s+\mu_t/\mu_b]$, where $\mu_i = \pm m_i$ and $\pm$ are the signs occuring in the Yukawa sector of the model. While $|k_2/k_1 \sin\omega| \leq m_b/m_t$ [7] there is an enhancement factor $m_t/m_b$ in the expressions for $\sin\sigma_{d,s}$ which thus can be as large as unity. Therefore, taking into account the present constraints on the right-handed particle masses it follows from Eqs (3), (4) that in the LRSM with spontaneous CP violation the phases $\delta_M^q$ can take any value from 0 to $2\pi$ and, consequently, the CP asymmetries in Eq. (2) can vary between -1 and 1.

Unfortunately the CP asymmetries in $B_s^0$ decays which are predicted to be very small in the SM and can easily show up the new physics cannot be studied in $B$ factories running on the $\Upsilon$ peak. $B_d^0$ decays, however, involve large CP asymmetries which are predicted with poor accuracy in the SM. The "benchmark" modes $B \to J/\psi K_S$ and $B \to \pi^+\pi^-$ measure $a_{CP} = \sin 2\beta$ and $a_{CP} = \sin 2\alpha$, respectively, where $\beta$ and $\alpha$ are the angles of the SM unitary triangle. The SM predictions for them are $0.3 \lesssim \sin 2\beta \lesssim 0.9$ and $|\sin\alpha| \leq 1$ [2]. Unless the experimental measurement $\beta_{exp} = \beta + \delta_M$ clearly lays outside the allowed region the new physics cannot be traced off. Moreover, since $\alpha$ gets modified as $\alpha_{exp} = \alpha - \delta_M$ then $\delta_M$ cancels out in $\alpha_{exp} + \beta_{exp}$ [13]. Therefore, finding new physics could rely only on the experimentally very challenging measurement of the third angle $\gamma_{exp}$.

On the other hand, it is known that in the SM the CP asymmetries in the theoretically clean decays $B_d \to J/\psi K_S$ ($b \to c\bar{c}s$) and $B_d \to \phi K_S$ ($b \to s\bar{s}s$) measure with high accuracy the same angle $\beta$. The uncertainty in the SM is estimated to be [3]

$$|\phi(B_d \to J/\psi K_S) - \phi(B_d \to \phi K_S)| \lesssim 0.04\,, \tag{5}$$

where $\phi = \phi_M + \phi_D$. Any deviation from this relation (which should be further tested as proposed in Ref. [14]) will be a clear indication of new physics. The decay $b \to c\bar{c}s$ is dominated by tree level $W_1$ exchange and the new physics contribution to it cannot be sizeable. However, the flavor changing decay $b \to s\bar{s}s$ is one-loop effect in the SM and can, therefore, be modified by new physics.

The flavor changing decay $b \to s\bar{s}s$ is induced by the QCD-, electroweak- and magnetic penguins. The dominant contribution comes from the QCD penguins with top quark in the loop. It is also known [15] that the electroweak penguins decrease about 30% the decay rate and we shall add their contribution to the QCD improved effective Hamiltonian. We start with the effective Hamiltonian due to the gluon exchange describing the decay $b \to s\bar{s}s$ at the scale $M_1$

$$H^0_{eff} = -\frac{G_F}{\sqrt{2}}\frac{\alpha_s}{\pi} V_L^{ts*} V_L^{tb} \left(\bar{s}\left[\Gamma_\mu^{LL} + \Gamma_\mu^{LR}\right] T^a b\right)\left(\bar{s}\gamma^\mu T^a s\right), \tag{6}$$

