Search for new physics in $\Delta S=2$ two-body (VV, PP, VP) decays of the B^- meson

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The $\Delta S=2$ $b\to\bar{d}ss$ transition proceeds via the box diagram in the standard model with a branching ratio calculated to be below 10^{-11}, thus providing an appropriate testing ground for physics beyond the standard model. We analyze the $\Delta S=2$ two-body $B^-\to K^{*-}\bar{K}^{*0}$, $B^-\to K^-\bar{K}^0$, $B^-\to K^{*-}\bar{K}^0$, and $B^-\to K^-\bar{K}^{*0}$ exclusive decays which are driven by the $b\to\bar{d}ss$ transition, both in the standard model and in several extensions of it. The models considered are the minimal supersymmetric models with and without $\mathcal{R}$-parity conservation and two Higgs doublet models. All four modes are found to have a branching ratio of the order of 10^{-13} in the standard model, while the expected branching ratio in the different extensions vary between 10^{-9} and 10^{-6}.

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The intensive search for physics beyond the standard model (SM) is performed nowdays in various areas of particle physics. Among these, rare B meson decays are suggested to give good opportunities for discovering new physics beyond the SM [1]. Recently, it has been suggested [2–4] to investigate effects of new physics possibly arising from $b\to ss\bar{d}$ or $b\to dd\bar{s}$ decays. As shown in Ref. [2], the $b\to ss\bar{d}$ transition is mediated in the standard model by the box diagram and its calculation results in a branching ratio of nearly 10^{-11}, the exact value depending on the relative unknown phase between t, c contributions in the box. The $b\to dd\bar{s}$ branching ratio is even smaller by a factor of 10^2, due to the relative $|V_{td}/V_{ts}|$ factor in the amplitudes. In Ref. [5] different scenarios were used in the analysis of the $b\to dd\bar{s}$ decay, which might be important in $B^{\pm}\to K\pi$ decays. The authors of Refs. [2,3] have calculated the $b\to ss\bar{d}$ transition in various extensions of the SM. It appears that for certain plausible values of the parameters, this decay may proceed with a branching ratio of 10^{-8} and 10^{-7} in the minimal supersymmetric standard model (MSSM) and in two Higgs doublet models [3].

Thus, decays related to the $b\to ss\bar{d}$ transition, which was calculated to be very rare in the standard model, provide a good opportunity for investigating beyond the standard model physics. In Ref. [2] it was suggested that the most suitable channels to see effects of the $b\to ss\bar{d}$ transition are the $B^-\to K^-K^-\pi^+$ or $\bar{B}^0\to K^-K^-\pi^+\pi^+$ decays. Moreover, when one considers supersymmetric models with $\mathcal{R}$-parity violating couplings, it turned out that the existing bounds on the involved couplings of the superpotential did not provide any constraint on the $b\to ss\bar{d}$ mode [2]. Recently, the OPAL Collaboration [6] has set bounds on these couplings from the establishment of an upper limit for the $B^-\to K^-K^-\pi^+$ decay $BR(B^-\to K^-K^-\pi^+)\leqslant 1.3\times 10^{-4}$. The long distance effects in $B^-\to K^-K^-\pi^+$ decay [4] have also been estimated recently and they have been found to be of the order 10^{-12}, comparable in size with the short-distance SM contribution, thus leaving this decay ''free'' for the search of new physics. Although it appears that $B^-\to K^-K^-\pi^+$ or $\bar{B}^0\to K^-K^-\pi^+\pi^+$ are very good candidates to search for the $\Delta S=2$ transitions, we investigate here another possibilty for the observation of the $b\to ss\bar{d}$ transition: the two body decays of B^-.

We consider the VV, VP, and PP states. Although in principle two body decays would appear to be simpler to analyze, there is the complication of $K^0-\bar{K}^0$ mixing. Hence one also needs a good estimate for the $b\to s\bar{s}d$ transitions as well. Nevertheless, not all the two-body states involve neutral K's and we shall return to this point in our summary. First, we proceed to describe the framework used in our analysis in which we concentrate on MSSM, with and without $\mathcal{R}$-parity and two Higgs doublet models as possible alternatives to the SM.

