Penguin-mediated exclusive hadronic weak B decays

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We estimate a number of exclusive two-body charmless decays of B^+ and B^- mesons. Some of these are mediated predominantly through one-loop gluon exchange, while others have a comparable or larger contribution from the doubly Cabibbo-suppressed tree diagrams. The rates for several decays are in an observable range and should test the standard model.

I. INTRODUCTION

Recently rare B-meson decays have been a subject of both theoretical and experimental interest. Penguin diagrams give rise to charmless B-meson decays.[1] The processes can be mediated by a photon or a gluon.[2] The photon-mediated process such as $B \to K^* \gamma$ (Ref. 3), and $B \to K^* l^+ l^-$ (Ref. 4) have been extensively discussed because of their clean signature. The corresponding processes at the quark level involving gluons, $b \to sg$ and $b \to s q\bar{q}$, have also been considered.[5] However, the experimental signature[6] for such a transition is a charmless exclusive mode[7] such as $B \to K\pi$, etc. The estimates for these processes involve matrix elements of four-quark operators, and these have not received much attention in the past. In this article we propose to study some of these charmless modes that have clean signatures. An added complication here is that charmless hadronic decays can also arise through the tree Hamiltonian with $b \to u$ transition. A careful study of the modes reveals some in which the penguin diagram clearly dominates, while in others the tree contribution can be significant.

The calculation proceeds in two steps. First we obtain the effective short-distance interaction including the one-loop gluon-mediated diagram.[7] We then use the factorization approximation to derive the hadronic matrix elements by saturating with a vacuum state in all possible ways. The resulting matrix elements involve quark bilinears between one meson and a vacuum and between two meson states. These are estimated using relativistic quark-model wave functions. Such a technique has been used extensively by Bauer, Stech, and Wirbel[8] for B and D nonleptonic decays and results are in good agreement with the experiment.

II. EFFECTIVE HAMILTONIAN

We shall first discuss the gluon-mediated penguin contribution. Dictated by gauge invariance, the effective flavor-changing neutral current J_μ contains two terms. The first, proportional to G_1, we call the charge radius, while the second one, proportional to G_2, is called the dipole moment operator:

$$J_\mu = (g_s/4\pi^2)(G_F/\sqrt{2})\bar{s}^{\,i}\tfrac{1}{2}\lambda_{ij} \times [G_1(q^2\gamma_\mu - q_\mu \not{q})L + iG_2\sigma_{\mu\nu}q^\nu(m_s L + m_b R)]bj \,. \tag{1}$$

Here $G_{1,2}$ are functions of Kobayashi-Maskawa angles and quark masses:

$$G_1 = \sum V_k G_1^k(x_k), \quad V_k = U_{bk}U_{ks}^*, \quad k = u, c, t, \ldots, \tag{2}$$

$$G_1^k = (x_k/12)(1/y_k + 13/y_k^2 - 6/y_k^3) + [2/3y_k - (x_k/6)(4/y_k^2 + 5/y_k^3 - 3/y_k^4)]\ln x_k \,, \tag{3}$$

where $x_k = m_k^2/m_W^2$ and $y_k = 1 - x_k$.

Note that when the gluon is on shell (i.e., $q^2=0$), the G_1 contribution vanishes. In the $q^2 \neq 0$ cases both terms participate. For a gluon-exchange diagram (i.e., for the process $b \to s\bar{q}q$) we find that the G_1 contribution dominates over G_2, and we neglect G_2 in what follows. We also evaluate G_1 at $q^2=0$, as in Eq. (3). At larger q^2, G_1 develops a small imaginary part, which is important for discussion of CP violation. In this article we shall restrict our attention only to the branching ratios.

From the current (1) via one-gluon-exchange diagram we find the following effective Hamiltonian (local 4-quark operator):

$$H_{\text{eff}}^P = \kappa(\bar{s}^{\,i}_L\gamma_\mu b^i_L \bar{q}^{\,j}_L\gamma_\mu q^j_L - 3\bar{s}^{\,i}_L\gamma_\mu b^j_L \bar{q}^{\,j}_L\gamma_\mu q^i_L + \bar{s}^{\,i}_L\gamma_\mu b^i_L \bar{q}^{\,j}_R\gamma_\mu q^j_R - 3\bar{s}^{\,i}_L\gamma_\mu b^j_L \bar{q}^{\,j}_R\gamma_\mu q^i_R) \,, \tag{4}$$

$$\kappa = (\alpha_s/6\pi)(G_F/\sqrt{2})G_1, \quad \alpha_s = g_s^2/4\pi \,. \tag{5}$$