where $\Gamma_\mu^{LL} = E_0(x_t)\gamma_\mu P_L + 2im_b/q^2 E_0'(x_t)\sigma_{\mu\nu}q^\nu P_R$, $\Gamma_\mu^{LR} = 2im_b/q^2 \tilde{E}_0'(x_t)[A^{tb}\sigma_{\mu\nu}q^\nu P_R + A^{ts*}\sigma_{\mu\nu}q^\nu P_L]$, and the $\Gamma_\mu^{LR}$ term describes the new dominant left-right contribution via the mixing angle $\xi$. Here $A^{tb} = \xi m_t/m_b V_R^{tb}/V_L^{tb} e^{i\omega} \equiv \xi m_t/m_b e^{i\sigma_1}$ and analogously $A^{ts} = \xi m_t/m_b V_R^{ts}/V_L^{ts} e^{i\omega} \equiv \xi m_t/m_b e^{i\sigma_2}$. Note that the phases $\sigma_{1,2}$ are independent and can take any value in the range $(0, 2\pi)$. The functions $E_0(x_t)$, $E_0'(x_t)$ and $\tilde{E}_0'(x_t)$ are Inami-Lim type functions [16] of $x_t = m_t^2/M_1^2$ and are given by $E_0(x_t) = -2/3 \ln x + x(18 - 11x - x^2)/(12(1-x)^3) + x^2(15 - 16x + 4x^2)/(6(1-x)^4) \ln x$, $E_0'(x_t) = x(2 + 5x - x^2)/(8(x-1)^3) - 3x^2/(4(x-1)^4) \ln x$, $\tilde{E}_0'(x_t) = -(4 + x + x^2)/(4(x-1)^2) + 3x/(2(x-1)^3) \ln x$. The left-right analog of $E_0'(x_t)$, $\tilde{E}_0'(x_t)$, is numerically about factor of four larger than the latter one. Together with the $m_t/m_b$ enhancement in $A^{tq}$ this practically overcomes the left-right suppression by small $\xi$.

To obtain reliable estimates for the CP asymmetries in $b \to s\bar{s}s$ induced modes in the LRSM we have to calculate the LO QCD corrections to Eq. (6). Using the operator product expansion to integrate out the heavy fields and calculating the Wilson coefficients $C_i$ in the leading logarithm approximation we run them with the renormalization group equations from the scale of $W_1$ down to the scale $\mu = m_b$ (since the contributions of $W_2, H_{1,2}^0$ are negligible we start immediately from the $W_1$ scale). Because the new physics appears only in the gluonic magnetic operators we can safely take over some well-known results from the SM studies. The effective Hamiltonian we work with is

$$H_{eff} = -\frac{G_F}{\sqrt{2}} V_L^{ts*} V_L^{tb} \left( \sum_{i=1}^{20} C_i(\mu) O_i(\mu) + \sum_{j=7}^{10} C_j^{ew}(\mu) O_j^{ew}(\mu) \right),$$

where we have explicitly separated the electroweak penguin operators (the second term) which to a good approximation will not receive any new contribution in the LRSM from the twenty operators which do mix with the gluonic and photonic magnetic operators. Due to the left-right symmetry the twenty operators split into two groups, $O_1$-$O_{10}$ and $O_1'$-$O_{10}'$, which can be obtained by $P_L \leftrightarrow P_R$ from each other. For the QCD penguin operators $O_1$-$O_6$, magnetic penguin operators $O_{7,8}$ as well as for the electroweak penguin operators $O_7^{ew}$-$O_{10}^{ew}$ we use the standard set of the operators from Ref. [1]. The new left-right operators $O_{9,10}$