The minimal supersymmetric extension of the standard model leads to the following effective Hamiltonian describing the $b\to ss\bar{d}$ transition [2,7]

$$\mathcal{H}=\tilde{C}_{MSSM}(\bar{s}\gamma^{\mu}d_L)(\bar{s}\gamma_{\mu}b_L), \tag{1}$$

where we have denoted

$$\tilde{C}_{MSSM}=-\frac{\alpha_s^2\delta_{12}^{d*}\delta_{23}^{d}}{216m_{\tilde{d}}^2}[24xf_6(x)+66\tilde{f}_6(x)] \tag{2}$$

with $x=m_{\tilde{g}}^2/m_{\tilde{d}}^2$, and the functions $f_6(x)$ and $\tilde{f}_6(x)$ are given in [7]. The couplings δ_{ij}^d parametrize the mixing between the down-type left-handed squarks. At the scale of b quark mass and by taking the existing upper limits on δ_{ij}^d from [7] and [2] the coupling $\tilde{C}_{MSSM}$ is estimated to be $|\tilde{C}_{MSSM}|\leqslant 1.2\times 10^{-9}$ GeV^{-2} for an average squark mass $m_{\tilde{d}}=500$ GeV and $x=8$, which leads to an inclusive branching ratio for $b\to ss\bar{d}$ of 2×10^{-7} [2]. The corresponding factor calculated in SM [2] is found to be

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$$C_{SM}=\frac{1}{2}\left\{\frac{G_F^2}{2\pi^2}m_W^2V_{tb}V_{ts}^*\left[V_{td}V_{ts}^*f\left(\frac{m_W^2}{m_t^2}\right)+V_{cd}V_{cs}^*\frac{m_c^2}{m_W^2}g\left(\frac{m_W^2}{m_t^2},\frac{m_c^2}{m_W^2}\right)\right]\right\} \tag{3}$$

with $f(x)$ and $g(x,y)$ given in [2]. Taking numerical values from [8], neglecting the CKM phases, one estimates $|C_{SM}| \simeq 4\times10^{-12}$ GeV^{-2}.

The authors of [2] have also investigated beyond MSSM cases by including R-parity violating interactions. The part of the superpotential which is relevant here is $W=\lambda'_{ijk}L_iQ_jd_k$, where i,j, and k are indices for the families and L,Q, and d are superfields for the lepton doublet, the quark doublet, and the down-type quark singlet, respectively. Using the notation of [9] and [2], the tree level Hamiltonian has the form

$$\mathcal{H}=-\sum_n \frac{f_{QCD}}{m^2_{\tilde{\nu}_n}}[\lambda'_{n32}\lambda'^*_{n21}(\overline{s_R}b_L)(\overline{s_L}d_R)+\lambda'_{n21}\lambda'^*_{n32}(\overline{s_R}d_L)(\overline{s_L}b_R)]. \tag{4}$$

The QCD corrections were found to be important for this transition [10]. For our purpose it suffices to follow [2] retaining the leading order QCD result $f_{QCD}\simeq 2$, for $m_{\tilde{\nu}}=100$ GeV.

Most recently an upper bound on the specific combination of couplings entering (4) has been obtained by OPAL from a search for the $B^-\to K^-K^-\pi^+$ decay [6]

$$\sum_n \sqrt{|\lambda'_{n32}\lambda'^*_{n21}|^2+|\lambda'_{n21}\lambda'^*_{n32}|^2}<10^{-4}. \tag{5}$$

Here we take the order of magnitude, while the OPAL result is 5.9×10^{-4} based on a rough estimate $\Gamma(B^-\to K^-K^-\pi^+)\simeq 1/4\ \Gamma(b\to ss\bar{d})$.

