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WEAK BOTTOM-QUARK DECAYS AND *QCD*

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QCD effects in weak *b*-quark decays are studied in the Kobayashi-Maskawa model using renormalization-group techniques to sum up leading logarithms. It is found that penguins are more stable to the variation of parameters than in the lowest-order estimate. Our analysis confirms that the Cabibbo pattern will be changed. It is estimated that for inclusive decays with no charmed quarks in the final states the ratio $\Gamma(B \to K + nX) / \Gamma(B \to (n+1) X)$ is approximately equal to 0.3—1.2 instead of being 5%.

1. *Introduction*

The existence of the bottom quark is now well established[1]. A question naturally arises how to incorporate the fifth quark in the standard Glashow-Weinberg-Salam (*GWS*) model[2]. Among the various schemes proposed[3] to answer this question, the most popular one is the six-quark model of Kobayashi and Maskawa (*KM*)[4]. This model is a *natural* generalization of the *GWS* model, with the top and bottom quarks in the left-handed weak doublet; it has three generalized Cabibbo-like angles ϑ_i, $i = 1, 2, 3$, and an arbitrary phase δ, the latter being responsible for *CP* violation. However, the angles and the phase are not determined by the model itself and it is important to know their values. Some constraints come, of course, from light-quark physics, e. g. from β decay, K_{e3} and strange-hyperon decays[5]. Recent experiments from the *CLEO*, *CUSB* and *JADE* collaborations[6,7]

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have proved to be very successful in ruling out most of the so-called topless models[8]. These experiments are, in fact, consistent with the KM model.

The question of the Cabibbo pattern in b-quark decays was studied some time ago by Peccei, Rückl and ourselves[9]. We found that in the pure hadronic sector strong interactions, described by perturbative QCD, might change the *natural* Cabibbo pattern of weak botom-quark decays through new QCD-induced operators, the so-called penguins. The same phenomena were studied independently by Eilam and Leveille[10], reaching similar conclusions. These operators were first invoked to explain the $\Delta I = \frac{1}{2}$ rule in weak decays of kaons and hyperons[11], but they also appeared in the study of other phenomena, such as parity violation in nuclei[12], or CP violation in K decays[13-15]. Since their role is somewhat masked by the uncertainties due to the possible nonperturbative effects and the unreliability of the use of perturbative QCD at the scale of light hadrons, it would be very useful to study these effects at the scale of b quarks, at which most of the difficulties mentioned above seem to disappear. In addition, the question of the Cabibbo pattern has been the subject of intensive theoretical studies[16,17], since the predictions concerning this pattern could be distinguished experimentally.

Therefore, to study the effects described in Refs. 9 and 10 in more detail, we use renormalization-group equations to sum up leading logarithms to all orders in strong interactions, and we compare these effects, to a certain extent, with penguin phenomena in light-hadron physics. We treat renormalization-group effects in a simple and general way, enabling the reader to apply the same technique in a trivial way to all other processes.

The plan of the paper is as follows. In Sect. 2 we discuss briefly the main properties of the Kobayashi-Maskawa model, the constraints on the KM angles and the Cabibbo pattern of nonleptonic weak decays. In sect. 3 we proceed with the calculation of QCD effects using renormalization-group techniques and discuss the choice of the renormalization point. Section 4 contains applications to different weak decays of bottom quarks and a comparison with adequate calculations in charmed and light-hadron physics. Finally, we discuss our results and draw conclusions.

2. *The Kobayashi-Maskawa six-quark model*

We start by discussing briefly the Kobayashi-Maskawa model[4]. The weak mixing-angle structure is given by the matrix

$$U = \begin{pmatrix} U_{ud} & U_{us} & U_{ub} \\ U_{cd} & U_{cs} & U_{cb} \\ U_{td} & U_{ts} & U_{tb} \end{pmatrix} = \tag{2.1}$$

$$= \begin{pmatrix} c_1 & s_1 c_3 & s_1 s_3 \\ -s_1 c_2 & c_1 c_2 c_3 - s_2 s_3\, e^{i\delta} & c_1 c_2 s_3 + s_2 c_3\, e^{i\delta} \\ -s_1 s_2 & c_1 s_2 c_3 + c_2 s_3\, e^{i\delta} & c_1 s_2 s_3 - c_2 c_3\, e^{i\delta} \end{pmatrix},$$

where c_i and s_i stand for cos ϑ_i and sin ϑ_i, respectively, and δ is an arbitrary phase responsible for CP violation. We use the convention from Ref. 1, with all ϑ_i lying in the first quadrant. The nonleptonic weak Hamiltonian is given by

$$H_w = \frac{G_F}{\sqrt{2}} J^+_\mu J^\mu \tag{2.2}$$

with

$$J_\mu = \overline{\Psi}^{up} \gamma_\mu (1 - \gamma_5) \, U \, \Psi^{down},$$
$$\Psi^{up} = \begin{pmatrix} u \\ c \\ t \end{pmatrix}, \quad \Psi^{down} = \begin{pmatrix} d \\ s \\ b \end{pmatrix}. \tag{2.3}$$

It is evident from Eq. (2.1) that the weak-transition structure is much richer than in the GWS four-quark model. There are two possibilities of studing this model. One possibility is to try to set constraints on the KM angles and phases from the known light-hadron physics[6,20,21]; the other is to try to give different predictions from the weak Hamiltonian that could be distinguished experimentally[16,17,20].