are [18] $O_9 = 4(m_b/m_c)(\bar{s}_\alpha\gamma^\mu P_L c_\beta)(\bar{c}_\beta\gamma_\mu P_R b_\alpha)$ and $O_{10} = 4(m_b/m_c)(\bar{s}_\alpha\gamma^\mu P_L c_\alpha)(\bar{c}_\beta\gamma_\mu P_R b_\beta)$. Keeping only the top and bottom quark masses to be non-vanishing, the matching conditions at $W_1$ scale are given as $C_2(M_1) = 1$, $C_7(M_1) = D_0'(x_t) + A^{tb}\tilde{D}_0'(x_t)$, $C_7'(M_1) = A^{ts*}\tilde{D}_0'(x_t)$, $C_8(M_1) = E_0'(x_t) + A^{tb}\tilde{E}_0'(x_t)$, $C_8'(M_1) = A^{ts*}\tilde{E}_0'(x_t)$ and the rest of the coefficients vanish. Here the SM function $D_0'(x_t)$ and its left-right analog $\tilde{D}_0'(x_t)$ are given by $D_0'(x_t) = x(7 - 5x - 8x^2)/(24(x-1)^3) - x^2(2-3x)/(4(x-1)^4)\ln x$, $\tilde{D}_0'(x_t) = (-20 + 31x - 5x^2)/(12(x-1)^2) + x(2-3x)/(2(x-1)^3)\ln x$ .

The $20\times 20$ anomalous dimension matrix decomposes into two identical $10\times 10$ submatrices. The SM $8\times 8$ submatrix of the latter one can be found in Ref. [17] and the rest of the entries have been calculated by Cho and Misiak in Ref. [18]. In the leading logarithm approximation the low energy Wilson coefficients for five flavors are given by $C_i(\mu = m_b) = \sum_{k,l}(S^{-1})_{ik}(\eta^{3\lambda_k/46})S_{kl}C_l(M_1)$, where the $\lambda_k$'s in the exponent of $\eta = \alpha_s(M_1)/\alpha_s(m_b)$ are the eigenvalues of the anomalous dimension matrix over $g^2/16\pi^2$ and the matrix $S$ contains the corresponding eigenvectors. We find

$$C_8(m_b) = \eta^{\frac{14}{23}}(E_0'(x_t) + A^{tb}\tilde{E}_0'(x_t)) + \sum_{i=1}^{5} h_i\eta^{p_i}\,, \tag{7}$$

$$C_8'(m_b) = \eta^{\frac{14}{23}}A^{ts*}\tilde{E}_0'(x_t)\,, \tag{8}$$

where $h_i$ = (0.8623, -0.9135, 0.0209, 0.0873, -0.0571) and $p_i$ = (14/23, 0.4086, 0.1456, -0.4230, -0.8994). We reproduced $C_7$ and $C_7'$ exactly as in Ref. [18] and $C_3$-$C_6$ numerically within 1% as in Ref. [1] and we shall not present them here.

Denoting $\langle O\rangle \equiv \langle K_S\phi|O|B\rangle$ the decay amplitude of $B\to K_S\phi$ can be written

$$\langle H_{eff}\rangle = -\frac{G_F}{\sqrt{2}}V_L^{tb}V_L^{ts*}\sum_{3-6,8,8',ew} C_i(\mu)\langle O_i(\mu)\rangle\,. \tag{9}$$

Contributions from $\langle O_7\rangle$ are suppressed by a factor of $\sqrt{3\alpha_s/\alpha}\approx 9.7$ if compared with $\langle O_8\rangle$ and therefore negligible. The hadronic matrix elements $\langle O_8\rangle$ and $\langle O_8'\rangle$ can be approximated to be of the form

$$\langle O_8\rangle = -\frac{2\alpha_s}{\pi}\frac{m_b}{q^2}\langle \bar{s}i\sigma_{\mu\nu}q^\mu P_R T^a b\bar{s}\gamma^\nu T^a s\rangle\,, \tag{10}$$