The decay $b\to ss\bar{d}$ has been investigated using two Higgs doublet models (THDM) as well [3]. These authors found that the charged Higgs box contribution in MSSM is negligible. On the other hand, THDM involving several neutral Higgs bosons [11] could have a more sizable contribution to these modes. The part of the effective Hamiltonian relevant in our case is the tree diagram exchanging the neutral Higgs bosons h (scalar) and A (pseudoscalar)

$$\mathcal{H}_{TH}=\frac{i}{2}\xi_{sb}\xi_{sd}\left[\frac{1}{m_h^2}(\bar{s}d)(\bar{s}b)-\frac{1}{m_A^2}(\bar{s}\gamma_5d)(\bar{s}\gamma_5b)\right], \tag{6}$$

with the coupling ξ_{ij} defined in [11] as a Yukawa coupling of the flavor changing neutral current (FCNC) transitions $d_i\leftrightarrow d_j$. In our estimation we use the bound $|\xi_{sb}\xi_{sd}|/m_H^2 >10^{-10}$ GeV^{-2}, $H=h,A$, which was obtained in [3] by using the Δm_K limit on ξ_{bd}/m_H and assuming $|\xi_{sb}/m_H| >10^{-3}$.

We proceed now to study the effect of Hamiltonians (1), (4), (6) on the various two body $\Delta S=2$ decays of charged B mesons. In order to calculate the matrix elements of the operators appearing in the effective Hamiltonian, we use the factorization approximation [12–14], which requires the knowledge of the matrix elements of the current operators or the density operators. Here we use the standard form factor representation [13,12] of the following matrix elements:

$$\langle P'(p')|\bar{q}_j\gamma^\mu q_i|P(p)\rangle = F_1(q^2)\left(p^\mu+p'^\mu-\frac{m_P^2-m_{P'}^2}{q^2}(p^\mu-p'^\mu)\right)+F_0(q^2)\frac{m_P^2-m_{P'}^2}{q^2}(p^\mu-p'^\mu), \tag{7}$$

where F_1 and F_0 contain the contribution of vector and scalar states respectively and $q^2=(p-p')^2$. Also, $F_1(0)=F_0(0)$ [13]. For these form factors, one usually assumes pole dominance [13,15]

$$F_1(q^2)=\frac{F_1(0)}{1-q^2/m_V^2};\quad F_0(q^2)=\frac{F_0(0)}{1-q^2/m_S^2} \tag{8}$$

and in order to simplify, we shall take $m_V=m_S$. The matrix element between pseudoscalar and vector meson is usually decomposed [14] as

$$\begin{aligned}\langle V(q,\epsilon_V)|\bar{q}_j\gamma^\mu(1-\gamma_5)q_i|P(p)\rangle &=\frac{2V(Q^2)}{m_P+m_V}\epsilon^{\mu\nu\alpha\beta}\epsilon^*_{V\nu}p_\alpha q_\beta\\ &+i\epsilon_V^*\cdot Q\frac{2m_V}{Q^2}Q^\mu[A_3(Q^2)-A_0(Q^2)]+i(m_P+m_V)\\ &\times\left[\epsilon_V^{\mu*}A_1(Q^2)-\frac{\epsilon_V^*\cdot Q}{(m_P+m_V)^2}(p+q)^\mu A_2(Q^2)\right],\end{aligned} \tag{9}$$

where $Q=p-q$,

$$A_3(Q^2)-\frac{m_H+m_V}{2m_V}A_1(Q^2)+\frac{m_H-m_V}{2m_V}A_2(Q^2)=0, \tag{10}$$

and $A_3(0)=A_0(0)$. For the vector and axial vector form factor we use again pole dominance [13,15], and relevant parameters are taken from [12,14] $F_0^{BK}(0)=0.38$, $A_0^{BK*}(0)=0.32$. For the calculations of the density operators we use the relations

$$\partial^\alpha(\bar{s}\gamma_\alpha b)=i(m_b-m_s)\bar{s}b, \tag{11}$$

$$\partial^\alpha(\bar{s}\gamma_\alpha\gamma_5 b)=i(m_b+m_s)\bar{s}\gamma_5 b. \tag{12}$$

We will use also the following decay constants:

$$\langle V(\epsilon_V,q)|\bar{q}_j\gamma^\mu q_i|0\rangle=\epsilon_\mu^*(q)g_V(q^2), \tag{13}$$

$$\langle P(q)|\bar{q}_j\gamma^\mu\gamma_5 q_i|0\rangle=if_Pq_\mu \tag{14}$$

with $f_K = 0.162$ GeV, $g_{K*} = 0.196$ GeV^2 [14]. Now we turn to the analysis of the specific modes.