Here q runs over all quark species, although only u, d, and s are relevant for our discussion. Charmless decays can also arise from the standard tree level interaction with $b \to u$ transition. This is given by

$$H_{\text{eff}}^T = \eta\bar{s}^{\,i}_L\gamma_\mu u^i_L \bar{u}^{\,j}_L\gamma_\mu b^j_L, \quad \eta = 4U_{ub}U_{us}^* G_F/\sqrt{2} \,. \tag{6}$$

Present limit on $U_{ub}/U_{uc} < 0.2$ gives rise to the constraint $\eta/\kappa < 2.5$, where we have used $\alpha_s(m_b^2) = 0.256$ and $G_1 = -5.91 U_{bc}U_{cs}^*$, which corresponds to $m_t = 80$ GeV.

The effects of the tree-level interaction could be large in general. We find however that there are certain modes of B decays in which the penguin diagram clearly dominate. In others the tree diagram is comparable or quite large. We evaluate the rates arising from these contributions independently and present our results in Table I. We shall define $U_{ub}/U_{bc}=0.1\zeta$, where $-2<\zeta<2$. The modes where we find that the penguin contribution clearly dominates are

$$\begin{aligned} B^+ \to\; & K^0\pi^+, K^{*+}\rho^0 ,\\ & K^{*0}(\pi^+,\rho^+) ,\\ & \phi(K^+,K^{*+}) ,\\ B^0 \to\; & \phi(K^0,K^{*0}) ,\\ & K^{*0}(\pi^0,\rho^0) . \end{aligned} \tag{7}$$

The penguin Hamiltonian is a pure $I=0$ operator. In contrast, the doubly Cabibbo-suppressed tree-level Hamiltonian is an equal mixture of $I=0$ and 1 operators. If U_{ub} turns out to be very small and the penguin diagram dominates all the two-body decays, then we would have $\Delta I=0$ rule. We find, assuming $\Delta I=0$, the following isospin relations between the decay modes:

$$\begin{aligned} \Gamma(B^+\to K^0\pi^+) &= \Gamma(B^0\to K^+\pi^-) \\ &= 2\Gamma(B^+\to K^+\pi^0) = 2\Gamma(B^0\to K^0\pi^0) ,\\ \Gamma(B^+\to K^+\phi) &= \Gamma(B^0\to K^0\phi) ,\\ \Gamma(B^+\to K^0\rho^+) &= \Gamma(B^0\to K^+\rho^-) \\ &= 2\Gamma(B^+\to K^+\rho^0) = 2\Gamma(B^0\to K^0\rho^0) ,\\ \Gamma(B^+\to K^{*0}\pi^+) &= \Gamma(B^0\to K^{*+}\pi^-) \\ &= 2\Gamma(B^+\to K^{*+}\pi^0) = 2\Gamma(B^0\to K^{*0}\pi^0) ,\\ \Gamma(B^+\to K^{*+}\phi) &= \Gamma(B^0\to K^{*0}\phi) ,\\ \Gamma(B^+\to K^{*0}\rho^+) &= \Gamma(B^0\to K^{*+}\rho^-) \\ &= 2\Gamma(B^+\to K^{*+}\rho^0) = 2\Gamma(B^0\to K^{*0}\rho^0) . \end{aligned} \tag{8}$$

Departure from these relations will clearly show the importance of doubly suppressed Cabibbo transitions.

III. FACTORIZATION APPROXIMATION

From experience we know that nonleptonic decays are extremely difficult to handle. For example, the $\Delta I=\frac{1}{2}$ rule in $K\to\pi\pi$ decays has not yet been understood in a satisfactory way. Enormous theoretical machinery has been applied to $K\to\pi\pi$ decays producing only up to 50% agreement with experiment. For energetic decays of heavy mesons (D,B) the situation is somewhat simpler. Here the direct generation of a final meson by a quark current is (probably) a good approximation.

According to the current-field identities, the currents are proportional to interpolating stable or quasistable hadron fields. The approximation now consists in taking, for one current of the current product, only the asymptotic part of the full hadronic field: i.e., its "in" or "out" field. Then the weak amplitude factorizes and is fully determined by the matrix elements of another current between two remaining hadron states. That is the reason why we call this approximation the factorization approximation.