and similarly for $\langle O_8' \rangle$, where the timelike gluon has produced $\bar{s}s$. Using factorization and the following parametrization for the hadronic matrix elements [19] $\langle\phi|\bar{s}\gamma_\mu s|0\rangle = f_\phi M_\phi \epsilon_\mu$, $\langle K|\bar{s}\gamma^\mu b|B\rangle = (q_+ - \Delta q_-)F_{BK}(q_-^2;1^-) + \Delta q_- F'_{BK}(q_-^2;0^+)$, where $f_\phi$=0.23 GeV, $q_\pm = q_B \pm q_K$ and $\Delta = (M_B^2 - M_K^2)/q_-^2$ one gets [15] $\langle O_3\rangle = \langle O_4\rangle = 4a/3$, $\langle O_5\rangle = a$, $\langle O_6\rangle = a/3$, $\langle O_7^{ew}\rangle = -a/2$, $\langle O_8^{ew}\rangle = -a/6$ and $\langle O_9^{ew}\rangle = \langle O_{10}^{ew}\rangle = -2a/3$, where $a = f_\phi M_\phi F_{BK}(M_\phi^2;1^-) q_+ \cdot \epsilon_\phi$. In the parametrization of [19] $F_{BK}(M_\phi^2;l) = F_{BK}(0)/(1 - M_\phi^2/M_{BK}^2(l))$, where $F_{BK}(0)$=0.38, $M_{BK}(0^+)$=5.8 GeV and $M_{BK}(1^-)$=5.4 GeV. The element $\langle O_8\rangle$ decomposes to $\langle O_8\rangle = -i4\alpha_s m_b/(9\pi q^2)\langle \bar{s}(\gamma^\mu q_s^\nu - \gamma^\nu q_s^\mu)P_L s \bar{s}\sigma_{\mu\nu} b - i2m_b \bar{s}\gamma^\mu P_L s \bar{s}\gamma_\mu P_L b\rangle$, where $q_s^\mu \approx q_\phi^\mu/2$. With factorization the new matrix element appearing is $\langle K|\bar{s}\sigma^{\mu\nu}b|B\rangle = i f_\sigma (q_+^\mu q_-^\nu - q_+^\nu q_-^\mu)$, where $f_\sigma = [(1+\Delta)F_{BK}(M_\phi^2;1^-) - \Delta F_{BK}(M_\phi^2;0^+)]/(4m_b) \approx F_{BK}(M_\phi^2;1^-)/(4m_b)$ is obtained using the heavy quark effective theory [20]. As a result we get $\langle O_8\rangle = -2\alpha_s/(9\pi) m_b^2/q^2 a[1 - M_\phi^2/(4m_b^2)]$, where the second term in brackets is negligibly small. The same result is valid also for $\langle O_8'\rangle$.

It has been shown that the" physical" range of $q^2$ in $B \to K_S\phi$ is $1/4 \lesssim q^2/m_b^2 \lesssim 1/2$ [21]. To be conservative we use $q^2 = m_b^2/2$, $\xi = 0.01$ and $m_b(M_1) = 2.8$ GeV [22] to estimate the possible effects of new physics. Numerically we obtain for the LO QCD improved amplitude $A \equiv \langle H_{eff}\rangle$

$$A = -\frac{G_F}{\sqrt{2}} V_L^{tb} V_L^{ts*}[-0.0154 + 0.0047(e^{i\sigma_1} + e^{-i\sigma_2})]a\,. \tag{11}$$

It is important to notice that in $B \to \phi K_S$ *both* phases $\sigma_{1,2}$ contribute to the CP asymmetry because only the hadronic matrix elements of the vector currents matter. This should be compared with $B \to X_s\gamma$ in which CP asymmetry is given only by the right-projected operators [23] and, consequently, the phase $\sigma_2$ does not contribute. Also, the major source of uncertainty in the decay rate, the hadronic matrix element $a$, cancels out in the CP asymmetry. The maximum effect is obtained if $\sigma_1 = -\sigma_2 = \pi/2 + \delta_D$. We get $(\bar{A}/A)_{max} = e^{1.3i}$ which implies $|\delta_D|_{max} = 0.65$. This result should be compared with Eq. (5) which implies that there could be a clear effect of the new physics. The maximum allowed difference of the CP asymmetries in $B \to J/\psi K_S$ and $B \to \phi K_S$ in the LRSM could thus be as large as

$|a_{CP}(B \to J/\psi K_S) - a_{CP}(B \to \phi K_S)|_{max} = 1$. If the difference of the phases in these two processes will be measured within 10% and if no difference will be seen then a new upper bound, $\xi \lesssim 0.002$, can be put on the left-right mixing angle for *large* phases which is stronger than the present limit for small phases $\xi \lesssim 0.0025$.