(a) $B^- \to K^{*-}\bar{K}^{*0}$ decay. For the analysis of pseudoscalar meson decay to two vector mesons it is convenient to use helicity formalism (see, e.g., [16]). We denote $\mathcal{O} = (\bar{s}\gamma^\mu(1-\gamma_5)d)\ (\bar{s}\gamma_\mu(1-\gamma_5)b)$, and then we use $\mathcal{H} = C\mathcal{O}$ with C being $1/4\tilde{C}_{MSSM}$, $1/4C_{SM}$. Using factorization and the definitions given above, one finds the following helicity amplitudes:

$$H_{00}(B^- \to K^{*-}\bar{K}^{*0}) = C g_{K*}(m_B + m_{K*}) \times [\alpha A_1^{BK*}(m_{K*}^2) - \beta A_2(m_{K*}^2)], \tag{15}$$

$$H_{\pm\pm}(B^- \to K^{*-}\bar{K}^{*0}) = C g_{K*}(m_B + m_{K*}) \times [\alpha A_1^{BK*}(m_{K*}^2) \mp \gamma V^{BK*}(m_{K*}^2)] \tag{16}$$

where $\alpha = (1-2r^2)/2r^2$, $\beta = k^2/[2r^2(1+r)^2]$, $\gamma = (1-4r^2)$ with $r = m_{K*}/m_B$, $k^2 = 1 + r^4 + t^4 - 2r^2 - 2t^2 - 2r^2t^2$. The decay width is then

$$\Gamma(B^- \to K^{*-}\bar{K}^{*0}) = \frac{|\vec{p}|}{8\pi m_B^2}[|H_{00}|^2 + |H_{++}|^2 + |H_{--}|^2].$$

Within MSSM model the branching ratio becomes $\leqslant 6.2\times 10^{-9}$, while SM gives this rate to be 6.8×10^{-14}. The $\mathcal{R}$-parity term described by the effective Hamiltonian (4) cannot be seen in this decay mode when factorization approach is used, since the density operator matrix element $\langle \bar{K}^{*0}|(\bar{s}d)|0\rangle$ vanishes. The two Higgs doublet model also cannot be tested in this mode due to the same reason.

(b) $B^- \to K^{*-}\bar{K}^0$ decay. The matrix element of the operator $\mathcal{O}$ is calculated to be $\langle \bar{K}^0(k_0)K^{*-}(k_-,\epsilon)|\mathcal{O}|B^-(p_B)\rangle = -2m_{K*}f_K A_0^{BK*}(m_{K*}^2)\epsilon^*\cdot k_0$. Denoting the decay amplitude by $\mathcal{A}$, one finds $\Sigma_{pol}|\mathcal{A}|^2 = |C|^2 f_K^2 |A_0(m_K^2)|^2 \lambda(m_B^2, m_K^2, m_{K*}^2)$ with the $\lambda(a,b,c) = a^2 + b^2 + c^2 - 2(ab + bc + ac)$. The branching ratio is straightforwardly found to be $BR(B^- \to K^{*-}\bar{K}^0)_{MSSM} \leqslant 1.6\times 10^{-9}$, which is comparable to the SM prediction of Ref. [12] for the $\Delta S = 0$ $B^- \to K^{*-}K^0$ decay, given as $BR(B^- \to K^{*-}K^0) = 1\times 10^{-9}$, 5×10^{-9}, 2×10^{-9} obtained for the number of colors $N_c = 2$, $N_c = 3$, $N_c = \infty$, respectively.

The SM calculation for the $\Delta S = 2$ transition leads to $BR(B^- \to K^{*-}\bar{K}^0)_{SM} = 1.7\times 10^{-14}$. The MSSM which includes $\mathcal{R}$ parity breaking terms can occur in this decay. The matrix element of the operator $\mathcal{O}_{\mathcal{R}} = (\bar{s}(1+\gamma_5)d)(\bar{s}(1-\gamma_5)b)$ can be found to be

$$\langle \bar{K}^0(k_0)K^{*-}(k_-,\epsilon)|\mathcal{O}_{\mathcal{R}}|B^-(p_B)\rangle = \frac{m_K^2 f_K}{(m_s+m_d)(m_s+m_b)}(2m_{K*}\epsilon^*\cdot k_0)A_0^{BK*}(m_K^2).$$

Taking the values of the quark masses as in [12] $m_b = 4.88$ GeV, $m_s = 122$ MeV, $m_d = 7.6$ MeV and using the bound given in Eq. (5) we obtain the estimation of the upper limit of the branching ratio $BR(B^- \to K^{*-}\bar{K}^0)_{\mathcal{R}}$ to be 4.4×10^{-8}. This limit can be raised to 1.5×10^{-6} for the upper bound on the couplings of 5.9×10^{-4} given in [6].