Note that in the replacement of interacting fields by asymptotic fields we neglect only the initial- or final-state interaction of the corresponding particles. In the case of B decays in Eq. (8) this is justified by a simple energy argument that one very heavy object decays into two light but very energetic objects whose interactions might be safely neglected.

Note also that $1/N_c$ expansion argument provides some theoretical justification for the factorization approximation, because factorization follows to leading order in a $1/N_c$ expansion.

Each one of the B decay modes [Eq. (7)] can receive three different contributions from each Hamiltonian. As an example we give the amplitude for $B^+\to K^+\pi^0$ obtained from $H^P_{\rm eff}$:

$$\begin{aligned} A^P[B^+(p)\to K^+(k)\pi^0(k')] &= L(\pi^0)\langle K^+|\bar{s}\gamma_\mu b|B^+\rangle\langle\pi^0|\bar{q}\gamma^\mu\gamma_5 q|0\rangle \\ &\quad + L(K^+)\langle K^+|\bar{s}\gamma_\mu\gamma_5 u|0\rangle[\langle\pi^0|\bar{u}\gamma^\mu b|B^+\rangle + (2m_K^2/m_b m_s)\langle\pi^0|\bar{u}\gamma^\mu b|B^+\rangle] \\ &\quad + L(B^+)\langle K^+\pi^0|\bar{s}\gamma_\mu u|0\rangle\langle 0|\bar{u}\gamma^\mu\gamma_5 b|B^+\rangle(1+2m_B^2/m_b m_s) , \end{aligned} \tag{9}$$

where the coefficients $L(\pi^0)$, $L(K^+)$, and $L(B^+)$ contain the coupling constant κ, color factors, flavor-symmetry factors, i.e., flavor-counting factors, and factors coming from Fierz transforming the operators (4). Note that the mass ratios $(2m_K^2/m_b m_s)$ in Eq. (9) come from Fierz transformation and application of the quark equation of motion on the left-right current $\times$ current parts of the effective weak Hamiltonian (4). The factors, proportional to $L(\pi^0)$ and $L(K^+)$ come from the quark decay diagrams, while the $L(B^+)$ corresponds to so-called annihi-

TABLE I. Branching ratios for charmless two-body decay modes of the B meson. Decay rates from the penguin and tree Hamiltonians are contributed separately. The calculation of Ref. 8 based on the penguin operator and experimental limits are also presented. Some of the zeros stand for extremely small rates.

Branching ratio mode		Ref. 8 (10^{-5})	Contribution from $H^T_{\rm eff}$ (10^{-5})	Contribution from $H^P_{\rm eff}$ (10^{-5})	Present experimental limit at 90% C.L. (10^{-5})
$B^+\to$	$K^+\pi^0$		$0.06\zeta^2$	0.53	
	$K^0\pi^+$		0	1.06	9.0
	$K^+\phi$	1.20	0	1.12	8.0
	$K^+\rho^0$		$0.01\zeta^2$	0	7.0
	$K^0\rho^+$		0	0	
	$K^{*+}\pi^0$		$0.05\zeta^2$	0.29	
	$K^{*0}\pi^+$		0	0.58	13.0
	$K^{*+}\phi$		0	3.12	
	$K^{*+}\rho^0$		0	0.62	
	$K^{*0}\rho^+$		0	1.24	
$B^0\to$	$K^0\pi^0$		0	0.53	
	$K^+\pi^-$		$0.06\zeta^2$	1.06	
	$K^0\phi$		0	1.12	49.0
	$K^0\rho^0$		$0.01\zeta^2$	0	58.0
	$K^+\rho^-$		0	0	
	$K^{*0}\pi^0$		0	0.29	
	$K^{*+}\pi^-$	0.44	$0.10\zeta^2$	0.58	
	$K^{*0}\phi$		0	3.12	44.0
	$K^{*0}\rho^0$		0	0.62	67.0
	$K^{*+}\rho^-$	0.78	$0.01\zeta^2$	1.24	

lation diagrams. These are different for each decay mode as indicated by the dependence on final-state meson. To obtain the amplitude for ($B\to K^*+$meson) decay modes one has to replace L with L^* in Eq. (9). We proceed with the following definitions of the coupling constants and Lorentz decomposition of the typical hadronic matrix elements:

$$\langle K^+(k)|\bar{s}\gamma_\mu\gamma_5 u|0\rangle = ik_\mu f_K,\quad f_K=1.2f_\pi=0.154\ \text{GeV}\ , \tag{10a}$$