Now, we turn to the nonleptonic weak Hamiltonian. In Table 1 we give explicitly the quark content of the product $J^+_\mu J^\mu$ from (2.2), the Cabibbo structure of quarks and their properties in respect to flavour. From the table we recognize the

TABLE 1

Flavour content	KM structure	Flavour selection rules
$(\bar{c}b)(\bar{d}u)$	$c_1 U_{cb}$	$\Delta C = 1, \Delta S = 0$ (open charm)
$(\bar{c}b)(\bar{s}c)$	$(c_1c_2c_3 - s_2s_2c_\delta) U_{cb}$	$\Delta C = 0, \Delta S = 1$ (hidden charm)
$(\bar{c}b)(\bar{d}c)$	$-s_1c_2 U_{cb}$	$\Delta C = 0, \Delta S = 0$ (hidden charm)
$(\bar{c}b)(\bar{s}u)$	$s_1c_3 U_{cb}$	$\Delta C = 1, \Delta S = 1$ (open charm)
$(\bar{u}b)(\bar{d}u)$	$c_1 U_{ub}$	$\Delta C = 0, \Delta S = 0$ (no charm)
$(\bar{u}b)(\bar{s}c)$	$(c_1c_2c_3 - s_2s_3c_\delta) U_{ub}$	$\Delta C = 1, \Delta S = 1$ (open charm)
$(\bar{u}b)(\bar{d}c)$	$-s_1c_2 U_{ub}$	$\Delta C = 0, \Delta S = 0$ (open charm)
$(\bar{u}b)(\bar{s}u)$	$s_1c_3 U_{ub}$	$\Delta C = 0, \Delta S = 1$ (no charm)

Flavour content, KM structure and flavour selection rules fo the product $J^+_\mu J^\mu$, with $U_{cb} = c_1c_2s_3 + s_2c_3c_\delta$, $U_{ub} = s_1s_3$.

simple hierarchy given by the KM angles. Assuming that all angles have roughly the same value, we obtain three levels, namely Cabibbo-allowed decays with one $\sin \vartheta$ suppression and two Cabibbo-suppressed levels with $\sin^2 \vartheta$ and $\sin^3 \vartheta$ suppressions, respectively.

It follows, for example, that $b \to c$ transitions are Cabibbo-favoured in comparison with $b \to u$ transitions, the latter being suppressed by at least $\sin \vartheta_1 \sin \vartheta_3$. Therefore, the study of the ratio

$$\frac{\Gamma(b \to u)}{\Gamma(b \to c)} \tag{2.4}$$

might give us useful information on the Cabibbo pattern of nonleptonic b decays[17,20].

This rather simple picture has been questioned by the study of QCD effects[9,10]. It has been found that penquin diagrams appear at two Cabibbo levels. Generally, they come only from the parts of $J_\mu^+ J^\mu$ which have the structure

$$(J_\mu^+ J^\mu)_l = \sum_j U_{jb} U_{jl}^* \bar{\psi}_l \gamma_\mu (1 - \gamma_5) \psi_j \bar{\psi}_j \gamma^\mu (1 - \gamma_5) b, \tag{2.5}$$

because this structure allows the quarks ψ_j to appear in the loops. It is obvious from the Hamiltonian structure in Table 1 that $\bar{\psi}_l$ would be either s or d quarks. This leads to new operators in the nonleptonic weak Hamiltonian which are of the form

$$O_l^{penguin} \sim \bar{\psi}_l \gamma_\mu (1 - \gamma_5) \lambda_a b \times \sum_j \bar{\psi}_j \gamma^\mu \lambda_a \psi_j. \tag{2.6}$$

The penguin operator with $\bar{\psi}_l = \bar{s}$, being proportional to $U_{tb} U_{ts}^* \sim \sin \vartheta$, is Cabibbo-favoured according to our classification and is the important one. In spite of the fact that the coefficient of this operator is of order α_s, i. e. rather small, it has been estimated[9] that its contribution will compete with the Cabibbo suppressed $b \to u$, $\Delta S = 0$ transitions and dominate the $b \to u$, $\Delta S = 1$ transitions! This means, for example[9], that inclusive decays with no c quark in the final state, such as $B \to K + n\pi$ and $B \to (n+1)\pi$ should appear at comparable rates, adding more kaons at the first Cabibbo-suppressed level.

In the next section we derive the nonleptonic weak Hamiltonian by summing up leading-logarithmic corrections to all orders in α_s. Since the operator (2.6) contains terms with both left-left and left-right helicities, these terms will be renormalized differently, in contrast to the lowest-order calculation in Ref. 3 where they had the same coefficient.

3. *QCD corrections to the weak Hamiltonian*

The Hamiltonian which generates leading penguins (LP) has the form

$$H^{LP} = \frac{G_F}{\sqrt{2}} \{\xi_1 (\bar{s}c)(\bar{c}b) + \xi_2 (\bar{s}u)(\bar{u}b) + \xi_3 (\bar{s}t)(\bar{t}b)\} \tag{3.1}$$

with

$$\xi_1 = U_{cb} U^*_{cs},\ \xi_2 = U_{ub} U^*_{us},\ \xi_3 = U_{tb} U^*_{ts}.$$

The structure of the Hamiltonian has the form $(V - A) \times (V - A)$. In order to simplify the calculation, we shall take only the parity-violating part $-(AV + VA)$. Final results will be valid for the $(V - A) \times (V - A)$ part, because the parity properties are not affected by strong interactions.

It is convenient to write the Hamiltonian in the basis $\psi = \{u, d, s, c, b, t\}$, using flavour matrices A, B, C,.... There are generally two types of operators:

$$O(A, B) = (\overline{\psi} \gamma_\mu \gamma_5 A\psi)(\overline{\psi} \gamma^\mu B\psi), \tag{3.2}$$

$$O(A\lambda, B\lambda) = (\overline{\psi} \gamma_\mu \gamma_5 A\lambda_a \psi)(\overline{\psi} \gamma^\mu B\lambda_a \psi)$$

with colour matrices λ_a, $a = 1, \ldots, 8$. Equation (3.1) can now be written as

$$H^{LP}_{PV} = -\frac{G_F}{\sqrt{2}} \sum_{i=1}^{3} \xi_i [O(A_i, B_i) + O(B_i, A_i)], \tag{3.3}$$

where each flavour matrix has only one nonvanishing element:

$$(A_1)_{45} = (A_2)_{15} = (A_3)_{65} = (B_1)_{34} = (B_2)_{31} = (B_3)_{36} = 1. \tag{3.4}$$