The two Higgs doublet model (6) gives for the amplitude of this decay

$$\mathcal{A}_{THDM}[B^-(p_B) \to \bar{K}^0(k_0)K^{*-}(k_-,\epsilon)] = \frac{i}{2}\frac{\xi_{sb}\xi_{sd}}{m_A^2}[2m_{K*}f_K A_0^{BK*}(m_K^2)\epsilon^*\cdot k_0] \times \frac{m_K^2 f_K}{(m_s+m_d)(m_s+m_b)}, \tag{17}$$

which gives for the limit $|\xi_{sb}\xi_{sd}|/m_H^2 > 10^{-10}$ GeV^{-2}, a branching ratio of the order 10^{-11}. Because of the specific combination of the products of the scalar (pseudoscalar) densities this is the only decay which has nonvanishing amplitude within the factorization assumption.

(c) $B^- \to K^-\bar{K}^{*0}$ decay. For this decay mode the matrix element of the operator $\mathcal{O}$ is determined to be

$$\langle \bar{K}^{*0}(k_0,\epsilon)K^-(k_-)|\mathcal{O}|B^-(p_B)\rangle = 2g_{K*}f_K F_1^{BK*}(m_{K*}^2)\epsilon^*\cdot k_- \tag{18}$$

giving the branching ratio in MSSM with an upper limit $BR(B^- \to K^-\bar{K}^{*0})_{MSSM} = 5.9\times 10^{-9}$ in comparison with SM result 6.5×10^{-14}.

The amplitude calculated in the MSSM including $\mathcal{R}$ breaking and the THDM vanishes, due to vanishing of the matrix element of the density operator for $\bar{K}^{*0}$ state.

(d) $B^- \to K^-\bar{K}^0$ decay. The matrix element of the operator $\mathcal{O}$ becomes in this case

$$\langle \bar{K}^0(k_0)K^-(k_-)|\mathcal{O}|B^-(p_B)\rangle = i f_K F_0^{BK}(m_K^2)(m_B^2 - m_K^2).$$

The multiplication with the corresponding $1/4\tilde{C}_{MSSM}$ gives the required amplitude $\tilde{\mathcal{A}}$. The decay width is then

$$\Gamma(B^- \to K^-\bar{K}^0) = (16\pi m_B^2)^{-1}\sqrt{m_B^2 - 4m_K^2}|\tilde{\mathcal{A}}|^2. \tag{19}$$

The branching ratio for MSSM is found to be $BR(B^- \to K^-\bar{K}^0)_{MSSM} \leqslant 2.3\times 10^{-9}$, in comparison with the 2.5×10^{-14} found in the SM. The matrix element of the R-parity breaking MSSM operator $\mathcal{O}^{(1)} = (\bar{s}\gamma_5 d)(\bar{s}b)$ is found to be

$$\langle K^-\bar{K}^0|\mathcal{O}^{(1)}|B^-\rangle = \langle \bar{K}^0|\bar{s}\gamma_5 d|0\rangle\langle K^-|\bar{s}b|B^-\rangle = -i\frac{m_K^2}{(m_s+m_d)(m_b-m_s)}f_K F_0^{BK}(m_K^2)(m_B^2 - m_K^2)$$

while the operator $(\bar{s}\gamma_5 b)(\bar{s}d)$ gives the same result with the opposite sign. The decay width is then

TABLE I. The predicted branching ratio for the $B^- \Delta S=2$ two-body decays calculated using the factorization approach within standard model (the first column), minimal supersymmetric standard model (the second column), minimal supersymmetric standard model extended by $\mathcal{R}$-parity breaking (the third column), and two Higgs doublet model (the fourth column). The values in columns two, three, and four are upper limits, as determined from present knowledge of upper limits for couplings involved.