$$\langle K^{*+}(k)|\bar{s}\gamma_\mu u|0\rangle = ig_{K^*}\epsilon_\mu(k),\quad g_{K^*}=g_{\rho^+}\ , \tag{10b}$$

$$\langle \rho^0(k')|\bar{u}\gamma_\mu u|0\rangle = ig\rho^0\epsilon_\mu(k'),\quad g\rho^0=\frac{1}{\sqrt{2}}g_{\rho^+}\ , \tag{10c}$$

$$\langle \phi(k')|\bar{s}\gamma_\mu s|0\rangle = ig_\phi\epsilon_\mu(k')\ , \tag{10d}$$

$$\langle K^+(k)|\bar{s}\gamma_\mu b|B^+(p)\rangle = (p+k)_\mu f^{(+)}k'^2) + k'_\mu f^{(-)}(k'^2),\quad k'=p-k\ , \tag{10e}$$

$$\begin{aligned}\langle K^{*+}(k)|\bar{s}\gamma_\mu(1-\gamma_5)b|B^+(p)\rangle &= i\epsilon_{\mu\nu\lambda\sigma}\epsilon^\nu(k)(p+k)^\lambda(p-k)^\sigma V(k'^2)\\ &\quad+\epsilon_\mu(k)(m^2_{K^*}-m^2_B)A_1(k'^2)\\ &\quad-(\epsilon\cdot k')(p+k)_\mu A_2(k'^2)\ ,\end{aligned} \tag{10f}$$

$$\langle K^+(k)\rho^0(k')|\bar{s}\gamma_\mu(1-\gamma_5)u|0\rangle = \langle \rho^0(k')|\bar{s}\gamma_\mu(1-\gamma_5)u|K^-(-k)\rangle\ . \tag{10g}$$

Any other hadronic matrix element needed to evaluate branching ratios of decays (8) can easily be obtained from the definitions (10). We shall use form factors at zero momentum squared obtained by Ref. 8 and assume that momentum dependence of the form factors $f^{(+)}(k'^2)$, $V(k'^2)$, etc., from Eq. (10), is well described by a single pole with mass $\approx m_B$, since masses of excited b-quark meson states $(0^+,1^+,1^-,\ldots)$ are very similar to m_B.

Explicit evaluation of a few modes for an illustrative purpose is in order.

A. $B^+\to K^0\pi^+$

Let us first consider the contribution due to $H^P_{\rm eff}$. For this decay $L(\pi^+)=0$ due to the flavor symmetry of the Hamiltonian (4). Nonvanishing contributions come from the $L(K^0)$ and $L(B^+)$ terms. Since the hadronic matrix elements from the annihilation diagram $L(B^+)$ is proportional to a very small quantity $[(m_K^2-m_\pi^2)f^{(+)}-m_B^2f^{(-)}]$ we shall neglect it in further calculations. The second term in (9) gives the branching ratio

$$B^P(B^+\to K^0\pi^+)=\alpha_1[(f_K/m_B)f^{(+)}_{\pi^+B^+}(m_K^2)]^2 \times\lambda^{1/2}_{\pi K}(1+2m_K^2/m_bm_s)^2\ , \tag{11}$$

where the $f^{(+)}$ form factor and phase-space factor λ are

$$f^{(+)}_{\pi^+B^+}(m_K^2)=f^{(+)}_{\pi^+B^+}(0)/(1-m_K^2/m_B^2)=0.336\ , \tag{12}$$

$$\lambda_{\pi K}=(1-m_\pi^2/m_B^2-m_K^2/m_B^2)^2-4m_\pi^2m_K^2/m_B^4\ . \tag{13}$$

The factor α_1 which will also occur in the other rates is defined by

$$\alpha_1=\tfrac{1}{54}(\alpha_s G_1/U_{bc})^2=0.04 . \tag{14}$$

Note here that the total B decay width is estimated from the leptonic branching ratio and is $\Gamma_B^{\rm tot}\approx 4|U_{bc}|^2(G_F^2 m_b^5/192\pi^3)$.

Now if we consider the tree Hamiltonian, we see that only the annihilation diagram proportional to $L(B^+)$ is present, giving extremely small contribution.