It is convenient to use flavour matrices because, once the corrections are found, they are valid, up to the trivial matrix multiplication, for any flavour matrices. This reflects the fact that QCD is flavour-independent. Furthermore, it is convenient to express the Hamiltonian in terms of the operators which are, at least in the flavour-symmetry limit, multiplicatively renormalizable[22]. One obtains

$$H^{LP}_{PV} = \sqrt{2}\, G_F\, [O'^{PV}_{189} + O'^{PV}_{405}], \tag{3.5}$$

where O_{189} and O_{405} transform under $SU(6)$ flavour as irreducible representations 189 and 405, respectively. They are given explicitly in terms of the operator $O(A_i, B_i)$ and $O(A_i\lambda, B_i\lambda)$ as

$$-2O'^{PV}_{189} = \frac{1}{3} \sum_{i=1}^{3} \xi_i O(A_i, B_i) - \frac{1}{4} \sum_{i=1}^{3} \xi_i O(A_i\lambda, B_i\lambda) + (A_i \leftrightarrow B_i),$$

$$-2O'^{PV}_{405} = \frac{2}{3} \sum_{i=1}^{3} \xi_i O(A_i, B_i) + \frac{1}{4} \sum_{i=1}^{3} \xi_i O(A_i\lambda, B_i\lambda) + (A_i \leftrightarrow B_i). \tag{3.6}$$

At this point we remind the reader that QCD modifies the weak Hamiltonian into the form[22]

$$H_w = \sqrt{2}\, G_F \sum_i c_i \left(\frac{M}{\mu}, \frac{M}{m_q}, \ldots, g\right) O_i, \tag{3.7}$$

where the sum extends over all operators including those which arise by mixing in the renormalization procedure. The calculation relies essentially on the assumption that the weak-transition matrix element might be factorized* into the *soft* part with small momenta of the order of quark masses and into the *hard* part dominated by the large momentum flow of the order of m_W. Then the soft part is described in terms of the matrix elements of renormalized local operators, whereas the Wilson coefficients c_i satisfy the renormalization-group (RG) equation with a well-known solution depending on the anomalous dimensions of the operator O_i. These coefficients also depend on the quark masses m_q the mass M of the W boson and the renormalization point μ. The RG equation is rather trivial as long as quark masses are small in comparison with μ. Since the matrix element of the weak Hamiltonian does not depend on μ, the μ dependence of the coefficients $c_i(\mu)$ in (3.7) has to cancel the μ dependence of the matrix elements of the operators O_i. The latter dependence is not known at all; μ should therefore be taken at the typical hadron scale $\mu \simeq \mu_0$, for which our calculations of the matrix elements have some sense[11-14]. This point will be discussed in more detail later.

Now, we turn to the anomalous dimensions. These can be evaluated by inserting the operators O_i in all possible irreducible one-loop Green functions, as shown in Fig. 1. These diagrams are usually renormalized at a symmetric point,

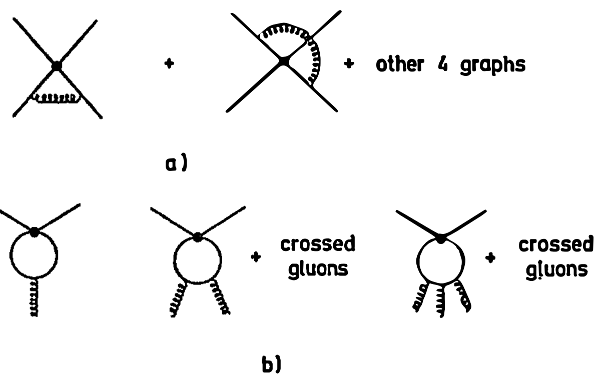

Fig. 1. Feynman diagrams for renormalization matrices. Heavy dots indicate operator insertion Full lines are quarks and wiggly lines are gluons. Set b are penguin diagrams.

$$p_i^2 = -\mu^2, \tag{3.8}$$

$$2p_i \cdot p_j = \frac{2}{3}\mu^2,$$

where p_i and p_j are momenta of the external legs. The renormalization matrix Z_{ij} which connects the unrenormalized operator O_i^0 with the renormalized operators O_j^{ren},

$$O_i^0 = Z_{ij} O_j^{ren}, \tag{3.9}$$

*The validity of this assumption has recently been studied and proved by Galić[23] for $Q\bar{Q}$ systems with heavy Q.

is easily calculated by using the general formulae given in Appendix A. From (3.9) one obtains the anomalous dimensions γ_{ij} in a trivial way,

$$\gamma_{ij} = \frac{\partial \ln Z_{ij}}{\partial \ln \mu}, \tag{3.10}$$

and one easily writes the solution for the coefficients c_i. However, functions, such as γ or β which enter the solution have to be evaluated for some fixed value of μ. Owing to flavour-symmetry breaking, the leading logarithmic parts of these functions depend crucially on the region from which μ is taken. Since μ should be chosen to be near a typical hadronic scale, which for b-quark decays is given by the b-quark mass (or even smaller), it follows that μ is smaller than the mass of the top quark, which appears in the loops of Fig. 1. Therefore, we first find the solution for μ in the range $m_W^2 > \mu^2 > m_t^2$, which, for $\mu = m_t$, is given by

$$H^{LP}(\mu = m_t) = \sqrt{2}\, G_F\, [\varkappa_t^{4/7}\, O'_{189} + \varkappa_t^{-2/7}\, O'_{405}],$$

$$\varkappa_t = \frac{\alpha_s(m_t^2)}{\alpha_s(m_W^2)} = 1 + \alpha_s(m_t^2)\frac{b}{4\pi}\ln\frac{m_W^2}{m_t^2}. \tag{3.11}$$

At this level, *QCD* corrections also generate penguin operators $O(C)$ of the form

$$O(C) = -\frac{2}{g}\,\overline{\psi}\,\gamma_\nu\,\gamma_5\, C\lambda_a\,\psi\,(\nabla_\mu F^{\mu\nu})_a, \tag{3.12}$$

where $F_a^{\mu\nu} = \partial^\mu A_a^\nu - \partial^\nu A_a^\mu + g f_{abc} A_b^\mu A_c^\nu$ is a gluonic tensor and ∇^μ is a gauge-covariant derivative. However, the coefficients of these operators are proportional to $\xi_1 + \xi_2 + \xi_3$, which is equal to zero owing to the *GIM* mechanism.