Note that the new effect in Eq. (11) is due to the LR contribution to the QCD magnetic penguins. This new contribution can also provide an answer to the enhancement of $b \to sg^*$ observed by CLEO [24].

Finally, let us consider the constraints on the LRSM coming from the decay $B \to X_s\gamma$. It is possible that due to the cancellation between the $LL$ and $LR$ contributions both the total rate $\Gamma$ and the CP asymmetry in this process can, within errors, correspond to the SM predictions [23]. If the SM predictions will be confirmed experimentally (the CP asymmetry in the SM is expected to be very small) this will constrain the phase $\sigma_1$ and the size of the $LR$ contribution to $B \to X_s\gamma$ but cannot probe the phase $\sigma_2$ which will still be a free parameter. Assuming $\Gamma_{LRSM}(B \to X_s\gamma)/\Gamma_{SM}(B \to X_s\gamma) = 1$ we obtain in the most conservative case for the decay $B \to \phi K_S$ in the LRSM that $(\bar{A}/A)_{max} = e^{0.47i}$ which means $|\delta_D|_{max} = 0.24$ and $|a_{CP}(B \to J/\psi K_S) - a_{CP}(B \to \phi K_S)|_{max} = 0.45$. Therefore, large observable effects are possible independently of the results in $B \to X_s\gamma$.

In conclusion, we show that the LRSM with spontaneous violation of CP can dramatically affect the time dependent CP asymmetries in $B^0_{d,s}$ decays. Due to the new contribution to the $B^0$-$\bar{B}^0$ mixing the CP asymmetries can vary from -1 to 1 in both $B^0_d$ and $B^0_s$ decays. Moreover, the $B$-meson decay amplitudes can receive significant new contributions as well. Most importantly for discovering the new physics in the $B$ factories, the CP asymmetries in $B \to J/\psi K_S$ and $B \to \phi K_S$ which are equal with high accuracy in the SM can differ from each other as much as unity in our model independently of the results in $B \to X_s\gamma$. Interestingly, the excess of $b \to sg^*$ observed by CLEO can also be explained by the LRSM.

We thank Y. Grossman for pointing out the interesting CP violation effects in the decay amplitudes and A. Pich and L. Silvestrini for discussions on the QCD corrections.

## REFERENCES

[1] A.J. Buras and R. Fleischer, TUM-HEP-275/97, hep-ph/9704376.

[2] Y. Nir and H.R. Quinn, Ann. Rev. Nucl. Part. Sci. **42**, 211 (1992); M. Gronau and D. London, Phys. Rev. **D 55**, 2845 (1997); Y. Grossman, Y. Nir and R. Rattazi, SLAC-PUB-7379, hep-ph/9701231; R. Fleischer, CERN-TH/97-241, hep-ph/9709291; Y. Nir, WIS-97/28/Sep-PH, hep-ph/9709301.

[3] Y. Grossman and M. Worah, Phys. Lett. **B 395**, 241 (1997); M. Worah, SLAC-PUB-7700, hep-ph/9711265.

[4] N.G. Deshpande, B. Dutta and S. Oh, Phys. Rev. Lett. **77**, 4499 (1996); M. Ciuchini et al., Phys. Rev. Lett. 79, 978 (1997); D. London and A. Soni, Phys. Lett. **B 407**, 61 (1997); A. Abd El-Hady and G. Valencia, Phys. Lett. **B 414**, 173 (1997); S.A. Abel, W.N. Cottingham and I.B. Whittingham, hep-ph/9803401.