B. $B^+\to K^+\pi^0$

The penguin contribution is now computed easily by using $\Delta I=0$ rule, and therefore

$$B^P(B^+\to K^+\pi^0)=\tfrac{1}{2}B^P(B^+\to K^0\pi^+) . \tag{15}$$

The tree contribution, neglecting the annihilation diagram, comes through $L(\pi^0)$ and $L(K^+)$ terms:

$$B^T(B^+\to K^+\pi^0)=(U_{ub}/U_{bc})^2(\pi/3m_B^3)^2\lambda_{\pi K}^{1/2}\times[f_\pi(m_B^2-m_K^2)f_{KB}^{(+)}(m_\pi^2)+3f_K(m_B^2-m_\pi^2)f_{\pi^+B^+}^{(+)}(m_K^2)]^2 . \tag{16}$$

Using definition[8] $f_{KB}^{(+)}(m_\pi^2)=f_{KB}^{(+)}(0)/(1-m_\pi^2/m_B^2)=0.379$, and $U_{ub}/U_{bc}=0.1\zeta$, we find the value

$$B^T(B^+\to K^+\pi^0)=0.06\times 10^{-5}\zeta^2 . \tag{17}$$

Similarly we find that for $B^0\to K^+\pi^-$, $K^{*+}\pi^-$ modes the tree contributions are potentially quite large, depending on the value of ζ^2.

C. $B^+\to K^0\rho^+$

Flavor symmetry gives $L(\rho^+)=0$. Terms proportional to $L(K^0)$ and $L(B^+)$ give, in principle, a nonvanishing contribution. Since we believe that this mode is important, as a test of our method we give some more details. First, according to Eq. (10) the amplitude is

$$A^P(B^+\to K^0\rho^+)=\tfrac{2}{3}\kappa[k\epsilon(k')]\{f_K(m_B^2-m_\rho^2)[A^1_{B^+\rho^+}(m_K^2)-A^2_{B^+\rho^+}(m_K^2)]+f_B(m_\rho^2-m_K^2)[A^1_{K^0\rho^+}(m_B^2)-A^2_{K^0\rho^+}(m_B^2)]\} . \tag{18}$$

The difference $[A^1-A^2]$ in the above expression is less than 0.1 in the model of Ref. 8. Thus the branching ratio is vanishingly small. From the tree Hamiltonian, we find $L(\rho^0)=L(K^0)=0$ and the contribution from the annihilation diagram is also small as is the penguin contribution. We thus find that this mode should not be found if our approximation scheme is correct. The same situation occurs for $B^0\to K^+\rho^-$ mode.

D. $B^+\to K^+\phi$

In all modes where ϕ is present, the tree Hamiltonian contribution is flavor forbidden because one needs three s quarks in the final state. From the penguin diagram in this case flavor symmetry gives $L(B^+)=L(K^+)=0$. The nonvanishing contribution comes from the term proportional to $L(\phi)$:

$$B^P(B^+\to K^+\phi)=\alpha_1(g_\phi/m_B m_\phi)^2[f_{KB}^{(+)}(m_\phi^2)]^2\lambda_{\phi K}^{3/2} , \tag{19}$$

$$\lambda_{\phi K}=(1-m_\phi^2/m_B^2-m_K^2/m_B^2)^2-4m_K^2m_\phi^2/m_B^4 ,\quad f_{KB}^{(+)}(m_\phi^2)=0.394 . \tag{20}$$

The couplings g_ϕ,g_ρ are determined from the $\phi,\rho\to e^+e^-$ experimental rates:

$$g_\phi^2=3m_\phi^3\Gamma(\phi\to e^+e^-)/(4\pi\alpha^2Q_s^2)=0.0575\ {\rm GeV}^4 , \tag{21}$$

$$g_{\rho^0}^2=3m_\rho^3\Gamma(\rho\to e^+e^-)/(4\pi\alpha^2)=0.0141\ {\rm GeV}^4 . \tag{22}$$

We then find

$$B^P(B^+\to K^+\phi)=1.50\times 10^{-5} . \tag{23}$$

E. $B^+\to K^{*0}\pi^+$

In this mode flavor symmetry and current conservation gives $L^*(\pi^+)=0$ and $L^*(B^+)=0$, respectively. The branching ratio from $L^*(K^{*0})$ is

$$B^P(B^+\to K^{*0}\pi^+)=\alpha_1(g_{K^*}/m_B m_{K^*})^2[f_{\pi^+B^+}^{(+)}(m_{K^*}^2)]^2\lambda_{\pi K^*}^{3/2} . \tag{24}$$

The g_{K^*} constant is determined from Eq. (10b). The other two unknown quantities from Eq. (24) are

$$\lambda_{\pi K^*}=(1-m_\pi^2/m_B^2-m_{K^*}^2/m_B^2)^2-4m_\pi^2m_{K^*}^2/m_B^4 ,\quad f_{\pi^+B^+}^{(+)}(m_{K^*}^2)=0.343 . \tag{25}$$

The tree Hamiltonian in this case gives a negligible amplitude.