Now, we proceed to the region $m_t^2 > \mu^2 > m_c^2$ and evaluate the Hamiltonian at the point μ. The external and internal momenta p_i flowing in the penguin diagram should be taken at the symmetric point $p_i^2 = -\mu^2 < m_t^2$. The diagram with a heavy quark t in the loop is suppressed. Its leading-logarithmic term has no $\ln \mu^2$ dependence and, according to (3.10), does not contribute to the anomalous dimension of the operators O_j. Furthermore, the coefficient of $O(C)$ is proportional to $\xi_1 + \xi_2 = -\xi_3$, and the *GIM* mechanism is no longer effective. The new operator $O(C)$ should also be inserted into all irreducible Green functions. However, one realizes that by using the equation of motion, one obtains the four-quark operator again,

$$O(C) = (\overline{\psi}\,\gamma_\mu\,\gamma_5\, C\lambda_a\,\psi)\sum_f(\overline{\psi}_f\,\gamma^\mu\,\lambda_a\,\psi_f) \equiv O(C\lambda, \lambda) \tag{3.13}$$

which can be easily renormalized by using our rules given in Appendix A. The matrix C has only one nonvanishing element, $(C)_{35} = 1$, giving in our case

$$\overline{s}\,\gamma_\mu\,\gamma_5\,\lambda_a\, b \sum_f(\overline{\psi}_f\,\gamma^\mu\,\lambda_a\,\psi_f). \tag{3.14}$$

This is, of course, only the PV part of the penguin operator. Its full Lorentz structure has the form $(V - A) \times V$, i. e. this operator contains two terms, one with left-left and the other with left-right helicities.

At this stage we choose the operator basis in such a way that the connection with the Shifman-Vainshtein-Zakharov (SVZ) operators[11] is transparent. We have to define the following operators in addition to O_{189} and O_{405}:

$$
\begin{aligned}
O_5'^{PV} &= \frac{1}{4} \xi_3 \, [O(C\lambda, \lambda) - O(\lambda, C\lambda)] = - \xi_3 \, O_5^{PV}, \\
O_6'^{PV} &= \frac{1}{4} \xi_3 \, [O(C, 1) - O(1, C)] = - \xi_3 \, O_6^{PV}, \\
O_7'^{PV} &= \frac{1}{4} \xi_3 \, [O(C, 1) + O(1, C)] = - \xi_3 \, O_7^{PV}, \\
O_8'^{PV} &= \frac{1}{4} \xi_3 \, [O(C\lambda, \lambda) + O(\lambda, C\lambda)] = - \xi_3 \, O_8^{PV}.
\end{aligned}
\tag{3.15}
$$

The operators O_5 and O_6 have mixed L—R helicities and correspond to the SVZ penguin operators. The last two operators, O_7 and O_8, can be expressed in terms of the operators O_{189} and O_{405} if we choose μ to be in the range $m_t^2 > \mu^2 > m_c^2$. In some cases, such as the SVZ $\Delta S = 1$ Hamiltonian, it is convenient to express O_{189} and O_{405} in terms of the operators O_1, O_2, O_3 and O_4 with definite $SU(3)$ properties. This explains the way we have labelled the operators (3.15).

The operators (3.15) together with O_{189} and O_{405} form a closed set, i. e. the renormalization procedure does not induce any new operator. However, they are not multiplicatively renormalizable and one needs to diagonalize the matrix Z_{ij}. As an example, by using the formulae given in Appendix A, one easily finds that the insertion of the operator O_5' in the diagrams of Fig. 1 leads to

$$
O_5'^{PV} \to \left(\frac{\alpha_s}{4\pi} \ln \mu\right) \left\{ \left(14 - \frac{2}{3} n_f\right) O_5'^{PV} + \frac{32}{3} O_6'^{PV} - \frac{2}{3} n_f \, O_8'^{PV} \right\}. \tag{3.16}
$$

The elements of the matrix Z_{ij} are simply the coefficients of the operators in (3.16). The complete matrix Z_{ij} could be obtained by using the same procedure. In Appendix A we give Z_{ij} explicitly together with the diagonalization matrices φ^{-1} and φ. A final expression for the Hamiltonian defined at the point μ is given in the form (3.7),

$$
H^{LP}(\mu) = \sqrt{2}\, G_F \sum_j c_j(\mu)\, O_j \tag{3.17}
$$

with

$$
c_j(\mu) = \sum_{i,r} \varkappa_t^{d_i} \varkappa^{d_r'} c_i^F (\varphi^{-1})_{ir} \varphi_{rj}. \tag{3.18}
$$

The sum over i, j, r extends over the operator set $\{O'_{189}, O'_{405}, O'_5, O'_6, O'_7, O'_8\}$. c_i^F are free-field values of the Wilson coefficients:

$$\begin{aligned} c_{189}^F &= 1, \\ c_{405}^F &= 1, \\ c_5^F = c_6^F = c_7^F &= c_8^F = 0. \end{aligned} \tag{3.19}$$

The last equation in (3.19) expresses the simple fact that the operators $O_5, \ldots, O_8$ are induced by *QCD* and do not exist in the original Hamiltonian. The exponents d_i and d'_i are proportional to the anomalous dimensions of the operator O_i and are given in Appendix A. $\varkappa_t$ is given by (3.11) and $\varkappa$ is

$$\varkappa = \frac{\alpha_s(\mu^2)}{\alpha_s(m_t^2)} = 1 + \alpha_s(\mu^2)\frac{b'}{4\pi}\ln\frac{m_t^2}{\mu^2}. \tag{3.20}$$

At this point we want to discuss our results and compare them with the all--order calculation of Ref. 18. We have taken the renormalization point μ from the region $m_t^2 > \mu^2 > m_c^2$, since μ should be chosen to be near a typical hadronic scale. This scale is governed by the heaviest-quark mass, being of the order of 4—5 GeV in our case. However, the b quark would decay to lighter quarks, and the momenta of the quarks entering the diagrams of Fig. 1 are not symmetrical at all, contrary to our condition (3.8). Therefore, it might be more appropriate to take μ to be some average smaller mass of the particles engaged, say, at 2—3 GeV, which is still larger than the mass of the charmed quark. In Ref. 18, the point μ was taken at the scale below the charm threshold, and we find it unreliable. Moreover, one more region, i. e. $m_c^2 > \mu^2$, had to be taken into account, which made the calculation of Ref. 18 too complicated. For such a small μ one cannot hope to match the μ dependence of the matrix elements of the operators O_i, as we have already discussed.