Decay	SM	MSSM	MSSM+$\mathcal{R}$	THDM
$B^- \to K^{*-}\bar{K}^{*0}$	6.9×10^{-14}	6.2×10^{-9}	–	–
$B^- \to K^{*-}\bar{K}^{0}$	1.7×10^{-14}	1.6×10^{-9}	$10^{-7}-10^{-6}$	10^{-11}
$B^- \to K^{-}\bar{K}^{*0}$	6.6×10^{-14}	5.9×10^{-9}	–	–
$B^- \to K^{-}\bar{K}^{0}$	2.5×10^{-14}	2.3×10^{-9}	$10^{-7}-10^{-6}$	–

$$\Gamma(B^- \to K^- \bar{K}^0)_{\mathcal{R}} = \frac{1}{16\pi m_B^2}\sqrt{m_B^2-4m_K^2}\,|\langle K^-\bar{K}^0|\mathcal{O}^{(1)}|B^-\rangle/4|^2 \times \frac{f_{QCD}^2}{m_{\tilde{\nu}}^4}\left(\sum_{i=n}|\lambda'_{n32}\lambda'^{*}_{n21}|^2+|\lambda'_{n21}\lambda'^{*}_{n32}|^2\right).$$

The constraint in Eq. (5) gives the bound 9.4×10^{-8}, while for the bound of 5.9×10^{-4} for the coupling constants (6) the rate $BR(B^- \to K^- \bar{K}^0)_{\mathcal{R}}$ can reach 3.3×10^{-6}.

The long distance effects are usually suppressed in the B meson decays. One might wonder if they are important in decays we consider here. We have estimated the tree level contribution of the $D(D^*)$ which then goes into $K(K^*)$ via weak annihilation. We found that these contributions give a branching ratio of the order 10^{-18} and therefore they can be safely neglected. One might think that the exchange of two intermediate states $D(D^*)$ and $K(K^*)$ can introduce certain long distance contributions. In decay $B \to$ ''D'' ''K'' $\to$ ''K'' ''K'' the first weak vertex arises from the decay $B \to$ ''D'' ''K'' and the second weak vertex (see e.g. [4]) can be generally obtained from the three body decays of $D \to KKK$. In Ref. [4] it was found that such contributions are also very small. Therefore, we are quite confident to suggest that the long distance effects are not important in the two body $\Delta S=2$ B decays.

Let us turn now to the possibility of detecting these decay modes. The $B^- \to K^{*-}\bar{K}^{*0}$ and $B^- \to K^-\bar{K}^{*0}$ modes have clean signatures of a $\Delta S=2$ transition and therefore these are the channels we recommend to look for. The other two modes we discussed, (b) and (d) have a $\bar{K}^0$ in the final states which complicates the possibilty of detection because of K^0-$\bar{K}^0$ mixing. Separating the desired amplitude requires the measurement of the decays of both K_S and K_L, since one can express [17]

$$\frac{\Gamma(B^- \to K^- K_S)-\Gamma(B^- \to K^- K_L)}{\Gamma(B^- \to K^- K_S)+\Gamma(B^- \to K^- K_L)} = \mathrm{Re}\ \eta(B^- \to KK^-),$$

where $\mathrm{Re}\ \eta(B^- \to KK^-) = A(B \to \bar{K}^0 K^-)/A(B \to K^0 K^-)$.

We summarize our results in Table I. The MSSM gives rates of the order 10^{-9} and 10^{-8}, while the $\mathcal{R}$-parity breaking terms in the MSSM can be seen only in the $B^- \to K^{*-}\bar{K}^0$ and $B^- \to K^-\bar{K}^0$ decay. These are the modes which as we mentioned are more difficult on the experimental side. The THDM model can give nonvanishing contribution only in the case of $B^- \to K^{*-}\bar{K}^0$ decay, with a rate too small to be seen. Thus, we conclude by stressing the possibilty of detecting physics beyond SM mainly in the $K^{*-}\bar{K}^{*0}$, $K^-\bar{K}^{*0}$ decays.

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