F. $B^+\to K^{*0}\rho^+$

The same arguments as in Sec. III D give here $L^*(\rho^+)=L^*(B^+)=0$ and

$$B^P(B^+\to K^{*0}\rho^+)=\alpha_1(g_{K^*}/m_B)^2\lambda_{K^*\rho}^{3/2}$$
$$\times\{2V_{\rho^+B^+}^2(m_{K^*}^2)+[3(1-m_\rho^2/m_B^2)^2/\lambda_{K^*\rho}+m_{K^*}^2/4m_\rho^2-1]A_{\rho^+B^+}^2(m_{K^*}^2)\}\ , \tag{26}$$

$$V_{\rho^+B^+}(m_{K^*}^2)=[V_{\rho^+B^+}(0)/(m_B+m_\rho)](1-m_{K^*}^2/m_B^2)^{-1}=0.0560\ \mathrm{GeV}^{-1}\ , \tag{27}$$

$$A_{\rho^+B^+}(m_{K^*}^2)=[A_{\rho^+B^+}(0)/(m_B+m_\rho)](1-m_{K^*}^2/m_B^2)^{-1}=0.0482\ \mathrm{GeV}^{-1}\ , \tag{28}$$

$$\lambda_{K^*\rho}=(1-m_{K^*}^2/m_B^2-m_\rho^2/m_B^2)^2-4m_{K^*}^2m_\rho^2/m_B^4\ . \tag{29}$$

Here as well as in the $K^{*+}\rho^0$ mode the tree contributions are extremely small.

G. $B^+\to K^{*+}\phi$

Here the same argument as in Sec. III C holds; i.e., flavor symmetry gives $L^*(K^{*+})=L^*(B^+)=0$. The branching ratio is

$$B^P(B^+\to K^{*+}\phi)=\alpha_1(g_\phi/m_B)^2\lambda_{K^*\phi}^{3/2}$$
$$\times\{2V_{K^*B}^2(m_\phi^2)$$
$$+[3(1-m_\phi^2/m_B^2)^2/\lambda_{K^*\phi}$$
$$+m_\phi^2/4m_{K^*}^2-1]A_{K^*B}^2(m_\phi^2)\} \tag{30}$$

$$V_{K^*B}(m_\phi^2)=[V_{K^*B}(0)/(m_B+m_{K^*})]$$
$$\times(1-m_\phi^2/m_B^2)^{-1}=0.0621\ \mathrm{GeV}^{-1}\ , \tag{31}$$

$$A_{K^*B}(m_\phi^2)=[A_{K^*B}(0)/(m_B+m_{K^*})]$$
$$\times(1-m_\phi^2/m_B^2)^{-1}=0.0552\ \mathrm{GeV}^{-1}\ , \tag{32}$$

$$\lambda_{K^*\phi}=(1-m_{K^*}^2/m_B^2-m_\phi^2/m_B^2)^2-4m_{K^*}^2m_\phi^2/m_B^4\ . \tag{33}$$

We have illustrated our method of calculation in the above examples. The results are presented in Table I.

IV. DISCUSSION

We have considered a large number of two-body charmless decay modes of the B meson. Although in principle they can receive contributions from the penguin or tree diagrams, we found a number of modes where one of the contributions clearly dominates. Many of the branching ratios are in the interesting range of a few times 10^{-5}, which should be accessible very soon, as seen from the present experimental limits. We further find a few modes that are highly suppressed, receiving a negligible contribution from both Hamiltonians. The lack of observation of these would test our method of computation. When the contributions are comparable, one, of course, should add the amplitudes and square. We have not done so because we do not know the magnitude and sign of U_{ub}, and there might be different final-state phases on these amplitudes. A clean test is provided by those decays where one of Hamiltonian dominates. In the case where the tree Hamiltonian dominates, one could even estimate U_{bu}. As our knowledge of these decays improves, we believe more ambitious methods might be used to calculate the amplitudes.

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