The simplicity of the calculation presented in this section is mostly due to the choice of the operators $O_5, \ldots, O_8$ given by (3.15). If one had been forced to proceed to the point $\mu < m_c^2$, one would have used the same set (3.15). Other choice of the operator basis would have led to a more complicated structure.

Moreover, the choice (3.15) and the flavour-matrix formalism enable one to perform similar calculations easily. For example, kaon and hyperon decays described by the $\Delta S = 1$ Hamiltonian could be obtained from our calculation by making the substitutions $b \to d$ and $U_{bs} \to U_{ds}$. Penguin coefficients would now be proportional to $\xi_3 = U_{ts} U_{td}$. A glance at the matrix U reveals that ξ_3 also contains a *CP*-violating part proportional to i sin δ. In this way, one immediately obtains the *CP*-violating $\Delta S = 1$ Hamiltonian defined at μ, $m_t^2 > \mu^2 > m_c^2$. However, the appropriate choice of μ would be $m_c^2 > \mu^2 > m_\mu^2$ and one is forced to find H_w defined at such μ. In this region one needs to renormalize the set $O'_5, \ldots O'_8$, since O'_{189} and O'_{405} would not contribute* to *CP* violation. The new renormalization matrix Z_{ij} and the diagonalization matrices φ' and φ'^{-1} may be found in the Appendix of Ref. 25, with Z' being identified as Z_B^{15}, φ' as U, etc. The matrices

*Strictly speaking, the terms in the Hamiltonian with $c\bar{c}$ and $t\bar{t}$ quarks do not contribute in the valence quark-model calculations, cf. Refs. 13—15.

Z_B^{15} and U in Ref. 27 were obtained for the $\Delta S = 0$ weak Hamiltonian. Part of this Hamiltonian coming from the product of neutral currents had, up to the flavour, the structure of the operators $O_5, \ldots, O_8$. Thus one immediately obtains the *CP*-violating Hamiltonian of Ref. 14. It is a suitable example of the advantage of using the general flavour matrix formalism.

4. *Decay properties*

In this section we want to study how the new operators $O_5, \ldots, O_8$ influence the inclusive — and exclusive-decay properties of the b quark. We express our Hamiltonian in terms of left — and right-handed quarks

$$\psi_L = \frac{1}{2}(1 - \gamma_5)\,\psi, \quad \psi_R = \frac{1}{2}(1 + \gamma_5)\,\psi. \tag{4.1}$$

Furthermore, we use the notation $O_{189}^{(n)} = \xi_n O_-^{(n)}$ and $O_{405}^{(n)} = \xi_n O_+^{(n)}$. In this way, the basis of the operators O_-, O_+, O_5, O_6, O_7 and O_8 corresponds exactly to the basis of the SVZ operators[11], apart from the flavour content. We obtain from (3.17)

$$H_W^{LP} = \sqrt{2}\, G_F \sum_{n=1}^{3} \xi_n \{c_-\, O_-^{(n)} + c_+\, O_+^{(n)}\} + \sqrt{2}\, G_F\, \xi_p \sum_{i=5}^{8} c_i\, O_i, \tag{4.2}$$

where ξ_1, ξ_2 and ξ_3 are given by (3.1) and $\xi_p = -U_{tb}\, U_{ts}^*$. The operators $O_\pm^{(n)}$ are given, for example for $n = 2$, by

$$\begin{aligned} O_-^{(2)} &= 2\left[\frac{1}{3}(\bar{u}_L \gamma_\mu b_L)(\bar{s}_L \gamma^\mu u_L) - \left(\bar{u}_L \gamma_\mu \frac{\lambda_a}{2} b_L\right)\left(\bar{s}\, \gamma^\mu \frac{\lambda_a}{2} u_L\right)\right], \\ O_+^{(2)} &= 2\left[\frac{2}{3}(\bar{u}_L \gamma_\mu b_L)(\bar{s}_L \gamma^\mu u_L) + \left(\bar{u}_L \gamma_\mu \frac{\lambda_a}{2} b_L\right)\left(\bar{s}_L \gamma^\mu \frac{\lambda_a}{2} u_L\right)\right]. \end{aligned} \tag{4.3}$$

The operators $O_\pm^{(1)}$ and $O_\pm^{(3)}$ are obtained from (4.3) simply by making the substitutions $u \to c$ and $u \to t$, respectively*. Using the relation.

$$(\lambda_a)^{ij}(\lambda_a)^{kl} = 2\,\delta^{il}\,\delta^{jk} - \frac{2}{3}\,\delta^{ij}\,\delta^{kl}, \tag{4.4}$$

one can express the operators (4.3) in the form

$$O_\pm^{(2)} = (\bar{u}_L \gamma_\mu b_L)(\bar{s}_L \gamma^\mu u_L) \pm (\bar{s}_L \gamma^\mu b_L)(\bar{u}_L \gamma^\mu u_L). \tag{4.5}$$

*The operator with the flavour structure $(\bar{s}b)\,(\bar{t}t)$ is of no interest for us.

This is precisely the form used by *SVZ*. However, for inclusive-decay rates it is more convenient to use the form (4.3). The operators $O_5, \dots O_8$ are given by

$$O_5 = (\bar{s}_L \gamma_\mu \lambda_a b_L) \sum_f \bar{\psi}^f_R \gamma^\mu \lambda_a \psi^f_R,$$

$$O_6 = (\bar{s}_L \gamma_\mu b_L) \sum_f \bar{\psi}^f_R \gamma^\mu \psi^f_R,$$

$$O_7 = (\bar{s}_L \gamma_\mu b_L) \sum_f \bar{\psi}^f_L \gamma^\mu \psi^f_L, \tag{4.6}$$

$$O_8 = (\bar{s}_L \gamma_\mu \lambda_a b_L) \sum_f \bar{\psi}^f_L \gamma^\mu \lambda_a \psi^f_L,$$

where the sum over f extends over quark flavours, $f = u, d, s, c$. The coefficients c^F_- and c^F_+ are normalized to unity according to (3.19) and otherwise are given by (3.18):

$$c_- = \varkappa_t^{4/7} \varkappa^{12/25}, \tag{4.7}$$

$$c_+ = \varkappa_t^{-2/7} \varkappa^{-6/25},$$

with $\varkappa_t$ and $\varkappa$ given by (3.11) and (3.20), respectively.