[5] J.C. Pati and A. Salam, Phys. Rev. **D 10**, 275 (1975); R.N. Mohapatra and J.C. Pati, Phys. Rev. **D 11**, 566 and 2558 (1975); G. Senjanović and R.N. Mohapatra, Phys. Rev. **D 12**, 1502 (1975); R.N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 44, 912 (1980) and Phys. Rev. **D 23**, 165 (1981).

[6] T.D. Lee, Phys. Rev. **D 8**, 1226 (1973); G. Senjanović, Nucl. Phys. **B153**, 334 (1979); G. Senjanović and P. Senjanović, Phys. Rev. **D 21**, 3253 (1980).

[7] D. Chang, Nucl. Phys. **B214**, 435 (1983); H. Harari and M. Leurer, Nucl. Phys. **B233**, 221 (1984); G. Ecker and W. Grimus, Nucl. Phys. **B258**, 328 (1985); M. Leurer, Nucl. Phys. **B266**, 147 (1986); J.-M. Frère et al., Phys. Rev. **D 46**, 337 (1992); G. Barenboim, J. Bernabéu and M. Raidal, Nucl. Phys. **B478**, 527 (1996).

[8] G. Ecker and W. Grimus, Z. Phys. **C 30**, 293 (1986); B. Barenboim, J. Bernabéu and M. Raidal, Nucl. Phys. **B511**, 577 (1998).

[9] G. Beall, M. Bander and A. Soni, Phys. Rev. Lett. **48**, 848 (1982).

[10] P. Langacker and S. Uma Sankar, Phys. Rev. **D 40**, 1569 (1989).

[11] G. Barenboim, J. Bernabéu, J. Prades and M. Raidal, Phys. Rev. **D 55**, 4213 (1997).

[12] A very strong bound, $M_H \gtrsim 50$ TeV, has been derived in the manifestly LRSM with hard CP violation in, M.E. Pospelov, Phys. Rev. **D 56**, 259 (1997). In the case of spontaneous CP breaking this bound is decreased by small $|k_2/k_1 \sin\omega| \leq m_b/m_t$ and becomes $M_H \gtrsim 6.4$ TeV.

[13] Y. Nir and D. Silverman, Nucl. Phys. **B345**, 301 (1990).

[14] Y. Grossman, G. Isidor and M. Worah, SLAC-PUB-7614, hep-ph/9708305.

[15] R. Fleischer, Z. Phys. **C 62**, 81 (1994).

[16] T. Inami and C. S. Lim, Prog. Theor. Phys. **65**, 297 (1981).

[17] M. Ciuchini et al., Phys. Lett. **B 316**, 127 (1993); Nucl. Phys. **B415**, 403 (1994).

[18] P. Cho and M. Misiak, Phys. Rev. **D 49**, 5894 (1994); K. Fujikawa and A. Yamada, Phys. Rev. **D 49**, 5890 (1994).

[19] M. Bauer, B. Stech and M. Wirbel, Z. Phys. **C 29**, 637 (1985); Z. Phys. **C 34**, 103 (1987).

[20] N. Isgur and M.B. Wise, Phys. Rev. **D 42**, 2388 (1990).

[21] J.-M. Gerard and W.-S. Hou, Phys. Rev. **D 43**, 2909 (1991); Phys. Lett. **B 253**, 478 (1991).

[22] G. Rodrigo, A. Santamaria and M. Bilenky, Phys. Rev. Lett. **79**, 193 (1997).

[23] D. Atwood, M. Gronau and A. Soni, Phys. Rev. Lett. **79**, 185 (1997); A.L. Kagan and M. Neubert, CERN-TH/98-1, hep-ph/9803368.

[24] D.H. Miller, CLEO Collaboration, talk in EPS-HEP97, Jerusalem, 1997; for theory overview see e.g., D. Atwood and A. Soni, hep-ph/9712243.