The values in (4.7) are precisely those which can be obtained in the limit of flavour symmetry. However, there also appear the operators O_7 and O_8 as a result of the operator mixing. Parts of O_7 and O_8 with the flavour structures $(\bar{s}b)(\bar{u}u)$ and $(\bar{s}b)(\bar{c}c)$ are linear combinations of the operators $O^{(2)}_\pm$ and $O^{(1)}_\pm$, respectively. Thus, as a result of mixing, the operators $O_\pm$ have changed their coefficients from the value given by (4.7). However, we prefer to keep the whole set (4.6) separate from the operators $O_\pm$ in order to see clearly the effect of the new operators.

The coefficients $c_5, \dots c_8$ can be obtained by using Eq. (3.18) with φ^{-1} and φ given in Appendix A. As can be seen from Eq. (3.18) the coefficients depend on the values of $\varkappa_t$ and $\varkappa$. These are functions of quark masses, *QCD* parameter Λ and renormalization point μ. We have fixed Λ and μ by choosing $\alpha_s = 1$ at 300 MeV and $m_c = 1.24$ GeV. Then, we vary the mass of the t quark in the range 20—60 GeV for μ in the range 2—4.5 GeV. Since the coefficients $c_5, \dots c_8$ vanish in the limit $\varkappa \to 1$, $\varkappa = \alpha_s(\mu^2) / \alpha_s(m_t^2)$, their values increase slowly with larger m_t and decrease slowly with larger μ. However, at reasonably large intervals, the values of, for example, c_5 stay in the range from -0.02 to -0.04. This is contrary to the $\Delta S = 1$ penguins of SVZ[11], for which the coefficients change by an order of magnitude for μ in the range 150—700 MeV. In our discussion we use, as the best guess, the values

$$c_- = 1.33,\ c_5 = -0.028,\ c_7 = 0.004, \tag{4.8}$$

$$c_+ = 0.867,\ c_6 = -0.005,\ c_8 = -0.024.$$

In light-hadron physics, $\mu \simeq m_\pi$, the values of c_-, c_+, c_5 and c_6 are typically 2.8, 0.6, −0.14 and −0.05, respectively. This difference with respect to (4.8) reflects the asymptotic freedom aspects of *QCD*.

The results obtained may be compared with the lowest-order calculation of Ref. 9. The coefficients c_5 and c_8 are equal and given by

$$c_5 = c_8 = -\frac{1}{3}\cdot\frac{\alpha_s(\mu^2)}{4\pi}\cdot\ln\left(\frac{m_t^2}{\mu^2}\right). \tag{4.9}$$

The values obtained are $-(0.015 \sim 0.05)$ for $m_t = 20$—60 GeV and $\mu = 2$—4.5 GeV. Thus, the all-order calculation (4.8) gives typical values for c_5 and c_8 which are less parameter-dependent. The coefficients c_6 and c_7 are equal to zero in the lowest-order calculation*.

To get a feeling of the strength of penguin effects, we shall look at their contribution to inclusive rates. We assume here that the decay rate for a B meson is given almost by the decay rate for a heavy b quark. Let us compare the typical rate coming from the operators $O_\pm^{(n)}$ with that coming, for example, from O_5. The calculation is straightforward[26] and gives (in units $\Gamma_0 = G_F^2 m_b^5/192\pi^3$)

$$\Gamma^{(n)} = 3\,\Gamma_0\,\xi_n^2\left\{\frac{1}{9}(2c_+ + c_-)^2 + \frac{2}{9}(c_+ - c_-)^2 + \frac{8}{9}\left(\frac{\xi_p}{\xi_n}\right)^2 c_5^2 n_f\right\}, \tag{4.10}$$

where the last two terms come from the $\lambda\lambda$ terms in $O_\pm^{(n)}$ and O_5, respectively. The first two terms add to give the well-known form $(2c_+^2 + c_-)^2$, i. e.

$$\Gamma^{(n)} = 3\,\Gamma_0\,\xi_n^2\left\{\frac{2c_+^2 + c_-^2}{3} + \frac{8}{9}\left(\frac{\xi_p}{\xi_n}\right)^2 c_5^2 n_f\right\}, \tag{4.11}$$

where n_f accounts for the fact that there are a few different flavour channels in (4.6) which we add incoherently.

Equation (4.11) is particularly suitable for studying the effects of the new operators inclusively. The value of the first term is $(2c_+^2 + c_-^2)/3 \simeq 1.1$ for the choice (4.8). One sees that the effect of *QCD* is rather modest, i. e. the enhancement in nonleptonic rate is of the order of 10%. It is clear that because of the smallness of penguin coefficients, the second term in (4.11) negligibly affects the rates for decays that include charmed quarks, i. e. those for which $\xi_n \simeq \xi_p$. On the other hand, the factor $(\xi_p/\xi_n)^2$ could be very large for other modes. This fact was pointed out in Ref. 9. Let us make a more detailed estimate. For the values (4.8), $c_5^2 \simeq (0.4$—$1.6) \times 10^{-3}$. However, since there are also a comparable contribution from O_8

*The difference between the lowest-order calculation and the all-order calculation is large for the coefficients c_- and c_+. For example, c_+ is even negative in the lowest order.

and small contributions from O_6 and O_7, the amplitude is enhanced at least by a factor of 2, so that effectively

$$c^2_{peng}\, n_f \simeq (0.5\text{—}2) \times 10^{-2}, \tag{4.12}$$

where we have incoherently added three noncharmed channels allowed by (4.6). We want to make a comparison between the contributions of the operators $O^{(2)}_{\pm}$ and those of the operators $O^{(4)}_{\pm}$ with the flavour structures $(\bar{u}b)(\bar{s}u)$ and $(\bar{u}b)(\bar{d}u)$, respectively. The former operator is maximally Cabibbo-suppressed and the latter is suppressed by roughly $\sin\vartheta_i$ in comparison with the Cabibbo-leading terms. Moreover, $O^{(2)}_{\pm}$ would always produce an odd number of kaons in the final state, contrary to the operator $O^{(4)}_{\pm}$. Therefore,

$$\frac{\Gamma^{(2)}(B \to K + nX)}{\Gamma^{(4)}(B \to (n+1)X)} = \left(\frac{\xi_2}{\xi_4}\right)^2 = c_3^2 \tan^2\vartheta_1 \simeq 5\%. \tag{4.13}$$

This estimate may be considered as an upper limit, because ϑ_1 is rather well known and $c_3^2 \leq 1$. Now, the important point is that the operators $O_5, \dots O_8$, like the operators $\widetilde{O}^{(2)}_{\pm}$, always produce an odd number of kaons because of their flavour structure (from the vacuum one obtains only $K\bar{K}$ pairs or φ). Their contribution is proportional to

$$\left(\frac{\xi_p}{\xi_4}\right)^2 = \frac{1}{s_1^2}\frac{c_2^2 c_3^2}{c_1^2}\left(c_2 c_\delta + c_1 c_3 \frac{s_2}{s_3}\right)^2 c_\delta^2 \simeq 60, \tag{4.14}$$

where we hawe used $c_i \simeq 1$ and $s_2 \simeq s_3$. It is important to note that the estimate (4.14) would remain true even if s_2 and s_3 were both very small. However, the opposite choice of δ, i. e. $\delta \simeq \pi$, would suppress both penguins and normal $b \to c$ transitions, leading to the dominance of $b \to u$ transitions. This seems to be unlikely, since current experiments favour $b \to c$ transitions. Therefore, by using (4.14) and (4.11), we make the estimate

$$\frac{\Gamma^{peng}(B \to K + nX)}{\Gamma^{(4)}(B \to (n+1)X)} \simeq 0.3\text{—}1.2, \tag{4.15}$$

i. e. it is of order 1 instead of being 5%. Furthermore, it is clear that the operators $O_5, \dots O_8$ should dominate the normal $\Gamma(B \to K + nX)$ rate. In fact, one obtains*

$$\frac{\Gamma^{peng}}{\Gamma^{(2)}} = \frac{1}{s_1^2}\left(\frac{c_1}{c_3}\right)^2\left(\frac{\xi_p}{\xi_4}\right)^2 \frac{8}{3}\frac{n_f\, c^2_{peng}}{2c_+^2 + c_-^2} \simeq 6\text{—}25. \tag{4.16}$$

*Our calculation presented in this paper was recently used by Eilam[27] to estimate the branching ratio $BR(B \to K + X)$ for very large m_t, $m_t \geq m_W$. For such values of m_t the contribution of the direct process $b \to s +$ gluon became sizeable and was added to the penguin contribution calculated in this paper. The $BR(B \to K + X)$ ratio was estimated to be 0.4 for $m_t = 20$ GeV, and up to 2% for $m_t = 2m_W$.

We have assumed in our calculations that weak b decays proceed via heavy--quark decays (spectator hypothesis). This rather simple picture has been questioned afer the discovery that D^+ and D^0 have different lifetimes. However, the mechanisms invoked in D decays[26], such as the interference of identical quarks[28] or soft-gluonic effects[29], are expected to be much less efficient in B decays[30]. The main reason is that all spectator effects scale as $|\varphi_M(O)|^2 m_Q^2$, whereas the quark decay rate scales as m_Q^5. The wave function at the origin, $\varphi_M(O)$, is related to the meson decay constant f_M via $|\varphi_M(O)|^2 \sim m_M f_M^2$. However, various estimates for f_D and f_B, such as those in the MIT bag model[31a], the potential model[31b], or from the QCD sum rules [31c], lead both f_D and f_B to be roughly of the order of f_π. Moreover, the experimental ratio for D lifetimes has drastically decreased, recent values being about 2—3[32].

We are aware of the difficulties in the experimental detection of these effects. Therefore, it was suggested in Ref. 9 to investigate the four-body channels $K3\pi$ and 4π. According to our analysis, both rates should be comparable. The background coming from sequential decays, such as $B \to D\pi \to K\pi\pi\pi$, was expected to be of the same order of magnitude. The main difficulty would be that the D contribution has to be separated out by invariant mass plots. Nevertheless, we find these measurements to be very important for the following reasons:

i) The effects described are pure QCD effects and the detection of penguin effects would provide a suitable test of QCD.

ii) The ratio estimated in (4.15) is at least order of magnitude larger than the ratio (4.13). Therefore, the effect is very large and clear.

iii) Penguin-like effects in $K \to 2\pi$ physics[11] are smaller (by a factor of 2.4) than needed to explain the $\Delta I = \frac{1}{2}$ rule. This should be ascribed to the failure of the use of perturbative QCD at low energy ($\mu \simeq m_\pi$) or some other mechanism might be responsible for the complete explanation[33]. Another penguin effect, $\varepsilon'/\varepsilon \sim 2 \times 10^{-3} - 2 \times 10^{-2}$ in K decays[13-15], is also difficult to measure. Current experiments[34] are expected to have a sensitivity of 2×10^{-3}. However, this sensitivity might not be sufficient since the presently favoured value $\Lambda \sim \sim 100$ MeV and smaler values for ϑ_2 would easily give ε'/ε to be $\sim 1 \times 10^{-3}$. (Cf. Table 3 in Ref. 14). Therefore, we believe that B decays should also be a suitable ground for testing penguin effects unambiguously.

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Appendix A

Here we present some formulae and expressions used in the text. More details may be found in Ref. 12.

In order to find the anomalous dimension γ_i of the operator O_i, one has to evaluate the renormalization matrix Z_i defined by (3.9). It is simply related to γ_i as

$$\gamma_i = \frac{\partial \ln Z_i}{\partial \ln \mu} = \frac{\alpha_s}{4\pi} \tilde{\gamma}_i. \tag{A. 1}$$

The operators $O(A, B)$ and $O(A\lambda, B\lambda)$ and the operator $O(C)$, defined in (3.2) and (3.12), respectively, should be inserted in the diagrams of Fig. 1. As a result of this procedure, one finds the following general expressions:

$$O(A, B) \to \frac{1}{2} \frac{\alpha_s}{4\pi} \ln \mu^2 \{-30(B\lambda, A\lambda) - \frac{1}{3} O(AB + BA)\},$$

$$O(A\lambda, B\lambda) \to \frac{1}{2} \frac{\alpha_s}{4\pi} \ln \mu^2 \left\{ -\frac{32}{3} O(B, A) - 5O(B\lambda, A\lambda) + 9O(A\lambda, B\lambda) + \right.$$

$$\left. + \frac{2}{9} O(AB + BA) - \frac{4}{3} \mathrm{Tr}\, B \cdot O(A) \right\}, \tag{A. 2}$$

$$O(C) \to \frac{1}{2} \frac{\alpha_s}{4\pi} \ln \mu^2 \left\{ \left[\frac{85}{9} - \frac{4}{3} n_f \right] O(C) - \frac{32}{2} O(1, C) - 5O(\lambda, C\lambda) \right\},$$

where n_f is the number of effective flavours.

If the operator O_i is not multiplicatively renormalizable, one has

$$O_i \to Z_{ij} O_j. \tag{A.3}$$

In this case, the matrix Z_{ij}, $\varphi Z \varphi^{-1} = \tilde{Z}$ should be diagonalized and the resulting Hamiltonian is given by (3.17) and (3.18).

As a result of the procedure discussed above, the matrix Z_{ij} obtained in the operator basis O_-, O_+, O_5, O_6, O_7 and O_8 is the following :

$$\ln Z = \frac{\alpha_s}{4\pi} \left(\frac{1}{2} \ln \mu^2 \right) \begin{bmatrix} 8 & 0 & -\frac{1}{3} & 0 & 0 & -\frac{1}{3} \\ 0 & -4 & -\frac{1}{3} & 0 & 0 & -\frac{1}{3} \\ 0 & 0 & \frac{34}{3} & \frac{32}{3} & 0 & -\frac{8}{3} \\ 0 & 0 & 3 & 0 & 0 & 0 \\ 0 & 0 & -\frac{2}{3} & 0 & 0 & -\frac{11}{3} \\ 0 & 0 & -\frac{20}{9} & 0 & -\frac{32}{3} & \frac{16}{9} \end{bmatrix} \tag{A. 4}$$

with the eigenvalues $\widetilde{Z}_i = \{8; -4; -5.782; \quad -2.195; 7.002; 14.086\}$, respectively. The exponents d'_i in (3.18) are given simply by

$$d'_i = \frac{\widetilde{\gamma}_i}{2b}, \tag{A. 5}$$

with $\widetilde{\gamma}_i$ defined by (A/1) and b given by

$$b = 11 - \frac{2}{3} n_f. \tag{A. 6}$$

Explicitly, $d_- = 0.57$, $d_+ = -0.29$, $d'_- = 0.48$, $d'_+ = -0.24$, $d'_5 = -0.35$, $d'_6 = -0.13$, $d'_7 = 0.42$, $d'_8 = 0.85$. Explicit values for the diagonalization matrices φ^{-1} and φ are the following:

$$\varphi^{-1} = \begin{bmatrix} 1. & 0 & 0.1290 & -0.0145 & -0.7852 & 0.2488 \\ 0 & 1. & 0.9979 & 0.0817 & 0.0712 & 0.0837 \\ 0 & 0 & 1. & -0.7317 & -0.5425 & -5.5486 \\ 0 & 0 & -0.5185 & 1. & -0.2325 & -1.1817 \\ 0 & 0 & 2.8683 & 0.2599 & 1. & 0.0021 \\ 0 & 0 & 4.3408 & 0.2887 & -1.8108 & 1. \end{bmatrix}, \tag{A. 7}$$

$$\varphi = \begin{bmatrix} 1. & 0 & 0 & 0 & 0.333 & -0.2497 \\ 0 & 1. & 0 & 0 & -0.2219 & -0.0832 \\ 0 & 0 & 0.0292 & -0.0538 & 0.1806 & 0.0979 \\ 0 & 0 & -0.1698 & 0.8251 & 0.1589 & 0.0327 \\ 0 & 0 & -0.0393 & -0.0598 & 0.4408 & -0.2894 \\ 0 & 0 & -0.1487 & -0.1127 & -0.0315 & 0.0416 \end{bmatrix}. \tag{A. 8}$$

Finally, we give explicit expressions for the coefficients c_5 and c_8. Thus, the reader could easily evaluate c_5 and c_8 for an arbitrary choice of $\varkappa_t$ and $\varkappa$:

$$\begin{aligned} c_5 = -\{&\varkappa_t^{0.57} [0.037\, \varkappa^{0.85} - 0.0308\, \varkappa^{0.42} - \\ &\qquad - 0.0024\, \varkappa^{-0.13} - 0.0038\, \varkappa^{-0.35}] + \\ &+ \varkappa_t^{-0.29} [0.0124\, \varkappa^{0.85} + 0.0028\, \varkappa^{0.42} \\ &\qquad + 0.014\, \varkappa^{-0.13} - 0.0292\, \varkappa^{-0.35}]\}, \\ c_8 = -\{&\varkappa^{0.57} [-0.01\, \varkappa^{0.85} - 0.227\, \varkappa^{0.42} - \\ &\qquad - 0.0125\, \varkappa^{-0.35} + 0.2497\, \varkappa^{0.48}] + \\ &+ \varkappa_t^{-0.29} [-0.0035\, \varkappa^{0.85} + 0.0207\, \varkappa^{0.42} - 0.0027\, \varkappa^{-0.13} - \\ &\qquad - 0.0979\, \varkappa^{-0.35} + 0.0834\, \varkappa^{-0.24}]\}. \end{aligned} \tag{A. 9}$$

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SLABI RASPADI B-KVARKOVA I KVANTNA KROMODINAMIKA

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Učinci kvantne kromodinamike u slabim raspadima *b*-kvarkova izučavaju se u Kobayashi-Maskawinom modelu uz uporabu tehnika renormalizacijske grupe da bi se zbrojili vodeći logaritmi. Nađeno je da su pingvinski grafovi stabilniji u odnosu na promjenu parametara nego u ocjeni za najniži red. Naša analiza potvrđuje da će se Cabibbova slika promijeniti. Ocijenjeno je da je za inkluzivne raspade bez šarmiranih kvarkova u konačnom stanju omjer $\Gamma(B \to K + nX) / \Gamma(B \to (n+1)X)$ približno jednak 0.3—1.2 umjesto 5%.